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A NEW APPROACH TO THE STUDY OF THE ECR HEATING
AND PARTICLE DYNAMICS IN THE PLASMA OF
ELECTRON CYCLOTRON RESONANCE ION SOURCES

PHD THESIS

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Introduction

Future particle accelerators require higher and higher currents of highly charged ions (HCI). In the case of the accelerators devoted to nuclear physics, the use of high currents and of highly charged ions extends significantly the number of low cross-section processes that can be studied with the available facilities and it may permit to fully exploit the potentiality of the forthcoming projects. High density plasmas and the ion beams that are generated can be also adapted to the use of industrial treatment of surfaces and of implantation devices, for which the stability, reproducibility, brightness and absence of beam halo are key issues. The ability to produce HCI is essential because the maximum energy of the accelerated particles depends on the charge over mass ratio (for circular machines it is proportional to the square of the charge state). The Electron Cyclotron Resonance Ion Sources (ECRIS) are nowadays the most effective devices that can feed the particle accelerators in a continuous and reliable way, providing high current beams of low and medium charge state ions and lower, but still important, beam current for highly charged ions.

In such equipments, a plasma is confined in a conductive cavity by strong magnetostatic fields and sustained by microwaves. Up to now a well established model has been used for ECRIS, following the experience made at INFN-LNS and other laboratories in Europe and US, but a remarkable enhancement is expected after the evidence that new coupling mechanisms and innovative techniques can be used. Some experimental results obtained at INFN-LNS shed some light over some unclear issues related to the modeling of the plasma used for these ion sources, as well as for ion sources based on non-resonant discharges (MDIS).

In ECRIS the ion beam is extracted from a plasma that can be created by microwaves with frequency between 2.45 GHz and 28 GHz with quite different magnetic field amplitude and structure, so that the plasma properties, as well as the ion beam properties, will be completely different.

Differences in ion beam parameters strongly depend on the plasma electron temperature T_e , on the ion lifetime τ_i , and on the electron density n_e ; the operation of ECRIS can be described in terms of such parameters, which determine the production of highly charged ion beams whose intensity may reach hundreds of $e\mu A$ of Ar^{12+} and up to some tens of $e\mu A$ of Xe^{30+} and higher charge states.

The creation of highly charged ions is a step by step process that requires long ion lifetimes; it can be obtained by means of an adequate magneto-static field (minimum-B structure) which traps both the ions and the electrons. The development of ECRIS during last twenty years has been driven by the improvement of the minimum-B magnetic trap, according to the High B mode concept, proposed in 1990 [1]. In particular, the need to achieve higher currents of HCI pushed the development of ECRIS towards the use of higher power and higher microwave frequency generators that require higher confining magnetic field, according to the Geller's scaling laws [2].

A classification in terms of ECRIS generation is commonly accepted according to the typical operational frequency. Second generation ECRIS are operated at $14 - 18\text{ GHz}$, whereas third generation ones work at 28 GHz or more. The use of higher frequency (28 to 37 GHz) makes the use of superconducting magnets for the confining trap mandatory. This trend is now deemed to saturate because of the technological limits of magnets and then we need to find an innovative approach to get additional and remarkable improvements, as requested by many users. Due to the technological problems, even the simple scaling of the microwave power is not anymore sustainable and for values above 10 kW the coupling efficiency is a key issue.

In order to overcome the limits imposed by the so called "Standard Model" [3], the thesis has been focused on the investigation of the ion and electron dynamics of ECR plasmas. Only a detailed investigation of the plasma physics regarding the processes involved with the plasma dynamics, and in particular with the wave-particle interaction, permits to further develop the ECRIS without any strong and expensive improvement of the technology. The activities here reported have been carried out in the framework of the INES experiment (INovative Ecr ion Sources), funded by the 5th National Committee of the INFN, and also in the framework of the international collaboration named ISIBHI (Ion Sources for Ion Beams of High Intensities), i.e. that project which aims to the production of highly charged ions with high currents. Such a collaboration is a part of the European initiative for the Nuclear Physics research, named EURONS (EUROpean Nuclear Science), funded by the 6th Framework Programme.

0. INTRODUCTION

The models and the clues presented hereinafter permit to explain many of the experimental results obtained in the last ten years, which could not find an adequate explanation up to now. Some of these explanations have been already published [4], [5], [6], others new and still under refereeing. Some other simulations permitted to formulate new ideas to design future experiments, that will be carried out in the frame of the HELIOS experiment (Hot Electron Layer in ecr Ion Sources) during the next three-four years, and to outline different methods of plasma ignition in ion sources. On this purpose, many international collaborations have been activated: in particular with the GSI, Darmstadt, the Kurchatov institute of Moscow, Russia, and with the Max Planck Institute for Plasma Physics of Greifswald, Germany. The theoretical efforts and the original experimental results obtained with the new 3rd generation ECRIS named VENUS [7], operating at Lawrence Berkeley Laboratory, as well as with the WEGA stellarator operating at the Max Planck Institute, have contributed to the comprehension of phenomena which may drive the ECRIS development towards a different microwave-plasma coupling. In particular, the Bernstein waves heating has been investigate at Max Planck Institute: sometime it has been observed that this mechanism may generate extremely hot electrons ($\sim MeV$), and the test on WEGA was important to understand when the formation of these particles can be avoided.

More in detail, the thesis reports in the first chapter the general features of plasmas, in particular in case of microwave heating. A special attention is devoted to the electromagnetic field propagation and to the stochastic heating in collisionless plasmas. The original contribution is mainly reported after the second chapter, where an overview of the different ion sources is given, together with the main peculiarities of the microwave discharge and ECR ion sources.

The third chapter concerns the experiments showing some new "tricks"¹ used to improve the sources performances, like the frequency tuning effect and the two frequency heating. The data about the VENUS experiment on the hot electrons production are here reported, with new results about the violation of the scaling laws for the magnetic field in case of very high magnetic field gradients. The chapter also reports the numerical simulations which describe many of the experimental results and propose alternative schemes of plasma heating. At first the simulations have been based on a single particle approach, then the same problem has been analyzed by

¹Historically, new techniques for the ECRIS plasma parameters enhancement have been called "tricks" in different workshop proceedings

means of a PIC code, able to self-consistently solve the Maxwell equations in a dense plasma. These simulations have been performed in collaboration with Dr. S. Chikin of the Kurtchatov Institute of Moscow. The single particle and PIC simulations are compared at the end of the chapter, and then both are compared with the experimental results, showing how the effective synergy between the experimental investigation and the theoretical study has been the "leitmotiv" of the entire thesis. The fifth chapter reports a series of experimental activities and of theoretical investigations regarding the possibility of plasma ignition by means of the Bernstein waves. The model presented in the chapter, which allows to design a new experiment on ECRIS-like systems, moves from the experiments carried out on the WEGA stellarator of Greifswald. In collaboration with the Dr. H. Laqua, a conceptual design of a new microwave injection scheme in ECRIS has been modeled. Finally, in the sixth chapter the conclusions are reported, and the relationship between the several approaches used in the plasma dynamics investigation is summarized, pointing out that the theoretical approach is consistent with the experimental results.

Chapter 1

Fundamentals of ECRIS plasma physics

*We want to show that without a minimum of
plasma science no progress is possible in ECRIS
and probably also in other ion source development.*

R. Geller

The plasma physics is a fundamental branch of modern science, as it finds a great number of applications in several fields of nature, from stars and universe to the issues linked to the production of energy by means of the nuclear fusion. The 98% of the matter of the entire universe is plasma, and only on earth this state of matter is so rare. In the next pages only some aspects directly linked with the activities carried out on ECR ion sources during the PhD course will be treated with some details. For more details see the textbooks [8] and [9].

1.1 Basic plasma parameters

Three fundamental parameters characterize the plasma generated in an Electron Cyclotron Resonance Ion Source:

1. the particle density n (measured in particles per cubic meter or cubic centimeter),
2. the temperature T of each species (usually measured in eV , where $1\text{ eV} = 11605\text{ }^\circ K$),

3. the steady state magnetic field B (measured in Tesla).

A host of subsidiary parameters (e.g., Debye length, Larmor radius, plasma frequency, cyclotron frequency, thermal velocity) can be derived from these three fundamental parameters. For partially-ionized plasmas, the fractional ionization and cross-sections of neutrals are also important.

1.2 Collisions in Plasmas

The collisions are of primary importance to determine the ionization rates in ion sources. They also regulate the electrons and ions lifetime inside magnetically confined plasmas. The number and the nature of the collisions is mainly determined by the electron temperature, the electron density and the background pressure of neutrals. Hence, each ion source, according to its own characteristics, will provide particular performances determined by the number of collisions per time unit.

In order to discuss about the main collisional phenomena occurring in ECR plasmas, it is convenient to introduce some physical quantities: the mean free path λ_{mfp} , the mean time between a collision and the following one τ , and finally the collision frequency ν ¹.

All these parameters can be calculated once known the cross section σ .

Generally in plasmas all the particles (electrons, ions and neutrals) follow some energy distribution function², then to determine ν we have to average over all the possible values of σv ³ [10]:

$$\nu = n\overline{\sigma v} \quad (1.2.1)$$

Figure 1.2.1 features a schematic overview of the main plasma collision processes; a classification has been done according to the number of colliding particles: **multiple collisions** are characteristics of plasmas, because they occur thanks to the *long range Coulomb interaction*. The **binary collisions** are similar to the collisions in gaseous systems; however in this case non-elastic and ionizing collisions occur because of the high energy content of the plasma particles, especially of electrons.

¹The analysis of the various collisions in plasmas will be carried out principally in terms of the collision frequency.

²In many cases it is convenient to use a Maxwellian-like distribution, for which a temperature can be determined. Anyway we will see that for low pressure and high temperature plasmas the energy distribution function does not follow a pure Maxwellian shape, but it can be viewed as a superposition of several plasma populations with different temperatures.

³We must average not only on v , but also on σ because in general $\sigma \equiv \sigma(v)$.

1. FUNDAMENTALS OF ECRIS PLASMA PHYSICS

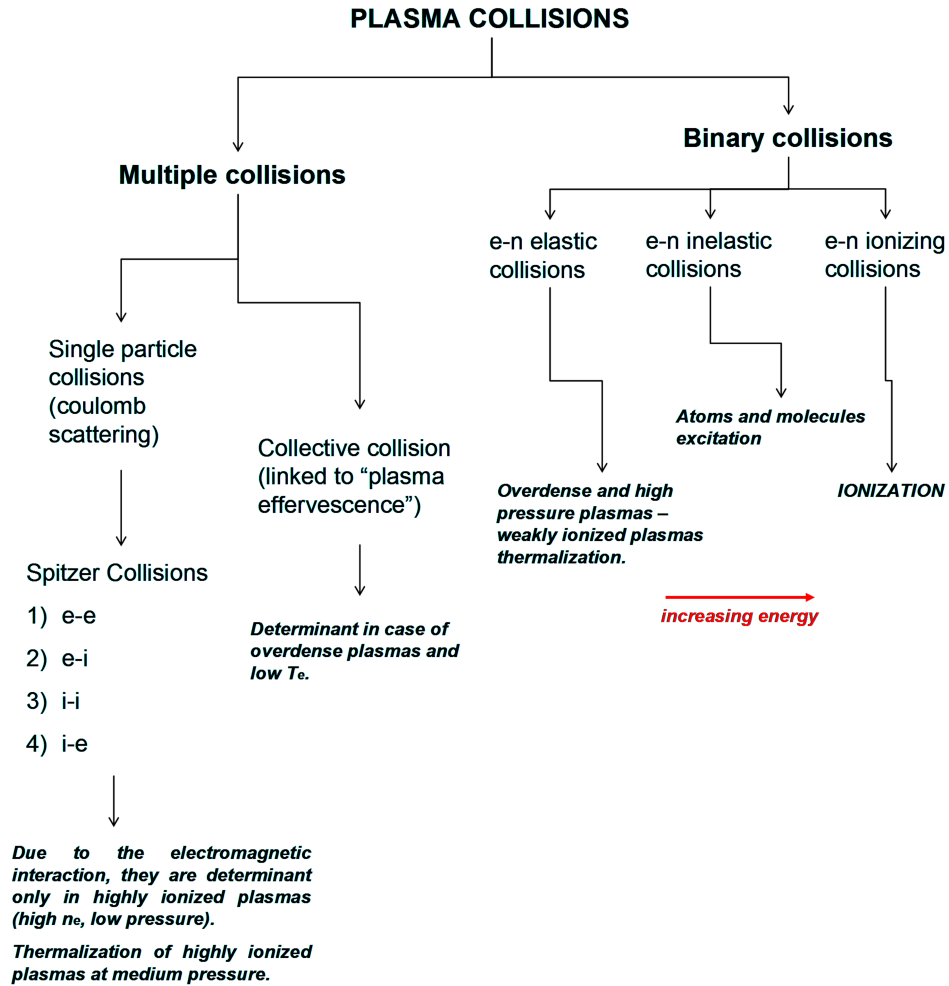


Figure 1.2.1: A schematic overview of the principal collision occurring in plasmas. the abbreviations *e*, *n*, *i* mean electrons, neutrals and ions respectively.

1.2.1 Multiple scattering: the Spitzer collisions

Note that in plasmas any electric field can be felt by charged particles only within a distance that is named **Debye Length** - Λ_D [8]:

$$\Lambda_D = \left(\frac{\epsilon_0 k T_e}{n e^2} \right)^{\frac{1}{2}} \quad (1.2.2)$$

which means that a proof particle feels the electromagnetic interaction due only to those plasma particles located inside the Debye sphere ⁴. A proof particle (electron or ion) moving in a highly ionized plasma experiences a lot of scattering collisions due to the electric field generated by the other particles located inside the Debye sphere.

We can define a cross section, σ_{90° , giving the probability of a multiple scattering with a final angle of 90° ⁵:

$$\sigma_{90^\circ} = 8\pi \left(\frac{z_1 z_2 e^2}{M v^2} \right)^2 \ln \left(\frac{\Lambda_D}{b_{min}} \right) \quad (1.2.3)$$

where z_1 and z_2 are the charges of the two particles, e is the electron charge, M is the mass of the colliding particle (if one considers the other particles much more massive) and the term $\frac{\Lambda_D}{b_{min}}$ is the so called Coulomb logarithm. This equation, except the logarithm, is practically the Rutherford formula for charged particles scattering (binary collision). As the Rutherford cross section is almost two order of magnitude lower than the (1.2.3), the 90° scattering due to a single collisions is much less probable than multiple deflections⁶.

A more precise representation of the multiple collisions can be given in the center of mass system, and the collision frequencies can be calculated by taking into account collisions among the two different plasma species: electrons and ions. Omitting, for sake of brevity, the mathematical description, we can use the final formulas to understand what is the role that each type

⁴The Debye sphere is the sphere with radius Λ_D .

⁵The quantities in the formula are expressed in the c.g.s. system

⁶This result was formerly guessed by Chandrasekhar.

1. FUNDAMENTALS OF ECRIS PLASMA PHYSICS

of collision plays in the plasma:

$$\nu_{90^\circ}^{ee} = \frac{1}{\tau_{sp}} = 5 \cdot 10^{-6} n \frac{\ln\left(\frac{\Lambda_D}{b}\right)}{T_e^{\frac{3}{2}}} \quad (1.2.4)$$

$$\nu_{90^\circ}^{ei} = \frac{1}{\tau_{sp}} \sim 2 \cdot 10^{-6} z n \frac{\ln\left(\frac{\Lambda_D}{b}\right)}{T_e^{\frac{3}{2}}} \quad (1.2.5)$$

$$\nu_{90^\circ}^{ii} = \frac{1}{\tau_{sp}} \sim z^4 \left(\frac{m_e}{m_i}\right)^{\frac{1}{2}} \left(\frac{T_e}{T_i}\right)^{\frac{3}{2}} \nu_{90^\circ}^{ee} \quad (1.2.6)$$

n_e is in cm^{-3} and T_e in eV .

Each frequency is usually called **characteristic Spitzer collision frequency**. Once given the collision frequency⁷ we can easily determine the time between two successive events: $\nu_{90^\circ}^{ee} \simeq \nu_{90^\circ}^{ei}$, and $\tau_{90^\circ}^{ee} \simeq \tau_{90^\circ}^{ei}$; on the contrary, if we consider ions colliding against electrons (i-e collisions), we have:

$$\tau_{90^\circ}^{ie} \simeq \left(\frac{m_i}{m_e}\right) \tau_{90^\circ}^{ei} \quad (1.2.7)$$

To investigate the thermalization of the plasma via collisions the energy transfer frequency must be analyzed. The amount of the energy transferred for particles of the same species (for example e-e or i-i) is about the half of the initial one, whereas it is equal to the $\frac{m_e}{m_i}$ fraction in case of e-i or i-e collisions. Then we can find the following relationship:

$$\tau_m^{ee} \sim \tau_{90^\circ}^{ee} \sim \tau_{90^\circ}^{ei} \quad (1.2.8)$$

$$\tau_m^{ii} \sim \tau_{90^\circ}^{ii} \sim \left(\frac{m_i}{m_e}\right)^{\frac{1}{2}} \tau_{90^\circ}^{ei} \quad (1.2.9)$$

$$\tau_m^{ei} \sim \tau_m^{ie} \sim \frac{m_i}{m_e} \tau_{90^\circ}^{ei} \quad (1.2.10)$$

According to the scheme 1.2.1, the Spitzer collisions play a role only in highly ionized⁸ and low pressure plasmas, where the internal electromagnetic interactions are predominant. Plasmas produced in ECR ion sources fulfill this condition and the Spitzer collisions play the main role to determine the

⁷We talk about the collision frequency for sake of simplicity, but strictly speaking the term 'collision' is uncorrect as we should speak about 90° scattering.

⁸Usually the expression "highly ionized plasmas" identifies that ones with an ionization fraction higher than 90%.

energy exchange processes. It can be noticed that $\tau_m^{ee} \sim \tau_{90^\circ}^{ee}$ and $\tau_m^{ei} \gg \tau_{90^\circ}^{ei}$; then the thermalization via collisions is much easier among the electrons than among electrons and ions. In addition, as the ion lifetime in ECRIS is shorter than the time required for the ion heating through the collisions, the ions remain cold, whereas the electrons increase their energy because of the external electromagnetic field. This result is of primary importance for ECRIS, as one of the most important quality parameters, the emittance, increases with the ion temperature.

Some numerical example help to better understand the above results. In ECR plasmas, for energies ranging from few eV ⁹ to some keV ¹⁰ the ν_{ee} varies from $10^5 \div 10^6 \text{ sec.}^{-1}$ ($T_e = 10 \text{ eV}$) to $10^2 \div 10^3 \text{ sec.}^{-1}$ ($T_e = 1 \text{ KeV}$); this means that the electron thermalization requires times on the order of few μs , in the first case, on the order of ms in the second case. Considering that $\frac{m_i}{m_e}$ is about $10^3 \div 10^4$, then the electron-ion thermalization time is much longer than the plasma lifetime, which is estimated to be on the order of some ms at maximum. This explains the large difference between the electron and the ion temperature. It is known that ECR plasmas are generated by means of electromagnetic waves resonating with electrons moving in a magnetostatic field. The electromagnetic wave frequency ($\omega_{RF}/2\pi$) is comparable with the Larmor frequency; for the magnetic field of few Tesla, $\omega_{RF}/2\pi$ must be on the order of many GHz . Therefore:

$$\omega_{RF} \gg \nu_{ee} \quad \text{and:} \quad \omega_{RF} \gg \nu_{ei} \quad (1.2.11)$$

i.e. electrons gyrating around the magnetic lines and subjected to the wave field are practically collisionless by the point of view of the Spitzer collisions. Here a delicate aspect of the ECR plasma dynamics arises: *how does the electron heating occur without collisions?* Only on timescales of several microseconds the electrons begin to collide each other, but simulations (see section 4.6 and 4.13) and experimental results demonstrate that an effective heating occurs even for shorter times (on the order of few hundreds of ns). In order to explain the electron heating without collisions we have to take into account some stochastic processes occurring during the electron acceleration due to the interaction with the electromagnetic wave.

⁹This temperature is typical for the ion sources used to produce lowly charged ion beams at high currents (up to tens of mA): these sources are often defined Microwave Discharge Ion Sources (described in section 2.2).

¹⁰For plasmas of ECRIS for highly charged ions.

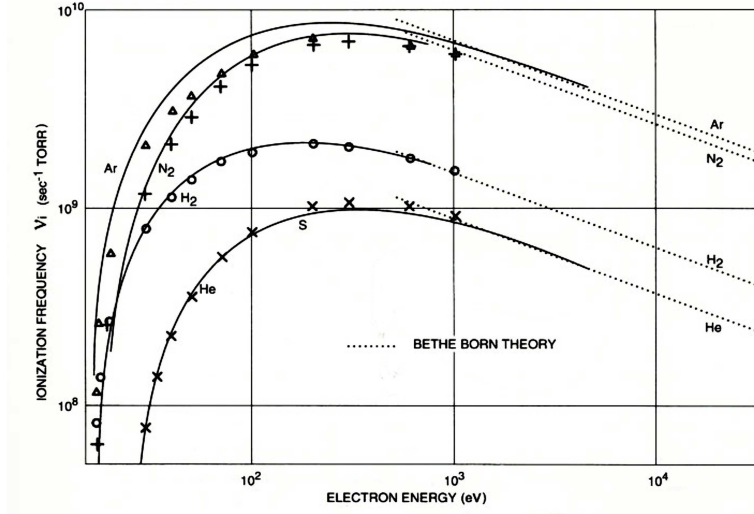


Figure 1.2.2: Ionization frequencies versus the electron energy. For high electron energies the experimental data have been fitted with the Bethe-Born theoretical formula. [10].

1.2.2 Binary collisions

The binary collisions in plasmas can be divided, according to figure 1.2.1, into two groups: the **electron-neutral collisions** (*elastic and inelastic*), and the **ionizing collisions**. The former play an important role especially in case of high pressure plasmas with a low degree of ionization and a low electron temperature¹¹. As featured in figure 1.2.1, these collisions provide the thermalization playing the same role of the Spitzer collisions in case of highly ionized and low pressure plasmas. As for the creation of multi-charged ions in ECRIS these collisions are not so important [10], then we focus our attention on the ionizing collisions.

The figure 1.2.2 shows the trend of the ionization frequency versus the electron energy for different ion species. It can be noticed that ν_i is parameterized by the background pressure; hence, to determine the real ionization rate, we have to multiply the frequency shown in figure by the operating pressure of our ion source. ν_i refers to all the possible ionizing events, i.e.

¹¹For high pressure we mean, in this case, $10^{-3} \div 10^{-4}$ mbar, i.e. the operating pressures of microwave discharge plasmas (MDP) (see section 2.2). Low temperature here means $T_e \leq 15$ eV, that also in this case is typical of MDP.

the production rates for all the charge states that can be produced by a single collision. It is clear from figure 1.2.2 that the ν_i decreases slowly at higher energy. If we take into account the usual maxwellian distribution function of the electrons, then we have to replace the electron energy in the x axis of figure 1.2.2 with the electron temperature. This means that a not negligible amount of ionizing collisions occur also below the energy threshold (i.e. the ionization potential) due to the electrons of the distribution's tail. For example, in microwave discharge plasmas the electron temperature is usually ~ 10 eV, but nevertheless the atoms can be easily ionized because of the tail electron component, whose energy is high enough to provide effective ionizing collisions.

In case of ECRIS, where the production of highly charged ions is of primary importance, a particular relationship among electron density, electron temperature, and ion confinement time has to be achieved. This discussion leads to the so called **multiply-charged ion (MI) production criterion**. The ionization down to inner shells of atoms required a **Step by Step** process, because the cross section of single ionization per collision are much higher than the multi-ionization.

The cross section for single ionization is on the order of 10^{-16} cm², but for multi-ionizations it rapidly decreases. To expel two electrons is one order of magnitude lower, to expel three electrons it is two order of magnitude lower, and so on. The time needed for the transition from charge state z_1 to charge state z_2 takes a time, on average [10]:

$$\tau_{z_1 z_2} = \frac{1}{[n_e \sigma_{z_1 z_2}(v_e) v_e]} \quad (1.2.12)$$

and averaging on a maxwellian distribution¹² we have:

$$\tau_z = \frac{1}{[n_e S_z(T_e)]} \quad (1.2.13)$$

$$S_z(T_e) = \langle \sigma_z(v_e) v_e \rangle(T_e) \quad (1.2.14)$$

only if the ion confinement time is longer than the time required for the given ionization (i.e. the time (1.2.13)) the transition $z \rightarrow z + 1$ takes place.

¹²We usually average the collision parameters on a maxwellian distribution. As said in the previous pages, the energy distributions of ECRIS plasmas differ from maxwellian trends, however they can be considered as superpositions of three different distributions, each one having its own temperature. Then, for sake of simplicity, we assume that the results obtained for a maxwellian distribution are valid for each plasma population, with the further consideration that the warm component is the most effective to produce highly charged ions. Since now we will use the parameters of the warm population to determine the ionizing frequencies of ECRIS plasmas.

1. FUNDAMENTALS OF ECRIS PLASMA PHYSICS

Then the condition to produce highly charged ions in plasmas can be written as:

$$n_e \tau_i \geq \frac{1}{[S_z(T_e)]} \quad (1.2.15)$$

In plasmas with long confinement times the step by step ionization provides highly charged ions. The equation (1.2.15) can be rewritten by substituting the $S_z(T_e)$ parameter, obtaining [10]:

$$\xi n_e \tau_i \geq 5 \cdot 10^4 (T_e^{opt})^{\frac{3}{2}} \quad (1.2.16)$$

where $\xi \equiv \sum_{j=1}^N q_j$ is the number of subshells in the atom outer shells.

The criterion which comes from (1.2.16) is of paramount importance in ECRIS, as it fixes the plasma operational parameters needed for good performance sources. Some numerical examples can be given about the required densities, confinement times and electron temperatures. In order to obtain completely stripped light ions, we need $T_e \simeq 5 \text{ eV}$, with $n_e \tau_i \geq 10^{10} \text{ cm}^{-3}$. The required electron temperature is given by the energy threshold of the given charge state that one would like to obtain.

$$T_e^{opt} \simeq 5W_{thr} \quad (1.2.17)$$

A more general picture of the criteria linked to the plasma parameters is given by the figure 1.2.3, which features what combinations of temperatures, densities and confinement times are able to produce different ion species with different charge states.

1.2.3 Three electron populations in ECR plasmas: the importance of the warm component

In ECRIS plasmas, the energy of the electrons ranges from a few eV to hundreds of keV or even 1 MeV [10]. A reasonable description is often given in terms of three different electron populations: cold, warm, and hot; the creation of the three populations can be explained by tacking into account the quasi-linear theory of the wave-particle interaction and the stochastic heating theory (see section 1.11). However, the existence of these populations for times longer than collisions is a consequence of the variation of the Spitzer collision frequency with the electron temperature (see equation (1.2.4)). A complete thermalization of the electrons is impossible, and three different electron populations are present at any time into the plasma. In

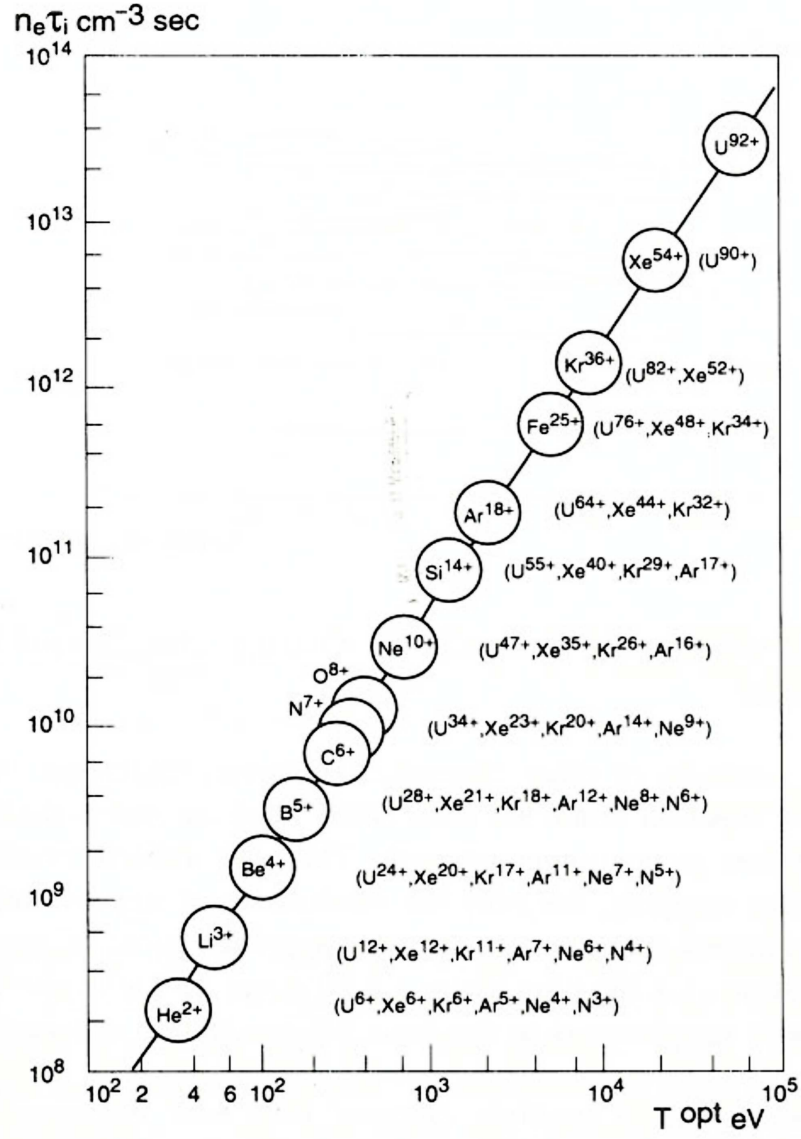


Figure 1.2.3: Golovanivsky's diagram of the $(n_e \tau_i)$ criteria for the production of highly charged ions. The ions enclosed by circles are completely stripped; some combinations of electron temperature, electron density and ion confinement time allow to produce completely stripped ions; inside brackets uncompletely stripped ions are shown: they can be produced with the corresponding plasma parameter of ions enclosed in circles. [10].

1. FUNDAMENTALS OF ECRIS PLASMA PHYSICS

particular, the cold population has temperatures in the range of $1 \div 100 \text{ eV}$, the warm in the range of $100 \text{ eV} \div 100 \text{ keV}$, and the hot above 100 keV [10]. Being the ionization cross section inversely proportional to the electron temperature, only the warm component of the electron energy distribution function has the right energy to ionize the atoms up to high charge states and, at the same time, a sufficient collision frequency for intense production of HCIs. The cold electrons cannot induce the ionization from the inner shells of atoms because of their insufficient energy, even if their collision frequency is the highest. It has been shown that the maximum ionization cross section is obtained at a given optimal temperature, as shown in figure 1.2.3. Therefore, the hot electrons have a too low cross section for the ionization process, although they have high energies, and in principle, they are able to create HCIs.

The ability to heat the electrons of the plasma in order to maximize the warm component and minimize the "hot" electrons is requested for high performance ECRIS. The advantage of the hot electrons is that they are almost collisionless and then well confined, stabilizing the plasma, but an excessive number is detrimental, as it increases the aging of insulators; in addition, for third generation ECRIS (that use superconducting magnets) the hot electrons increase the heat load of the cryostat, making difficult the operation of the sources at high power rate. These phenomena have been already observed in 2000 with the SERSE source [11] at INFN-LNS, and in 2005-2006 with the VENUS source at LBNL [7].

Then the importance of a sort of electron energy distribution function "tailoring" is evident, and the definition of the methods to change the EEDF-Electron Energy Distribution Function has been one of the main goals of the PhD activity.

1.2.4 Recombination effects

The charge state of highly charged ions may decrease because of some electron-ion recombination process. The recombination mechanisms are of three different types:

- *radiative recombination*: free electrons are captured by ions with the following emission of a photon;
- *dielectronic recombination*: it is a three steps process, where an electron firstly excite an ion, then it is captured, and finally a photon is emitted;

- *charge exchange recombination*: highly charged ions, colliding with lowly charged species and/or with neutrals, capture some electrons from the electronic cloud of the other particle, thus decreasing their own charge state;

Only the charge exchange recombination play a remarkable role in ECRIS [10]. Because of its own nature, the recombination rates must depend on the background pressure: the higher is the number of the neutral or lowly charged ions, the higher is the probability of charge exchange collisions. For high performance ion sources the charge exchange collision time must be longer than the ionization time for a given charge state, i.e. $\tau_{exch} > \tau_i \geq \tau_{z \rightarrow +1}$. This implies that:

$$\frac{n_0}{n_e} \leq 7 \cdot 10^3 \xi \frac{\sqrt{A}}{z \left[T_e^{opt} \right]^{\frac{3}{2}}} \quad (1.2.18)$$

where n_0 is the density of neutrals, and A is the atomic mass number.

For example, in order to obtain Argon fully stripped ions (Ar^{18+}) at $n_e \simeq 10^{12} \text{ cm}^{-3}$, we need an operative pressure of $\sim 10^{-7} \text{ mbar}$, corresponding to a number of neutrals per cubic centimeter $\sim 10^9 \text{ cm}^{-3}$. These numbers are typical of many modern ECRIS like SERSE, where the minimum operative pressure is 10^{-8} mbar , but it usually works at $1 \div 4 \cdot 10^{-7} \text{ mbar}$.

1.2.5 Summary and additional remarks

A more general description of the collision dynamics in plasmas can be given by looking to the figure 1.2.4.

It features that the momentum transfer processes in plasmas changes as the electron energy increases: they are mainly governed by the electron-neutral collisions at low temperatures (some eV , or few tens of eV), then by long-range Coulomb collisions (or Spitzer collisions). At very high energies (higher than 1 keV) the ionization collisions prevail. For these ones, the high energy electrons lose a small fraction of their energy (used to remove electrons from the ion's shells). As the energy increases, the scattering angle diminishes and also the transferred momentum decreases, thus leading to a non randomization of electron velocities. This means that the higher is the electron energy, the lower is the probability to reduce their energy by means of collisions. This leads to the formation of the so-called *electron runaway*; without any friction due to collisions of various type, a group of electrons may become decoupled for the Maxwellian bulk plasma (colder) and then

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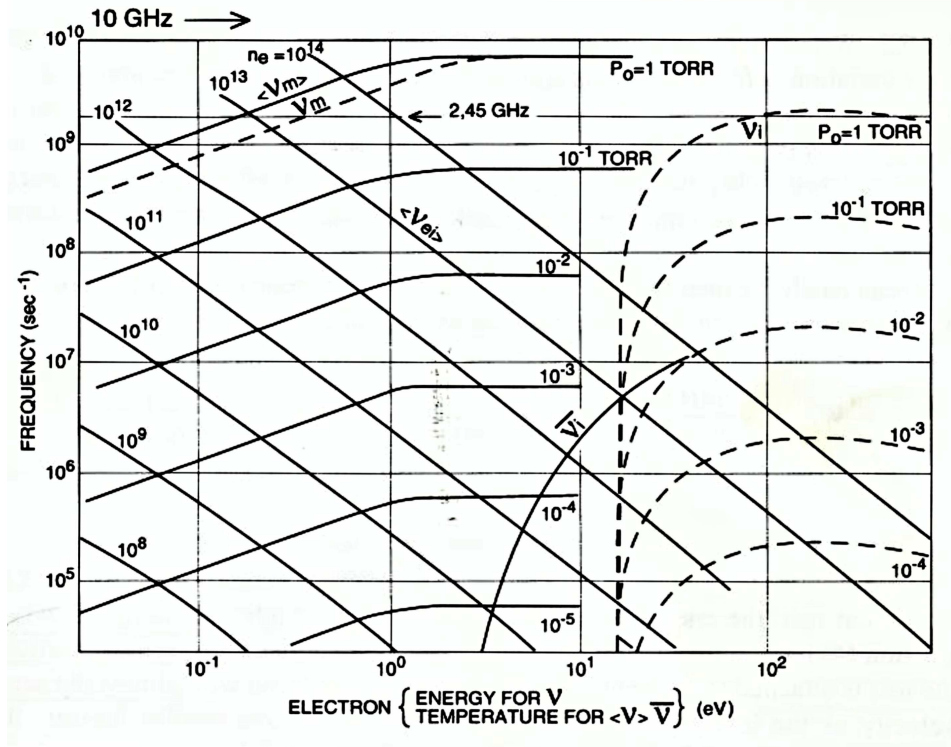


Figure 1.2.4: Electron-ion momentum transfer frequency ν_{ei} with respect to the electron energy or temperature for singly charged ions and for different electron densities as a parameter. The electron-neutral momentum transfer, and the ionization frequency ν_{ion} are shown for comparison (gas H_2) [10].

some instabilities may arise (beam plasma instability). It is not clear at this moment whether such electrons could be thermalized, in such a way to constitute the hot tail of maxwellian-like distribution. Several experimental results suggested that no separated groups of electrons are present in very hot plasmas [7]; nevertheless, the number of electrons extremely hot (with energies above 500 keV) is incredibly high [12].

1.3 The single particle, kinetic and fluidodynamics approaches

The most simple approach to the study of the plasmas is based on the single particle dynamics. Many properties of plasmas, from the confinement to the energy absorption during the wave particle interaction, and also many aspects of the electromagnetic field propagation can be described in terms of the single particle approach.

The magnetic confinement in Simple Mirrors, minimum-B and toroidal configurations can be studied, at first approximation, by means of a single particle approach. In section 1.6 some highlights on particle trapping are given, principally based on a single particle description (**"orbits theory"**).

The simplest equation describing the electromagnetic field propagation is the following one [13]:

$$m_e \frac{d\vec{v}}{dt} = q_e \vec{E} - m_e \vec{v} \omega_{eff} + q_e \vec{v} \times \vec{B}_0 \quad (1.3.1)$$

this equation takes into account the effects on the single particle motion, in a non relativistic approximation, of the wave electric field¹³, of the magnetostatic field and of the collisions.

As in case of plasmas produced in ECR ion sources, the electrons easily reach energies of the order of some hundreds of *keV*, or also *MeV*, the relativistic effects must be taken into account, and the complete single particle equation of motion becomes [14]:

$$\frac{d\vec{v}}{dt} = \frac{q}{m_0 \gamma} \left[\vec{E} + \vec{v} \times \vec{B} - \frac{\vec{v} \cdot \vec{E}}{c^2} \vec{v} \right] \quad (1.3.2)$$

where $\vec{E}$ is the electric field of the wave, and $\vec{B}$ is the sum of the wave magnetic field and of the magnetostatic field used for the plasma confinement. In section 4.7 we solve numerically this equation for several configurations of electromagnetic field filling the plasma cavity, and also for different sets of waves with different frequencies.

1.3.1 The kinetic approach

The most detailed approach to the plasma dynamics is the kinetic theory, which is based on non-equilibrium thermodynamics. Actually, an exhaustive

¹³In a non relativistic framework the force due to the magnetic field is much lower than the electric force.

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approach requires to know all the particle positions and velocities, the fields in which these particles are located and also the reciprocal interaction, or the correlations between pair or groups of particles.

From considerations based on statistical mechanics it comes out that a given system of N particles can be described by the function [15]:

$$\rho(p, q, t) p = p_1, \dots, p_{3N}; \quad q = q_1, \dots, q_{3N} \quad (1.3.3)$$

where p and q are the momenta and the positions of all the N particles¹⁴. In presence of conservative forces the **Liouville theorem** is valid, i.e. **the volume occupied by the system in the phase space is constant for different times**. This means that, in the phase space, the system evolves like an incompressible fluid. Starting from the Liouville theorem, the time evolution of the system can be determined, i.e. by solving the equation $d\rho/dt = 0$. This implies that we should solve equations with $6N$ variables, with N comparable with the Avogadro's number. In order to reduce the complexity of the problem, the so called *reduced densities* ρ_S can be introduced, each one depending on $1 \leq 6S < 6N$ variables. Each ρ_S gives the probability to find S particles in a given volume element of the phase space. Their temporal evolution can be determined, but unfortunately each equation depends on the density of the next order: i.e. the equation for the ρ_S is coupled to the ρ_{S+1} . The most simple reduced density is ρ_1 , which describe the so called single particle distribution. The complete set of equations for the reduced densities is named *BBGKY hierarchy*¹⁵ [15] and it can be truncated to the second equation, or even to the first one, if some approximations are done. In particular we can imagine that the particles are uncorrelated, and that they move in presence of a mean field generated by all the other particles; for this reason it takes the name of *mean field approximation*. Particles are uncorrelated if the system can be considered *collisionless*. If we suppose that it is formed by charged particles (ions and electrons), i.e. if we have a plasma, and if the mean field is the electromagnetic field, then the equation for ρ_1 (or f_1)¹⁶ will be [9]:

$$\frac{\partial f_\alpha}{\partial t} + \vec{v} \cdot \vec{\nabla} f_\alpha + \frac{q_\alpha}{m_\alpha} \left(\vec{E} + \frac{\vec{v} \times \vec{B}}{c} \right) \cdot \vec{\nabla}_{\vec{v}} f_\alpha = 0 \quad (1.3.4)$$

¹⁴The function ρ is the density in the phase space: the expression $\rho d^{3N}p d^{3N}q$ gives the number of particles in the hypervolume $d^{3N}p d^{3N}q$.

¹⁵This formal theoretical result is named after Bogoliubov, Born, Green, Kirkwood, and Yvon.

¹⁶Since now, instead of ρ we will use the distribution function f . Multiplying f by the average density of the system, $\bar{n}$, we obtain the number of particles in a given small volume element of the phase space.

f_α is the distribution function for the single particle, and the $\vec{E}$ and the $\vec{B}$ fields can be determined by:

$$\vec{\nabla} \cdot \vec{E} = 4\pi \sum_{\alpha} \bar{n}_{\alpha} q_{\alpha} \int f_{\alpha} d\vec{v} + 4\pi \rho_{ext} \quad (1.3.5)$$

$$\vec{\nabla} \times \vec{B} = \frac{1}{c} \frac{\partial \vec{E}}{\partial t} + \frac{4\pi}{c} \sum_{\alpha} \bar{n}_{\alpha} q_{\alpha} \int \vec{v} f_{\alpha} d\vec{v} + \frac{4\pi}{c} \vec{J}_{ext} \quad (1.3.6)$$

$$\vec{\nabla} \times \vec{E} = -\frac{1}{c} \frac{\partial \vec{B}}{\partial t} \quad (1.3.7)$$

the subscript α indicates both ions and electrons. The equations (1.3.4) and (1.3.5) describe the stationary plasma states, the plasma waves, the instabilities and other short time¹⁷ effects of plasma dynamics in terms of the distribution function $f(\vec{r}, \vec{v}, t)$.

The equation (1.3.4) is often named *collisionless Boltzmann equation or Vlasov equation* [9]. By making the integrals $\int f d\vec{v}$, $\int \vec{v} f d\vec{v}$, $\int v^2 f d\vec{v}$, and so on, we can determine the equation of continuity, the equation for the momentum and the equation for the energy of a fluid respectively. Therefore *the plasma fluid equations can be derived from the Vlasov equation*.

Although the plasma looks like a gas, it behaves more like a fluid, and the fluid mechanics is usually a powerful method to describe it. By means of the fluidodynamics we can describe the plasma dynamics without a precise knowledge of all the particle positions and velocities, but in terms of macroscopic parameters like the temperature, the density, the pressure, etc. In this sense the fluid approach is simpler than the Vlasov and the collisions can be included by considering the momentum exchange between electrons and ions. However all the microstructures due to particle effects cannot be taken into account by means of this approach. Collisionless absorption of electromagnetic waves, and other aspects concerning instabilities due to velocity-distribution inhomogeneities, remain out from the fluidodynamics possibilities.

¹⁷Here short times mean times shorter than the ones required for collisions $\tau \ll \tau_{coll}$. Such an approximation will be used many times in the most important simulations reported in this thesis.

1.3.2 Plasmas as fluids

The general equation which describe the momentum transport in a plasma subjected to electric and magnetic fields is [8]:

$$mn \left[\frac{\partial \vec{u}}{\partial t} + (\vec{u} \cdot \vec{\nabla}) \vec{u} \right] = qn \left(\vec{E} + \vec{u} \times \vec{B} \right) - \vec{\nabla} p \quad (1.3.8)$$

in the "one fluid approximation", i.e. when the fluid equations are written for a mixture of ions and electrons and the average mass density and velocity are calculated. $\vec{u}$ is the velocity of a fluid element, and p is the kinetic pressure. We can do some approximations which simplify considerably the model by supposing that the particles move with velocities lower than c (as in case of the single particle approximation). If we take into account the plasma quasi-neutrality (i.e. $n_e \simeq n_i$), and we neglect the ratio m_e/m_i , we can consider only large scale and low-frequency phenomena. We finally consider that the system remain isotropic at all times. Hence the set of the so-called **magneto-hydro-dynamics** equations becomes [9]:

$$\frac{\partial \rho_m}{\partial t} + \vec{\nabla} \cdot \rho_m \vec{V} = 0 \quad (1.3.9)$$

$$\rho_m \frac{\partial \vec{V}}{\partial t} = \frac{\vec{J} \times \vec{B}}{c} - \nabla p \quad (1.3.10)$$

$$\vec{E} + \frac{\vec{V} \times \vec{B}}{c} = \eta \vec{J} \quad (1.3.11)$$

$$\vec{\nabla} \times \vec{E} = -\frac{1}{c} \frac{\partial \vec{B}}{\partial t} \quad (1.3.12)$$

$$\vec{\nabla} \times \vec{B} = \frac{4\pi \vec{J}}{c} \quad (1.3.13)$$

where the (1.3.9) is the mass continuity equation, the (1.3.10) is the momentum transport equation, the (1.3.11) is the Ohm's law and the (1.3.12) and the (1.3.13) are the Maxwell equations. From the momentum transport equation comes that, for plasmas in a stationary state:

$$\vec{\nabla} p = \vec{J} \times \vec{B} \quad (1.3.14)$$

The equation (1.3.14) states that the currents and the magnetic field lines lie on isobaric surfaces, and the vector $\vec{J}$ and $\vec{B}$ are orthogonal each other. Furthermore, the equation (1.3.9), by making use of the equation:

$$\vec{\nabla} \times \vec{B} = \mu_0 \vec{j} \quad (1.3.15)$$

becomes:

$$\vec{\nabla} \left(p + \frac{B^2}{2\mu_0} \right) = \frac{1}{\mu_0} (\vec{B} \cdot \vec{\nabla}) \vec{B} \quad (1.3.16)$$

Assuming that $\vec{B}$ weakly varies along its direction, the second term is vanishing, so we can write [8]:

$$p + \frac{B^2}{2\mu_0} = \text{const.} \quad (1.3.17)$$

The term $\frac{B^2}{2\mu_0}$ is the magnetic pressure. The sum of the kinetic pressure and of the magnetic pressure is constant, then *in a plasma with a density gradient the magnetic pressure will be high where the particle pressure is low (i.e. in the plasma periphery), and it will be low where the plasma density is high*. This effect is due to the so-called *diamagnetic current*, which arises as a consequence of the pressure gradient. In a cylindrically shaped plasma, this current has the direction shown in figure 1.3.1 and it balances the pressure gradient by creating an internal magnetic field, in order to fulfill the equation (1.3.14).

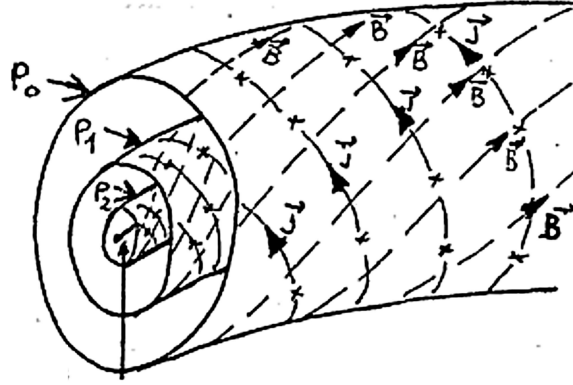


Figure 1.3.1: Current and magnetic field directions on the isobaric surface in case of a $\vec{\nabla}p \neq 0$ term in a cylindrically shaped plasma.

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The intensity of the diamagnetic current is [8]:

$$\vec{j}_\perp = \frac{\vec{B} \times \vec{\nabla} p}{B^2} = (KT_i + KT_e) \frac{\vec{B} \times \vec{\nabla} n}{B^2} \quad (1.3.18)$$

the intensity of the diamagnetic effect can be evaluated by making the ratio between the two terms of the equation (1.3.17). Such a ratio is usually indicated as the β **parameter**:

$$\beta \equiv \frac{\sum nkT}{\frac{B^2}{2\mu_0}} \quad (1.3.19)$$

Generally, β can be used to determine the quality of the plasma confinement:

- a) $\beta \ll 1$ then the magnetic pressure is much higher than the plasma kinetic pressure; this is the condition for an optimal confinement;
- b) $\beta > 1$ the plasma pressure is higher than that due to the magnetic field, then the plasma cannot be magnetically confined;
- c) $\beta = 1$ the diamagnetic effect generate an internal magnetic field exactly equal to the external one and two regions exist: one where *only the magnetic field is present*, the other where *only the plasma, without magnetic field, exists*.

It can be demonstrated that only in case of $\beta \ll 1$ the plasma is well confined, otherwise the equilibrium is unstable.

Working on the equation (1.3.11) it comes out that [16]:

$$\frac{\partial \vec{B}}{\partial t} = \frac{c^2}{4\pi\sigma} \nabla^2 \vec{B} + \vec{\nabla} \times (\vec{V} \times \vec{B}) \quad (1.3.20)$$

this equation is often called *magnetohydrodynamic equation*, as it describes the time evolution of the magnetic field¹⁸. If the plasma resistivity σ differs from zero, and the plasma is stationary, then the (1.3.20) reduces to:

$$\frac{\partial \vec{B}}{\partial t} = \frac{c^2}{4\pi\sigma} \nabla^2 \vec{B} \quad (1.3.21)$$

¹⁸Using the Maxwell equations the electric field has been eliminated from the equation, then the fields evolution can be studied only in terms of the magnetic field.

that is formally the equation of diffusion. It means that in resistive plasmas the magnetic field diffuses through the plasma on a timescale:

$$\tau_{diff} = \frac{L^2}{\eta} = \frac{4\pi\sigma L^2}{c^2} \quad (1.3.22)$$

Note that:

$$\eta \propto \frac{\nu_{coll}}{n_e} \quad (1.3.23)$$

i.e. the plasma resistivity linearly depends on the collision frequency.

For magnetically confined and resistive plasmas, the (1.3.21) says that the magnetic field diffuses into the plasma and the plasma diffuses through the magnetic field. Considering a perfectly conductor plasma the equation (1.3.20) becomes:

$$\frac{\partial \vec{B}}{\partial t} = \vec{\nabla} \times (\vec{V} \times \vec{B}) \quad (1.3.24)$$

from this equation it can be inferred that *the magnetic flux concatenated to the plasma must be conserved*. This means that *the magnetic field lines are "glued" to the plasma*. If one twists a plasma column, then also the magnetic field lines will be twisted, if one distorts a plasma column, then also the magnetic lines will be distorted.

1.3.3 Plasma diffusivity

The plasma diffusivity needs to be defined as it influences the particles lifetime. In neutral gases the collisions make slower the diffusion, as it can be determined from the Fick equation:

$$\Gamma = -D\vec{\nabla}n \quad (1.3.25)$$

where Γ is the particles flux, and D is the so called *diffusivity coefficient* given by:

$$D \equiv \frac{KT}{m\nu} \quad (1.3.26)$$

with ν collision frequency. For magnetized plasmas the diffusion is facilitated by the collisions, because they help the particles to leave the field lines around which they gyrate at the Larmor frequency. The classical diffusion coefficient along the perpendicular direction to the magnetic field is given by [8]:

$$D_{\perp} \propto \frac{1}{B^2} \quad (1.3.27)$$

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In some experiments the diffusion coefficient does not follow the above proportionality, but it is given by the formula:

$$D_{\perp} \propto \frac{1}{B} \quad (1.3.28)$$

this expression was derived by Bohm, considering that in some cases the particle diffusion is driven principally by plasma fluctuations and turbulence [10].

In fully or at least highly ionized plasmas the **electrons cannot diffuse without ions** and viceversa¹⁹. This property is usually called **Ambipolar diffusion** and its diffusion coefficient is given by [8]:

$$D_a \equiv \frac{\mu_i D_e + \mu_e D_i}{\mu_i + \mu_e} \quad (1.3.29)$$

where $\mu \equiv \frac{|q|D}{kT}$ is the so called "mobility". By virtue of its definition, the diffusion coefficient is mostly determined by the slower species, usually ions.

1.4 Waves in plasmas

The wave propagation in plasmas is quite a complex argument, as it involves many aspects of the plasma physics and the mathematical approach is not so easy. Waves can be studied by different point of views, especially using the kinetic or the fluidodynamical models. The most simple approach is the single particle, which however gives information about the principal modes propagating in the plasma. Here we limit our analysis to the general issues, with the aim to determine the cutoffs and resonances of each mode propagating in a magnetized plasma.

Plasma waves are always associated with a time and space varying electric field, and the propagation properties of such fields are determined by the dielectric properties of the plasma, which in turns depend on the steady-state values of magnetic or electric fields. A plasma may be both inhomogeneous and anisotropic, and this affects its dielectric properties.

The self-generated electrostatic plasma oscillations are crucial to determine the electromagnetic field propagation. These oscillations arise in a unmagnetized plasma²⁰ because of the electric fields acting among the plasma

¹⁹To distinguish between ionized gases and plasmas usually the quasineutrality condition is applied: plasmas must be globally neutrals, except for small regions whose dimension is given by the so-called Debye length Λ_D ($\Lambda_D \ll L_{plasma}$). This property is due to the electric field acting inside the plasma as a consequence of the ions-electrons interaction.

²⁰They also exist in magnetized plasmas, but in that case, as we will see, we have to distinguish between electrostatic waves along and across the magnetic field lines.

particles. If a group of electrons is displaced from its original position, an electric field arises which tends to retain the electrons to their initial position. This mechanism generates oscillations which in a *cold plasma approximation*, i.e. $T_e = T_i \sim 0$, do not propagate into the plasma²¹.

They can be found mathematically by perturbing the fluid equations²² (1.3.9)÷(1.3.13) with time varying harmonic electric and magnetic fields:

$$\vec{E}_1 = \hat{E}_1 e^{-i\omega t} \quad (1.4.1)$$

$$\vec{B}_1 = \hat{B}_1 e^{-i\omega t} \quad (1.4.2)$$

Assuming that the perturbing fields have small amplitudes, the fluid equations can be linearized and the Maxwell equation for B_1 gives [9]:

$$\vec{\nabla} \times \hat{B}_1 = -\frac{i\omega}{c} \left[1 - \frac{\omega_p^2 (1 + m_e Z/m_i)}{\omega^2} \right] \hat{E}_1 \quad (1.4.3)$$

considering that $m_e \ll m_i$, we can write for the dielectric constant:

$$\epsilon(\omega) \sim 1 - \frac{\omega_p^2}{\omega^2} \quad (1.4.4)$$

$\omega_p^2 = \frac{4\pi n_e e^2}{m_e}$, or $\omega_p^2 = \frac{n_e e^2}{\epsilon_0 m_e}$ in MKS units, is the so-called plasma frequency, and it can be demonstrated that electromechanical oscillations exist inside the plasma at this frequency. Note that a corresponding frequency can be determined also for ions.

The same results can be obtained by means of the kinetic theory. As a natural consequence of the kinetic approach it comes out that in *warm plasmas* the electrostatic plasma oscillations propagate into the plasma and they are weakly damped even in absence of collisions. In fact it can be demonstrated [9] that the dispersion relation becomes:

$$\omega_r^2 = \omega_p^2 (1 + 3k^2 \lambda_D^2) \quad (1.4.5)$$

$$\omega_i = -\sqrt{\frac{\pi}{8}} \frac{\omega_p}{|k^3 \Lambda_{De}^3|} e^{\left(\frac{1}{2k^2 \Lambda_{De}^2} + \frac{3}{2} \right)} \quad (1.4.6)$$

²¹Their group velocity is zero and they are dispersionless.

²²Actually, in order to calculate the plasma oscillations both for ions and electrons we should take into account two different sets of equations: the first one for ions and the second for electrons. Then we should assume a *two fluid model*.

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These weakly damped normal modes are generally called **Langmuir oscillations**, but they are also called plasma waves, plasma oscillations or electrostatic waves. The damping mechanism of such waves is predicted by the Vlasov (kinetic) theory, but it is completely outside the scope of fluid theory. In addition, it is completely collisionless. The electrostatic potential damps according to:

$$\phi_1(t) \propto e^{\omega_i t} \quad (1.4.7)$$

where:

$$\omega_i \propto \frac{\partial f}{\partial v} \quad (1.4.8)$$

with f energy distribution function, and v electron velocity. The decrease of the ϕ_1 amplitude with time was firstly calculated by Landau, and for this reason this mechanism is named **Landau Damping**. As it can be easily seen from 1.4.8, the damping is proportional to the slope of the particles distribution function decrease with the velocity. Then the Landau damping is a resonant phenomenon proportional to the slope of the distribution function, for those particles moving with a thermal velocity near to the wave phase velocity. The damping is usually explained by noting that for distribution decreasing with increasing v , there are more particles moving slightly slower than the wave than particles moving slightly faster. If the slower particles are accelerated by the waves, then the wave's energy will reduce, and the wave damps. As the number of slower electrons is higher than faster, then summing over a great number of electrons it results that the wave damps and the electrons are globally heated.

By looking again to the expression (1.4.4) some further considerations can be done. In particular note that electromagnetic waves with frequencies lower than ω_p cannot propagate into the plasma ²³. In other words *in homogeneous and non magnetized plasmas the density cannot exceed the so called cutoff value given by:*

$$n_{cutoff} = 4\pi^2 \frac{m\epsilon_0}{e^2} f_p^2 \quad (1.4.9)$$

We can express the above formula in a more convenient form:

$$n_{cutoff} = 1.26 \cdot 10^{-8} f_p^2 \quad (1.4.10)$$

If we apply a magnetostatic field to the plasma, then it will become an anisotropic medium for the electromagnetic waves propagation. Necessarily

²³Otherwise the dielectric constant becomes negative.

the dielectric constant will transform in a tensor ($\bar{\epsilon}$), as the field propagation will depend on the direction of the wave injection into the plasma. $\bar{\epsilon}$ can be determined in the cold plasma approximation and in the high frequency limit (neglecting the ions response to the electromagnetic field because of their high inertia), by means of the single particle approach. Starting from the equation (1.3.1), by adding the term connected with the magnetostatic field, oriented, for sake of simplicity, along the $\hat{z}$ axis, one can find the $\vec{k}$ vectors of waves propagating in magnetized plasmas [13]. It comes out that $\vec{k}$ depends on the angle formed by the incoming wave with the magnetic field direction. In addition, two different constants of propagation exist for the wave [13]:

$$k'_\theta = \frac{\omega}{c} \left[1 - \frac{X(1+iZ-X)}{(1+iZ)(1+iZ-X) - \frac{1}{2}Y_T^2 + \sqrt{\frac{1}{4}Y_T^4 + Y_L^2(1+iZ-X)^2}} \right]^{\frac{1}{2}} \quad (1.4.11)$$

$$k''_\theta = \frac{\omega}{c} \left[1 - \frac{X(1+iZ-X)}{(1+iZ)(1+iZ-X) - \frac{1}{2}Y_T^2 - \sqrt{\frac{1}{4}Y_T^4 + Y_L^2(1+iZ-X)^2}} \right]^{\frac{1}{2}} \quad (1.4.12)$$

similarly to the case of electromagnetic waves propagating in non isotropic crystals. Here $X = \left(\frac{\omega_p}{\omega}\right)^2$, $Y = \left(-\frac{\omega_c}{\omega}\right)$, and $Z = \left(\frac{\omega_{eff}}{\omega}\right)$, where ω is the wave frequency, ω_c is the cyclotron frequency and ω_{eff} is the collision frequency. The imaginary parts of $\vec{k}$ is connected with the wave collisional damping, which however appears unrealistic in ECRIS plasmas. Neglecting the collisions, for a wave incidence angle of 0° we obtain:

$$k'_0 = \frac{\omega}{c} \sqrt{1 - \frac{X}{1+Y}} = \frac{\omega}{c} \sqrt{1 - \frac{\omega_p^2}{\omega(\omega - \omega_c)}} \quad (1.4.13)$$

$$k''_0 = \frac{\omega}{c} \sqrt{1 - \frac{X}{1-Y}} = \frac{\omega}{c} \sqrt{1 - \frac{\omega_p^2}{\omega(\omega + \omega_c)}} \quad (1.4.14)$$

if the wave incidence angle is 90° we obtain:

$$k'_{\frac{\pi}{2}} = \frac{\omega}{c} \sqrt{1 - X} = \frac{\omega}{c} \sqrt{1 - \frac{\omega_p^2}{\omega^2}} \quad (1.4.15)$$

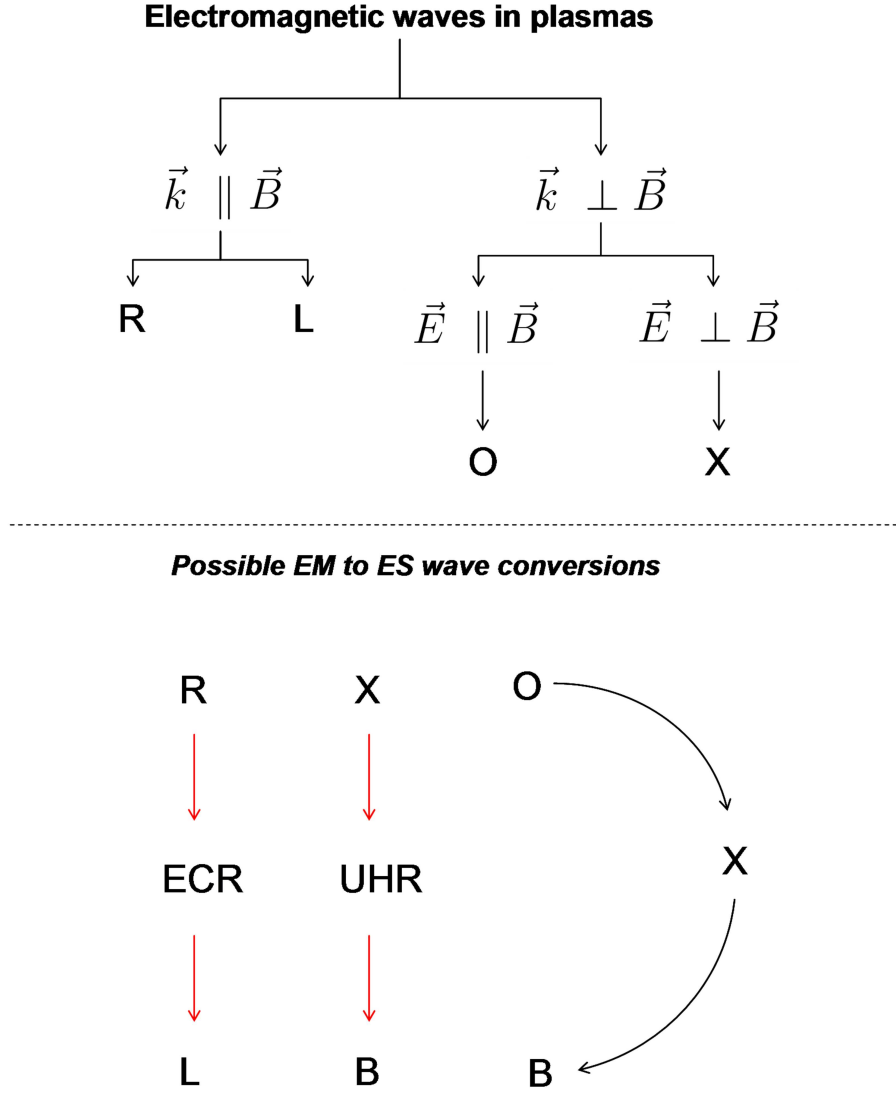


Figure 1.4.1: Scheme of the wave propagation in anisotropic plasmas. According to the mutual orientation of the $\vec{k}$, $\vec{E}$ and $\vec{B}$, the last being the magnetostatic field, the waves have different characteristics and are named R , L , O and X . Modes may convert each other at the cutoffs or resonances (R waves may convert in Langmuir waves at the ECR [17], X may directly or via OXB convert in Bernstein waves at the UHR [18]).

$$k''_{\pi} = \frac{\omega}{c} \sqrt{1 - \frac{\frac{\omega_p^2}{\omega^2}}{1 - \frac{\omega_g^2}{(\omega^2 - \omega_p^2)}}} \quad (1.4.16)$$

Hence for a given arbitrary direction, defined by θ , we have two waves characterized by k'_θ and k''_θ . As a function of X , k'_θ is more similar to a wave propagating in an isotropic plasma, then the associated wave will be called **ordinary wave**; the wave with k''_θ is named **extraordinary wave**. Considering the two extreme conditions (0° and 90°) we actually have four waves and in plasma physics another classification prevails, according to the scheme shown in figure 1.4.1. Usually, the waves propagating along the magnetic field direction are named ***R* and *L* waves**, according to their polarization²⁴. The waves propagating along the perpendicular direction to $\vec{B}_0$ are named ***O* and *X***, according to the orientation of the wave electric field with respect to the magnetostatic field direction.

vectors orientation	refraction index	wave type
$\vec{B}_0 = 0$	$\omega^2 = \omega_p^2 + k^2 c^2$	light waves
$\vec{k} \perp \vec{B}_0, \vec{E} \parallel \vec{B}_0$	$\frac{c^2 k^2}{\omega^2} = 1 - \frac{\omega_p^2}{\omega^2}$	<i>O</i> wave
$\vec{k} \perp \vec{B}_0, \vec{E} \perp \vec{B}_0$	$\frac{c^2 k^2}{\omega^2} = 1 - \frac{\omega_p^2 \omega^2 - \omega_p^2}{\omega^2 \omega^2 - \omega_h^2}$	<i>X</i> wave
$\vec{k} \parallel \vec{B}_0$	$\frac{c^2 k^2}{\omega^2} = 1 - \frac{\omega_p^2 / \omega^2}{1 - (\omega_c / \omega)}$	<i>R</i> wave
$\vec{k} \parallel \vec{B}_0$	$\frac{c^2 k^2}{\omega^2} = 1 - \frac{\omega_p^2 / \omega^2}{1 + (\omega_c / \omega)}$	<i>L</i> wave

Table 1.1: Cutoffs and resonances of the waves propagating inside the plasmas [8].

The *O*-mode is characterized by $\vec{k} \perp \vec{B}_0$ and $\vec{E} \parallel \vec{B}_0$. If $\vec{k} \perp \vec{B}_0$ and $\vec{E} \perp \vec{B}_0$, whatever will be the direction of $\vec{E}$, the wave will be called *X* mode.

Let us consider the cutoffs and resonances of the *R*, *L*, *O* and *X* waves. For the *R* waves, the refraction index $\frac{c^2 k^2}{\omega^2}$ becomes infinite at $\omega = \omega_c$. In

²⁴As it can be intuited by the denomination, the *R* wave is a right hand polarized wave, the *L* a left hand polarized wave

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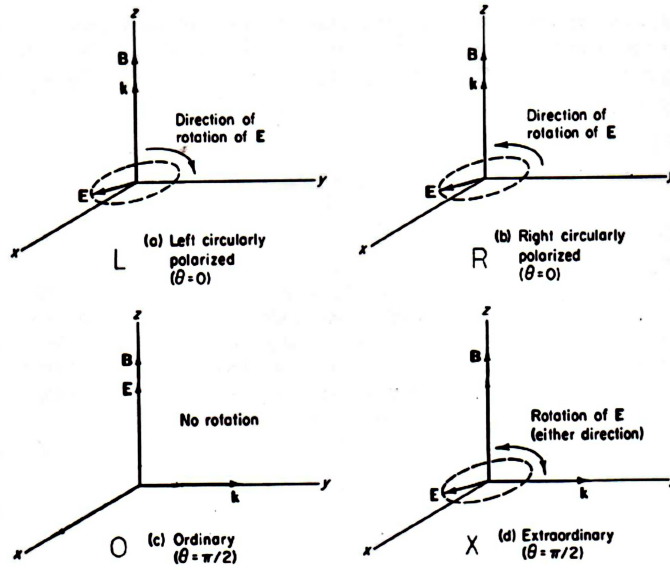


Figure 1.4.2: Diagram showing the possible orientations of the electric field with respect to the magnetostatic field, and also the possible polarizations in waves propagating in magnetized plasmas [10].

this case the microwave frequency is equal to the Larmor frequency, hence **the wave is in resonance with the cyclotron motion of the electrons**. This resonance is of primary importance for ECRIS and takes the name **Electron Cyclotron Resonance**. More generally the cutoffs and the resonances of the various propagation modes can be summarized by reporting the different $\frac{c^2 k^2}{\omega^2}$ formulas obtained from the (1.4.11) and (1.4.12). A schematic summary of the main electron waves²⁵ is given in table 1.1.

In the table 1.1 $\omega_h = \sqrt{\omega_c^2 + \omega_p^2}$ is the so called *upper hybrid frequency*. This frequency practically corresponds to the frequency of the Langmuir oscillations in non magnetized plasmas. Electrostatic electron waves across $\vec{B}$ have this frequency, while those along $\vec{B}$ have the usual $\omega = \omega_p$. The X

²⁵With the statement *electron waves* the high frequency waves are usually indicated, i.e. the electromagnetic waves. For these waves the ions can be considered as an immobile medium, as the wave frequency is too high to have some ions' response to the electromagnetic field. It is usually assumed that ions ensure the quasi-neutrality, but they do not take part to the electromagnetic wave propagation.

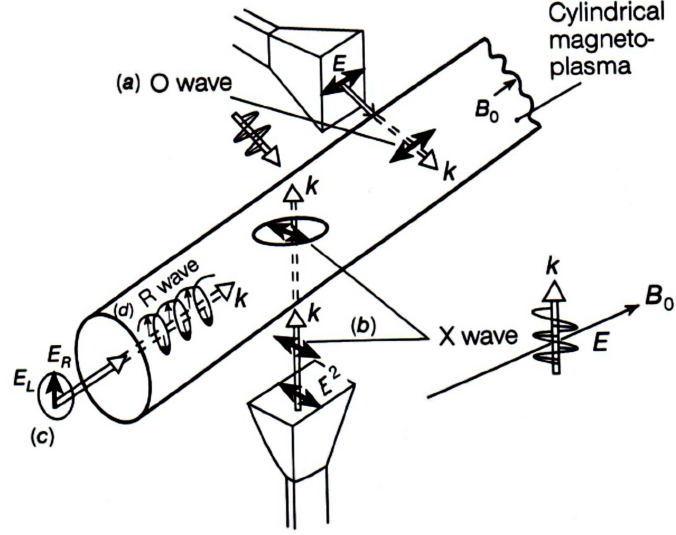


Figure 1.4.3: Schematic representation of the possible ways to inject electromagnetic waves into a column of a magnetized plasma. The different modes can be excited, according to the location of the waveguide providing the electromagnetic field, and also according to the wave polarization inside the waveguide. [10].

mode has a resonance at the upper hybrid frequency and it occurs if:

$$\omega^2 = \omega_p^2 + \omega_c^2 = \omega_h^2 \quad (1.4.17)$$

The L wave does not suffer any resonance, as well as the O mode. Instead they suffer a cutoff that can be determined by the equations reported in table 1.1 when the index of refraction $\frac{c^2 k^2}{\omega^2}$ goes to zero. The figure 1.4.3 shows schematically the direction of propagation of the four modes in a cylindrically shaped plasma column. A powerful method to visualize the cutoffs and the resonances of the various modes is to plot them on the so called CMA diagram²⁶ [8]. This diagram is reported in figure 1.4.4.

When the propagation occurs at a given angle with respect to the magnetic field, the phase velocities change with the angle. Some of the modes

²⁶The CMA diagram takes the name from Clemmow, Mullaly and Allis who proposed the graphical representation of the many cutoffs and resonances of electromagnetic waves propagating in a cold plasma.

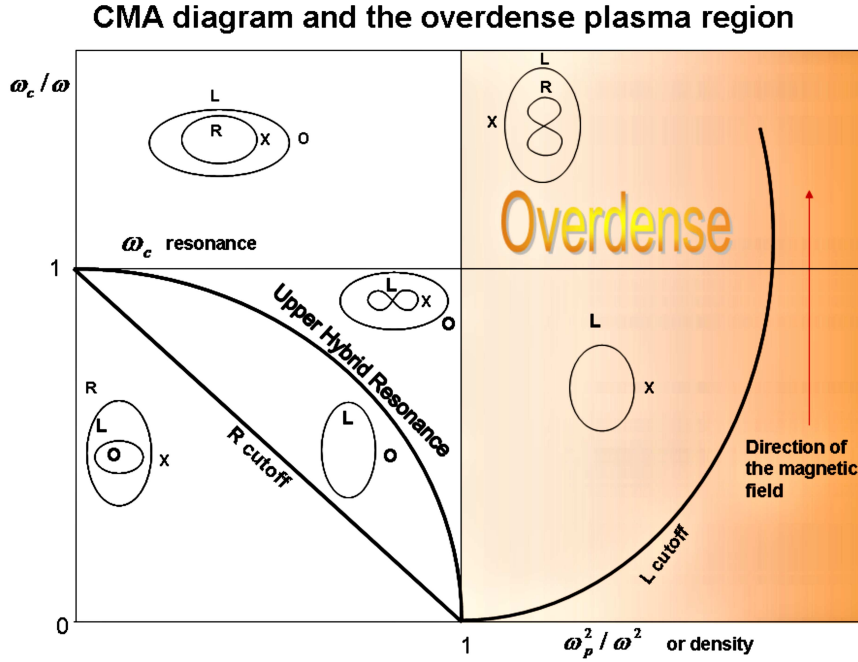


Figure 1.4.4: A detailed view of the Clemmow-Mullaly-Allis diagram showing the region of the so-called overdense plasma, i.e. that region where the plasma density exceeds the cutoff of the ordinary (O) mode.

listed above with $\vec{k} \parallel \vec{B}_0$ or $\vec{k} \perp \vec{B}_0$ change continuously into each other. Such a complicated picture is greatly simplified by the CMA diagram, which however works only in case of cold plasma approximation ($T_e = T_i = 0$). Any finite temperature modification implies a so great number of complications that the diagram cannot be further employed.

The diagram can be viewed as a plot of the ω_c/ω versus ω_p^2/ω^2 , or equivalently a plot of the magnetic field versus the plasma density. For a given frequency ω , any experimental situation characterized by ω_c (i.e. the magnetic field) and by ω_p (i.e. the plasma density) is denoted by a point on the graph. The total plane is divided into several zones, and the boundaries of each zone are the cutoffs and the resonances mentioned above. For example, considering an X wave propagating from a region of high magnetic field, inside a plasma with a fixed value of density, the value of the magnetic field where the UHR occurs can be easily determined on the graph.

The vertical line at $\omega_p^2/\omega^2 = 1$ is the so called O cutoff, and it corresponds

to the cutoff density in an unmagnetized plasma. Above this value the plasma will be **overdense**. The small diagrams in each region delimited by the cutoffs and the resonances (i.e. the eight-shaped or ellipsoidal small diagrams) indicate not only which wave propagates, but also the qualitative variation of the phase velocity with the angle (considering that the magnetic field is directed along the vertical direction).

The kinetic theory provides more information about the wave propagation than the macroscopic fluid and the single particle approaches. These two models apply only to low temperature plasmas, or for particle velocities near to the fluid element one. *The Landau damping, the ion waves and all that waves at harmonics of the cyclotron frequency depend on plasma temperature, thus requiring the kinetic approach to be described.*

1.5 Bernstein waves

Electron Bernstein waves (EBW or BW) are electrostatic waves propagating across the magnetic field in a hot plasma. Then the cold plasma approximation cannot be used, and the fluid and single particle approach fail. They are so named after Ira B. Bernstein, who described these modes in 1958 [19]. The Bernstein waves are sustained by electron cyclotron motion and can be thought as fronts of electron rarefaction and compression perpendicular to the magnetic field with wavelength λ of the order of electron Larmor radius ρ . Clearly, a finite Larmor radius should be considered. In a Maxwellian plasma this also implies that the temperature be finite. In fact in a cold magnetized plasma the electrons are tight to the field lines and the charge separation (and space-charge waves) is possible only along the field²⁷. At high enough temperatures the electron gyroradius $\rho = mv_T/(eB)$ (where v_T is the electron thermal velocity and the other notations are standard) becomes comparable with the wavelength λ . The EBW occurs when the electron gyrophase is organized in such a way that the space charge distribution has maxima and minima, which are periodically distributed perpendicularly to external magnetic field lines [18]. If additionally the wave frequency ω is of the order of the cyclotron frequency ω_c , the electrons can move in phase with the local wave field and "participate" in wave motion, in the following sense: by moving in response to the wave field, the electrons alter the local concentration of electric charge, giving their own contribution to the electric field [20]. This means that the electrons gyromotions cannot be random as

²⁷The existence of electrostatic waves in unmagnetized plasmas can be proved by means of the fluid theory in the framework of the cold plasma approximation, as done before.

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in the absence of the wave, otherwise the wave cannot propagate. We need some regular bunching of the electron charge in order to facilitate the wave propagation.

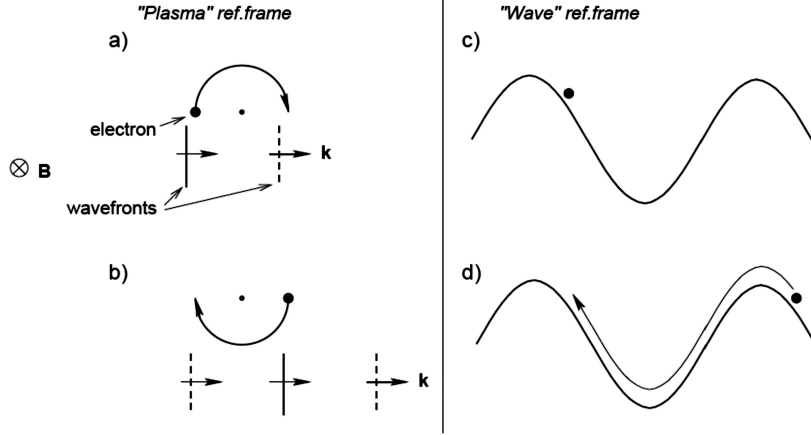


Figure 1.5.1: Description of the wave-particle interaction at the basis of the BW propagation. a) and b): half cycles of electrons gyrating in a magnetic field and in presence of an electrostatic wave (solid wavefront: $E > 0$, dashed: $E < 0$). c) and d) wave-particle interaction in a frame moving with the wave potential. [20].

In order to be constructive, the contribution due to the plasma electrons must clearly have proper frequency, wavelength and phase. Supposing that the electrons oscillate at nearly the same frequency as the wave and that $\rho = \lambda/4$, the electron experiences a constant field during the first half of the gyromotion, as it is clearly shown in figure 1.5.1a), whereas it feels a rapidly varying field in the second half cycle: in fact the electron moves in the opposite direction with respect to the wave front; it moves back for a space equal to $\lambda/2 = 2\rho$, while the wave moves forth for a space equal to $\lambda/2$ (see figure 1.5.1b)). Then during the second half cycle the electron experiences the entire wave, with a net electric field equal to zero. This means that only in the first half cycle some energy exchange may occur between the wave and the particle; a mechanism like the Landau damping occurs: those electrons slightly slower than the wave are accelerated by the wave front, the faster ones are decelerated. Summing over many electrons the final result is the *phase organization*, or *phase bunching*, shown in figure 1.5.2. *The gyrophases are no longer random as they were before the impinging of the wave, but they are organized depending on the gyrocenter positions, in such a way that*

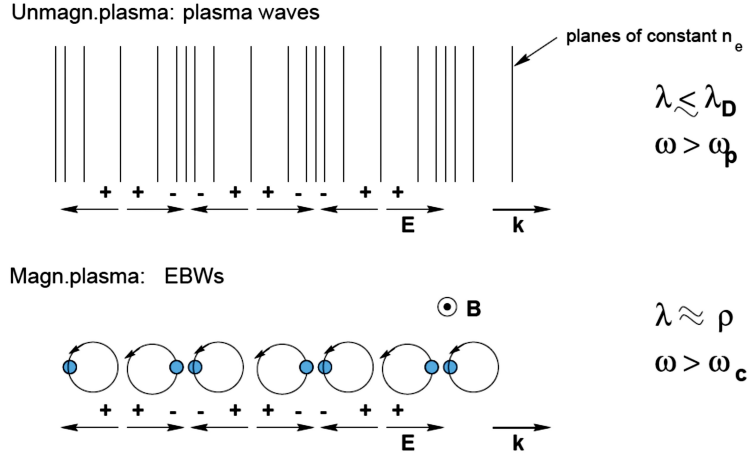


Figure 1.5.2: Description of the charge bunching in plasmas without and with a magnetic field. In case of magnetized plasmas the BW propagate if the electrons gyrate with a particular organization around the magnetic field lines. [21].

electrons form propagating planes of net charge accumulation-rarefaction. The discussion done so far has been applied to resonating electrons (i.e. $\omega \simeq \omega_c$, $\lambda = 4\rho$), but other coherent motions are possible more generally at cyclotron harmonics $\omega \simeq n\omega_c$.

In the real world the situation is different: not all the electrons gyrate with the same Larmor radius around the field lines, but a Maxwellian electron energy distribution function (EEDF) can be guessed²⁸. Hence the electrons have different gyroradius and the above discussion is not strictly valid. If the wave phase velocity is much higher than the thermal velocity, i.e. $\omega/k \gg v_{th}$, then the electrons will experience many wave periods even in the first half cycle of their gyromotion. This means that the previous condition cancels any possibility for BW propagation, as they cannot be sustained by the plasma electrons. On the contrary, as long as $\omega/k < 3v_{th}$, the "bulk plasma" populated by a great amount of electrons has the right velocity to sustain the BW propagation.

One of the most important properties of EBW is that they can propagate in overdense plasmas without any cutoff. As long as $\omega \geq \omega_c$, the BW propagate in overdense plasmas as $\omega_c > \omega_p$, then the electrons are still able to transfer information for the wave propagation outside any

²⁸About the applicability of such hypothesis see section 1.2.

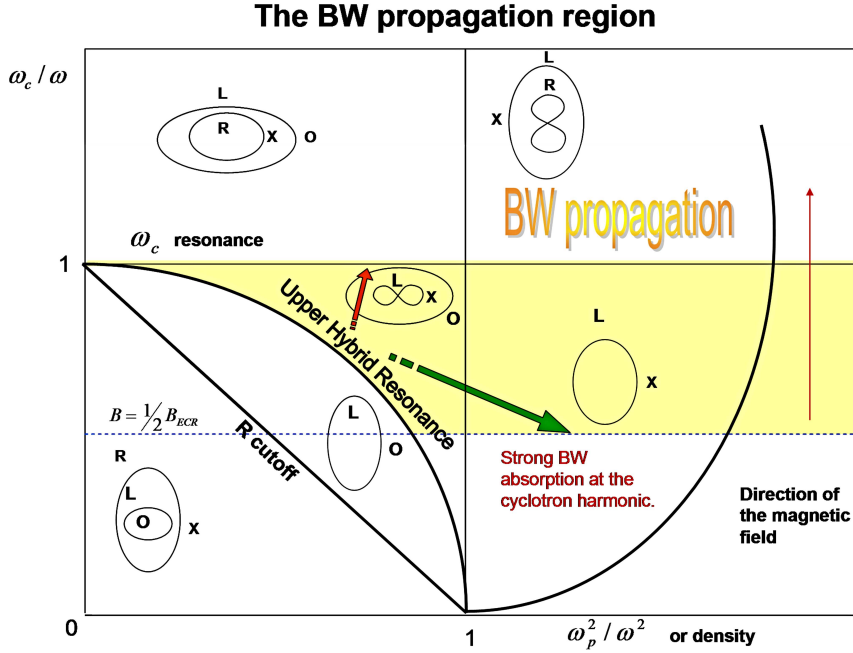


Figure 1.5.3: Representation on the CMA diagram of the region of possible propagation for the Bernstein waves.

Debye sphere²⁹.

Because of their electrostatic nature, **the Bernstein waves cannot propagate outside the plasma, and their boundary is constituted by the UHR layer** [18]. Figure 1.5.3 shows the regions, in the CMA diagram, where the BW can propagate.

The BW are undamped except in the vicinity of the cyclotron harmonics [18]. In fact for $\omega \sim n\omega_c$ their refractive index has an imaginary part. A possible region of BW absorption is shown in figure 1.5.3, where the possible second harmonic absorption has been indicated.

It is important to underline that *O* and *X* modes can be determined either in a fluidodynamics and in a kinetic frame. However, by applying the kinetic theory, it can be demonstrated that the *O* mode is a pure electromagnetic wave, whereas *X* has a small electrostatic component. Near the UHR layer the *X* mode also adopts an electrostatic character. Its refractive

²⁹The Debye sphere has a dimension fixed by the density, and it becomes smaller and smaller as the density increases.

index becomes large and the longitudinal electric field dominates. Then near the UHR the X and EBW-modes coincide; the X -wave is coupled into an EBW and vice versa. In a full wave view the X -wave expresses a longitudinal polarization which becomes dominant when the X -wave reaches the UHR [18]. In the vicinity of the UHR the phase velocity is strongly reduced until it approximates the thermal velocity. Here the wavelength reduces down the order of the electron gyro radius and an EBW is excited³⁰. Figure 1.5.3 features that once it is created near the UHR layer, the BW propagate according to the possible direction indicated on the graph and they can exist everywhere in the shadowed region [18].

Hence **to generate BW in plasmas an X mode is required, together with the presence of the UHR layer**. This result will be largely discussed in chapter 5, where some possible techniques to provide BW in ECRIS plasma will be presented.

1.6 Simple mirror and minimum-B configuration for the plasma trapping

The confinement of charged particles, and then of the plasma, has a primary importance in many field of physics, and in particular in the field of nuclear fusion. In order to fulfill the Lawson [22] criterion, a plasma must be sustained for a long time inside the confining structure, with high densities and temperatures. In case of ECRIS, and in general of all the ion sources, in order to satisfy the relation (1.2.16) the ions must live inside the plasma for time long enough for a high number of collisions.

The most simple device for plasma confinement was early investigated by Fermi and it is named **Simple Mirror**³¹. In this magnetic configuration the field is provided by two solenoids with coinciding axes, located to definite distances each other. Figure 1.6.1 features the shape of the field lines.

The principal ingredients for charged particles confinement are here summarized [8]:

1. **a magnetic field gradient $\vec{\nabla}_{\parallel} B$ is needed, directed along the direction of the field itself;**

³⁰Note that all the condition required to sustain the BW propagation are in this way satisfied

³¹Fermi proposed the simple mirror configuration, with sources of the magnetic fields moving towards the center of the system, as a possible mechanism involved in the generation of cosmic rays.

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2. *even in the simple cases, as for Simple Mirror configurations, **the field** maybe axisymmetric but it **must have a radial component**;*
3. ***the velocity vector of the charged particle must form a particular angle with the field lines.** A boundary exists for this angle, and for smaller angles the particle is not confined anymore.*

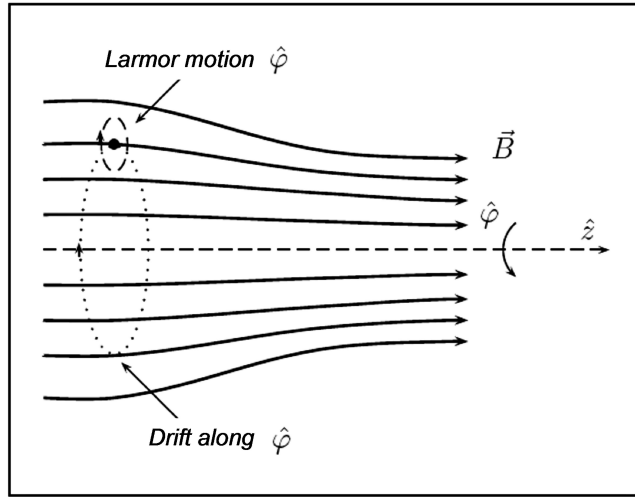


Figure 1.6.1: Magnetic field lines produced in a simple mirror configuration. The density of the field lines increases going to right, and this generates the gradient which allows the confinement of the charged particles. [8].

In figure 1.6.1 only one of the two mirrors is displayed, and the Larmor motions and the drift around the axis are shown. It can be seen that the field strongly depends on the z coordinate; in addition, we can neglect the dependence on φ , i.e. $B_\varphi = 0$. The B_r component cannot be neglected and it plays a fundamental role for trapping.

All the field components must satisfy the Maxwell equation: $\vec{\nabla} \cdot \vec{B} = 0$.

The simplest profile to trap charged particles satisfying this condition is [23]:

$$\begin{aligned} B_z &= B_0 + B_1 z^2 \\ B_r &= -r z B_1 \end{aligned} \tag{1.6.1}$$

with constant B_0 and B_1 . This approximation will be used in the next

1.6 Simple mirror and minimum-B configuration for the plasma trapping

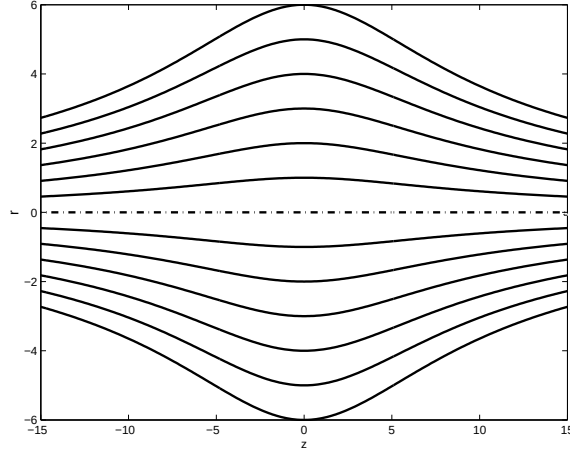


Figure 1.6.2: Force lines of a simple mirror like configuration with a parabolic approximation of the z field component. [25].

chapters where simulations about the particles motions inside magnetic traps will be discussed.

The figure 1.6.2 shows the force lines calculated by means of MATLAB [24] on the basis of the (1.6.1), taking $B_0 = 0.57$ and $B_1 = 80$.

The particle motion evolves by keeping the guiding center of the spiral-like trajectory "glued" to a fixed field line; slow drifts also affect the motion, the most important one around the system axis³². Transversally, i.e. along the perpendicular direction with respect to the mirror axis, the particles are automatically confined because they gyrate around the magnetic field lines³³. From the Lorentz force comes that the component along the z axis can be written as [8]:

$$F_z = \frac{1}{2} q v_\phi r \left(\frac{\partial B_z}{\partial z} \right) \quad (1.6.2)$$

and by replacing B_r with the second expression of (1.6.1) we have:

$$F_z = q v_\phi r z B_1 \quad (1.6.3)$$

³²About the drift motions of charged particles in magnetized plasmas see [8]. In general, drift are generated by electric fields, by the magnetic field lines curvature and by magnetic gradients. A force is determined and its effect on the guiding center motion can be evaluated: $\vec{v}_f = \frac{\vec{F} \times \vec{B}}{B^2}$, where $\vec{v}_f$ is the drift velocity, $\vec{F}$ is the particular force which arises because of field lines curvature or gradients, and B is the magnetic field.

³³This statement is true only in a single particle approach. Considering the plasma as a fluid, and then applying the MHD, we will see in section 1.6.3 that the simple mirror structure cannot guarantee the plasma stability and then a long time confinement.

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Then along z a force arises and it is able to confine the particles if some conditions are satisfied.

By averaging F_z on a gyro-period, considering a particle moving along the mirror axis we obtain [8]:

$$\bar{F}_z = -\frac{mv_{\perp}^2}{B} B_1 z \quad (1.6.4)$$

then the force is directed in the opposite direction with respect to the particle motion. We can also define the so-called **magnetic moment** of the charged particle as [8]:

$$\mu \equiv \frac{1}{2} \frac{mv_{\perp}^2}{B} \quad (1.6.5)$$

Then the force along z becomes:

$$\bar{F}_z = -\mu \left(\frac{\partial B}{\partial z} \right) \quad (1.6.6)$$

generalizing for a whatever gradient along the particle motion:

$$\vec{F}_{\parallel} = -\mu \frac{\partial B}{\partial s} = -\mu \vec{\nabla}_{\parallel} B \quad (1.6.7)$$

where $d\vec{s}$ is a line element along $\vec{B}$ direction.

The magnetic moment μ is an adiabatic invariant for the particle motion³⁴. The confinement of charged particles in mirror-like devices can be studied in terms of μ invariance. *The breakdown of the adiabatic invariance of μ expels the particle from the confinement.* The adiabatic invariance is valid as long as the B_z component slowly varies with z during a gyro-period; this condition can be written in the following manner:

$$\frac{r_L}{L} \ll 1 \quad (1.6.8)$$

i.e. the Larmor radius r_L must be much smaller than L , the characteristic length for the B variation³⁵.

The mirror effect in terms of the μ invariance can be explained as it follows: consider a proof particle moving from the center to the periphery

³⁴To read more about the adiabatic invariants in plasmas see [8]

³⁵This properties of the μ invariance is of primary importance also for ECRIS physics, as it will be demonstrated in section 4.13. If such condition is not satisfied non adiabatic effect begin to play their role, leading to the expulsion of the particle from the confinement.

1.6 Simple mirror and minimum- B configuration for the plasma trapping

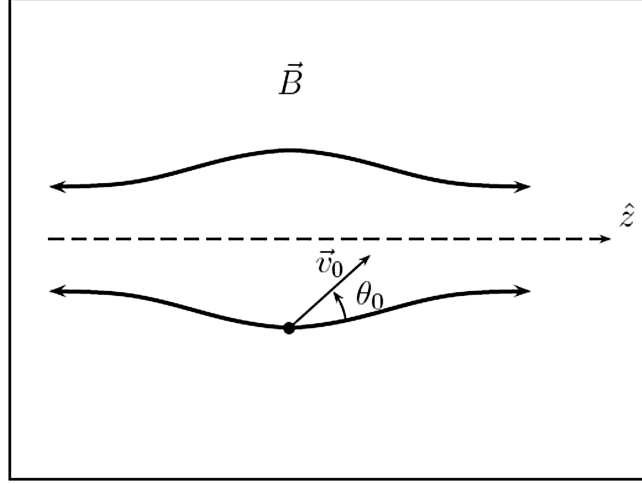


Figure 1.6.3: Schematic representation of the geometrical meaning of the θ_0 angle.

of the magnetic bottle³⁶; it sees a growing magnetic field; as μ must be kept constant the perpendicular component of the velocity $v_{\perp}$ must increase because of the B increase. As the magnetic field does not make any work on the system, then the kinetic energy remains constant. Hence, as $v_{\perp}$ increases, the $v_{\parallel}$ decreases. If B is sufficiently high at the periphery, there will be a mirror point where $v_{\parallel} = 0$ and the particle is reflected.

The ability to confine the particles by means of a magnetic bottle depends on the so called **loss cone aperture** [10]. In fact it can be easily understood from the above discussion that particles with $v_{\perp} = 0$ do not feel any mirror force, as they have $\mu = 0$; anyway $v_{\perp} \neq 0$ does not mean that the particle is confined, because particles with $\frac{v_{\perp}}{v_{\parallel}}$ very small on the midplane escape from the trapping if B_m at the system boundary is not sufficiently high to make $v_{\parallel}$ null. A minimum angle exists at the midplane (see figure 1.6.3), formed by the velocity vector and the magnetic lines, below which particles are lost because the maximum field cannot guarantee the mirroring.

The smaller is θ , the stronger must be B in order to have the mirroring effect. A mathematical relation can be found to link this angle and the **mirror ratio**, i.e. the ratio between the maximum and the minimum field

³⁶The Simple Mirror systems are often named "magnetic bottles".

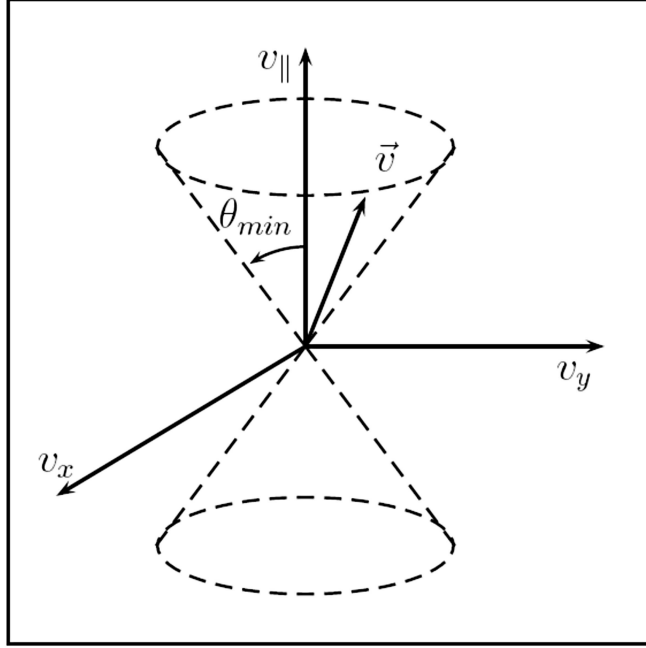


Figure 1.6.4: Schematic representation of the loss cone in the velocity space; θ_{min} indicates the cone aperture.

in the system:

$$\sin^2 \theta_{min} = \frac{B_0}{B_m} \equiv \frac{1}{R_m} \quad (1.6.9)$$

where R_m is the mirror ratio $\frac{B_m}{B_0}$. The equation (1.6.9) allows to determine the so called **loss cone**, i.e. that region in the velocity space ($\theta < \theta_{min}$) where particles that cannot be mirrored lie. The figure 1.6.4 clearly shown the main characteristics of the loss cone.

1.6.1 Non-adiabatic effects on μ

As defined in the previous section, μ is not an integral of the particle motion, but only an adiabatic invariant, which remains constant as long as some conditions are satisfied. Several attempts to describe the possible non-adiabatic evolution of the particle dynamics have been carried out in the past; we will follow the theory of Dubinina [26] which starts from the condition of the adiabatic motion of a charged particle in a non uniform magnetic field. Let us recall the condition for the adiabaticity:

$$\left(\frac{r_L}{L}\right) \ll 1 \quad (1.6.10)$$

Then we can introduce the following theorem:

Theorem 1 *A critical parameter exists: $\left(\frac{r_L}{L}\right)_1 \ll 1$ for which, if $\left(\frac{r_L}{L}\right) < \left(\frac{r_L}{L}\right)_1$ the particle motion is stable and the particle will remain indefinitely inside the magnetic trap, otherwise if: $\left(\frac{r_L}{L}\right) > \left(\frac{r_L}{L}\right)_1$ the particle leaves the plasma immediately³⁷*

then the parameter $\frac{r_L}{L}$ can be named **adiabaticity parameter**. Let us express this quantity in terms of the magnetic field gradient:

$$\frac{r_L}{L} = r_L \frac{\vec{\nabla} B}{B} \quad (1.6.11)$$

If the conditions required by the theorem 1 are fulfilled, then the trapping is quite stable, the lifetime of the particle inside the trap is very long, and it can be expelled from the plasma only by means of collisions³⁸. If the "stable trapping" conditions are not satisfied, the magnetic moment varies during the particle motion, and after N reflections the particle will leave the trap.

The confinement time depends on the adiabatic parameter according to the equation [26]:

$$\tau = Ae^{\left(\frac{B}{\frac{r_L}{L}}\right)} \quad (1.6.12)$$

where A and B are constant. The closer we are to the adiabaticity condition, the longer will be τ . Taking the limit $\frac{r_L}{L} \rightarrow 0$ then $\tau \rightarrow 0$.

1.6.2 Resonances in particle trapping

In the previous pages we said that the dynamics of a particle in a simple mirror system is characterized by two different motions: a fast rotation around the magnetic field lines (the Larmor rotation), and a slower bouncing between the two magnetic mirrors (with frequency ω_b). It is possible to

³⁷Immediately means, in this case, in a time comparable to that required to cross the magnetic trap.

³⁸The Spitzer collisions are able to change the direction of the particle velocity, then they can inject electrons into the loss cone.

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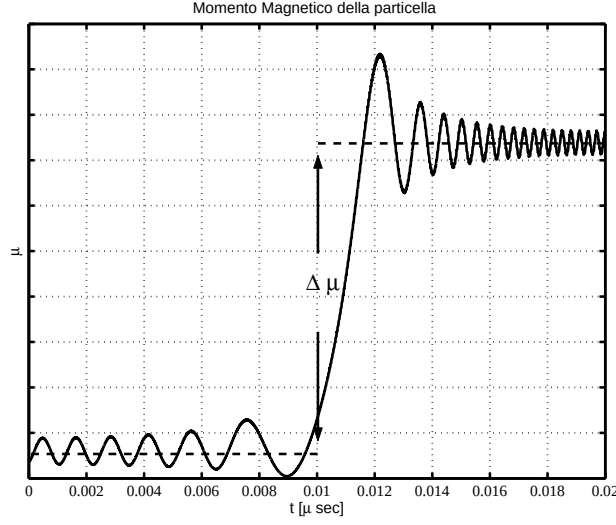


Figure 1.6.5: Trend of the magnetic moment μ for a charged particle moving in a simple mirror magnetic field [25].

demonstrate that possible resonances in a simple system constituted by a linear oscillator with a sinusoidal modulation may lead to uncontrollable variations of the adiabatic invariants [27].

Generally, if in these systems it occurs that:

$$n = \frac{2\bar{\omega}}{\Omega_m} \quad (1.6.13)$$

where $\bar{\omega}$ is the oscillator frequency and Ω_m the modulation frequency, then a parametric resonance occurs, and the oscillator dynamics changes abruptly and any further evolution of the system cannot be exactly foreseen.

In case of Simple Mirrors, we can replace Ω_m with the bouncing frequency ω_b , and the $\bar{\omega}$ with the Larmor frequency ω_L . This means that some possible resonances between the bouncing frequency and the Larmor frequency may lead to parametric resonances able to strongly change the value of μ , thus breaking its adiabatic invariance.

Many simulations were carried out [25] in order to investigate the μ dynamics in several magnetic field configurations. It was determined that the time variation of μ is characterized by many small oscillations between the mirror point and the midplane of the trap, and by a rapid and wider variation when the particle crosses the midplane, as shown in figure 1.6.5.

1. we indicate with $\Delta\tilde{\mu}$ the small oscillations occurring at the Larmor

1.6 Simple mirror and minimum-B configuration for the plasma trapping

frequency³⁹;

2. we indicate with $\Delta\mu_n$ the *jump* of μ at the n crossing of the midplane.

To avoid any possible breakdown of the μ invariance due to the small oscillations we need [27]:

$$\frac{\Delta\tilde{\mu}}{\mu} \sim \frac{\omega_b}{\omega_L} \sim \epsilon \ll 1 \quad (1.6.14)$$

and the dimensionless parameter ϵ is again the so-called ***adiabatic parameter***.

About the μ jump at the midplane we have:

$$\frac{\Delta\mu_1}{\mu} \simeq e^{-\frac{1}{\epsilon_a}} \quad (1.6.15)$$

In [25], simulations put in evidence that in some proof configurations of the magnetic field the confinement is not possible at all because of parametric resonances.

1.6.3 The MHD stability of magnetically confined plasmas

The simple mirror configuration allows to confine those particles obeying to the adiabatic invariance of the magnetic moment. Gibson, Jordan and Lauer [28] proved experimentally that single charged particles remained trapped in a magnetic mirror according to the prediction of the adiabatic theory. The experiment confirmed that single particles are satisfactorily trapped in a magnetic mirror field but the trapping of a single particle in some field configuration does not mean that a plasma has a stable equilibrium. Not only the losses induced by the collisions, which can change the velocity vector direction putting it inside the loss cone, but also other instabilities may arise increasing the particles losses. Because of the magnetic confinement, the simple mirror plasmas have a non isotropic energy distribution function. This leads to some instabilities in the velocity space, the so-called *loss cone instability*. However, some hydrodynamical instabilities may be more dangerous. The most important instability involving simple mirror plasmas is the so called *flute instability*. For magnetically confined plasmas

³⁹See [27] for the formula which describes the small oscillations of the μ parameter. The oscillations at the Larmor frequency are due to the perturbation of the particle motion given by the weak variation of the transversal magnetic field during the gyromotion.

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it is equivalent to the well known Rayleigh-Taylor instability, which arises in neutral fluids and also in plasmas subjected to the gravitational force. Is the magnetic field lines curvature which triggers the instability. In figure 1.6.6 a schematic representation of the mechanism at the basis of the flute instability is reported.

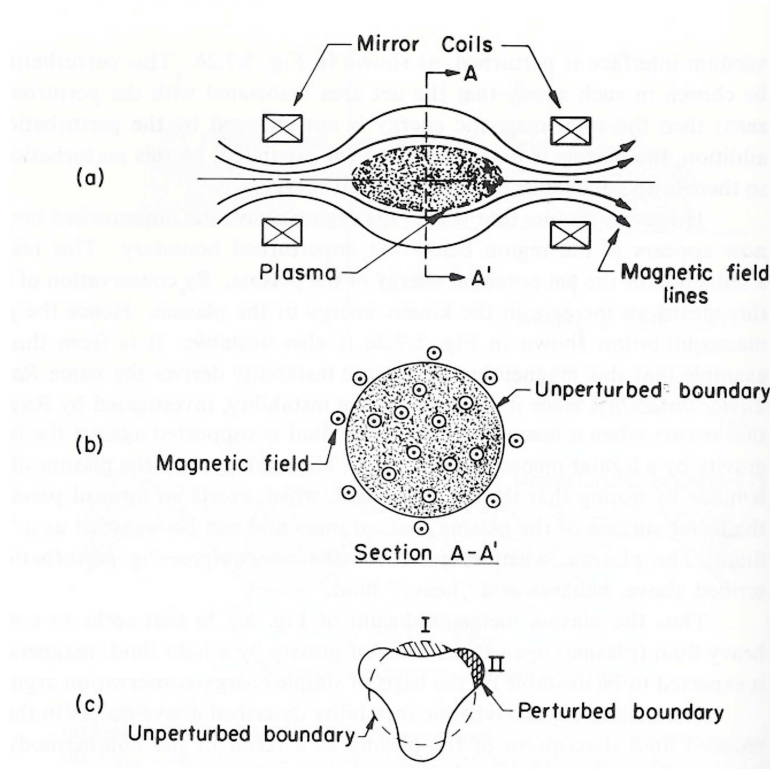


Figure 1.6.6: Description of the possible instability arising in a plasma confined by a magnetic field. (a) The plasma confined in the simple mirror configuration. (b) The midsection view of the plasma before the perturbation. (c) The same view of (b), but here plasma has been perturbed and the figure shown the first development of the "flutes". [9].

It can be seen that by perturbing the plasma boundary some flutes arise which change the energy balance of the plasma-field system. In particular the magnetic field energy remains unchanged (only some distortion of the field lines may change the magnetic energy), whereas the plasma energy increases. If the magnetic field lines have the curvature shown in figure, the perturbation will develop with time, producing a flux of plasma in the

1.6 Simple mirror and minimum- B configuration for the plasma trapping

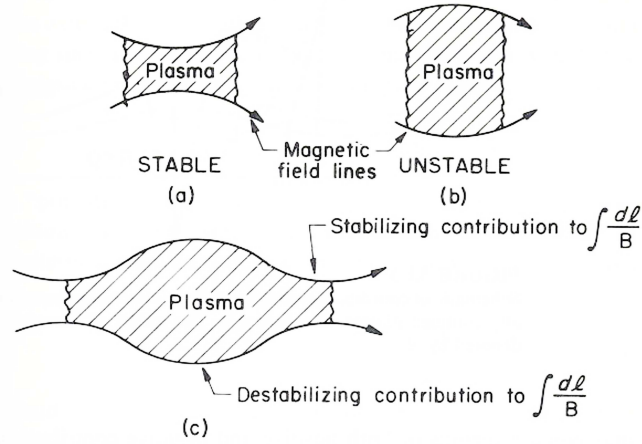


Figure 1.6.7: Schematic representation of the stable and unstable field line shapes in magnetically confined plasmas. [9].

perpendicular direction with respect to the mirror axis.

Figure 1.6.7 clearly shows that only if the magnetic field lines curve into the plasma, the confinement will be stable. This result gives us some practical information to how realize optimal plasma traps that are magnetically stable. The figure 1.6.8 shows the cusp field configuration: in this case the magnetic field lines curve into the plasma everywhere, and this ensures a stable confinement by the point of view of the flute instability. However the magnetic field lines are shaped in such a way that two large loss cones appear in the midplane, thus reducing the plasma lifetime (the plasma is lost because of particle scattering into four loss cones).

1.6.4 The B -min magnetic configuration

In 1962 Ioffe [29] first reported on plasma-confinement experiments in a magnetic field configuration that had the desirable feature that the magnetic field increased in every direction away from the plasma boundary, and that did not have the undesirable feature of a region where the magnetic field went to zero inside the plasma (as in the case of cusp configurations). Such a configuration can be obtained as a superposition of two different magnetic fields: one created by two solenoids (simple mirror), and the other created by six conductors surrounding the plasma chamber (an hexapole). In this way the magnetic field increases in every direction going from the plasma

1. FUNDAMENTALS OF ECRIS PLASMA PHYSICS

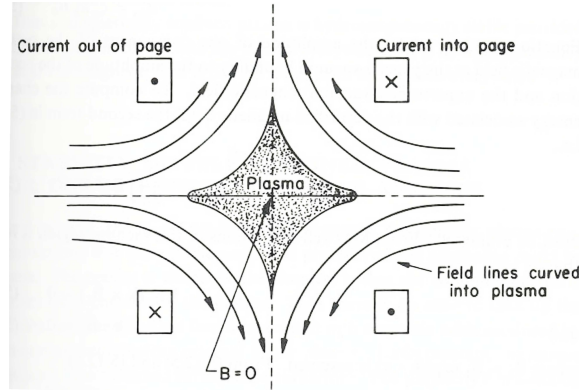


Figure 1.6.8: Cusp magnetic field configuration: the magnetic lines are curved into the plasma everywhere. [9].

center to the periphery. In fact the field produced by the hexapole increases along the plasma radius, and it is constant along the plasma axis. This stabilizes radially, being the simple mirror able to confine the plasma axially. This configuration can also be thought as a multi-mirror device: then the particle is reflected in many points traveling along the field lines. Ioffe and his co-workers at Kurchatov Institute of Moskow used a device known as the PR-5 (Probkotron-5) to demonstrate that the hydromagnetic instability could be suppressed with a minimum-B configuration. The figure 1.6.9 schematically shows the structure of such a field. More details about the structure of the minimum-B field are discussed in section 1.6.4. Note that the geometrical locus where B is constant are **egg-shaped surfaces**. As there are no axial symmetries, the μ adiabatic invariance cannot be invoked to study the particles confinement. Another adiabatic invariant can be used [10], [8]:

$$J = \oint v_{\parallel} dl \quad (1.6.16)$$

J essentially determines the field lines length between two reflection points. The J invariance implies that even after the reflection the particle will move along the same field line, or at least along a field line with the same J value. Field lines with the same J value define a surface over which particles with fixed values of $\frac{W}{\mu}$ will move.

Then the plasma trap obeys to the condition that the *constant* J surfaces do not intercept the plasma chamber walls.

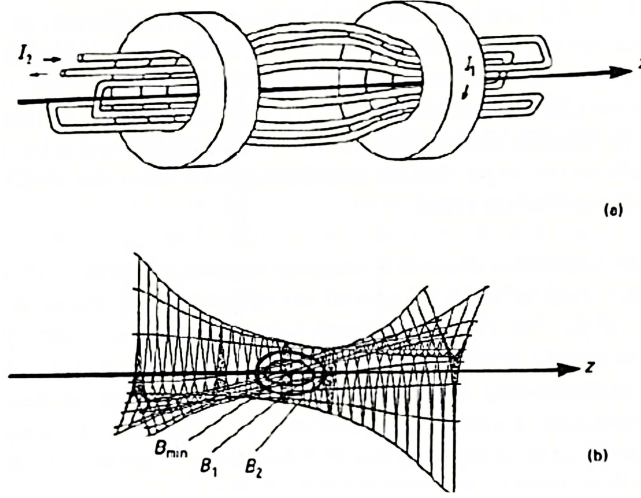


Figure 1.6.9: Magnetic system (a) and magnetic field structure obtained by the superposition of the field produced by two solenoids and an hexapole (minimum-B field) (b). [10].

1.7 Ion confinement in ECR plasmas

The present model to describe the ion confinement in magnetically confined plasmas is based on the creation of a negative potential well in the plasma core [10]. In the section 1.2 we demonstrate that for plasmas without turbulence and instabilities confined in a *B-min* configuration the electron losses are due only to the Spitzer scattering into the loss cone. For the Spitzer collisions we can write:

$$\nu_{ee} \simeq \frac{2\nu_{ei}}{z} \quad (1.7.1)$$

$$\tau_{ii} \neq \tau_{ei} \quad (1.7.2)$$

Then the scattering time is completely different for ions and electrons, and also the time needed to each species to leave the plasma. The scattered particles leave the plasma in a short time, thus determining the violation of the plasma quasi-neutrality. Without any additional mechanism, the plasma tends to become globally electrically charged. However the quasi-neutrality must be preserved, then an ambipolar diffusion process arises, as described in section 1.3.3. The slower and the faster species can be determined by analyzing the collisions times. The slower species determine the sign of the plasma potential, i.e. the potential which arises to accelerate

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the slower species and to accelerate the faster. Tonks and Langmuir formerly described the properties of plasmas located inside absorbing boundaries⁴⁰: the preservation of the quasi-neutrality means that $n_e \simeq Zn_i$, with Z mean ion charge state, and $j_e \simeq j_i$; then the electron and the ion densities must remain the same, along with the currents of particles leaving the plasma. Starting from this assumption it can be affirmed that in ECRIS plasma halo the electrons⁴¹, having $\frac{\tau_{ei}}{\tau_{ii}} < 1$, leave the plasma more easily than ions, thus creating a positive potential at the plasma boundary which accelerates ions and tends to retain the escaping electrons. On the contrary, into the plasma core, the electrons are less collisional, because their temperature is very high. Then they are well confined and generate a hot electrons cloud which retains the ions. The negative potential dip can be indicated as $\Delta\phi$:

$$\Delta\phi \propto -\frac{T_i}{ze} \quad (1.7.3)$$

Assuming that the bouncing time of ions inside the plasma trap is τ_b we obtain:

$$\tau_b \simeq \frac{B_{max}}{B_{min}} \left[\frac{l_p}{\langle v_i \rangle} \right] \quad (1.7.4)$$

$$\tau_i \simeq \tau_b e^{\frac{Z\Delta\phi}{KT_i}} \quad (1.7.5)$$

with $Z = ze$, l_p plasma length and $\langle v_i \rangle$ mean ion velocity. The final situation can be visualized in figure 1.7.1, where the complete shape of the plasma potential is shown.

According to this model electrons at high energies ($T_e > 10 \text{ keV}$) provide a stabilizing effect for the plasma because they practically confine the ions. In order to increase the ion lifetime, we have to increase the potential depth.

This model for the ion confinement has been put in crisis by some preliminary results obtained by our group in collaboration with the Kurchatov Institute of Moscow. In section 4.13.8 some preliminary results about a possible alternative scheme of ion confinement will be presented.

⁴⁰For absorbing boundaries we usually mean the walls of a metallic plasma chamber.

⁴¹Because of their higher mobility, due to the low plasma halo temperature and its high density, the electrons mobility is larger than ions in the plasma periphery, then they tend to leave the plasma with an higher rate than ions.

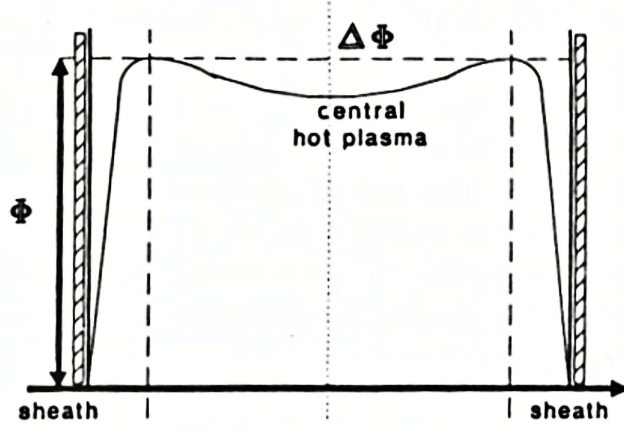


Figure 1.7.1: Plasma potential profile with the hollow shape providing the ion confinement as a consequence of the high energy electrons trapping. [10].

1.8 Particles confinement times in ECR plasmas

The first attempt to determine the particle confinement times is to take into account the breakdown of the adiabatic invariance of μ . However, in section 4.13 the deconfinement of the electrons due to such a reason has been investigated: it results also from experimental results [29] that charged particles inside a Simple mirror configuration are lost before any adiabatic invariance breakdown arises. The main reason which leads to particles losses are the long range coulomb collisions (the Spitzer collisions). Particles are injected into the loss cone in the velocity space, and it can be possible to derive a qualitative expression to determine the confinement time [10]:

$$\tau_{conf} \sim \tau_{90^\circ} \ln \left(\frac{B_{max}}{B_{min}} \right) \quad (1.8.1)$$

In this case τ_{90° indicates the scattering time of the charged particle to be analyzed. Anyway, as said in section 1.6.3, the MHD instabilities play the main role in case of simple mirrors, than the particles are lost mainly because of the flute instability. The RF diffusion also plays a fundamental role, as demonstrated in section 4.13 and in [30]. Reasonable estimations of the electron lifetime in plasma magnetically confined by Simple Mirrors give: $\tau_i \sim \tau_e \leq 10^{-5} \text{ sec}$. For the hot electrons ($E > 10 \text{ KeV}$), for which the cross section of Spitzer collisions is very low, we obtain that $\tau_e^{hot} > 10^{-3} \text{ sec}$.

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Ioffe demonstrated in his former experiment on B-min configurations that the confinement times are one hundred times higher than simple mirrors, as we have: $\tau \simeq 10^{-3}$ sec even for low energy electrons. If some conditions on MHD stability are satisfied, a confinement times up to 70 msec were possible. The main advantage in case of stabilized B-min configurations, i.e. for structures where the fluid instabilities have been suppressed and also any turbulence, is that the particles losses can be described only in terms of scattering inside the loss cones. This scattering is practically due to the Spitzer collisions.

The crucial issue of the plasma confinement is the ion trapping. Electrons are magnetized, then they are confined by means of the magnetic field structure. Ion are collisional in the plasma core (i.e. their collisional frequency, ν_i , is higher than the Larmor frequency), then they cannot be confined by the magnetic field. The ion confinement, at first approximation, occurs because of the ambipolar potential: electrons are strongly confined by the magnetic field; the ions are confined as a consequence of the electron confinement.

Whatever the ion confinement system be, the global plasma confinement depends strongly on the magnetic field shape.

Roughly we can write: $\tau_i \sim \frac{B_{max}}{B_{min}}$. Another way to investigate the plasma confinement is to use the diffusion approach: let us define *the particle diffusion time* t_D . This time takes into account many scattering phenomena between the plasma and the confining chamber. As introduced in section 1.3.3 t_D strictly depends on B and we can define it as the time needed to the particle to reach the plasma chamber walls.

Hence, strictly speaking, t_D represents a timescale for loss of confinement. We also have $t_D > \tau_{conf}$, with τ_{conf} given by the relation (1.8.1), as we have $t_D \simeq A\tau_{conf}$, with $A \sim 10$. The time for the energy transfer to the chamber walls is related to t_D :

$$T_{EN} \propto t_D \quad (1.8.2)$$

More generally, in order to evaluate the particle lifetime in a magnetized plasmas, we should take into account the different diffusion coefficients with respect to the magnetic field lines: the diffusion along the field lines is by far different with respect to the diffusion across the field. Then a detailed analysis requires the knowledge of the local structure of the magnetic field.

1.9 From Geller's scaling laws to the High-B mode concept

The ECR ion sources have their roots in plasma fusion developments in the late 1960s. The use of the plasma heating by means of the ECR with the aim to produce highly charged ions was suggested, for the first time, in 1970. The first operating ECR ion source able to produce multicharged ions (MI) was the MAFIOS source of the R. Geller group at CEA Grenoble, in 1972 [31]. In the same year the ECR-heating applied to ion sources was also used by K. Wiesemann in Germany [32]. However, room temperature solenoids were used, and their application demonstrated that simple mirror configurations were not able to guarantee long plasma lifetimes. Although the other source parameters were quite good (keV electrons, plasma densities on the order of $10^{12} cm^{-3}$ and pressures lower than $10^{-5} mbar$) their performances were negatively affected by the short ion lifetime, which limited the charge state distribution (CSD) to low charged ions (most abundant species were double charged ions).

It was quite soon understood that the key issue was the plasma lifetime. According to the calculations described in [30] and also reported in section 4.13 it can be demonstrated that once the plasma stationary state has been achieved, the self-generated plasma potential inject electrons inside the mirror loss cone. Both these effects become more clear by looking to figure 1.9.1.

According to what discussed in section 1.3.3 the plasma potential arises because of the ambipolar diffusion. Then it accelerates ions because the electrons are usually more mobile, and they escape with higher rates from the confinement region. Regardless to the shape of the plasma potential in the central region of the plasma⁴² near the plasma edge the potential is positive and accelerates the ions [30]. The region where this potential is applied coincides with the shape of the plasma. In simple mirrors, the constant density and temperature zones coincide with the constant B zones. These surfaces are shown in a two dimensional representation in figure 1.9.1. Let us assume that they are the edge surfaces of the plasma. The green arrows shown in figure indicate the direction of the ambipolar accelerating (for ions) electric fields. The ions are accelerated mainly towards the mirror

⁴²Remember the two conflicting model for ion confinement. The electrostatic model has been discussed in section 1.7, while the dynamical ion confinement will be discuss in section 4.13.8, as it is a completely new approach to this question and it constitutes a part of the original work carried out during the PhD period.

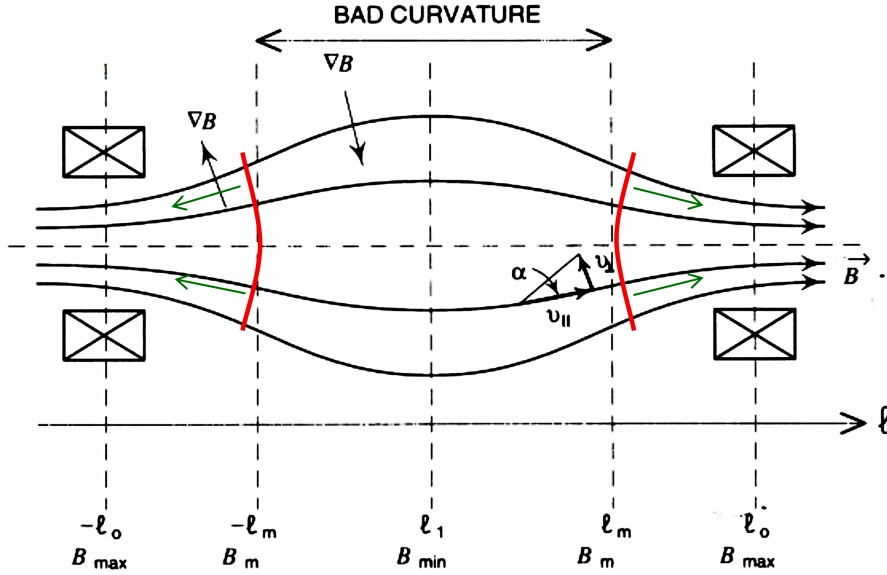


Figure 1.9.1: Field lines in a simple mirror configuration. The zone of the bad field lines curvature is pointed out together with the plasma emitting particle surface.

axis, and also the electrons will follow that direction. This acceleration inject both the particles inside the loss cone, thus leading to a rapid expulsion of plasma particles from the confinement region [33].

In addition to this mechanism, the bad central curvature of the magnetic field is detrimental for the plasma confinement. In figure the direction of the magnetic gradients is shown. Because of these gradients the plasma is pushed away from the confinement in the central region of the trap (the force is directed outward), and it is well confined laterally. This effect leads the plasma to escape transversally and it reduces considerably the particles lifetime.

The simple mirror is not the best structure to confine plasmas for the production of highly charged ions, because the relation (1.2.16) is not satisfied and the quality criteria (see section 2.1) are not completely fulfilled.

On the contrary, minimum-B configurations provide higher charge states. Figure 1.9.2 shows the shape of the plasma accelerating layer (as determined by means of the simulations reported in section 4.13.8) and of the magnetic field lines. In this case the lines curvature produces a stabilizing effect on

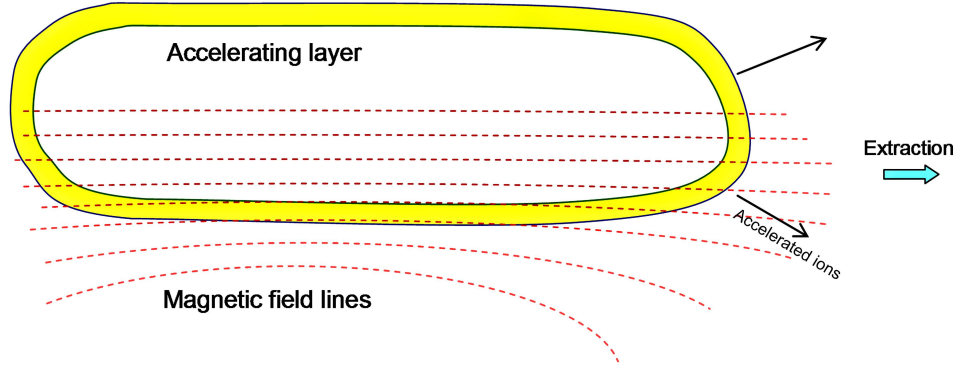


Figure 1.9.2: Field lines and plasma accelerating layer in a minimum B configuration. The black arrows indicate the ion and electron acceleration direction. A considerable amount of plasma is accelerated outside from the loss cone. This property of minimum-B configuration contributes to increase the plasma lifetime.

the plasma, and the accelerating ambipolar electric field is mainly directed outward with respect to the loss cone in proximity of the extraction [33]. This helps to remove particle from the loss cone and provides an additional confining effect which leads to longer ion lifetimes, then to the production of higher charge states. The strength of the magnetic field maxima determines the weight of this additional confinement.

1.9.1 Scaling Laws for the microwave frequency

In section 1.4 the different cutoffs and resonances of waves propagating in plasmas have been investigated. It has been shown that in an homogeneous and unmagnetized plasma the cutoff frequency assumes the simplest mathematical form, and it is the same of the $O - mode$ in case of magnetized plasmas. Although in the cylindrical cavities of ECRIS more than a single mode is present, so that all the possible cutoffs should be taken into account, the simplest picture to determine a relationship between the microwave frequency and the maximum achievable plasma density considers the $O - mode$ cutoff frequency, so that, writing it in terms of maximum density as a function of the field frequency, we obtain:

$$n_{cutoff} = \epsilon_0 m_e \frac{\omega^2}{e^2} \quad (1.9.1)$$

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More generally, all the possible electromagnetic modes propagating in plasmas (R, L, O, X) suffer a cutoff. Then the maximum density cannot overcome some fixed values. Only electrostatic modes are able to propagate in plasmas of whatever density (see section 1.5 on Bernstein modes).

The main consequence which comes from the equation 1.9.1 is that by increasing the electromagnetic field frequency heating the plasma the electron density can be increased. In magnetized plasmas the relationships between density and frequency may change, but the dependence of the extracted current on the square of microwave frequency has been demonstrated in many experiments. Such a dependence comes from the (1.9.1) by virtue of the 2.1.2.

On the basis of these considerations and of the experimental results, twenty years ago R. Geller proposed the scaling laws [2] which have been for long time a guideline for the ECRIS community:

$$q_{opt} \propto \log B^{\frac{3}{2}} \quad (1.9.2)$$

$$I^{q+} \propto f^2 M_i^{-1} \quad (1.9.3)$$

where q_{opt} is the optimum charge state, B is the peak field of the magnetic trap, f is the microwave frequency, I^{q+} is the intensity of the charge state q and M_i is the mass of ions.

These formulas were verified with several ECRIS and particularly with SC-ECRIS at Michigan State University and with SERSE at INFN-LNS [34]. In spite of the microwave frequency increase, in order to fulfill the (1.9.2), it was clear that the MINIMAFIOS source [35] produced at 16.6 GHz a mean charge state lower than the SC-ECRIS operating at MSU⁴³. The reason was found in the confinement, that for the SC-ECRIS was much better than for MINIMAFIOS, and permitted increase of the microwave frequency to be exploited.

In section 1.6.4 it has been pointed out that the minimum-B configuration is able to ensure the MHD stability, but no numbers were given about the relative values of the confining magnetic field, i.e. about the mirror ratio at the extraction and injection side, and also in the radial direction.

In 1990 the so-called **High-B Mode** concept was proposed [1], stating that *only the optimization, in terms of mirror ratios, of the confining magnetic field leads to the exploitation of the electron density as expected from the frequency scaling, thus leading to an effective increase of the sources per-*

⁴³The MINIMAFIOS source was able to extract CSD with $\langle q \rangle = 9$, whereas the SC-ECRIS produced $\langle q \rangle = 12$.

formances. Then such a principle does not conflict with the Geller laws, but it limits their applications to well confined plasmas.

The main advantage of the HBM concerns the possibility to give some numbers about the actual configuration of the magnetic field. These results have been obtained by virtue of MHD considerations. First, note that in section 1.6.3 we determined that a stable MHD stability can be reached only if the magnetic pressure is higher than the plasma pressure⁴⁴. By estimating the plasma pressure in ECRIS, and comparing it with the magnetic pressure (which depends on the maximum field), we can understand whether or not a given magnetic configuration is stable according to the MHD criteria for stability⁴⁵.

We report the expression for the β parameter:

$$\beta = \frac{n_e k T_e}{\frac{B^2}{2\mu_0}} \quad (1.9.4)$$

well confined plasmas are characterized by a very low β parameter:

$$0.005 < \beta < 0.01 \quad (1.9.5)$$

the plasma electron density can be roughly estimated by means of the (1.9.1):

$$n_e \simeq n_{cutoff} = \frac{m_e}{e^2} \epsilon_0 \omega_{RF}^2 = \frac{\epsilon_0}{m_e} B_{ECR}^2 \quad (1.9.6)$$

with:

$$\omega_{RF} = \frac{q B_{ECR}}{m} \quad (1.9.7)$$

Hence, by substituting the (1.9.6) into the (1.9.4), and considering the equation (1.9.5):

$$\frac{\epsilon_0}{m_e} B_{ECR}^2 K T_e < 0.01 \frac{B^2}{2\mu_0} \quad (1.9.8)$$

$$\left(\frac{B}{B_{ECR}} \right)^2 > 100 K T_e 2 \frac{\mu_0 \epsilon_0}{m_e} \quad (1.9.9)$$

we also have:

$$\mu_0 \epsilon_0 = \frac{1}{c^2} \quad ; \quad m_e c^2 = 511 \text{ KeV} \quad (1.9.10)$$

⁴⁴For plasma pressure we mean the usual kinetic pressure due to the particle temperature and density

⁴⁵Note that magnetically confined plasmas can be considered MHD-stable only if the β parameter is appreciably smaller than one.

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and then we obtain:

$$\left(\frac{B}{B_{ECR}}\right)^2 > \frac{200KT_e}{511 \text{ KeV}} \quad (1.9.11)$$

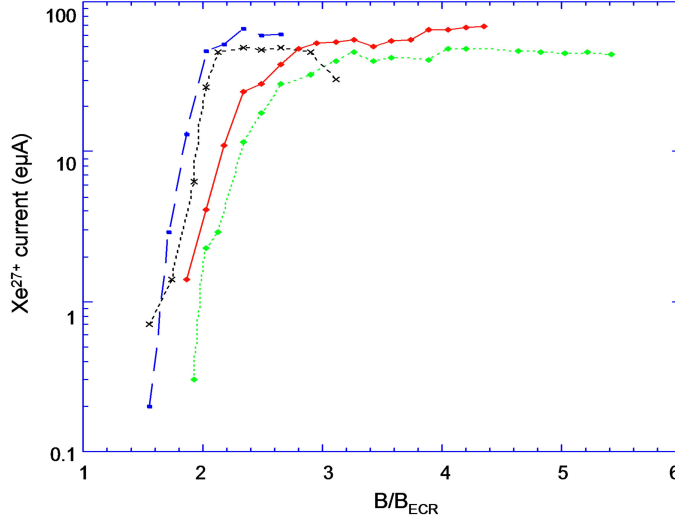


Figure 1.9.3: Experimental confirmation of the HBM concept. The extracted current increases as the ratio B/B_{ECR} increases, but it saturates around $B/B_{ECR} = 2$, in perfect agreement with the equation (1.9.12) in case of $T_e \simeq 10 \text{ keV}$.

For instance, if $KT_e = 10 \text{ KeV}$, then:

$$\frac{B}{B_{ECR}} > 2 \quad (1.9.12)$$

These results have been confirmed by a great amount of the experimental data. One of them is shown in figure 1.9.3, showing the trend of the extracted current from the SERSE source of the INFN-LNS for different values of the radial mirror ratio $\frac{B}{B_{ECR}}$. Note that the extracted current rapidly

increases, then tends to saturate when $\frac{B}{B_{ECR}} > 2$, confirming that the electron density reaches its maximum value only when the plasma stability condition is established.

More precise scaling laws can be determined by taking into account all the typical values of the magnetic field in the plasma trap: the injection, the extraction and the radial field, all of them compared with the field at ECR

and with that at the bottom of the configuration (usually named B_{min} or B_0).

Finally, note that the relation (1.9.12) strictly links the magnetic configuration to the used frequency (B_{ECR} is fixed by the microwave frequency). Hence, *to effectively improve the ECRIS performances one has to increase both the microwave frequency and the magnetic field, in particular the mirror ratios, of the plasma trap.*

Then the so-called "ECRIS Standard Model" [3], which has traced the road for the development of electron cyclotron resonance ion sources in the last twenty years, is based on experimental evidences and adds the Geller's scaling laws and the High-B mode concept. This model have permitted an average increase of about one order of magnitude per decade in the performance of ECR ion sources since the time of pioneering experiment of R. Geller at CEA Grenoble [2]. However nowadays such a trend is limited not only by the technological limits of the magnets, but also by the experimental evidence of hot electron population formation that becomes more important for high plasma density (i.e. for operational frequency of 28 GHz and higher) and for microwave power above 3 kW.

1.10 Modeling of plasma heated by microwaves

The study of the possible mechanisms that are at the basis of the ECRH started in the 50s, with regard to the thermonuclear fusion devices. Since that time, the models have considered the ECRH as an acceleration of individual noncorrelated electrons during the simultaneous actions of a magnetostatic field and microwaves in the case that the so-called ECR condition is satisfied, i.e., $\omega_{RF} = \frac{qB}{m}$. This model was formalized by Lieberman and Lichtenberg, who proposed a stochastic heating theory [36], [37]. Then in the next pages an overview of the main theoretical aspects of the electron heating will be given, even if more details will be presented in sections where the simulations results and the consideration about applications for ECRIS will be done.

Beside the single particle approach in the investigation of the electron heating in plasmas, the collective aspects have been also studied, formulating some alternative models of plasma heating driven by turbulence. Some general aspects of these models will be presented in the following pages.

1.11 The single particle theory of the electron heating: basic principles

A simplified way to describe the particle acceleration in presence of the ECR resonance is the Delcroix theory [38]. According to this model a sinusoidal electric field can be considered:

$$\vec{E} = E \cos \omega t \quad (1.11.1)$$

Two different conditions can be investigated, depending on the velocity along the magnetic field line ($V_{\parallel}$).

$V_{initial} = 0$: the equation for the total electron velocity is:

$$\vec{V} = \frac{q}{m} \vec{E}_+ t \quad (1.11.2)$$

and the Larmor radius:

$$\vec{r}_c = \frac{\vec{E}_+}{B} t \quad (1.11.3)$$

i.e. the velocity and the Larmor radius increase linearly with the time and *no phase correlations exist between the electron motion and the wave field*. E_+ is the electric field of the wave polarized in the same direction of the electron gyromotion.

$V_{initial} \neq 0$: the superposition between the Larmor motion given by $\vec{\omega}_c \times \vec{r}_c$ and the electric field rotation $\frac{q}{m} \vec{E}_+ t$ will depend on the relative phase. It can be demonstrated that:

$$V_{\perp}(t) = \frac{q}{m} \vec{E}_+ t_0 \sqrt{\left[1 + \frac{2t}{t_0} \cos \varphi + \frac{t^2}{t_0^2}\right]} \quad (1.11.4)$$

$$t_0 = \frac{m}{q} \frac{V_{\perp}(0)}{E_+} \quad (1.11.5)$$

The figure 1.11.1 shows the trend of the $V_{\perp}(t)$ versus the time. It is clear that for $\varphi = \pi$, the electron is firstly decelerated and then accelerated when $t > t_0$. As $t > t_0$ all the electrons are accelerated proportionally to the time and regardless of φ . After some time one has $V_{\perp} \gg V_{\parallel 0}$, and:

$$\vec{V}_{\perp} = \frac{q}{m} E_+ t = gc\omega t \quad (1.11.6)$$

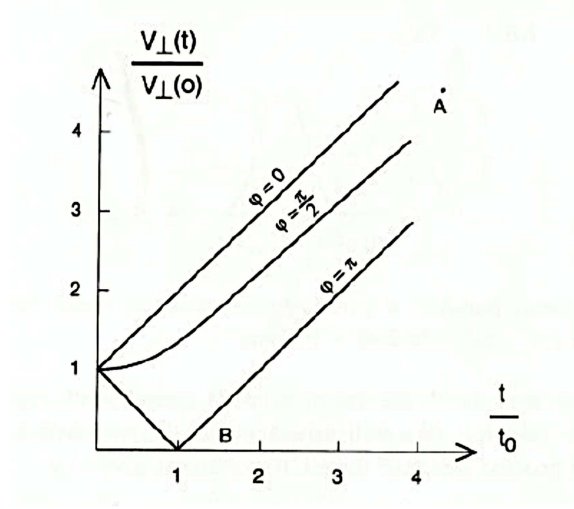


Figure 1.11.1: Trend of the electron velocity versus the time. At the beginning the energy gain depends on the relative phase; then, after a given time, all the electrons are accelerated.

for the energy along the perpendicular direction (the so called reduced energy) we obtain:

$$\frac{W_{\perp}}{m_0 c^2} = \frac{1}{2} g^2 \omega^2 t^2 \quad (1.11.7)$$

where no mass changes due to relativity are taken into account and g is the dimensionless quantity:

$$g = \frac{qE}{mc\omega} \quad (1.11.8)$$

Then the electron is accelerated as the electric field be continuous.

The Delcroix theory takes into account only the the electron acceleration, but it does not provide any result about the electron heating.

Considering the relativistic effects on the mass change due to the increase of the electron energy we have:

$$\omega_{RF} = \frac{eB}{\gamma m_0} = \omega_c \quad (1.11.9)$$

the relativistic change of the electron mass implies a variation of the wave-particle phase which leads to bad phase conditions when the mass increase

is appreciable. At that time the electron will be decelerated until the phase condition will be again favorable to the acceleration. If the magnetic field has no gradients, this process occurs indefinitely. The maximum attainable energy is: [10]:

$$W_{max} (KeV) = 18.8 \left[\frac{E}{B} \right]^{\frac{2}{3}} \propto mc^2 g^{\frac{2}{3}} \quad (1.11.10)$$

where E is expressed in $\left[\frac{kV}{cm} \right]$ and B in $[kGs]$.

Usually the time needed for the pulsation is longer than that really spent by the electrons in proximity of the ECR, as usually $\vec{\nabla} B \neq 0$.

1.11.1 ECR in presence of a magnetic gradient

The most powerful theory to explain the ECR in presence of a magnetic field gradient was proposed by Canobbio in 60s [39]. He considered collisionless plasmas and defined some dimensionless quantities that are useful to simplify the treatment. Else than g , defined in the previous section, the other two are:

$$\delta = \frac{1}{|B_0|} \left[\frac{d|B_0|}{dZ} \right]; \quad Z = \frac{\omega_c}{\omega} = \frac{eB}{m\omega} \quad (1.11.11)$$

In order to take into account the phase relationship between the resonant electron and the wave electric field, Canobbio proposed to expand the phase in series of powers:

$$\varphi(\tau) = Ag^2 \frac{\tau^3}{3!} + B\delta g^2 \frac{\tau^4}{4!} + C\delta^2 g^2 \frac{\tau^5}{5!} + D\delta^2 g^4 \frac{\tau^6}{6!} + \dots \quad (1.11.12)$$

with $A, B, C, D, \dots$ constants, and $\tau = \omega t$.

The main effects on the ECR dynamics due to the non uniformity of the magnetic field are the following:

- (a) The electron spends a finite time to cross the ECR region, and its final energy will depend on the magnetic gradient;
- (b) the phase relationship between the gyromotion and the electric field is not fixed for all the times; φ is fixed only for a limited number of gyro-orbits;
- (c) in case of very low field gradients, because of relativistic effect the electron energy $W_{\perp}$ may pulsate as reported above.

1.11 The single particle theory of the electron heating: basic principles

Generally, in the Canobbio theory the electron is considered to be at rest at the beginning of the resonance. This means that the time needed to go out of phase can be calculated by opportunely solving the equation of motion.

It is convenient to divide our treatment into two parts: one considers **gentle magnetic field gradients**, the other one considers **strong gradients**. First of all the knowledge of the B variation with z is mandatory⁴⁶.

The non uniformity of the B field implies that φ is not constant. The thickness of the resonance zone is calculated as that time needed to φ to become $\frac{\pi}{2}$. Z_R is the ECR slab thickness, and it is a function of g , τ , δ , N , where N is the plasma refraction index.

The gentle gradient regime can be mathematically defined according to the following condition:

$$\delta < 3.8g^{\frac{2}{3}} \quad (1.11.13)$$

In this case we consider only the first term of the expression (1.11.12), as δ is very small:

$$\varphi(\tau) \sim g^2 \frac{\tau^3}{3!} \quad (1.11.14)$$

indicating with $\psi(\tau)$:

$$\psi(\tau) = \int_0^\tau e^{i\varphi(\tau)} d\tau \quad (1.11.15)$$

we can calculate the maximum energy that an electron gains:

$$W_{\perp max} \sim \frac{1}{2}mc^2 g^2 \left| \int_0^\infty e^{i\varphi(\tau)} d\tau \right|^2 \quad (1.11.16)$$

i.e.:

$$W_{\perp max} \simeq me^2 g^{\frac{2}{3}} = 2mc^2 \left(\frac{eE}{m\omega_0 c} \right)^{\frac{2}{3}} \quad (1.11.17)$$

or:

$$W_{\perp max} (eV) = 1.5 \cdot 10^9 \left[\frac{E}{\omega_0} \right]^{\frac{2}{3}} \quad (1.11.18)$$

where E is expressed in $\left[\frac{V}{cm} \right]$. $W_{\perp max}$ does not depend on δ , i.e. on the magnetic gradient. If $\vec{\nabla}B$ tends to zero some initial stages of the energy pulsation may occur.

⁴⁶We consider the unidimensional problem where the magnetic field applies only along the z direction.

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For more clearness it is convenient to give some numerical examples. Suppose to work with an ECRIS operating at 18 GHz , with a magnetic field at the resonance of $B_r = 0.643 \text{ T}$. Suppose also that the magnetic field gradient, expressed in terms of the Canobbio's dimensionless parameters, is: $\delta = 10^{-2}$, as $\frac{\Delta B}{\Delta z} \simeq 0.7 \frac{T}{m}$. We will have: $\omega \simeq 1.13 \cdot 10^{11}$. First of all let us verify if the condition (1.11.13) be satisfied. Note that this condition depends not only on the B gradient, but also on the wave electric field amplitude at the resonance. In order to achieve a gentle gradient regime we need, in terms of the electric field:

$$E > \left[\frac{\delta}{3.8} \right]^{\frac{3}{2}} \frac{m\omega c}{e} \quad (1.11.19)$$

that for numbers above reported gives: $E > 260 \frac{V}{cm}$. This value appears quite low for modern ECRIS; for example assuming to have $E \sim 10^3 \text{ V cm}^{-1}$: $W_{\perp max} \sim 6.4 \cdot 10^3 \text{ eV}$.

The conclusions of the Canobbio theory in case of strong magnetic field gradients are completely different. Mathematically, the strong gradient regime can be expressed with the following relation:

$$\delta > 3.8g^{\frac{2}{3}} \quad (1.11.20)$$

In this case the dominant term in the expansion of $\varphi(\tau)$ is τ^5 :

$$\varphi(\tau) \simeq \frac{g^2 \delta^2 \tau^5}{5!} \quad (1.11.21)$$

φ becomes $\frac{\pi}{2}$ when:

$$\tau = \tau_R = \omega t_R \quad (1.11.22)$$

then the electron reaches the maximum energy when:

$$\varphi(\tau) = \frac{\pi}{2} \simeq \frac{g^2 \delta^2 \tau_R^5}{120} \quad (1.11.23)$$

and the ECR crossing time is given by:

$$\omega t_R = \tau_R = \left[\frac{\pi}{2} \frac{120}{g^2 \delta^2} \right]^{\frac{1}{5}} \quad (1.11.24)$$

Finally, considering the (1.11.7), with $m_0 = m$ to take into account the relativistic variation of the electron mass, we obtain:

$$W_{\perp} \sim 2.9mc^2 \left(\frac{g^3}{\delta^2} \right)^{\frac{2}{5}} \quad (1.11.25)$$

having replaced t in the (1.11.7) with the ECR crossing time t_R .

Differently than for the gentle gradient case, now $W_{\perp}$ depends not only on $\left[\frac{eE}{m\omega_0 c} \right]^{\frac{2}{5}}$, but also on the magnetic field gradient $\left(\frac{1}{\delta^{\frac{4}{5}}} \right)$. Then $W_{\perp}$ is smaller than the energy obtainable in case of gentle gradients.

In this case the electron motion is only weakly relativistic. It is easy to see that to τ_R corresponds a length Z_R which represents a sort of ECR thickness. Z_R can be calculated according to:

$$Z_R \simeq 1.4 \left[\frac{g^2}{\delta^3} \right]^{\frac{1}{5}} \simeq 1.4 \left(\frac{c}{\omega} \right) \left[\frac{\left(\frac{eE}{mc\omega} \right)^2}{\left(\frac{c}{\omega} \right) \left(\frac{\Delta B}{\Delta z} \right) \left(\frac{1}{B_{ECR}} \right)^3} \right]^{\frac{1}{5}} \quad (1.11.26)$$

usually Z_R is on the order of some mm for many ECRIS.

The maximum energies that can be obtained in case of strong gradient are about one third of the gentle gradient case. This difference is due to the mirror force $\mu \vec{\nabla} B$, as the higher is the B gradient the shorter is the time that each electron spends in proximity of the ECR zone.

1.11.2 ECR-Stochastic Heating

As ECR plasmas can be considered collisionless, then the particle heating cannot be collisional. Some other mechanisms must be invoked, and the stochastic heating mechanism of collisionless plasma may be a solution to the problem.

In principle, to obtain highly charged ions we do not need any global electron heating. Localized groups of energetic electrons may be able to provide high ionization rates and then good ECRIS performances.

In real situations, the plasma is confined by a magnetic trap; both in simple mirror or in multi mirror configurations the particles go forth and back along the magnetic field lines, as they are reflected at the two or more mirror points. Then the electron passes through the resonance zone many times. To some extent, the charged particle dynamics in presence of a magnetostatic field and of an electromagnetic wave can be considered adiabatic.

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However, as shown in section 1.6.2, abrupt variations of some physical variables may break the invariance by introducing some stochastic dynamics. The presence of the electron cyclotron resonance provides these non adiabatic effects, thus leading to an effective stochastic collisionless heating of plasma electrons. If the electrons remember their phase relatively to the wave after each ECR zone crossing, after few passages the energy increase should stop, as the phase correlation becomes bad for any further electron acceleration, and the all the electrons should have the same energy.

Actually, the build up of a EEDF typical of heated plasmas is usually observed experimentally, proving that non adiabatic effects play a remarkable role. The main consequence of the adiabatic motion breakdown is that any electron forgets its phase relationship with respect to the wave rotating electric field after each passage through the resonance. Hence the multiple resonance crossings are needed for the electron heating up to high energies.

In order to correctly describe the heating process we have to distinguish between the initial electron acceleration, due to the wave electric field, from the subsequent heating which occurs as the adiabaticity is violated. A simple picture can be proposed, where these complicated processes are qualitatively described [33].

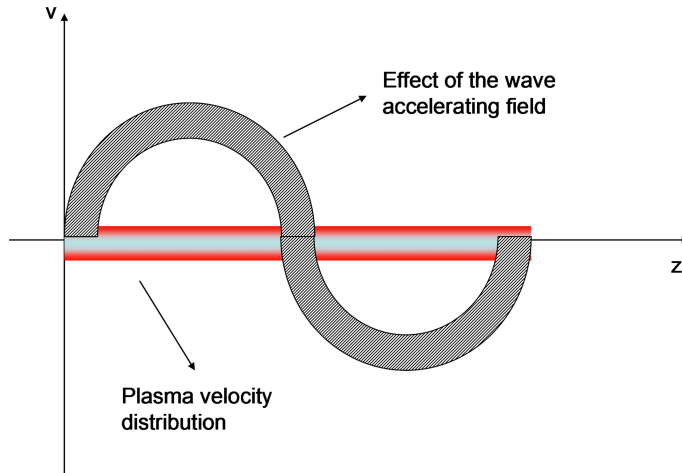


Figure 1.11.2: Simplified picture of a plasma wave interaction in the phase space. In a one dimensional model, the plasma extends along the z axis with a given distribution of velocities.

In figure 1.11.2 a simple picture of the plasma-wave interaction is shown.

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The figure represents the one dimensional plasma in the phase space: along the vertical axis the plasma velocity is reported, and the colors of the plasma slab represent the different plasma electrons energies; note also that the plasma, in this one dimensional model, extends along the z direction. The electrons far away from the z axis have a higher kinetic energy. Then the electron temperature is strictly connected with the vertical thickness of the slab. A wave carrying an electric field of a given amplitude interacts with the plasma. For sake of simplicity we suppose that the wavelength is comparable with the plasma dimension. The wave field accelerates the plasma particles according to the sign of the electric field. In a frame traveling with the wave there are no time variations, and the plasma organizes itself as it is shown in figure: where the electric field is stronger the electrons gain more energy, otherwise they remain in proximity of the z axis. At the beginning of the process no variations of the slab area occur: the interaction is completely adiabatic, then non conservative effects are not taken into account. The plasma retains its temperature, because the slab thickness remains unchanged, but nevertheless a great amount of electrons are strongly accelerated. Such a process may be enough to provide a great number of energetic electrons able to ionize atoms, even if no electron heating is present. However, collisions (in case of collisional plasmas), or stochastic effects lead to the breakdown of the adiabatic invariance and to the particle spread in the phase space.

The figure 1.11.3 shows that after some time the stochastic effects begin to play a role in the plasma-wave interaction dynamics. This leads to the enlargement of the plasma area in the phase space, being this a sign of an effective heating.

As the electron lifetimes in ECRIS are on the order of $msec$ the non-adiabatic effects arise in times shorter than $msec$. Reasonable estimations and numerical simulations give a time on the order of some $nsec$ to observe non these adiabatic effects (see section 4.13 for more details). Hence the collisionless plasma heating is completely determined by the stochasticity, as demonstrated also by Liebermann and Lichtenberg [37].

In order to better understand how the heating develops in multi-passage interactions, it is better to introduce the *separatrix* concept. Particles which resonantly interact with waves can be considered as trapped by the wave potential. The resonant condition is achieved if the well known equation: $\omega = kv_e$ is satisfied. Let us consider a frame moving with the wave velocity: particles moving with a similar velocity can oscillate into the wave

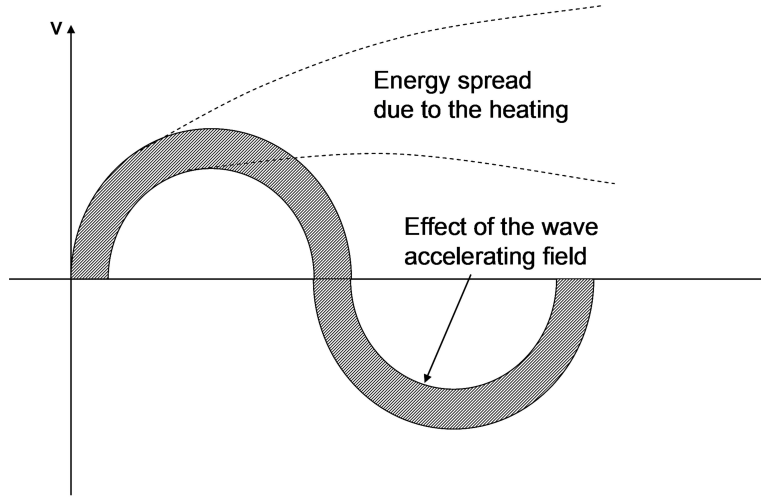


Figure 1.11.3: Spread of the electron velocity in the phase space due to the activation of the stochastic effects.

potential⁴⁷ $\varphi(z)$. Then the electron moves with a velocity v_e , while the wave potential, indicated with $\varphi(z)$:

$$\varphi = \varphi_0 \sin(\omega t - kz) \quad (1.11.27)$$

is stationary and sinusoidal, as the frame moves with v_φ . Because of the energy conservation we can write:

$$\frac{1}{2}mv_e^2 + e\varphi(z) = \epsilon = \text{const.} \quad (1.11.28)$$

then the velocity v of the electron is:

$$v_e = \pm \sqrt{\frac{2\epsilon - 2e\varphi(z)}{m}} \quad (1.11.29)$$

this velocity can be plotted in the phase space $z - v_e$.

Particles at the bottom of the wave potential (i.e. perfectly resonant with the wave) stay exactly at the center of each separatrix, and their trajectory in the phase space is just a point (their velocity with respect to the wave is zero, then they must lie on the z axis). If an electron is moved from

⁴⁷As before, we consider a one dimensional model to explain the particle trapping in the wave field at the resonance.

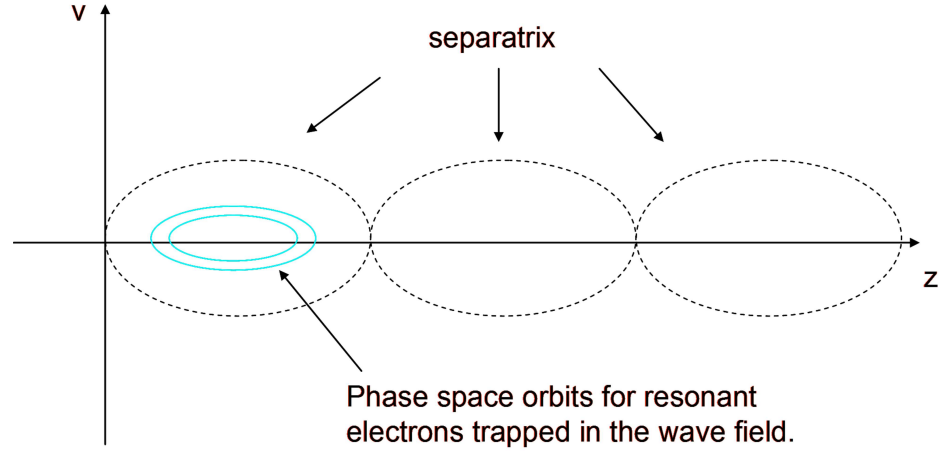


Figure 1.11.4: Definition of the separatrix concept on the phase space. The trajectories of electrons trapped by the wave in resonantly conditions are also shown.

the bottom of the potential, it will begin to oscillate around the equilibrium point, similarly to mechanical oscillators with small amplitude oscillations around the equilibrium. As long as the perturbation is small enough, the motion around the minimum will be harmonic. In the phase space, these motions are represented as ellipsoidal trajectories, as shown in figure 1.11.4. The electron velocity may be higher and higher, then the ellipsoidal trajectories become bigger and bigger. Actually, if the electron energy begins appreciable with respect to $\varphi(z)$, the motion is no longer harmonic and non linear effects arise. If the electron stays on the top of the $\varphi(z)$ potential, then the electron velocity will be zero according to the equation (1.11.29). In this case the electron has the same velocity of the wave, but as it stays on top of the wave potential, its equilibrium is not stable, and chaotic motions may arise. When this condition is satisfied the phase space separates in different zones. Inside each zone the particle is trapped by the wave. The figure 1.11.4 features the separatrix of a sinusoidal wave interacting with a single particle. This theory, named **pendulum model**, explains easily how the non linearity and the stochastic effects can play a fundamental role in the collisionless wave-particle interaction. That electrons whose velocity differ considerably from the wave phase velocity are not trapped and their velocity is slightly perturbed by the wave.

Let us discuss now the heating in ECRIS through electron multi passages

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across the ECR zones. An easier treatment is possible if we consider a Simple Mirror configuration, because in that case a one dimensional model can be applied and the bouncing of the electrons inside the trap are subjected only to the potential μB . Then the equation for the z variables becomes:

$$z = z_0 \cos(\omega_b t + \psi_0) \quad (1.11.30)$$

where ω_b is the bouncing frequency. The perpendicular component of the wave electric field can be simply written as:

$$E_{\perp}(t, z, r, \theta) = E_{\perp 0} \cos(\omega t - kz + \theta_{\perp}) \quad (1.11.31)$$

where $\theta_{\perp}$ is the initial perpendicular phase. By substituting the equation (1.11.30) into the (1.11.31) we obtain [33]:

$$\begin{aligned} E_{\perp}(t, z, r, \theta) &= E_{\perp 0} \cos(\omega t + \theta_{\perp} - kz_0 \cos(\omega_b t + \psi_0)) = \\ &= \sum_n A e^{i[(\omega - \omega_c)t + \theta_{\perp} + \theta_b]} \end{aligned} \quad (1.11.32)$$

where:

$$\theta_b = \omega_b t + \psi_0 \quad (1.11.33)$$

Then the equation (1.11.32) takes into account the cyclotron frequency, the frequency of the wave and the bouncing frequency. The Fourier expansion puts in evidence that the composition of the bounce and cyclotron frequencies leads to an effective multi-waves interaction: i.e. we may imagine the single particle interacting with a couple of waves at the resonance. Then each wave-particle interaction has a proper separatrix and many zones like that shown in figure 1.11.4 are present in the phase space⁴⁸. The figure 1.11.5 shows the structure of the several separatrix in the phase space in proximity of the ECR zone.

It is clear that all the separatrices are overlapped in $p_{\perp}$, being the distance between them proportional to ω_b . The electrons can be trapped into the lowest separatrix, as it intercepts the plasma slab. According to the electron kinetic energy and phase the motion will evolve into the several separatrices in a complicated manner, taking into account also non linearities in the wave-particle interactions. Because of the separatrices overlapping, the electron perpendicular momentum (then its energy) will increase more and more passing from a separatrix to the subsequent one, as schematically shown in figure. In the real situation the horizontal separatrix extension

⁴⁸Here we talk about the perpendicular phase space, i.e. that one formed by $p_{\perp} - \theta_{\perp}$.

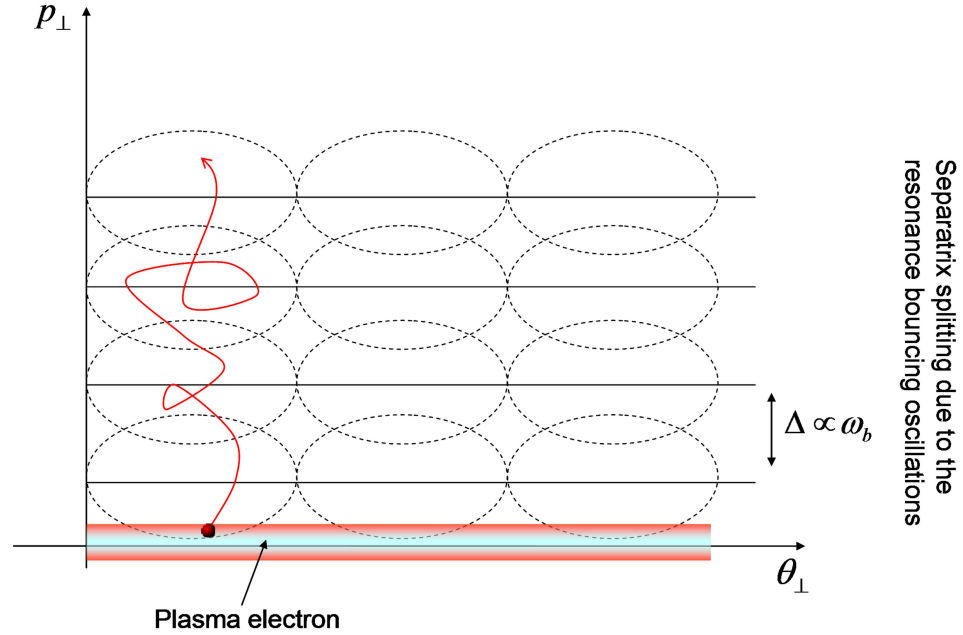


Figure 1.11.5: Separatrix overlapping occurring in proximity of the ECR zones.

at the resonance is limited, then the electron will escape after a fixed time from the interaction zone. But because of its bouncing motion, it will pass again through the resonance. As the former interaction was regulated by a stochastic motion in the phase space, it will arrive again to the ECR with a completely random phase, thus being able to gain more energy. This process goes on for many passages through the resonance. However, the higher is the electron energy, the higher is ω_b . When the bouncing frequency becomes comparable with the Larmor frequency, then the separatrix split, and the electron cannot pass through. This effect corresponds to the activation of some adiabatic invariants: the electron heating stops because the phase randomization process, strictly connected to the separatrix overlapping, does not work anymore.

A more quantitative treatment of the heating process starts from the Canobbio parameters. Liebermann and Lichtenberg formulated the theory to calculate the maximum value of the electron energy for particles bouncing inside the mirror trap. Being t_e the time effectively spent by the electron in

1. FUNDAMENTALS OF ECRIS PLASMA PHYSICS

the ECR region:

$$t_e = \frac{0.71}{\omega} \left(\frac{2\omega}{\delta v_{\perp 0}} \right)^{\frac{2}{3}} \quad (1.11.34)$$

where the perpendicular velocity $v_{\perp}$ is proportional to:

$$v_{\perp} \propto \frac{E}{\delta^{\frac{2}{3}}} \quad (1.11.35)$$

i.e.:

$$v_{\perp} \propto \frac{\sqrt{P}}{\delta^{\frac{2}{3}}} \quad (1.11.36)$$

by virtue of the (1.11.34), we have that t_e does not depend on $v_{\parallel R}$, and it is inversely proportional to $v_{\perp R}$, i.e. the higher is the transversal velocity, the less effective the particle acceleration will be.

For multiple passages through the ECR region, $v_{\perp}$ evolves according to:

$$v_{\perp}^2 = v_{\perp 0}^2 + v^2 + 2(vv_{\perp 0} \cos \theta) \quad (1.11.37)$$

where v is the maximum velocity gained for each passage, and θ a function of the phase. $vv_{\perp 0} \cos \theta$ is the stochastic term. After some passages, the separatrix splitting works and the heating stops. Defining with W_S the maximum value of the energy that can be obtained through the stochastic heating, we have [36], [37]:

$$W_S = 0.34eEL \sqrt{\left(\frac{\bar{t}}{\tau_S} \right)} \quad (1.11.38)$$

with:

$$\tau_S \sim 4\pi \frac{m}{e} \frac{1}{B} \frac{L^2}{l^2} \quad (1.11.39)$$

that is the time necessary to have a phase shift of 2π ; L , as usual, is the characteristic length of the mirror field.

$$\bar{t} = \sqrt{\left(\frac{mL}{eE} \right)} \left(1 + \frac{l^2}{L^2} \right)^{\frac{5}{4}} \quad (1.11.40)$$

in correspondence of the resonance position $x = l$. In this conditions the electron is mirrored just inside the ECR zone [36], [37].

The maximum energy achievable by the electrons will be:

$$W_b \sim 5W_S \quad (1.11.41)$$

with W_b that is usually named **absolute stochastic barrier**.

To give some estimation of W_b , we can replace with realistic values into the equations (1.11.38)÷(1.11.41), thus obtaining:

$$W_S \sim 2E^{\frac{3}{4}} \left[\frac{V}{m} \right] \propto P_{RF}^{\frac{3}{8}} \quad (1.11.42)$$

with E and P_{RF} wave's electric field and power. As for ECRIS we have:

$$0.1 \leq E \leq 10 \frac{kV}{cm} \quad (1.11.43)$$

the absolute stochastic barrier will lie in the range:

$$10 \text{ keV} \leq W_b \leq 270 \text{ keV} \quad (1.11.44)$$

The estimation of the maximum energy achievable by ECRIS plasma electrons is realistic if compared with experiments. This is a remarkable result, as the single particle theory is somehow oversimplified. However such calculations fail for modern ECRIS, like VENUS, which feature the presence of electrons with energies well above the stochastic barrier [40], [12], [41].

It can be noticed by the separatrix treatment that the heating may continue if the separatrix splitting does not occur, i.e. if additional perturbations to the particle motion change the separatrix amplitudes or provide additional separatrix, in order to have an overlapping. In this way the phase randomization should continue even above the stochastic barrier. The possible sources of further randomization of the electron-wave phase relationship will be largely discussed in section 4.13.5, thanks to numerical simulations.

1.12 Turbulent heating

As stated by Golovanivsky in [17], the physics of ECRIS may be looked by an entirely different point of view: the single particle approach is not sufficient to completely describe the ECRIS plasma dynamics. Many phenomena are involved in ECRIS plasma dynamics: some of them can be treated by means of the single particle theory, others can be described by the non-linear plasma waves theory, sometime turbulences and instabilities must be taken into account. In the most cases a mix of all these approaches is necessary. In order to understand the process that can be involved in the "ultra-hot" electrons creations, alternative schemes of plasma heating have been investigated:

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1. Stochastic heating of plasmas with nonlinear ECR.
2. Non linear mode conversion and HEL-Hot Electron Layers formation.
3. Turbulent heating.

The Liebermann and Lichtenberg theory predicts that the electron energy grows until the electron motion in phase space becomes periodic. Hence the kinetic energy of the electron will also be a periodic function of time and thus cannot increase indefinitely; the maximum value of the electron energy is the "absolute stochastic barrier". These results are valid as long as a quasi-linear electron-wave interaction can be considered. For very high power densities and for high electron densities one can assume that, instead of organized electromagnetic wave - electron interactions, the electrons encounters large-amplitude plasma waves with parametric decay. In such a medium, characterized by fluctuating fields, the electrons can randomly walk in phase space. A further randomization of electron-wave phase can occur and the electron energy can overcome the quasi-linear adiabatic barrier; this process is equivalent to a further overlapping of the separatrixes areas. No quantitative information are available about this mechanism, but it should be able to qualitatively explain a part of the experimental results featuring electron heating above the stochastic barrier. For the point 2 above we have to outline the mechanism which is at the basis of the HEL creation. It is important to underline that HEL creation has been observed so far only where the microwaves are injected into the plasma along the perpendicular direction with respect to the magnetic field [42]. This injection scheme encourages the X-wave creation inside the plasma. These waves travel inside the plasma till they approach the Upper Hybrid Resonance zone (see section 1.5). There the electromagnetic (EM) - electrostatic (ES) waves conversion can occur, with the following generation of two ES waves: UH-waves and B (Bernstein) waves. The nascent UH-waves are absorbed in a thin area near the UHR, then in that regions hot electrons can be created (with energies $\sim 1\text{ MeV}$), as experimentally observed). Being the X-wave reflected when it encounters the UHR layer, a short tunneling occurs, with a skin depth that varies from some mm to 1 cm (for lower frequencies, as 2.45 GHz). Inside the skin depth the longitudinal component of the electric field can excite longitudinal periodical oscillations of the space charge which are able to excite B-waves. The electrons should be heated by the dissipation of these waves inside the plasma, but it is difficult to determine how the dissipation occurs. In collisionless plasmas in fact, the Landau damping is usually the most probable mechanism of plasma heating. But Bernstein demonstrated that

the Landau damping doesn't occur in the transverse direction with respect to the magnetic field. This means that another collisionless heating mechanism has to be taken into account, as the one proposed by Segdeev and Shapiro [43]. However no quantitative estimations can be done by applying this model, due to the great amount of independent variables involved in its equations, and semi-empirical estimations are used. In [42] the measurements carried out on a 2.45 *GHz* ECRIS feature the presence of electrons of some tens of *keV*, or even 100 *keV* in HEL. It can be estimated that for higher frequencies, powers and magnetic fields, as in the case of VENUS, electrons with *MeV* energies are generated.

Finally, we have to investigate the third possible heating process, i.e. the "turbulent heating" [10]. It is based on the so-called LCC-Langmuir Caviton Collapse [17]. In ECRIS the microwaves are usually injected along the magnetic field. This means that the great amount of the electromagnetic energy is associated to the R and L waves, which propagate inside the plasma with different $\vec{k}$ vectors and opposite circular polarization. Only the R waves are able to get resonances within the magnetized plasma and it can suffer an EM-ES conversion like the X-wave. In this case an electrostatic wave is generated and it runs off from the the conversion area and transports the energy in the plasma core. The zone in which the mode conversion occurs depends on the ratio $\frac{n_e}{n_{cutoff}}$, and in ECRIS for multicharged ions, in which such a ratio may be roughly equal to 1, it takes place just a few millimeters before the ECR surface. Then a considerable fraction of the electromagnetic energy is converted into electrostatic waves, named Langmuir waves (L-waves). As the mode conversion takes place in proximity of the electromagnetic wave resonance, the energy density associated to the ES wave is very high and non-linear phenomena may occur. The non-linear dynamics of L-waves has been deeply studied and many quantitative calculations are available [17]. In physics of non-linear waves in plasmas the waves themselves are considered as a collection of quanta with energy $\hbar\omega$ and momentum $\hbar\vec{k}$. In the case of L-waves these quanta are named "plasmons". The major aspect of the L-plasmons gas in a plasma is that it is unstable so it suffers a series of non-linear transformations. In conditions of density usually found in ECRIS (i.e. $\frac{n_e}{n_{cutoff}} \simeq 1$) the plasmon gas undergoes the so-called parametric decay. This process fills the plasma core of a so-called "plasmon condensate", i.e. plasmons which are not absorbed by the plasma because of a different velocities of the electrons with respect to the plasmon phase velocity. After the creation of the "plasmon condensate" a second non-linear phenomenon occurs. The transformations that were involved in the parametric decay are

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due to instabilities of each individual plasmon. On the contrary, the instability of the plasmon condensate has a collective character: it essentially fragments the L-plasmon uniform gas so that the energy of the electrostatic oscillations is no longer uniformly distributed over the ECRIS plasma core, but it is rather concentrated in a plurality of localized bubbles. The plasma density inside the bubbles is lowered, and the initially homogeneous plasma density is thus spatially modulated by the bubbles (in fact such an instability is named "modulation instability"). Then the wave energy is trapped in a plurality of density cavities, the so-called "cavitons", which move inside the plasma core. The time needed to get this stage of L-waves evolution, named "cavitation", can be calculated and it is about 50 *nsec*. This time is shorter than the typical collision time for the great amount of ECRIS for multicharged ions, and this means that the non-linear evolution of the L-waves quickly occurs. Another picture of the cavitons can be given: they are small cavities with lower density, filled by plasmons. If some plasmons are freely travelling in the plasma core, they are immediately trapped in the cavitons when they encounter them. Some numerical values can be given about the cavitons properties and the energy that can be acquired by the electrons when the last step of the non linear L-wave evolution occurs, i.e. the Langmuir Caviton Collapse (LCC). The electric field generated inside the caviton can be calculated by means of the following equation:

$$E_{CAV} = \sqrt{\frac{2\delta n_e k T_e 1.6 \cdot 10^{-19}}{\epsilon_0}} \quad (1.12.1)$$

where kT_e is the electron energy in *eV*, δn_e is the fluctuation of the electron density inside the caviton, and ϵ_0 the dielectric constant. The density fluctuation can be calculated by means of the following equation:

$$\frac{\delta n_e}{n_e} = \frac{\bar{w}}{n_e k T_e 1.6 \cdot 10^{-19}} \quad (1.12.2)$$

where $\bar{w}$ is the energy density of the incident electromagnetic wave. Once that the ratio $\frac{n_e}{n_{cutoff}}$ is determined, the electric field inside the caviton can be calculated and then the energy density. This will be the energy distributed among the electrons after the LCC. In fact, due to the plasmon capture inside the cavitations, the plasmon accumulation makes the electric field higher and higher and the electron density lower and lower, until a limit value is reached. Then the caviton implodes and disappears, and its energy is transferred to the plasma electrons as kinetic energy. This is a crucial

point for the determination of the kinetic energy in average associated to each electron. The energy accumulation during the plasmon capture fill the entire volume of the caviton, whose diameter can be calculated as it follows:

$$\Delta = 5.5r_d \sqrt{\left(\frac{n_e k T_e 1.6 \cdot 10^{-19}}{\bar{w}} \right)} \quad (1.12.3)$$

During the caviton formation, its radius grows and in particular it becomes much larger than the Debye length. During the caviton collapse, the plasmons are compressed in a volume which becomes smaller and smaller, but no energy transfer occurs among the cavitons and the electrons. This means that the plasmons are still not absorbed by the plasma unless the caviton diameter becomes comparable to the Debye length, when they begin to be absorbed by means of the Landau damping. Hence the energy stored inside the caviton volume is distributed to the electrons inside a smaller volume comparable to the Debye sphere, or slightly larger.

Hence a simplified picture of the LCC can be guessed [41]:

- a) during the first part of the caviton collapse the nascent Landau damping transfers a part of the energy to a great amount of electrons (the ones contained in a volume of about two or three Debye spheres); these electrons get energies in the order of several tens or a few hundreds of keV;
- b) As the collapse goes on, the remaining stored energy is distributed among few electrons, contained in a smaller volume, comparable, or even smaller, than the Debye sphere.

Even a small fraction of the total energy, stored in the caviton before the collapse, allows the electrons to reach energy of hundreds of *keV* or *MeV*. This model may explain why in modern ECRIS very hot electrons are created, with energy well above that one predicted by the single particle theory. The LCC, anyway, requires that turbulences are active at any time inside the plasma. As stated by Geller [10] it can be thought that the turbulence is the final state of the plasma instabilities evolution. No turbulence occurs when instabilities are absent. The so-called "Russian school", which takes into account the quasi-particles (i.e. plasmons, cavitons, solitons, and so on), supposes that some turbulence is present in all plasmas, even for those usually considered quiescent as no noise is observed. In section 4.13 the main consequences of the processes discussed so far will be applied to the analysis of the VENUS results [41].

Chapter 2

Ion Sources for modern accelerators

2.1 Sources for the production of multiply charged ion beams

The construction of an accelerator facility requires many technological efforts and, first of all, a complete physical knowledge of those processes involved in the ion beam production, transport and acceleration. The first stage of an acceleration apparatus is constituted by the ion source. In general ions are generated by plasmas. Plasmas can be produced in several manners in laboratories, and they differ each other essentially for the electron temperature, electron density and plasma lifetime. Generally, plasmas can be created by means of electrical discharges under vacuum, by means of electron beams passing through the neutral gases or, in a more complicated manner, by means of electromagnetic waves interacting with gases or vapors, in presence of a well shaped magnetic field.

Ion sources based on this last method usually produce plasmas by means of the Electron Cyclotron Resonance and they are frequently used in modern accelerator facilities, as they allow to reach charge states and extracted currents relatively higher than other devices. In this chapter a brief description of other ion sources based on different principles will be also given; this will permit to compare the advantages and disadvantages of the different sources and to understand why the future development of ECRIS is of primary importance for the forthcoming accelerator facilities like SPIRAL2, FAIR, LHC, RIA, etc. ([44], [45], [46]).

First of all it is the case to summarize the main characteristics that are

required to satisfy the accelerator requests [47]:

- highly charged ions must be produced, especially for *cyclotrons*, for which the maximum accelerating energy is proportional to the square of the charge state.
- intense extracted current, especially for experiments involving low cross section reaction measurements; high currents help to considerably reduce the acquisition times.
- for specific applications high reliability, high efficiency and a good emittance are required.

The characterization of the different types of ion sources can be done by taking into account some fundamental **quality factors**. In particular, according to what discussed in section 1.2 about the step by step ionization process, the ion confinement time will be a crucial parameter to determine the maximum achievable charge state. In addition, the number of available electrons is important, because it is clearly related to the number of electrons able to ionize the atoms down to inner shells. Hence we have:

$$\langle q \rangle \propto n_e \tau_i \quad (2.1.1)$$

where $n_e \tau_i$ is usually defined as Q and it is the most important quality parameter.

Each source will have a specific capability to increase the electron density or the ion confinement time, then to maximize the Q parameter. The electron temperature also plays a fundamental role: in section 1.2 it has been demonstrated that a given charge state can be produced only with a proper T_e .

It is obvious that a too long confinement time limits the amount of ions extracted from the system per time unit, thus reducing the expected current. Hence well confined plasmas, if the electron density is not sufficiently high, may be not so good system to produce high currents. Mathematically we have:

$$I \propto \frac{n_e}{\tau_i} \quad (2.1.2)$$

By looking to the equation 2.1.2 we notice that the extracted current is inversely proportional to the ion confinement time (it is the contrary of the mean charge state), whereas it is proportional to the electron density (as in case of Q). This means that the role of n_e is crucial if one wants to maximize

2. ION SOURCES FOR MODERN ACCELERATORS

both the current and the charge states. These simple considerations on the main proportionalities of current and charge states are decisive to drive the research in the field of ion sources. Once known the processes which are able to affect density, temperature and plasma lifetime, then techniques and method to apply on ion sources can be developed.

The total background pressure completes the set of quality parameters. To minimize the charge exchange phenomena, according to the description of section 1.2.4), we need that the number of neutrals be sufficiently low, as stated by the equation (1.2.18). Finally, in order to perfectly match the ion sources with the accelerators the beam quality must be high, especially in terms of beam brightness.

Other additional characteristics of ion sources are important, as the reliability, the lifetime of the system, the maintenance and the last, but not the least, the cost (for building and for ordinary operations).

The principal ion sources able to produce multi-charged ions are:

1. **PIG**, i.e. *Penning Ion Sources*;
2. **EBIS**: i.e. *Electron Beam Ion Sources*;
3. **LIS**: i.e. *Laser Ion Sources*;
4. **ECRIS**: i.e. *Electron Cyclotron Ion Sources*.

On the basis of the above discussion the comparison of the data (both for currents and for charge states) plays in favor of ECRIS, which in fact are the most used in the laboratories for nuclear and particle physics spread over the world.

2.1.1 PIG sources

The PIG sources [48] are composed of two cathodes positioned at the extremity of an hollow anode. A gas is injected into the device and the ionization is provided by the electrons emitted from the cathode and traveling inside the anode. Because of their negative potential, the cathode repels the electrons, which oscillate many times inside the device, thus ionizing the atoms up to high charge states. In this way it is possible to obtain high ionization efficiency and then highly charged ions.

A strong magnetic field is used to further increase the electrons' path. The electron beam is focused by the field in a restricted region inside the device, thus leading to the formation of a negative potential well which is

2.1 Sources for the production of multiply charged ion beams

able to confine the ions. The improvement of the confinement immediately leads to the increase of the maximum obtainable charge state.

Very high electron densities ($n_e \leq 10^{13} \text{ cm}^{-3}$) and low operative background pressures are reached, so that multiply charged ions can be created efficiently. The ion confinement time can be estimated by means of [48]:

$$\tau_i = \frac{r^2 B}{T_e} \quad (2.1.3)$$

where r is the plasma radius, B is the magnetic field in Tesla and T_e the electron temperature expressed in eV. Considering the usual values of the magnetic field, electron temperature and plasma radius, the value of τ_i lies in the range $10 - 100 \mu\text{sec}$. The average charge states are for example 8^+ for Xenon, but for this element charge states up to 16^+ have been produced. The figure 2.1.1 features the scheme of a PIG source.

In the following some of the most important parameters characterizing the PIG sources are reported:

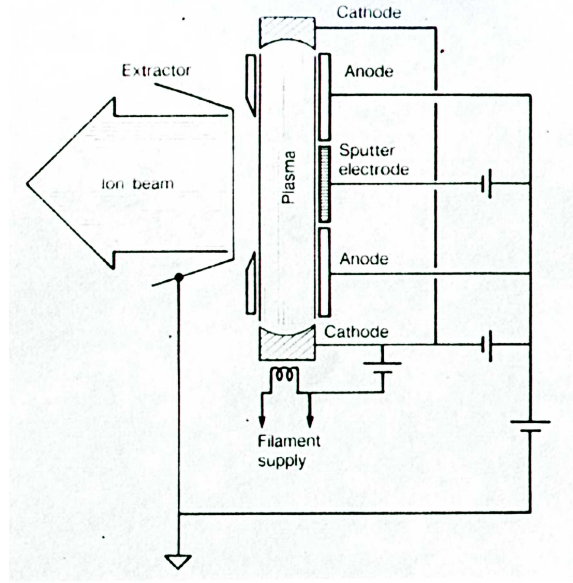


Figure 2.1.1: Scheme of a typical PIG source.

- The magnetic field lies in the range $0,1T - 1T$ and can be produced by means of room temperature magnets. Usually uniform magnetic field are used, but some attempts of mirror-like fields have been done.

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- The potential difference between the anode and the cathodes lies in the range $100 - 1000\text{ V}$.
- the extracted current is on the order of $\frac{10}{A} \div \frac{100}{A} \frac{mA}{cm^2}$, where A is the ion's mass; the current depends on the ion mass and on the anode's surface.

In the following some advantages and disadvantages of PIG sources are summarized:

- **ADVANTAGES:** intense current of multi-charged ions; quite easy methods to produce metallic ions; possibility to operate as an internal source for cyclotrons.
- **DISADVANTAGES:** beam instability; poor quality of the extracted beam; the beam quality varies during the source operations; short lifetime of the components, especially when the source is optimized for charge states as $12^+ \div 14^+$.

2.1.2 Electron Beam Ion Sources

The EBIS sources (Electron Beam Ion Source) [48] have some characteristics similar to the PIG ones: an electron beam is focused by means of a strong magnetic field in the central part of the the solenoid generating the field, thus creating a plasma with ions confined by the potential well due to the focused electron beam. A simplified picture of this mechanism is shown in figure 2.1.2. The atoms filling the plasma chamber are ionized by the electron beams. Then ions are confined by the potential well, and their confinement time increases as their charge state increases.

Cylindrical electrodes surrounding the device can be used to increase the potential well, thus increasing the ion confinement time, producing higher charge state. The solenoids for plasma confinement are shown in figure 2.1.2 along with the electron beam and the extraction system.

The EBIS are largely used as injectors for Synchrotron of completely stripped light ions (Dubna, Saclay, Stoccolma), or of highly charged heavy ions up to 54^+ .

For very high charge states the maxima currents are not higher than some nA (e.g. JINR - Dubna, Russia for X_e^{54+} currents do not exceed one nA). Some experimental tests carried out at the Los Alamos National Laboratories, USA, have demonstrated that the creation of totally stripped

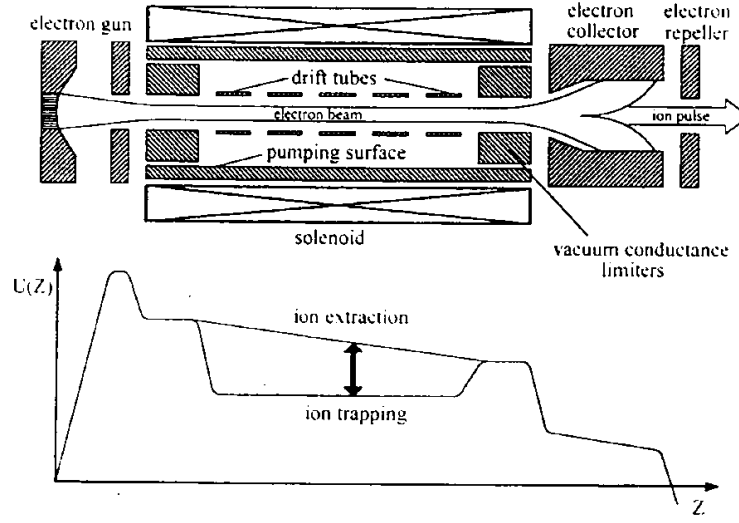


Figure 2.1.2: Scheme of a typical EBIS.

Uranium ions is possible, but the extracted current is on the order of few tens of ions per second, i.e. far from the requirements of accelerator facilities and their use is limited to atomic physics experiments. While in PIG sources the quality factor is enhanced by the electron density n_e , being τ_i modest, in the case of EBIS the density of the electrons and the ion confinement time are both high, thus permitting to reach high charge states. Conversely the current is very low, as it is proportional to n_e/τ_i .

Finally we can summarize the main advantage and disadvantage of the EBIS:

- **ADVANTAGES:** high charge states can be produced; the beam quality is quite good because of the low temperature of ions (the emittance is negatively affected by high temperature of ions); large possibility to vary the pulse width (μs beams).
- **DISADVANTAGES:** low currents, high realization costs, very complex device, not operator-friendly.

2.1.3 Laser Ion Sources (LIS)

The Laser Ion Sources [48] are based on the interaction of focused laser beams with solid targets. The figure 2.1.3 shows in particular the scheme of the LNS's laser source; the system is composed by different parts:

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- the laser beam generator;
- the optical system for laser focalization;
- the vacuum chamber where the target is located;
- several diagnostics tools especially for ion properties analysis (TOF - Time of Flight measurements, IEA - Ion Energy Analyzers, etc.);
- pumping system for vacuum.

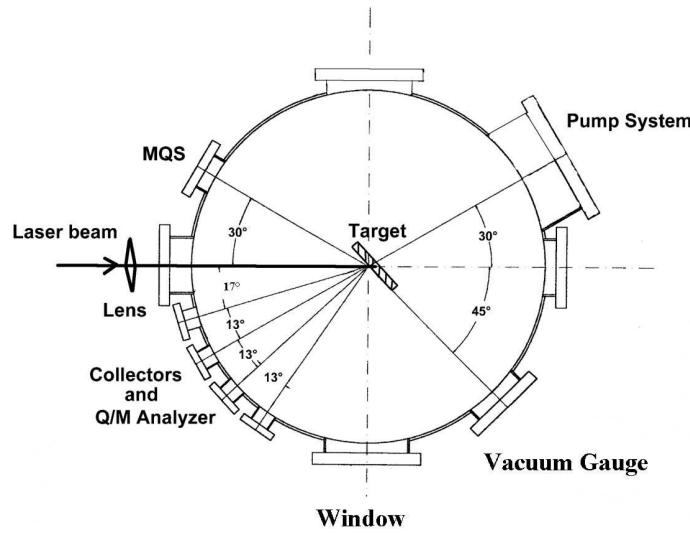


Figure 2.1.3: Schematic view of a LIS source.

If the power density of the laser beam is high enough, a rapid vaporization of the target material occurs, and a *plasma plume* is generated expanding at supersonic velocities (higher than 10^4 m/sec). The figure 2.1.4 shows the expanding plume during a recent experiment carried out at LNS [49] to characterize, by means of a Langmuir probe, the plasma plume in terms of temperature and density, for different times after the laser pulse. This figure helps to understand what are plasma dimensions: the probe was put 1.5 cm far from the Tantalum target, and the plume front reach the metallic tip (as shown in the picture) after one μs approximately. The transverse plasma dimension is comparable with the probe length, then it is about 1 cm .

The electron temperature in the plasma may range from tens to hundreds of eV as in the case of INFN-LNS laser.

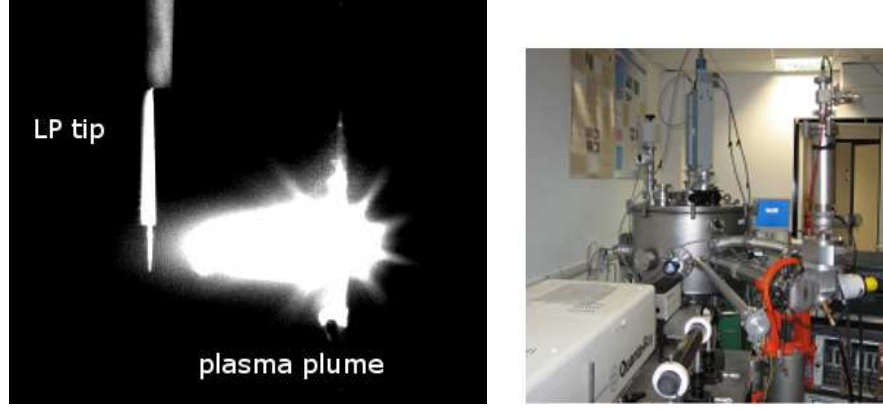


Figure 2.1.4: Plasma plume in a LIS source and the experimental apparatus for Langmuir probe measurements of laser generated plasma properties.

(a Q-switched Nd:YAG laser, operating at 1064 nm wavelength, with 9 ns pulse duration and 1 – 900 mJ energy), that permits to get charge states as high as 10+, to tens of keV in the case of the most powerful LIS, getting charge states above 50+. Usually the laser beam is focused through a convergent lens on a metal target (Ta, Ag, Cu, etc.) placed inside a vacuum chamber at 10^{-6} mbar. The laser beam incidence angle is $< 90^\circ$ with respect to the target surface and the spot diameter is 1 mm. The plasma plume is emitted along the normal to the irradiated target whatever the direction of laser beam is with respect to the normal direction.

Typically the laser energy used by LIS varies from 1 J to 50 J, with pulse lengths up to 1 μ sec and repetition rate up to 1 Hz for CO₂ lasers. The electron temperature T_e , and the charge states distribution depend on the laser properties as:

$$T_e \propto (P\bar{Z})^{\frac{2}{7}} \quad (2.1.4)$$

where P is the laser power.

The main limitation to the use of LIS in the accelerator facilities is linked to the low repetition rate. This is mainly due to the laser pumping, which requires a lot of time especially in case of high power - high intensity lasers. Anyway, as mentioned above, LIS are powerful methods to investigate

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Charge State	Pow. Dens.	Current.	Pulse time.
$1 \div 2$	$10^9 \frac{W}{cm^2}$	$10 \div 100 \text{ mA}$	$20 \text{ } \mu s$
$z \geq 10^+$	$\geq 10^{12} \frac{W}{cm^2}$	$0.5 \div 10 \text{ mA}$	$5 \div 10 \text{ } \mu s$

Table 2.1: *Characteristic performances of Laser Ion Sources.*

non equilibrium plasma physics, and they find applications in many field of fundamental and applied physics (recently applications to archeometry were developed).

The most important application of laser sources in the accelerator physics is connected to the possibility that the expanding plumes may work as accelerators including either the source and the accelerator. It has been demonstrated that a strong electric field can be generated inside the plasma plume, and these field can accelerate the ions up to MeV energies in case of laser at high intensity. The possibility to obtain small accelerators thanks to non equilibrium expanding plasmas is under investigation in many laboratories, and also at LNS some interesting results have been obtained with low intensity lasers, accelerating ions up to 10 keV [50]. If the plasmas produced by laser ablation are used as ion source the performances summarized in table 2.1 can be reached [51].

2.2 Overdense plasmas in an off-resonance microwave discharge

Since the middle of '90 a series of "Microwave Discharge based" Ion Sources (MDIS) have been designed and built at LNS to fulfill the request of intense beams of single charged ions for several applications: from the exotic beam production to the nuclear waste transmutation. The MIDAS ion source [52] was the first MDIS designed and built at LNS. It was a compact microwave discharge source which exploited the principle of the off-resonance discharge to generate an overdense low-temperature plasma, which is an effective tool to generate $1+$ ion beams with high efficiency, short delay time and low energy spread.

The prototype of the MIDAS source was tested in 1993-95 but the results were not acceptable with respect to the efficiency request, mainly because of the low RF power coupled to the plasma. Some improvements were done on such a source in 1997-1999. After, a completely new MDIS was designed and named TRIPS [53] (TRasco Intense Proton Source, see fig. 2.2.1), used

in the framework of the TRASCO (TRAsmutazione SCORie, i.e. nuclear waste transmutation). The aim of the TRASCO project was to verify the conceptual design of an accelerator driven system (ADS) for nuclear waste transmutation and to develop the related technologies, according to the Rubbia proposal. This proposal was based on a high current continuous wave proton accelerator driving a subcritical reactor to transmute nuclear wastes. The program consists of two main parts: the accelerator and the subcritical system. The injector of the proposed 30 *mA* proton linac consists of an 80 *kV* ECR source and an RFQ which accelerates the proton beam up to 5 *MeV*. TRIPS has been operational since the fall of 2000 at LNS [54] and in 2005 was installed at Laboratori Nazionali Legnaro (INFN-LNL).

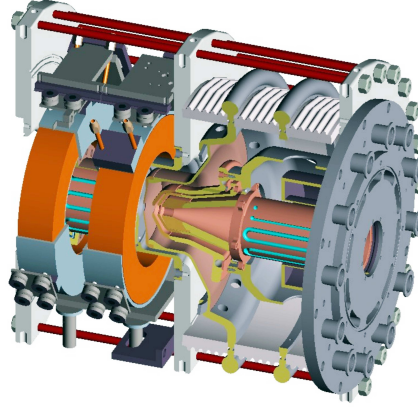


Figure 2.2.1: Design of the TRIPS source for the TRASCO project.

It was able to reach a proton current of 35 *mA* for an operating voltage of 80 *kV* with a rms normalized emittance lower than $0.2 \pi \text{ mm mrad}$, with availability close to 99.8%.

After MIDAS1 and 2 and TRIPS projects, two other devices have been built at INFN-LNS. The first one is the VIS source [55] (see figure 2.2.2), which is under installation; it is a permanent magnet version of TRIPS following the same principles. It will operate at similar plasma conditions, but with permanent magnet which ensure more compactness. In addition, some improvements in the extraction system have been made to operate up to 80 *kV* without a bulky high voltage platform, producing multi-*mA* beams of protons and H_2^+ .

The second one is a Plasma Reactor (see figure 2.2.3) [56] used for en-

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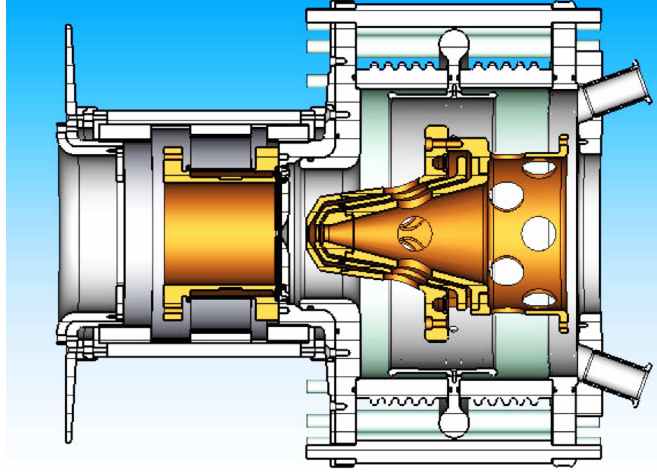


Figure 2.2.2: Design of the VIS source.

vironmental applications. Actually it is not an ion source, as ions are not extracted from the plasma chamber, but it has quite the same magnetic field magnitude and structure, microwave power and plasma chamber dimensions of the TRIPS and VIS ones. Only the operative pressure is considerably higher with respect to the TRIPS and VIS ones. The plasma reactor has been used as a testbench to provide information about the electromagnetic field propagation in MDIS plasmas.

MDIS are largely diffused also in several other laboratories and in industry, as they allow to produce high current intensities of lowly charged ions. However an exhaustive theory able to explain the plasma ignition for such sources has not been inferred so far. It is important to underline that in many of the MDIS operating at this time the microwave injection occurs along the longitudinal direction with respect to the plasma chamber axis. Different theories have been proposed to explain the plasma ignition in MDIS and they are largely discussed in [10].

Some sources operate at magnetic field below the ECR everywhere. In these cases the microwave absorption occurs at harmonic fields equal to $\frac{1}{2}$ or $\frac{1}{3}$ of B_{ECR} . The incident electromagnetic wave absorption can be described by means of a collective approach: the single particle motion is disturbed by fluctuating fields and by the plasma effervescence with high frequencies (the frequency of such oscillations should be comparable with the microwave frequency). Then the electron orbits around the magnetic

2.2 Overdense plasmas in an off-resonance microwave discharge

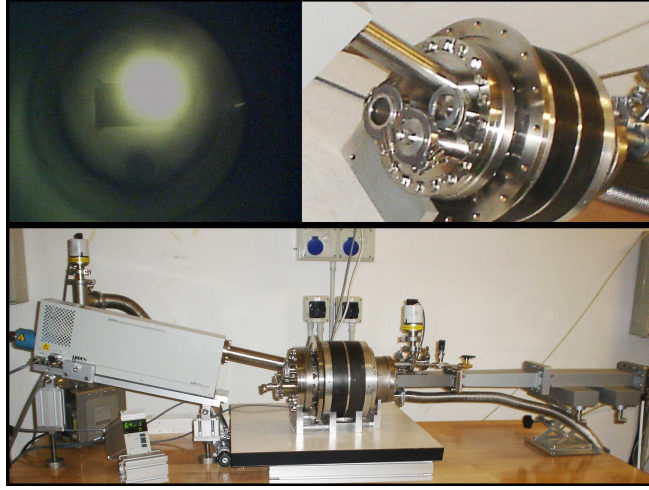


Figure 2.2.3: The plasma reactor installed at LNS and used for environmental application (toxic molecules dissociation). A panoramic view is shown, together with some details of one of the two flanges and an image of the plasma obtained through a CCD camera.

field lines are perturbed and harmonic waves at $Y\omega_{ce}$ (where ω_{ce} is the electron Larmor frequency) are generated. If the Larmor orbit present an Y deviation from the normal gyro-frequency, with $Y \sim \frac{\omega_p}{\omega_{ce}}$, **Bernstein waves (BW) are excited**. The BW are able to penetrate in the warm plasma core without any cutoff. Then, the most simple picture that one can give about this phenomenon is that a fraction of the incident electromagnetic energy is coupled to an internal wave and the conversion occurs where $B = B_{ECR}/Y$. This means that some values of B exist for which such conversion has the maximum efficiency and the plasma reach the maximum density. An experiment demonstrating the density optimization by varying B is reported in [see also Popov et al.]. It is important to underline that a plasma created in such a way is turbulent and any attempt to confine it may be meaningless.

Another theory applies to the cases of off resonance wave-plasma interaction. Some theories [57]-[58] affirm that resonances between the electromagnetic waves and the plasma electrons can occur in off-resonance magnetic fields and for high densities. However in this case a first ignition due to the single particle ECR heating should be necessary. This theory is also not so clear as it is hard to fit the experimental data with the theoretical ones: the theory would require to know exactly the direction of the incident

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electromagnetic wave, the angle between the k vector of this wave with the magnetic field for each position and so on. This theory looks practically inapplicable for the case of ion sources.

The third theory takes into account the off-resonance plasma effervescence. At the basis of the single particle theory of the ECR heating there is the assumption that the electron motion remains undisturbed (in the sense of collision's kicks) during the wave-particle interaction. However in the case of dense and turbulent plasmas a strong effervescence may arise, leading to a collective scattering frequency for the single electron that we name ν_{ω_p} . This frequency strongly depends on the plasma frequency ω_p . When $\nu_{\omega_p} > \omega_{ce}$. Then the electron gyro-motion is strongly perturbed by the collective collisions. Hence, during the ECR acceleration, the electron may go out of phase with respect to the wave because of such collective collisions. This means that the equivalent gyro-period (a period calculated by taking into account the frictions due to the collisions) is higher than in the collisionless case. In order to recover the synchronization with the rotating wave electric field, the electron needs to go faster and faster, then to gyrate around a magnetic field line whose field values is higher than B_{ECR} . Once the synchronization is recovered, the acceleration goes on. Then the off-resonance magnetic field allows the acceleration even for those electrons which go out of phase at the ECR zone.

The fourth theory is somehow similar to the previous one but it considers the electron-neutral collisions instead of the collective collisions. However it is valid only if the electron-neutrals collision frequency is comparable with the microwave frequency. This occurs when the neutral pressure is higher than 10^{-1} Torr, but in many cases this condition is not satisfied, as MDIS usually operate at lower pressure. Then also the collisional absorption theory is valid as long as we replace the electron-neutral collision frequency with the collective frequency ω_p .

From the several pictures proposed above it comes out that a complete comprehension of the physical phenomena involved in overdense plasma formation has not been achieved up to now. However it is important to underline that the collisional absorption (considering the plasma effervescence and the collective scatterings) play a role in the creation of such plasmas. It is reasonable to state that a lot of power is transferred to the plasma after the conversion of the electromagnetic wave in an internal wave (as proposed in the first theory reported above), whereas the remaining fraction is absorbed by means of a "collisional absorption" due to the plasma effervescence. Hence the possibility to built a MDIS at 28 GHz instead of the

usual 2.45 GHz should take into account both the mechanisms. The possible construction of a 28 GHz source based on principles of MDIS have been investigated during the last period of the PhD activities, in the frame of the collaboration with the Max Planck Institute of Greifswald (Germany). New ideas about future sources operating at high frequency and able to obtain high currents of lowly charged ions will be discussed later.

Note that $\nu_{\omega_p} \sim F(\omega_p/2\pi)$ [10], where F is a dimensionless parameter strictly related to the plasma effervescence. In the case of weakly over-dense plasmas $F \sim 10^{-1}$ or lower; this means that, up to densities of $10^{12} - 10^{13} \text{ cm}^{-3}$, ν_{ω_p} may be lower than ω_{RF} , thus reducing in a considerable manner the collisional absorption efficiency. The magnetic field trap of a future 28 GHz source has to fit a magnetic field around $\frac{1}{2}$ or $\frac{1}{3}$ of B_{ECR} thus permitting to optimize the internal wave conversion and the following plasma heating. In addition, a strongly non-homogeneous plasma may be created if such heating prevails, as for off-resonance magnetic field profiles the creation of over-dense plasmas often occurs in proximity of the microwave window. In order to avoid such a complication (which leads to a strong reduction of the ion source performances, reducing the ionization volume) it may be better to feed the plasma with microwaves injected along the perpendicular direction with respect to the plasma chamber axis. In this case another fraction of power may be provided through the microwave injection along the chamber axis, in order to provide a first ignition of the plasma by means of an ECR discharge. In these cases Golovanivsky et al. [42] have observed the creation of hot electron layers whose energies reach several hundreds of keV . These layers were located just in the regions in which the BW absorption may occurs. Then such a system may permit to study fundamental aspects of wave-plasma interaction, and it is of primary importance also for applications in X-Ray Sources (XRS). In fact, by introducing tungsten plates inside the plasma (in correspondence of the HELs) a strong production of keV X-rays has been observed.

2.3 ECRIS - Electron Cyclotron Resonance Ion Sources

The ECR ion sources are characterized by a complex technology, as the use of superconducting magnets and of high frequency microwave generators has become frequent. The SERSE source (Superconducting EcR ion SourceE), operating at INFN-LNS, since 1997 has been a breakthrough towards the highest performances. Many of the simulations and of the experimental data

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presented in this thesis are based on this source. In the following chapter also the CAPRICE [59] and the VENUS sources will be described, but the general aspects of all these sources can be summarized as in figure 2.3.1.

We know that in ECR ion sources the Electron Cyclotron Resonance ensures a resonant absorption of electromagnetic energy by the plasma electrons. At the beginning few free electrons, spiralizing around the magnetic field lines, are accelerated by the electric field of the injected microwaves. The electron-atom collisions allow to obtain multiply-charged ions, and very high charge states can be reached because of quite long confinement times ($\sim 10^{-3} \div 10^{-2}$ sec).

In the following pages the ECRIS characteristics will be investigated in details, anyway here we can summarize the advantages and disadvantages of such devices:

- **ADVANTAGES:** intense currents even for highly charged ions, long source lifetime, high stability, high reproducibility, capability to produce cw and pulsed beams;
- **DISADVANTAGES:** high power consumption especially in case of variable current magnets, expensive and complex technology both for magnets and microwave components, high ripple for highly charged ion beams (up to $\sim 10\%$), long conditioning times, some difficulties for the production of metallic beams.

2.3.1 Subsystems and technology

The different subsystems which compose the ion sources are shown in fig. 2.3.1. Hereinafter they will be described in detail. The SERSE design [60] is shown in figure as an example, putting in evidence also the cryostat for the superconducting magnets.

2.3.2 The magnetic confinement system

The magnetic system in the case of SERSE consists of three solenoids for the axial confinement (in the two external solenoids the current flows in the same direction, while in the middle solenoids it flows in the opposite direction). Then the dynamics of the plasma electrons along the plasma chamber axis can be described on the basis of the simple mirror trap (see section 1.6). In order to obtain the minimum-B configuration a sextupole encloses the plasma chamber. The figure 2.3.3 features the design of a typical magnetic structure of a recent ECRIS (the MS-ECRIS [61] at GSI,

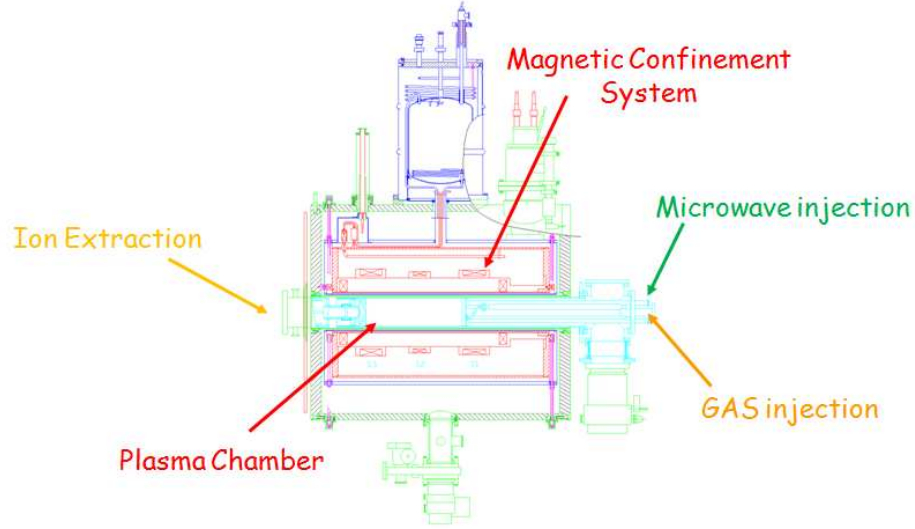


Figure 2.3.1: Main subsystems of the ECRIS.

Darmstadt). Three of the six coils of the hexapole are visible, together with the three solenoids for the axial trap. A part of the cryostat which envelopes the superconducting coils is also shown.

The technological problems related to the magnetic systems used for ECRIS are not simple to solve, because of the very high values of the magnetic field needed for the source operations. As stated above, in modern ECRIS superconducting magnets are required. The strength of the magnetic field at a fixed microwave frequency is:

$$B_{ECR} = 2\pi \frac{m}{q} f = 0.03577 f [GHz] \quad (2.3.1)$$

Hence if $f = 18 GHz \Rightarrow B_{ECR} \simeq 0.65 T$. Obviously this value of the magnetic field regards the resonance, but to ensure high stability and long confinement times the magnetic field must be considerably higher (see section 1.9), therefore superconducting magnets are required. In third generation ECRIS (like the VENUS source) magnetic fields up to $3.7 T$ are used. Superconducting magnets give also the possibility to operate with variable magnetic field which is highly appreciated by the ECRIS scientists in order to optimize the source performances. In general, any type of multipoles may be used for the radial confinement, and it is known that the higher is the

2. ION SOURCES FOR MODERN ACCELERATORS

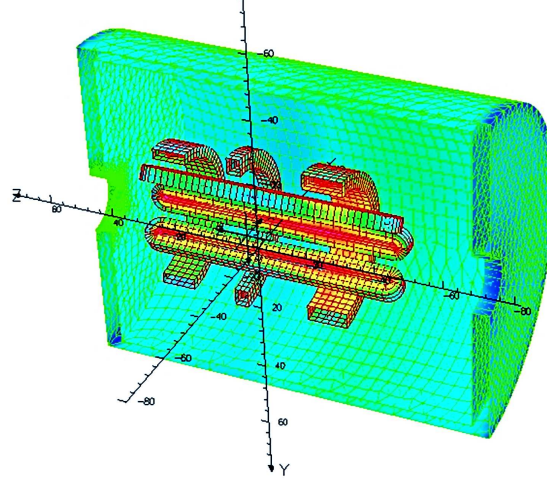


Figure 2.3.2: General design of the magnetic field structure of a modern ECRIS; the magnetic system is characterized by three solenoids and an hexapole; an additional solenoid is located outside the source and it is used for beam focusing (see below for more details).

multipole order, the larger the loss area of the magnetic structure will be. Larger areas mean higher ion currents and better beam uniformity, but on the other side this lowers the ion lifetime and then the mean ion charge state. In order to reach a compromise between the high extraction currents without any reduction of the ion lifetime the hexapole is usually employed. The loss area of the hexapole is shown in figure 2.3.4.

The figure shows the loss area shape for three different magnetic multipoles (octupole, hexapole and quadrupole) and the radial trend of the magnetic field for each magnet, featuring that the field grows from the center to the periphery. Let us define with B_m and r_m the magnetic field (on the multipole walls) and the radius of the multipole and r_C the plasma chamber radius; then the magnetic field on the plasma chamber walls is given by:

$$B_C = B_m \left(\frac{r_C}{r_m} \right)^{\frac{n}{2} - 1} \quad (2.3.2)$$

where n corresponds to the multipole order. Note that the higher is the multipole order, the lower will be the value of the magnetic field at the chamber

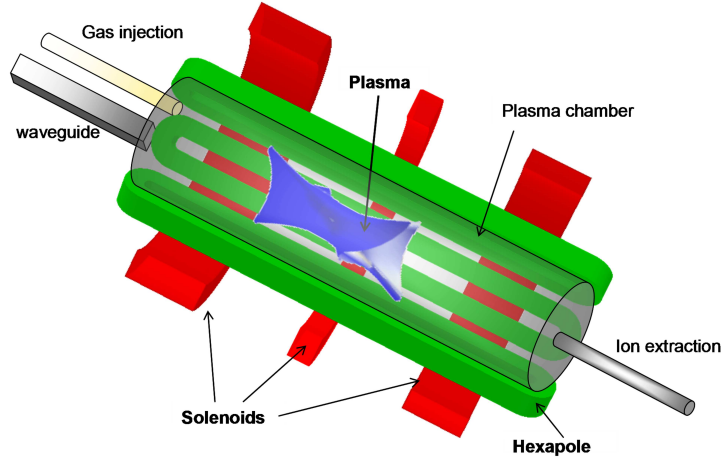


Figure 2.3.3: Location of the main ECRIS devices (magnets, microwave injection, gas injection, ion extraction, plasma chamber) together with the six-cusp plasma generated in the central region of the plasma chamber.

walls, that is detrimental for the confinement. Once again the hexapole is the best compromise between the maximum value of the magnetic field and the the loss area. The figure 2.3.5 shows one of coils of the MS-ECRIS source and the complete hexapole. Its maximum hexapolar field is equal to 2.7 T (the maximum available for ECRIS currently operating in the world). The SERSE hexapolar field is 1.55 T , VENUS and SECRAL [62] have a maximum hexapolar field of 2.1 and 1.9 T respectively.

2.3.3 The microwave injection system

The microwave injection system consists essentially of one or more waveguides (in cases when the sources operate with the two frequency heating) with the corresponding ports located in the injection flange of the plasma chamber. As an example, the figure 2.3.6 shows the injection flange with the location of the microwave ports for the SERSE source.

The insulation of the chamber is ensured by a DC break located in proximity of the injection flange. When two waveguide ports are present, they are perpendicular one each other to decouple the two channels (in the case of the SERSE source a -40 dB factor is obtained).

2. ION SOURCES FOR MODERN ACCELERATORS

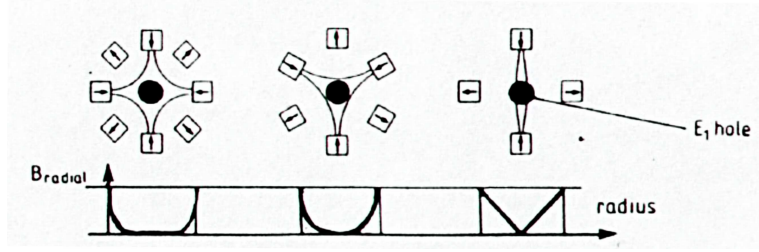


Figure 2.3.4: Loss area and trend of the radial magnetic field in case of multipolar magnets.

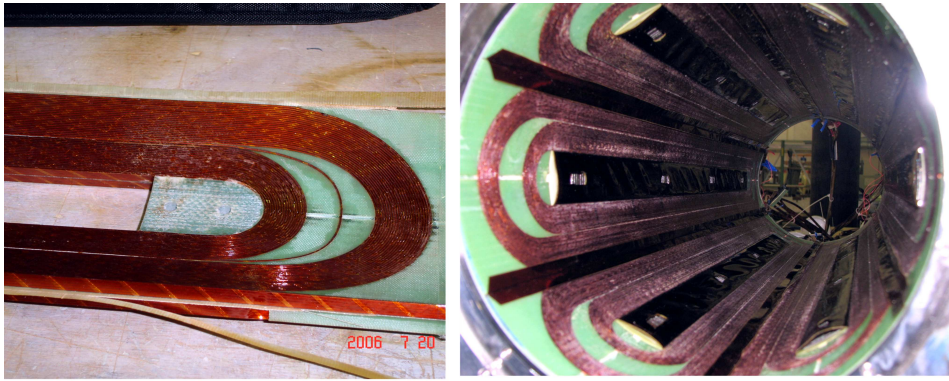


Figure 2.3.5: Superconducting coils and hexapole assembly for the MS-ECRIS source.

Different **microwave generators** are used: *klystron generators* and *TWT (Travelling Wave Tubes) amplifiers* up to 18 GHz . For 3^{rd} generation sources the gyrotron-based generators are used, with frequencies of 28 GHz or higher.

2.3.4 The plasma chamber

The chambers used for ECRIS plasma creation usually are cylindrical cavities with stainless steel or aluminum walls. As an example, the SERSE chamber is aluminum made with a radius of 6.5 cm and a length of 45 cm . The one of the MS-ECRIS source will be slightly bigger, with a chamber radius of 7.5 cm and 50 cm of length. In the chamber injection flange the *gas injection system* is located, together with the holes for vacuum pump-

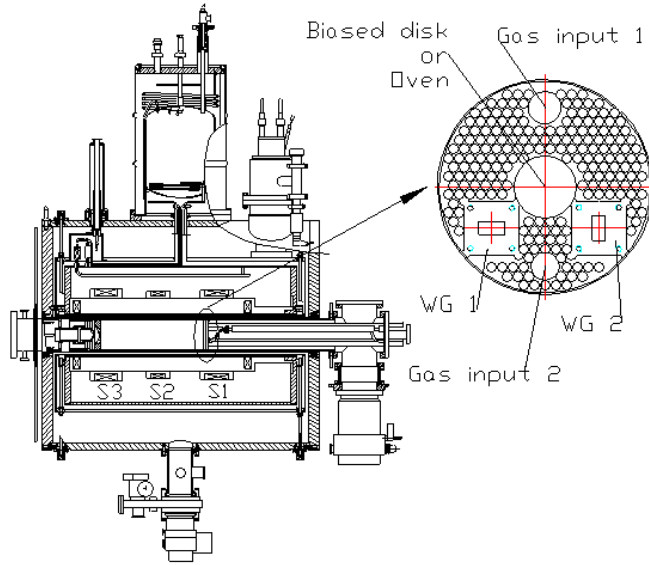


Figure 2.3.6: Detail of the injection flange of the SERSE source with the two microwave ports and a sketch of the microwave generators-plasma chamber connections via waveguides.

ing and the microwave ports. The extraction flange is characterized by the *extraction system*. The chamber has an hole used for the extraction of the ions. The chamber in fact is usually biased up to high voltages (e.g. 25 kV for SERSE) and a complicated series of electrodes are located outside the chamber, near the hole, in order to ensure the extraction of high currents beams with low emittance. Details about the extraction system of ECRIS can be found in [10].

2.3.5 Additional Devices

Ancillary equipment of LEBT (Low Energy Beam Transfer line) is not essential to produce the plasma but instead for the beam formation and transport. A *solenoid* or multiplet of quadrupoles are employed after the extraction system, in order to focus the extracted beam. *Faraday cups* are used to measure the beam current. A 90° magnetic dipole is used as charge over mass particle selector, and it permits to determine the Charge State Distribution (CSD) by means of the ramp of the magnetic field. The low operative pressure is obtained by means of rotative and turbo-molecular pumps. A pressure

2. ION SOURCES FOR MODERN ACCELERATORS

down to $2 \cdot 10^{-8}$ to $8 \cdot 10^{-8}$ mbar can be reached in the case of SERSE.

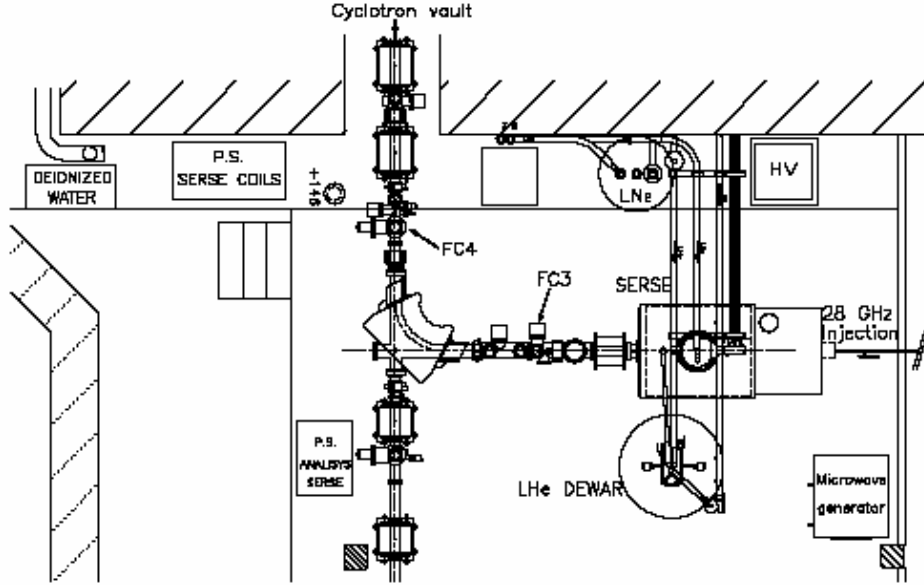


Figure 2.3.7: Design of the beam line of the SERSE source towards the superconducting cyclotron of INFN-LNS. The position of the two faraday cups is shown: FC3 is located before the analysis magnet, measuring the total extracted current; the FC4 is located after the 90° magnet in order to detect the current of each charge state. The position of the microwave generators is also shown.

The ECRIS are able to obtain also metallic ion beams when vapors of metallic elements are injected in the plasma. The *ovens* can operate with different principles: some of them are resistive, others are inductively heated and are able to reach a temperature higher than 2000 °C. The fig. 2.3.8 shows the inductively heated oven constructed for MS-ECRIS.

Some problems arise when metals with high fusion temperatures have to be used (like tungsten, tantalum, etc.). In these cases metal sample sputtering or laser ablation can be used. On the injection side, along with the oven, the gas input and the waveguides, the *bias disk* [63]-[64] is used to increase the electron density; it consists of a metallic and cylindrically shaped electrode biased at variable voltage (from few to hundreds volts); its position inside the cavity can be varied, but usually it stays outside the plasma core, i.e. that region of the plasma inside the resonance surface. This device permits to replace the electrons lost because of confinement leakages and to decrease the plasma potential which affects the ion confinement time.

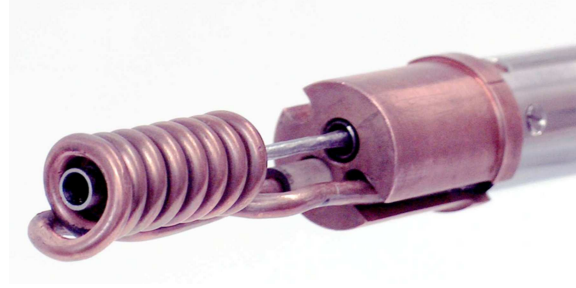


Figure 2.3.8: Inductively heated oven of the MS-ECRIS source.

Finally, in modern ECRIS with superconducting magnets a cryostat containing liquid helium is needed. The SERSE cryostat is 1310 *mm* long and it has 1000 *mm* of radius, the MS-ECRIS one is 1347 *mm* long, 1100 *mm* of radius. The figure 2.3.9 shows the design of the MS-ECRIS cryostat with the location of the plasma chamber. Liquid Helium is provided by an external liquifier (for SERSE) or by cryocoolers attached to the cryostat (for VENUS and MS-ECRIS).

Due to the complexity of the ECRIS systems and to the many parameters to control (other than for safety problems), a remote control panel is necessary, which allows to vary the microwave generators power (and frequency where possible), to change the magnetic field profile (in case of variable magnetic field), the amount of gas, the bias disk and the extraction system voltages.

The figure 2.3.7 shows the layout of the SERSE source at LNS, including the LEBT line and the faraday cups.

2.4 R&D activities devoted to plasma-based ion sources for INFN-LNS accelerators

The development of different types of ion sources has been one of the most valuable assets for the Laboratori Nazionali del Sud, strongly supporting the upcoming needs of nuclear physics research during last twenty years and of multidisciplinary physics in the most recent years.

The activity of the team started in **1983**, with standard negative ion sources for the Tandem, but it was only at the end of Eighties that the R&D activity started, aimed to the design of the injectors for the **K-800**

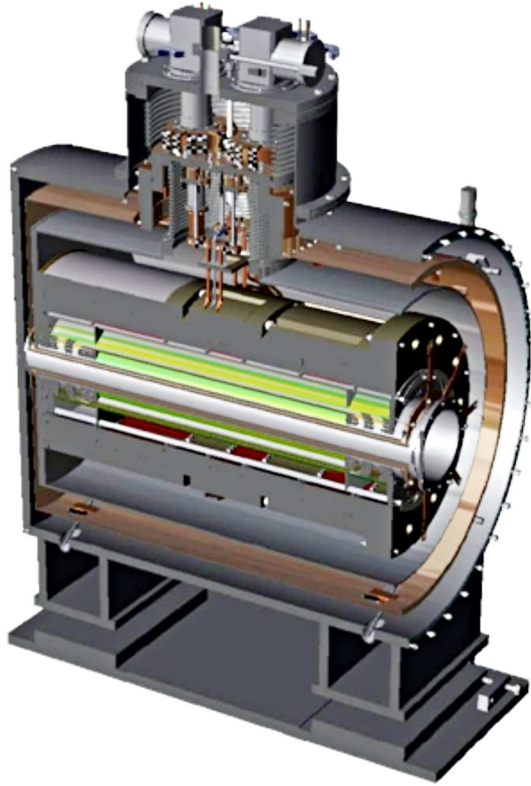


Figure 2.3.9: The design of the MS-ECRIS cryostat.

superconducting cyclotron.

In 1990 the proposition of the **High B mode (HBM)** was the trigger for the approval of the **superconducting SERSE source in 1991** [60]. The demonstration of HBM at MSU-NSCL came only in 1993 [34].

The construction of SERSE began in 1994; the source was completed in 1997 and *installed in June 1998* at LNS.

In 1997 the construction of a second conventional ECRIS, named **CAE-SAR** [65], was funded; the source was completed at the end of 1998 and installed in 1999 at LNS. The efficient ionization of atoms, for the **Radioactive Ion Beam Facility EXCYT** under construction, was studied with the microwave discharge ion source **MIDAS**, *designed in 1993* and built in the following years.

The intense proton source for the **TRASCO/ADS project** (named

TRIPS) was the first spin-off of the MIDAS activity. TRIPS was completed in 2003. The commissioning of the compact version VIS is under way. The **plasma reactor** was designed at the end of 2006, built and tested during the 2007. A development for higher frequency plasma reactor based on the old superconducting magnets of the SERSE source is in the conceptual phase.

Even if different R&D was developed in the course of last years, the major commitment remains the study and design of innovative ECR ion sources. In fact the *study of ECRIS operations at higher frequency than 18 GHz* was carried out in the frame of a RTD programme funded through the 5th Framework Programme of European Union and a following programme named **INES (Innovative ECRIS)** was mainly based on the study of intense beam production for highly charged heavy ions; a test for the *coupling of a 28 GHz – 10 kW gyrotron-based generator to the SERSE source* was carried out and the study of the scaling laws was done at the different frequencies (14, 18, 28 GHz) permitted by this superconducting ECR source, at that time the unique tool able to perform such experiment and still now the only one in Europe [66]. A beam current close to the one needed for the next generation of accelerators was obtained (above 0.5 mA for charge states between 20 and 25+ for Xenon), and the data essential for the design of 3rd generation ECRIS were collected, thus completing the preparation of the **GyroSERSE project**, conceived in 1994 and proposed in 1999.

In 2003 the **ISIBHI collaboration** was established within the EURONS proposal for the 6th Framework Programme and the GyroSERSE project was adopted for the design of the prototype **MS-ECRIS** [61]. The construction was funded for the period **2005-2008** with the contribution of INFN, GSI, GANIL and other European Laboratories. The source commissioning is under way at GSI, Darmstadt.

The activities concerning the ECR ion sources have been also applied to *medical applications* with the construction of the ion sources for the **CNAO accelerators (Italian Center for Oncological Hadrotherapy in Pavia)** under the responsibility of LNS Ion Sources R& D team.

Other plasma-based ion sources have been developed during the last ten years, as the hybrid ion source (ECRIS + LIS) and different types of Laser Ion Sources [50], [51] but their description is not essential for the argument described in this PhD thesis and it will be skipped for sake of brevity.

Chapter 3

Experimental equipments

A general overview of a typical ECRIS system has been given in the previous chapter. Several ECR sources and other experimental equipments have been used during the PhD course, then in the next pages they will be described in detail. They have permitted to obtain the results discussed in the following chapters, whose analysis has been of primary important for the ideas and the conclusions reported in this thesis.

3.1 Standing waves measurements in a plasma reactor

The plasma reactor was designed at INFN-LNS to dissociate complex molecules (in particular environmental applications).

It practically works like a MDIS, with a plasma that can be considered "weakly ionized" and strongly collisional because of low electron temperatures ($T_e \sim 1 \div 10 \text{ eV}$) and high pressures ($0.1 \div 0.5 \text{ mbar}$), as discussed in section 2.2. The plasma chamber is a stainless steel cylinder 276 mm long, with a radius of 67.8 mm . A *NdFeB* permanent magnet generates an off-resonance magnetic field along the plasma chamber axis (with a maximum of 0.1 T on axis). The high electron-neutral collision frequency determines a high rate of molecules fragmentation, thus increasing the dissociation fractions. For safety reasons, the reactor efficiency has been tested by using non dangerous gas, in particular n-hexane (C_6H_{14}) and cyclehexane (C_6H_{12}) [56] but the final aim of the experiment is the dissociation of dangerous complex molecules (e.g. dioxine). In the injection side of the reactor the microwave injection aperture, the flanges for the pumping and for the gas injection are located. In the opposite side $DN 40$ and $DN 25$ flanges are used to con-

3.1 Standing waves measurements in a plasma reactor

nect the mass spectrometer (if necessary), the optical window (to be used for plasma observation and for optical spectroscopy), the microwave probe and the Langmuir probe. The microwaves are generated by a Magnetron operating at 2.45GHz with a 300W of maximum power and they are injected into the plasma chamber through a WR284 (72.136mm width and 34.036mm height) rectangular waveguide placed on the cavity axis that operates in the TE_{10} dominant mode. A rotative pump is used for vacuum (pressure down to 10^{-3} mbar are usually obtained). Several measurements have been carried out in order to determine the optimum operating conditions in terms of microwave power and gas pressure [56] and in 2005 n-hexane and cyclo-hexane have been successful dissociated [56]. Now this set-up has been used as a testbench to investigate the standing wave formation inside the resonant cavity even in presence of plasma. Preliminary studies about the coupling between microwave generators and ECRIS have been carried out at INFN-LNS for the SERSE superconducting ion source. Several measurements were performed by means of a vectorial network analyzer (NA) able to operate up to 50GHz [67], in order to characterize the system including the plasma chamber (in vacuum) and the microwaves line in terms of S matrix. The same NA has been used for the plasma reactor and the S matrix has been characterized with and without plasma. In our experiment we fixed the microwave feeding the plasma at 2.45GHz and analysed the plasma properties in the $1 - 3\text{GHz}$ frequency range.

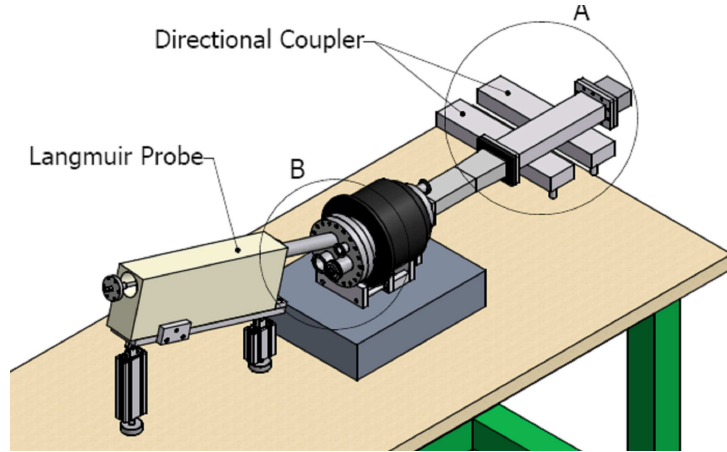


Figure 3.1.1: Schematic view of the plasma reactor and of the microwave injection system.

3. EXPERIMENTAL EQUIPMENTS

A general view of the plasma reactor is shown in the previous chapter, while a schematic view of the experimental set-up is shown in fig. 3.1.1. This figure shows in particular the directional coupler, which allows to determine the forward and reflected power for different plasma conditions. The results of these measurements are reported in [68] and they feature a complicated mechanism of wave plasma interaction which strongly depends on the background pressure and on the RF power. A coaxial connector is also shown in figure 3.1.1; it permitted to determine the S_{11} parameter in presence of the plasma (a series of attenuator were used to reduce the output power) and to investigate the best condition of wave-plasma coupling thanks to a microwave probe that can be connected to it. For S_{11} parameter characterization the plasma was ignited through the waveguide and the NA was connected to the coaxial connector. By analyzing the S_{ij} parameters it is possible to characterize the modes that exist inside the plasma chamber. In fact the frequencies for which the minimum values of the S_{11} occur represent the excited modes. Measurements of scattering parameters have been carried out in parallel with Langmuir Probe measurements of the plasma parameters, i.e. electron temperature and density, ion temperature and density. Main features of the Langmuir Probe diagnostics are discussed in [49].

The Langmuir probe used for our experiment consists of a 4 mm length of tungsten wire, with a diameter of 150 μm (the metallic wire protrudes from a cylindrical alumina sheath). An apposite motor allows to move the probe inside the plasma chamber with an accuracy of 1 mm.

The Langmuir probe is located so that it can penetrates inside the cavity for its entire length, forming an angle of 14° with respect to the chamber axis, as shown in fig. 3.1.2. An HIDEN commercial Langmuir probe has been used [69], and it was "user friendly" especially for the data analysis: the HIDEN software allows to easily determine the main plasma parameters by fitting the characteristic $I - V$ curve.

3.2 The Frequency Tuning with the SUPERNANO- NOGAN ECRIS

In 2005 and 2006 a series of tests were carried out on the SUPERNANO-NOGAN source [70] operating at 14.5 GHz, in order to verify the dependence of the extracted current for a particular charge state on the microwave frequency (i.e. the so-called frequency tuning effect). SUPERNANO-NOGAN is an ECR source using permanent magnets for the radial and longitudinal magnetic fields, so the total electrical power is extremely reduced. The RF power

3.3 X-ray measurements on the VENUS ion source

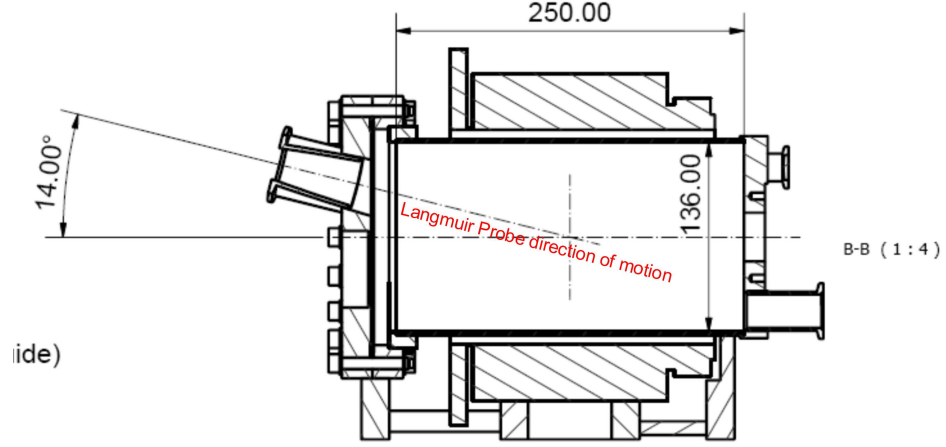


Figure 3.1.2: Design of the plasma reactor (lateral view) with the direction of motion of the Langmuir probe.

needed varies from 200 W to 500 W at 14.5 GHz . The maximum extraction voltage is 35 kV . This ion source can produce C^{4+} beams with intensities well suited especially for hadrontherapy facilities [71].

In order to verify the performances variations with the frequency, a TWT amplifier was used, able to vary the output frequency from 13.75 to 14.75 GHz . The experiment was carried out in the frequency range 14.440 – 14.530 GHz . The extracted ion current was measured, as usual, by means of a Faraday cup. A 90° magnetic dipole was employed as a charge over mass selector. Results obtained during this experiment are reported in section 4.2.

3.3 X-ray measurements on the VENUS ion source

The VENUS ECR ion source has been designed for operation at 28 GHz following scaling laws for ECR ion sources [2], [34], [72], [73]. The VENUS source is the first ECR ion source with a sufficiently strong magnetic field to operate at optimum values of magnetic field for 28 GHz fulfilling the scaling laws [7]. The higher plasma density, probably close to the cutoff $n_{cutoff} \sim 10^{13} \text{ cm}^{-3}$, permits to obtain the highest currents for highly charged heavy ions up to date. The choice of 28 GHz as the heating frequency with a B_{ECR} of 1 T implied several design consequences and challenges. Figure

3. EXPERIMENTAL EQUIPMENTS

3.3.1 [7] shows the mechanical layout of the VENUS Ion Source with its beam transport system. Table 3.1 summarizes the principal characteristics of the Source.

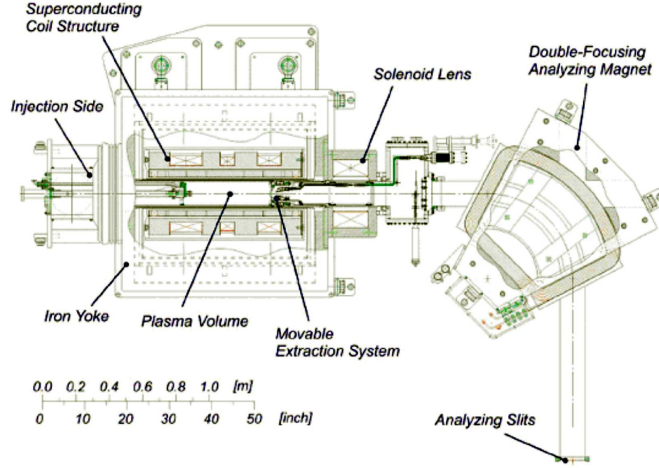


Figure 3.3.1: Mechanical layout of the VENUS Ion Source with its beam transport system.

One of the most important goals of the VENUS Ion Source is the production of high current for the highly charged heavy ion beams. For these ions typically intensities needed are $0.5 - 10 \text{ p}\mu\text{A}$. To achieve such intensities on the target in the experimental areas, for instance with xenon, 200 to 300 $e\mu\text{A}$ of Xe^{27+} current out of the ion source is needed.

It is important to notice that the values of currents and the extracted beams of VENUS are much higher than the ones obtained by any other ECRIS (including SECRA and SERSE)

3.3.1 The magnetic field of VENUS

In order to fulfill the requests presented in the previous paragraph an appropriate magnetic system has been designed for VENUS. It can operate at

3.3 X-ray measurements on the VENUS ion source

ID of the plasma chamber	15 cm
Mirror field on axis	4.0 T at injection, 3.0 T at extraction
Mirror-mirror spacing	50 cm
Central field	0.2 – 1.0 T (variable)
Min. radial field, plasma wall	2.0 T

Table 3.1: Typical values of magnetic field and characteristic dimensions of the plasma chamber for VENUS.

18 and 28 GHz or both (in the Two Frequency Heating mode) with axial magnetic fields up to 4 T at the injection side and 3 T at the extraction side and a wide range of B_{min} fields between the magnetic mirrors [7]. In addition the sextupole for the radial field generation can operate up to 2 T (i.e. twice B_{ECR}) at 28 GHz .

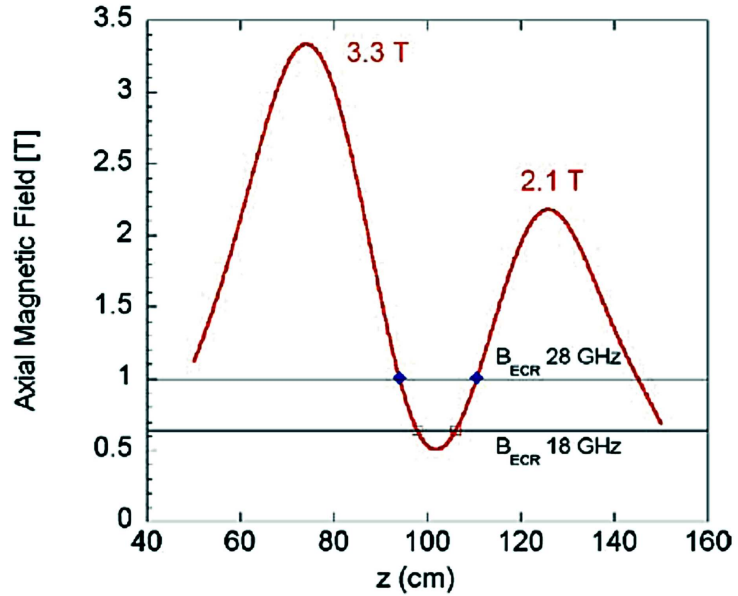


Figure 3.3.2: Typical VENUS axial magnetic field profile.

The range and flexibility of the magnetic field system and of the microwave generators allows to study the effect of different magnetic fields, micro-

3. EXPERIMENTAL EQUIPMENTS

wave power and frequencies on the ECR plasma, hence on the ECR ion sources performances. On this purpose the fig. 3.3.2 [7] shows the axial profile of the magnetic field (the field generated by the three axial solenoids) with typical operating values. For sake of simplicity, all the tests described in the following have been carried out in single frequency mode at 28 GHz .

Finally it is to be remarked that the design and construction of this source has been done by the ECR ion source team of LBNL, but they have largely based their activity on the studies and results of the Ion Source Team of LNS and of CEA, Grenoble [ref GyroSERSE, 94-2000].

3.3.2 Description of the X-Ray measuring device

The measurements of X-ray emitted by the VENUS sources have been carried out by placing on the axis a NaI phototube, which is particularly suitable for the detection of X-rays in the range of tens and hundreds keV . Fig. 3.3.3 shows the scheme of X-ray detection. The X-ray signal was collimated by a $5''$ (12.7 cm long) heavy met collimator with a square groove of 1 mm by 1 mm along the lengths. In addition the collimator was surrounded by lead shielding blocks minimizing the leakage of the X-ray radiation around the collimator set-up. The collimator was aligned with the extraction hole of the VENUS ECR ion source. The NaI phototube detector was placed about $15''$ away and was surrounded by lead shielding blocks to minimize the background coming from other X-ray sources in the vicinity (e.g. the AECR ion source). The energy calibration was done by using a ^{207}Bi source, with characteristic gamma lines at 569 and 1063 keV . The electronics used for this detector were the following: High voltage (Ortec) 1000 V Amplifier (as specified by the manufacturer): coarse gain 1 k , fine gain " 510 ", shaping time 0.5 ms . TFA unit (Ortec): coarse gain 20 , fine gain " 2 o'clock ", time constant " 100 ". Discriminator unit. GDG (Gate & delay Generator): delay " 110 " (1.1 ms). The signal was processed with a MCA-8000A Pocket Multichannel Analyzer. This set-up was used for the experiments here described. Anyway as severe problems occurred with dead time (pile up in the detector) and also with the stability of the PM tube in terms of gain, a different setup was prepared for the future experiments. In fact the gain shifts take to change of energy calibration and they were observed because of temperature changes and high count rates. In addition, the lead shielding around the X-ray port was insufficient so that high energy X-rays could leak around the main collimator. All these systematic errors did not change significantly many of the results and in the most of case they just lowered the amount of detected X-rays. In the cases with higher counting rate, anyway, the pile-up

3.4 Frequency tuning affecting the beam shape: experimental set-up

may change the slope of the curves and for this reasons it was decided to improve the detection set-up. The results described below are not affected by the systematic errors.

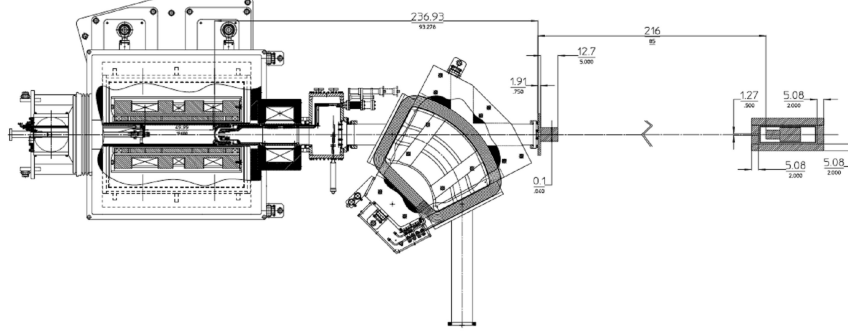


Figure 3.3.3: The VENUS source and the X-ray detection system (on the right) after the changes.

The new set-up is shown in fig. 3.3.3. To reduce the dead time the detector was moved further away. The shielding was improved so that the signal is effectively collimated by the heavy met collimator. With the current set-up, only a few counts of background were obtained, so the shielding is adequate. A new phototube was used so that more stable detector will be available in the future experiments.

3.4 Frequency tuning affecting the beam shape: experimental set-up

In the next chapter we will describe some experimental results concerning the frequency tuning effect, as a great part of the theoretical studies carried out during the PhD course have been devoted to the understanding of this effect. One of the most important results, that is also a further proof of the strong influence of the frequency variation on the ECRIS plasma dynamics (both on electrons and ions), has been obtained with the CAPRICE ion sources at GSI testbench by ramping the RF frequency in a narrow range (± 40 MHz) around 14.5 GHz and by looking to the ion output and the beam shape [6]. The CAPRICE sources are compact sources, developed by the historical Grenoble ion source group, which are able to produce large currents of intermediate charge state ions such as O^{6+} , Ar^{11+} , Kr^{17+} , etc.

3. EXPERIMENTAL EQUIPMENTS

and which are still adopted by different laboratories all over the world. A detailed description of these sources can be found elsewhere [74], [75], [76], [59] and here we will briefly remind some main characteristics (see fig. 3.4.1).

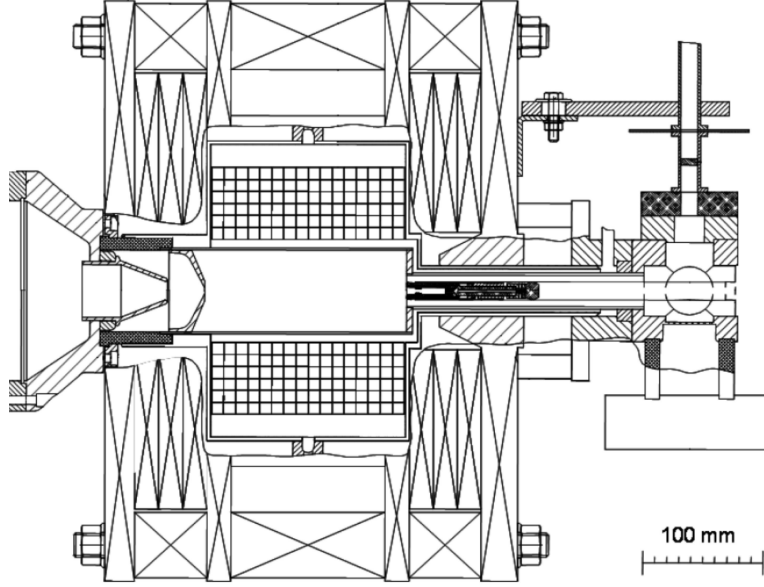


Figure 3.4.1: Design of the CAPRICE source installed and operating at GSI, Darmstadt.

It can be observed that the electromagnetic waves are fed to the plasma chamber by means of a coaxial transmission line working in the TEM mode [76].

During the tests the source was fed by a microwave generator based on a 8 channel klystron, all of them tuned at 14.5 GHz central frequency and with a -1 dB bandwidth of about 100 MHz . The microwave generator provided waves in the frequency range: $14.5\text{ GHz} \pm 40\text{ MHz}$. The frequency sweep has been obtained by replacing the local oscillator driving the klystron amplifier with a HP8620C analogic sweeping generator.

The plasma chamber is a cylindrical cavity 64 mm in diameter and 181 mm long (considering also the length due to the shape of the extraction electrode, see fig. 3.4.1). Theoretically it comes out that only five modes can be excited in this cavity, coherently with the span of the frequency sweep. Some of them, even in the adjacent frequencies, have a structure presenting the minimum electric field on chamber axis. This fact is due to

the large chamber dimensions with respect to the wavelength of the feeding microwave (the first resonant frequency is at 2.8 GHz) and therefore at our operating frequencies higher order modes may be excited and part of the power may be coupled to them. However, in presence of plasma, being it a dispersive medium, the electric field present in the plasma chamber for a given frequency may be due to the contributions of different modes present simultaneously and overlapped.

Target viewers and pepper-pot measurement devices have been used to investigate the beam shape modifications depending on the beamline and source settings. Three viewers (VT1, VT2 and VT3) are located just after the extraction (before the focusing solenoid), after the solenoid and in the dipole focus point respectively. All of them are tilted by 45° to permit the visualization of the beam profile by means of a CCD camera through a quartz window. Different observations have been done by varying the source parameters (magnetic field, gas pressure, extraction voltage, microwave frequency and power) and the beamline settings (repeller voltage, solenoid current, quadrupole current and dipole current) outlining the importance of each parameter in the optimization of the ion beams supplied to particle accelerators.

Helium and Argon were used for plasma creation. The forward and reflected power were precisely measured by using two *ROHDE&SCHWARTZ* power sensors connected to the microwave line through high directivity waveguide reflectometers.

3.5 The WEGA stellarator and the experimental apparatus

WEGA is a classical mid-size stellarator, whose toroidal vacuum chamber has a major radius of $R = 0.72\text{ m}$ and minor radius of $r_0 = 0.19\text{ m}$ [21]. Hydrogen, Helium and Argon are used as working gas. The pulse duration varies from a few seconds up to a few minutes depending on the magnetic field strength and hence the dissipated energy in the field coils. Short shots are used for high magnetic field experiments, the long time shots may be performed with magnetic field strength up to 0.34 T . In this case the an operation of WEGA in steady regime is possible. Figure 3.5.1 describes the structure of the stellarator magnets. It can be seen that three different types of coil exist. In particular, the toroidal coils generate the field for the toroidal confinement, while the helical fields generate the poloidal field, that is fundamental to obtain the stabilization of the plasma confinement both

3. EXPERIMENTAL EQUIPMENTS

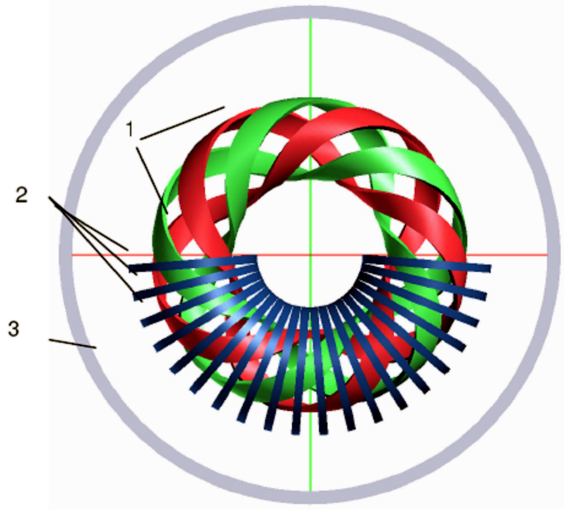


Figure 3.5.1: Design of the WEGA stellarator showing the helical coils (1), the toroidal field coils (2), and the vertical field coils (3).

in terms of the single particle and fluid approach.

The figure 3.5.2 shows the different tools used for the plasma diagnostics. In addition, during our experiment, a soft X-rays detector was used, in order to detect X-rays emitted by plasma electrons via bremsstrahlung with energies 500 eV or more. The plasma is ignited and heated by ECR with a maximum power of $20 + 6\text{ kW}$ coming from two separate magnetrons operating at the frequency of 2.45 GHz . The microwaves are launched in the equatorial plane of the torus through A-ports at two positions shifted of 36° , as shown in figure 3.5.2.

Because of the low heating frequency, our experiment has been carried out with $B \sim 0.0875\text{ T}$ where 0.0875 T is the electron cyclotron resonant field for the 2.45 GHz heating frequency. A detailed description of the wave injection system used on WEGA is given in figure 3.5.3. A general view of the WEGA Stellarator is given in figure 3.5.4.

3.5 The WEGA stellarator and the experimental apparatus

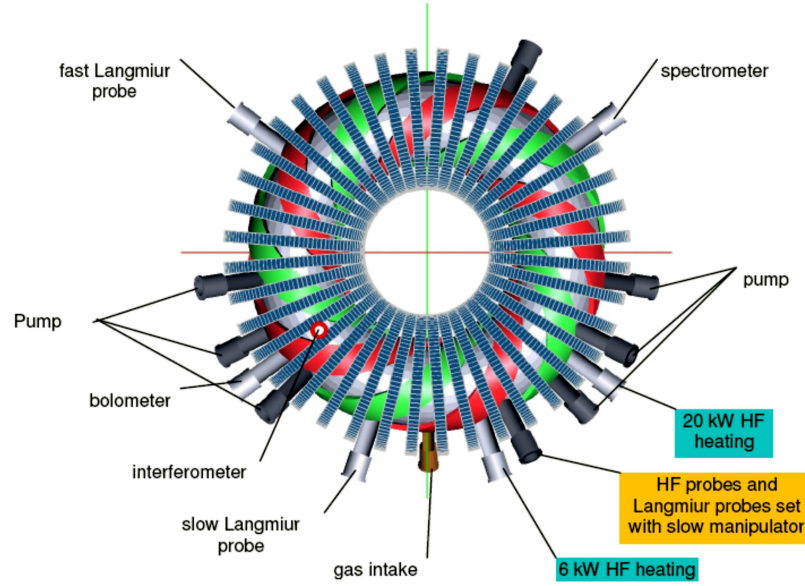


Figure 3.5.2: Locations of the various diagnostics tools and of the two ports for the microwave injection.

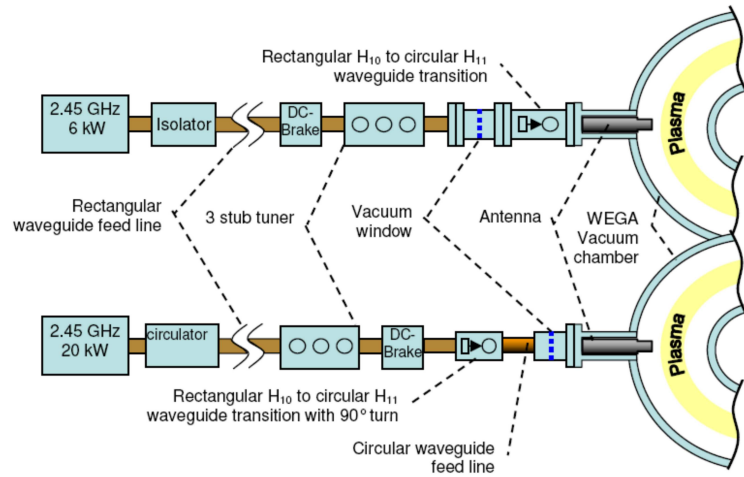


Figure 3.5.3: Description of the microwave injection system used on the WEGA stellarator for the 6 kW magnetron (upper) and for the 20 kW (lower).

3. EXPERIMENTAL EQUIPMENTS

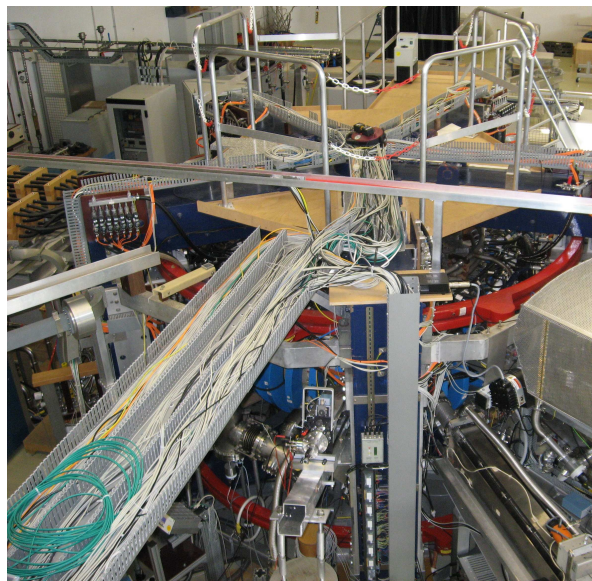


Figure 3.5.4: General view of the WEGA system.

3.5 The WEGA stellarator and the experimental apparatus

Chapter 4

Microwave coupling to ECRIS plasmas

4.1 Overcoming the ECR standard model

Up to now the so-called "ECRIS Standard Model" has been proven by many experiments carried out with different sources, operating at different frequencies. The main rules were confirmed by experiments performed at MSU-NSCL in 1993-94 and in 1995 and they can be summarized as it follows [77]:

- a) the radial magnetic field value at the plasma chamber wall must be $B_{rad} \geq 2B_{ECR}$;
- b) the axial magnetic field value at injection must be $B_{inj} \simeq 3B_{ECR}$ or more;
- c) the axial magnetic field value at extraction must be about $B_{ext} \simeq B_{rad}$;
- d) the axial magnetic field value at minimum must be in the range $0.30 < \frac{B_{min}}{B_{rad}} < 0.45$.
- e) the optimum power increases with the volume of the plasma and with the square of the frequency.

At power, magnetic fields and microwave frequencies used up to now all the ECRIS obey to the Standard Model: the extracted current strongly increases as the microwave frequency increases, but only the increase of mirror ratio can exploit the performances, making effective the increase of electron density with frequency.

Then according to the Standard Model the ECRIS development is strictly linked to the improvements of superconducting magnets and of the microwave generators technology. The scaling law concerning the microwave power and the ECR heating process (ECR-H) was not confirmed in all the cases because the relation between the power and the magnetic field is not so simple. Some authors studied the RF coupling to the plasma in terms of the maximum power rate per unit volume and of its relationship with the beam intensity produced by different ECR ion sources [78], but this description is not satisfactory. The cavity design, and the microwaves injection geometry, are of primary importance for a high RF energy transmission to the plasma chamber. Note that the problem of the wave energy transmission into the plasma must be divided into two parts: the first is connected to the microwave generator-waveguide-plasma chamber coupling, while the second one regards the wave-plasma interaction. Both these two aspects are described in this thesis and play a notable role for future improvements of ECRIS performances. The standard model does not allow to argue any prediction about the mechanism at the basis of electron and ion dynamics for different profiles of magnetic field and for slightly different microwave frequencies. Some preliminary experimental results obtained with III generation ECRIS (VENUS) put in evidence the limits of Standard Model, in particular regarding the magnetic field scalings. We expect that the technological limits imposed by this model can be overcome by means of a better understanding of plasma physics, in particular of the phenomenon that is at the basis of the ECR plasma creation, i.e. the ECR-heating.

Many experiments have been carried out to verify the possibility of different plasma heating schemes out of the framework traced by the Standard Model. Principally they follow three different roadmaps: a first series of experiments has been devoted to the study of the variation of ECRIS performances with slight variations of the microwave frequency. The second one regards the possibility to operate with more than one frequency for plasma heating (usually two) or by using broad microwave spectra. The third method actually is a mix of the previous ones: two or more frequencies are used for plasma heating, but at least one is provided by a broadband microwave generator (like a TWT) which allow to combine the effects of frequency tuning and of multi-frequency heating.

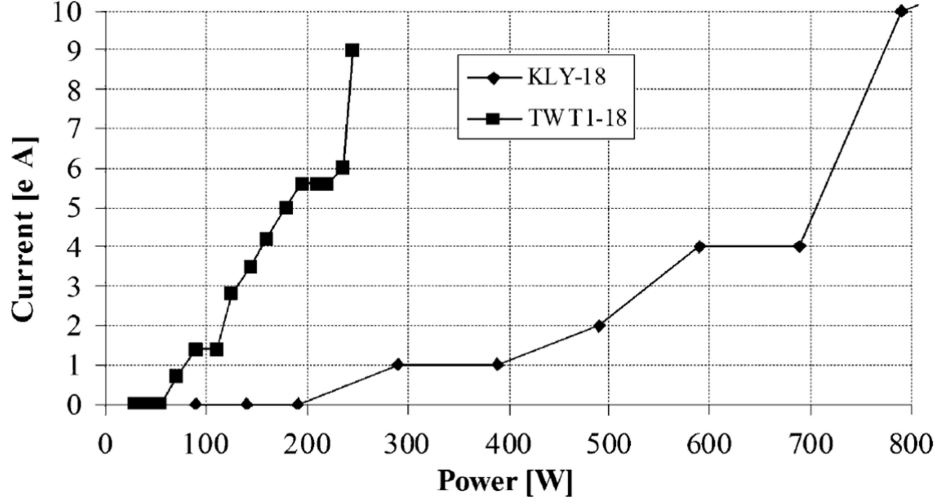


Figure 4.2.1: Comparison between trends of O^{8+} at 18 GHz for klystron (up to 800 W) and TWT1 operating in the same range of frequency.

4.2 The frequency tuning effect

Many experiments in the last years have shown that significant improvements of ECRIS performances (both in terms of total extracted current and highly charged ions production) are obtainable by slightly varying the microwave frequency in the case of SFH (Single Frequency Heating) [79], [6], that is defined as a "frequency tuning" effect. It was known that a large increase of the frequency ($\sim GHz$) increases the electron density (see eq. (1.9.2)) and improves the ECRIS performances because of the increase of the cutoff density, but Since 2001 several experiments have demonstrated that even slight variations of the pumping wave frequency ($\sim MHz$) may lead to strong variation of the extracted current.

The first evidence was given by the different performances observed for the SERSE and CAESAR ion sources when fed by a klystron based or a travelling-wave-tube (TWT) based generator [80], [81], [82] either at 14 and 18 GHz .

Other interesting results came from experiments performed at ORNL and at JYFL [83].

Figure 4.2.1 features that a remarkable increase in the production of O^{8+} can be obtained by using the TWT instead of Klystrons; in particular

the current of O^{8+} obtained with the TWT at 240 W is obtained with the Klystron at 800 W, i.e. a power three times higher. It is important to underline that the two generators operated at two different but very close frequencies; in particular the Klystron was operating at 18.0 GHz, while the TWT at 17.9 GHz (both the amplifiers were fed by DRO-type oscillators). Such difference of performances was initially explained as a greater frequency dispersion of the TWT, but a series of measurements carried out with the two microwave generators, by means of a Spectrum Analyzer, pointed out that the spectrum of the emitted radiation of a TWT is similar to the one of the Klystron [84]. Then the only difference between the two generators is the output frequency, with the further possibility for the TWT to vary the emitted frequency, thus optimizing the source performances.

Fig. 4.2.2 shows the current for C^{4+} obtained with the SUPERNANOGAN ion source versus the microwave frequency, keeping unchanged the power and all the other source parameters. The figure features strong fluctuations of the extracted current in a frequency span of about 90 MHz.

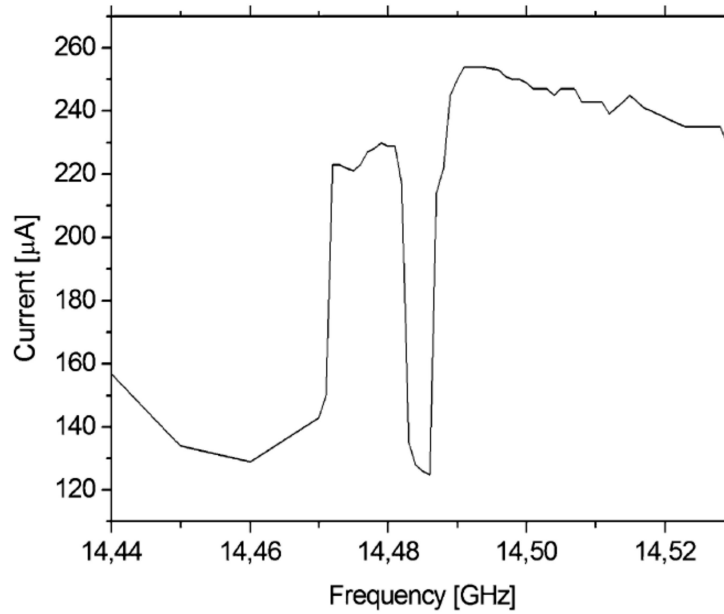


Figure 4.2.2: Trend of the analyzed C^{4+} current for the SUPERNANOGAN ECR Ion Source versus the RF frequency.

The frequency was changed in the range 14.44 to 14.53 GHz with a

4. MICROWAVE COUPLING TO ECRIS PLASMAS

step of 1 MHz and it was observed that changes of a few MHz changed the C^{4+} current even of 50%. Further experiments including the final validation of the frequency tuning effect have been carried out with the SUPERNANO-GAN source of CNAO. The effect was remarkable for this source, as it permitted to increase by 30% the best performance for C^{4+} and by 50% for H^{3+} (these ions are particularly requested for medical applications); an additional increase of reliability and availability was also registered.

Interesting results came from the GSI experiment [6] (see section 3.4). Figure 4.2.3 reports the shape evolution of a Helium beam recorded on the first viewer related to an acquisition lasting 150 sec (500 W of microwave power with standard magnetic field profile and injection pressure of $4.3 \cdot 10^{-6}$ $mbar$).

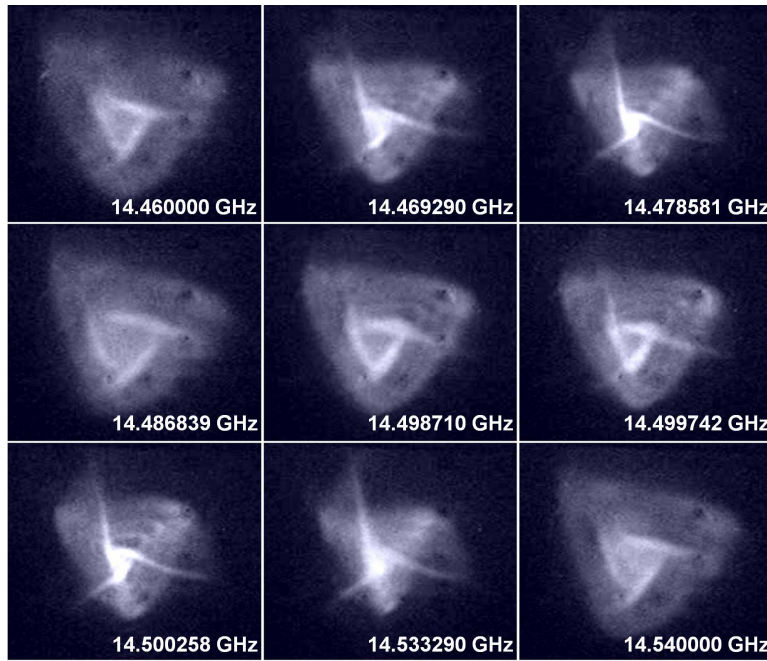


Figure 4.2.3: Variation of the extracted beam shape for different microwaves frequency.

Figure 4.2.3 shows that the beam intensity distribution is inhomogeneous and that this distribution changes with the microwave frequency, while keeping all the other parameters constants.

By keeping the source in these conditions and looking to the analyzed

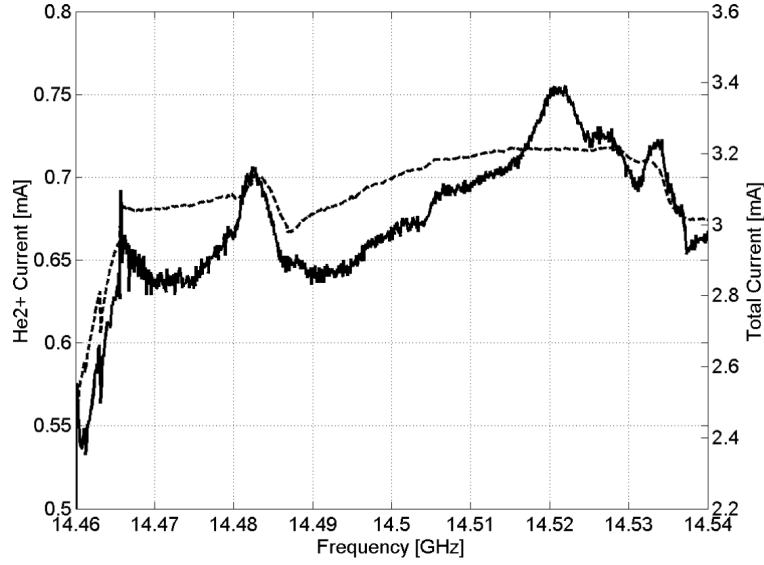


Figure 4.2.4: Trend of the He^{2+} extracted current (full line) for different microwave frequencies. The total extracted current is also shown (dashed line).

beams after the dipole magnet, it can be observed a change of the He^{2+} beam current during the sweep together with a change of the total current extracted by the ion source (fig. 4.2.4).

The current increase for some values of the microwave frequency is about the 50%, as in fig. 4.2.2. The total extracted current also varies with the frequency, but with a lower percentage with respect to the He^{2+} current, because different ion species are differently affected by the plasma parameters changes.

The experiment was repeated by positioning the viewers in different position of the beam line, but in any case beam transport settings were kept constant: the effect of the frequency change was clearly visible also in these cases.

Similar observations have been done for Argon beams, for which the ion current evolution of the Ar^{8+} charge state during some frequency sweeps (each one 150 sec long) was investigated. The measured trends of the extracted current were repeatable in each sweep. Measurements of the S_{11} scattering parameter were carried out in parallel. The S_{11} parameter is connected with the reflection coefficient, and the measurements have shown that for some frequencies the amount of the reflected power increases, thus

demonstrating that some resonances are present inside the cavity with or without plasma. In addition, fluctuations in the extracted power occurs also in case of quasi-constant reflection coefficient.

4.3 Standing waves in dense plasmas

The explanation of the data above presented is based on the assumption that the frequency tuning changes the electromagnetic field distribution inside the resonator in terms of its distribution over the resonance surface, i.e. where the wave-electron energy transfer takes place. This assumption requires that even in case of plasma filling the cavity the resonant modes persist, i.e. the formation of standing waves is still possible. In [85] and [25] it was shown that the plasma chamber of an ECRIS works as a high-Q resonator in vacuum, but up to now no confirmations about the presence of standing waves (and then of resonant modes) are available in literature for ECR-generated plasmas.

The determination of the resonant frequency shift and the measurement of the quality factor of the cavity Q allow often to calculate the plasma density and the collision frequency [86], [87]. In particular Adgur has given a complete picture of the mode pattern existing in a cylindrical cavity, including modes that are not directly related to the ones existing in an empty cavity. In this case the cavity is cylindrically shaped and it contains a central plasma column. The tests here described have been carried out with the device described in section 3.1, that is somehow different. In fact the plasma completely fills the cavity and the magnetic field makes it a non-isotropic medium for the electromagnetic field propagation. Additional hints come from Asmussen [88], who studied the interaction of microwaves and cavity in a plasma source, and he observed that the plasma density increases less than the incident power. Only a small fraction of the additional incident power is absorbed into the plasma and the remaining part is reflected from the cavity. This is a fundamental problem when attempting to produce high-density plasma inside a microwave cavity and it can be solved by varying the excitation frequency while the plasma is RF sustained or by to continuously varying the cavity dimensions. This model explains the saturation of electron density with the pressure. A more similar situation to the MDIS devices was investigated by Shohet et al. [89], as he assumed a multiple-shell model considering a cylindrical plasma column of variable radial density filling entirely the plasma chamber.

As mentioned before, the main difference between the MDIS and other

ECR plasma based devices above mentioned is connected with the presence of the magnetic field. It is not so easy to determine a priori whether the ECR destroys or not the mode structure. The question can be solved by means of the properties of the stochastic ECR heating and of other additional hypothesis, as it will be more clear in section 4.13.

In different ways we made some attempts to determine theoretically the shift of resonance frequencies for isotropic and homogenous plasmas. However the experimental measurement of the plasma parameters, i.e. the electron density and the electron temperature, have put in evidence that they have a radial and axial variation (the plasma density is usually higher in proximity of the wave guide feeding the cavity, and the plasma radius is comparable with the chamber's one), then it is not easy to link the electron density and the mode shift with simple proportionalities, but semi-empirical laws can be obtained. Finally, microwave discharge plasma sources are quite different devices with respect to the ECRIS, as for the plasma energy content, but the plasma densities are quite similar and the experimental results that will be shown in the following may be partially valid also for wave pattern formation in ECRIS.

4.3.1 Experimental Results: electromagnetic waves propagation in a under vacuum resonating cavity

A theoretical calculation of the resonant modes in vacuum have been done in parallel with the experiment by solving the eigenvalue equations for the electromagnetic field inside the plasma chamber and by using an oversimplified model of the cavity (it has been considered empty and completely closed). As outlined before the plasma chamber is a cylindrical (stainless steel) cavity and the electromagnetic waves are provided through a WR284 rectangular waveguide. In a cavity with radius a and length l , the possible frequencies for TM and TE modes are given by the formula [13]:

$$\omega^2 = c\sqrt{\frac{r^2\pi^2}{l^2} + h^2} \quad (4.3.1)$$

where r is an integer, l is the length of the plasma chamber and h is the eigenvalue obtained by solving the following eigenvalue equations:

$$\nabla_t^2 H_{z0} + h^2 H_{z0} = 0 \quad (4.3.2)$$

$$\nabla_t^2 E_{z0} + h^2 E_{z0} = 0 \quad (4.3.3)$$

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Number	Mode	Calculated resonance frequency [GHz]
1	$TE_{1\ 1\ 1}$	1.4049
2	$TE_{1\ 1\ 2}$	1.6908
3	$TM_{0\ 1\ 0}$	1.6924
4	$TM_{0\ 1\ 1}$	1.7774
5	$TM_{0\ 1\ 2}$	2.0110
6	$TE_{1\ 1\ 3}$	2.0817
7	$TE_{2\ 1\ 1}$	2.2169
8	$TM_{0\ 1\ 3}$	2.3492
9	$TE_{2\ 1\ 2}$	2.4082
10	$TE_{1\ 1\ 4}$	2.5295
11	$TM_{1\ 1\ 0}$	2.6965
12	$TE_{2\ 1\ 3}$	2.6971
13	$TE_{0\ 1\ 1}$	2.7507
14	$TM_{1\ 1\ 1}$	2.7507
15	$TM_{0\ 1\ 4}$	2.7538
16	$TE_{0\ 1\ 2}$	2.9071
17	$TM_{1\ 1\ 2}$	2.9071

Table 4.1: Resonant modes numerically calculated for the plasma reactor chamber.

The modes excitation depends on the frequency of electromagnetic wave launched in the plasma chamber and on the matching between the plasma chamber and the waveguide. As shown in section 3.1 the coaxial probe can be used to detect the S_{11} parameter or to excite the plasma instead of the waveguide. In order to enhance the matching impedance the probe was inserted into the cavity through the $DN40$ flange, located in the opposite direction with respect to the waveguide location and about 30 mm off-axis. Two different probe lengths have been used: the probe A that is 20 mm long and the probe B that is 30 mm long, having both 3 mm of diameter. Other than the calculated modes, the table 4.2 reports the resonance frequencies measured thanks to the two probes described before and to the directional coupler connected to the $WR284$.

Note that many of the calculated modes are excited, but minor differences exist with the measured frequencies because of the poor parallelism of the plasma chamber dimensions that is the source of errors. Some modes

Number	Mode	Calc. f [GHz]	Meas. f probe A [GHz]	Meas. f probe B [GHz]	Meas. f WR284 [GHz]
1	$TE_{1\ 1\ 1}$	1.4049	1.3978	1.3978	
2	$TE_{1\ 1\ 2}$	1.6910	1.6731	1.4006	
3	$TM_{0\ 1\ 0}$	1.6924	1.6897	1.6734	
4	$TM_{0\ 1\ 1}$	1.7774		1.6897	1.7534
5	$TM_{0\ 1\ 2}$	2.0110		1.8175	1.9822
6	$TE_{1\ 1\ 3}$	2.0817	2.0838	2.0844	
7	$TE_{2\ 1\ 1}$	2.2169	2.2091	2.2097	2.2089
8	$TM_{0\ 1\ 3}$	2.3492	2.3687	2.4020	
9	$TE_{2\ 1\ 2}$	2.4082		2.4041	2.4033
10	$TE_{1\ 1\ 4}$	2.5295	2.5288	2.5291	2.5287
11	$TM_{1\ 1\ 0}$	2.6965	2.6925	2.6938	2.6925
12	$TE_{2\ 1\ 3}$	2.6971	2.6975	2.6981	2.6975
13	$TE_{0\ 1\ 1}$	2.7507			
14	$TM_{1\ 1\ 1}$	2.7507	2.7198	2.7400	
15	$TM_{0\ 1\ 4}$	2.7538	2.8191		2.8191
16	$TE_{0\ 1\ 2}$	2.9071	2.8956	2.8956	
17	$TM_{1\ 1\ 2}$	2.9071			

Table 4.2: Comparison between theoretical and measured modes through the two probes A and B, and the waveguide WR284.

cannot be excited by one of the probe or by the WR284 according to the matching impedance to the cylindrical cavity. In case of plasma parameters measurements (carried out by means of the Langmuir probe) a 7 mm diameter alumina cylinder penetrates inside the cavity up to 9 mm off axis with a slope of 14° with respect to the plasma chamber axis. The plasma characterization has been performed by positioning the Langmuir probe at $z = 22$ cm inside the cavity. The measurements have demonstrated that the Langmuir probe alumina cylinder affects the reflection coefficient measured at the coaxial connector, as shown in fig. 4.3.1 (b). A similar behavior was observed when the measurements were performed through the waveguide, and this means that the amount of power which enters into the cavity depends on the LP position (the relative variation of the reflection coefficient reached values on the order of 10 – 15%).

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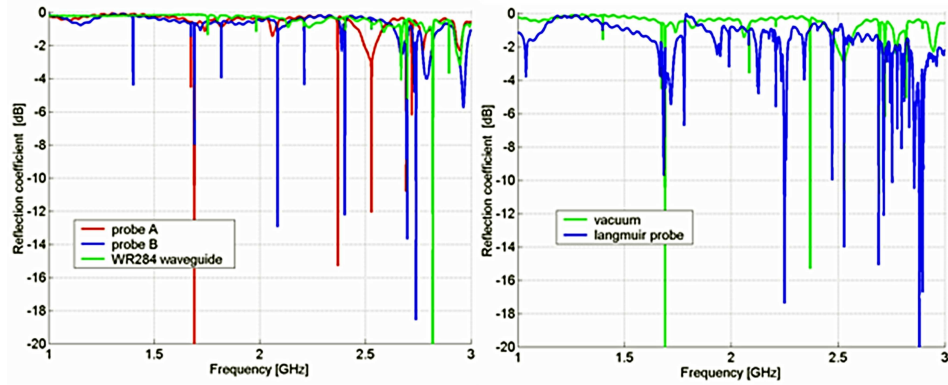


Figure 4.3.1: (a) S_{11} parameter obtained through the probe A, the probe B and the WR284 waveguide. (b) Comparison between the S_{11} parameter measured in case of vacuum (green line) and in case of Langmuir Probe inserted into the chamber.

4.3.2 Modes detection in presence of plasma

A detailed characterization of the scattering parameter S_{11} has been carried out in vacuum at the beginning, then in presence of plasma generated at different power and pressure. At first it is important to demonstrate that the plasma creation does not destroy all the resonant modes; then the shift of the resonant frequencies with the RF power and the pressure have been investigated. In order to verify how the S_{11} structure changes during the plasma creation some frames of the network analyzer (NA) display images have been acquired (with time interval of 0.15 sec) and two of them are shown in fig. 4.3.2.

The figure features that during the plasma ignition the S_{11} structure strongly changes because of the different coupling between magnetron and reactor (as a consequence of plasma formation); clear resonances are still visible. Some resonances disappear as the plasma ignition goes on, whereas some others appear at different frequencies. Modes existing only in vacuum have been indicated with red circles, that ones existing with and without plasma (a frequency shift occurs in this case) have been indicated with yellow circles, and in green the only mode that has been observed only in presence of the plasma.

The shift of modes can be correlated with the electron density measured by means of a Langmuir probe, even if the Langmuir probe strongly perturbs the magnetron- cavity coupling. With the Langmuir Probe into the chamber

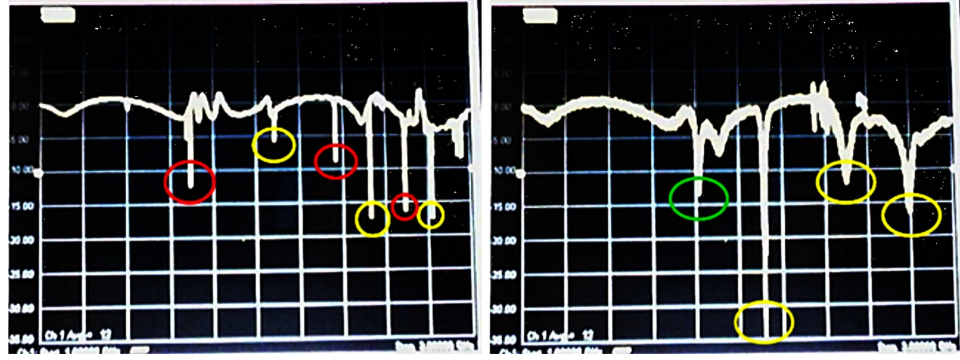


Figure 4.3.2: S_{11} parameter in vacuum (left) and after the plasma ignition (right). In red we have indicated the modes existing only in vacuum, in yellow the modes existing with and without plasma, finally in green the mode existing only in case of plasma filling the activity has been indicated.

two modes clearly shift towards higher values of frequency after the plasma ignition, thus demonstrating that some modes are preserved even after to the plasma creation. In order to investigate in details the frequency shift, the frequencies 2.249 and 2.896 GHz , corresponding to resonant modes in the cavity without plasma, have been chosen. These frequencies correspond to modes in vacuum, but in presence of the Langmuir probe we cannot argue if they are TE or TM modes because they probably have an hybrid nature [86]. The frequency shift was observed for different values of power and pressure. The microwave power range used for the experiment was $30 \div 70$ W and the gas pressure was variable between 0.1 and 0.9 $mbar$.

The modes frequency shift for modes at 2.249 and 2.896 GHz (in vacuum) in case of cavity with the plasma is shown in fig. 4.3.4. The peaks vary with frequency; the wave coupling changes as a consequence of the RF power and chamber pressure variations, which affect the electron density. The coupling changes are put in evidence by the variation of the peak height, which can change up to 10 dB by varying the power or the pressure.

The shift has been determined by means of the following equation:

$$\frac{\Delta f}{f} = \frac{f_{on} - f_{off}}{f_{off}} \cdot 100 \quad (4.3.4)$$

where f_{on} and f_{off} are the frequencies with and without plasmas. From the figure 4.3.4 it can be observed that the shift decreases for higher value of

4. MICROWAVE COUPLING TO ECRIS PLASMAS

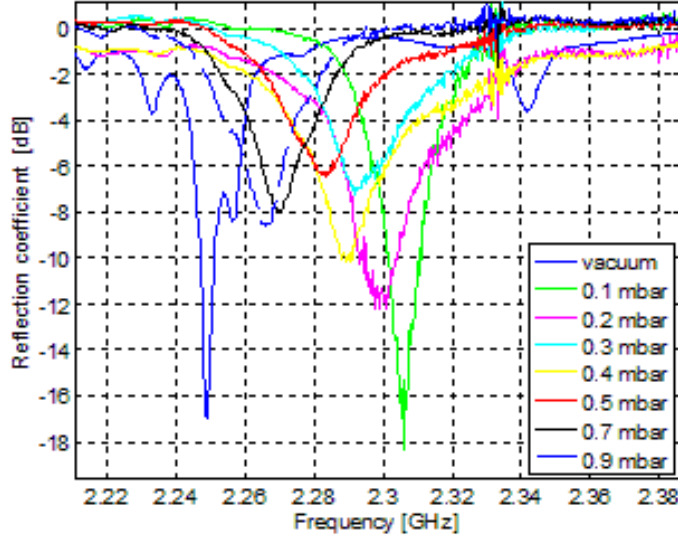


Figure 4.3.3: Shift of the resonant mode at 2.249 GHz for different values of pressures at 50 W of microwave power.

pressure¹, i.e. it is linked to the plasma operative parameters, in particular to the variation of the electron density.

4.3.3 Plasma parameters characterization

During the interaction with the plasma the wave can be weaker or stronger absorbed or reflected at the plasma boundary if the pressure and/or the power are changed. Also the electromagnetic field structure inside the plasma chamber is strongly biased by the electron density distribution. The measurement of the electron density, in a fixed position of the plasma chamber, gives a not complete characterization of the plasma properties. On the other side, it is not possible to perform measurements in different regions of the cavity, as the variations of the probe position lead to notable variation in the structure of the S_{11} parameter. Then we must focus our attention on partial measurements carried out in a position where the density is less affected by phenomena which lead to strong changes of the plasma properties. The figures 4.3.5 and 4.3.6 clearly correlate the electron density and

¹Analogous measurements at constant pressure and different powers demonstrate that the shift is greater for higher value of power.

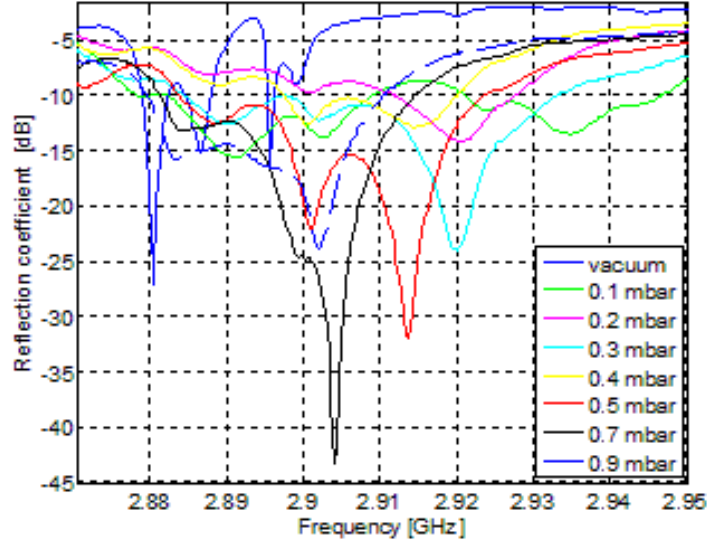


Figure 4.3.4: Shift of the resonant mode at 2.896 GHz on the right) for different values of pressures at 50 W of microwave power.

the frequency shift. The frequency shift decreases with the density, as expected from calculations in case of homogeneous and isotropic plasma filling cylindrical cavities.

It is also evident that the shift for 2.249 GHz it is higher than for 2.896 GHz . The trend of the frequency shift for the two modes with respect to the dimensionless plasma density η is shown in figure 4.3.6. The dimensionless density is defined as:

$$\eta = \frac{\omega_p^2}{\omega^2} = \frac{ne^2}{m\epsilon_0\omega^2} \quad (4.3.5)$$

The two curves shown in fig. 4.3.7 have been fitted with functions of type:

$$y = ax^b \quad (4.3.6)$$

For the mode at 2.249 GHz the $b \sim 0.5$, while for the other mode $b \sim 0.4$. The formula (4.3.6) features that the for both the modes no saturation occurs.

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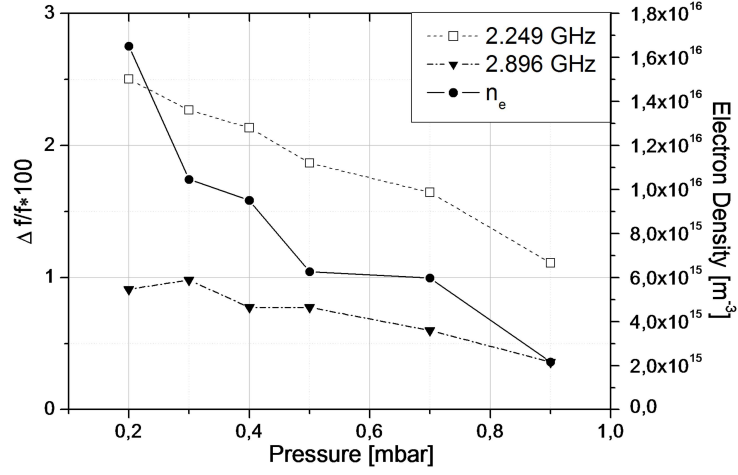


Figure 4.3.5: Comparison between the frequency shift and the electron density trends for different background pressures at $P_{RF} = 70 \text{ W}$.

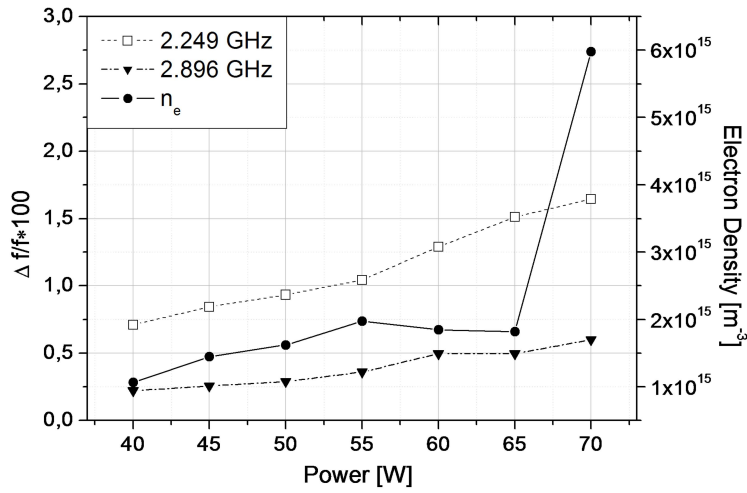


Figure 4.3.6: Comparison between the frequency shift and the electron density trends for different microwave power at 0.7 mbar of background pressure.

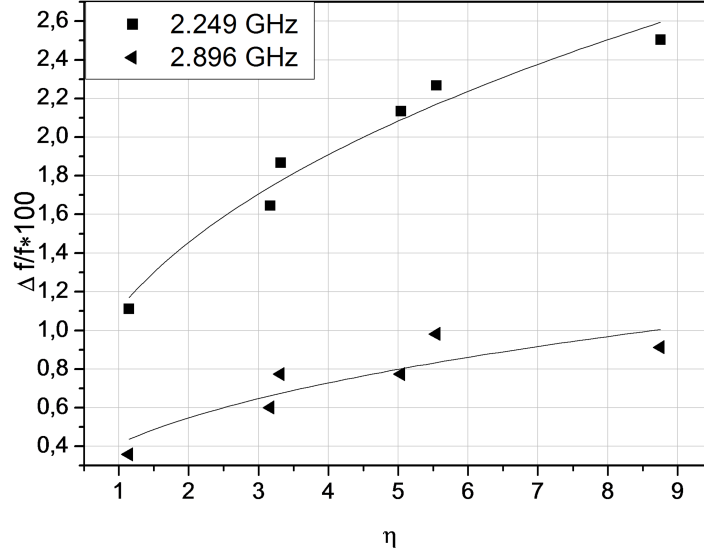


Figure 4.3.7: Trend of the frequency shift with respect to the electron density for the two modes under investigation.

4.3.4 Discussion

The experimental data here reported feature a maximum shift below 3%, i.e. ~ 30 MHz. Calculations performed for the vacuum chamber demonstrate that the spacing between subsequent modes increases as the frequency increases. In cavities operating around 2.45 GHz the mean spacing is about 20 MHz [90] and it decreases down to few hundreds of kHz for cavities operating at 28 GHz. Then the frequency shift occurring in the plasma reactor probably do not change considerably the mode structure, but the same shift at higher frequencies (like in case of ECRIS) may lead to strong changes in the excited modes patterns over the resonance surface. The strong change in the standing wave formation may lead to unstable operations which may be avoided by slightly changing the microwave frequency.

In case of plasmas excited by microwaves at 2.45 GHz if one increases the injected power the absorbed fraction becomes lower and lower and a sort of density saturation occurs. In order to operate in stable conditions, the plasma density adjusts itself as the power increases and it tends to an asymptotic value. The frequency tuning, and/or the variation of the cavity length (which leads to similar effects), usually improves the wave-plasma

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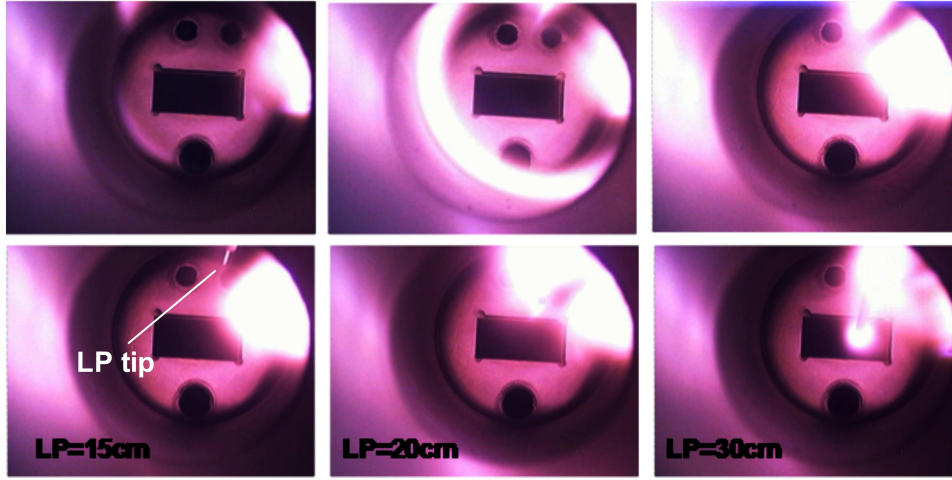


Figure 4.3.8: Frames showing the variation of the plasma shape variation for the LP moving inside the plasma chamber. In addition, in this case the plasma has been generated by means of a metallic probe inserted into one of the flanges located on the side of the plasma chamber, thus forming a sort of coaxial-like excitation.

energy transfer and it maximizes the plasma density.

A particular attention must be devoted to the possible excitation of the plasma through the coaxial connector which permitted to carry out the measurement of S_{11} parameter in case of plasma filling the cavity. As demonstrated in tab. 4.2 the length of the internal probe determines the coupling properties of the coaxial system formed by the injection flange and the internal probe itself. It is possible to cut the probe to a length of ~ 130 mm obtaining an excitation mode near to the magnetron output frequency. In such a way it was possible to excite the plasma with just 5–6 W of RF power instead of 20–30 W, usually required for the waveguide excitation. This result demonstrates once again that the coupling of the microwaves is essential for plasma excitation, as it determines the effective amount of power that can be spent for electron heating.

It was also observed that the plasma shape changed a lot for different position of the Langmuir probe inside the cavity. Also the plasma properties (electron temperature and density, etc.) change remarkably when the LP moves into the plasma. The change of the plasma shape is shown in figure 4.3.8; in that case the plasma is generated through the coaxial-like microwave injection.

The strong variation of the reflected power when the LP moves inside the cavity has a similar explanation. The reflected power measured through the directional coupler located in the injection waveguide fluctuates when the LP moves inside the plasma cavity, and it strongly increases when the probe approaches the microwave injection flange. A series of simulations carried out to investigate the electromagnetic field propagation inside the plasma chamber have demonstrated that the resonant modes present very pronounced maxima in proximity of the microwave injection flange (impedance matching devices are usually employed to have the maximum field in the center of the plasma chamber) and some small side peaks far from the microwave window. Then we may link the structure of the excited modes with the trend of the reflection coefficient by affirming that the perturbation induced by the LP is more pronounced where the electric field is higher.

The experiment demonstrated that the standing waves formation occurs inside the cavity also in case of high density plasmas and in presence of two ECR zones. The same measurements will be performed in future on the resonant cavities of high performance ECRIS, in order to confirm the results even in case of higher frequencies. Theoretical analysis on the mode formation must be extended to the presence of complex plasma structures: variation of density along the plasma axis and radius, presence of large gradients of magnetic field, etc. have to be taken into account.

4.4 The Two Frequency Heating Effect

Since 1994, the so called TFH has been used [91], [92] to improve the HCI production by feeding the plasma with two electromagnetic waves at different frequencies instead of one. In some cases, even three or more close frequencies or a white noise generator (WNG) [83] have been used. The TFH demonstrated to be a powerful method: e.g., in the case in [91] for ^{238}U , it increased the production of higher charge states (from 35+ to 39+) by a factor from two to four and shifted the peak charge state from 33+ to 36+. Unfortunately, neither the relationship between the two frequencies nor the respective power was univocally determined. In fact, any source features a different set of parameters and the optimization is done empirically, just by looking to the maximization of beam current. Several qualitative explanations have been given about this phenomenon, all related to the increase of the average electron temperature T_e and the ionization rate by assuming that the crossing of two resonance surfaces helps the electron to gain more energy. This simple picture does not explain the reason of the rel-

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evant changes in the charge state distribution (CSD²) for different couples of frequencies (even for the case of minor changes, let us say few megahertz over 14 or 18 *GHz*), that can be explained nowadays by the frequency tuning effect. It is important to underline that, even in the case of TFH applied to many existing sources, a TWT is often used, the other being a klystron-based generator. The choice of TWT allows the experimentalists to slightly vary the second frequency: as an example, a strong variation of ECRIS performances has been observed in TFH operations in [92]. It was observed, for example for the production of O^{7+} , that 60 *W* emitted by the TWT at the optimum frequency gives the same effect as 300 *W* from the fixed-frequency klystron. The maximum current is obtained by means of klystron at 427 *W* and TWT at 62 *W* ($I = 66 \text{ e}\mu\text{A}$) operating at the same time; in order to obtain the same current, they needed $P_{RF} = 800 \text{ W}$ by the klystron in SFH. Furthermore, in the case of TFH, the current increases almost 20% (from 57 to 66 *eμA*) when the TWT emitted frequency shifts from 11.06 to 10.85 *GHz*, being constant the klystron and TWT emitted power. Then, the TFH is an effective method to increase the extracted current from ECRIS, but it can be fully exploited only by means of the frequency tuning. Several measurements have been carried out with the SERSE ion source. The TFH has been used either for operation at 14 or 18 *GHz*, with a clear advantage with respect to SFH. It was not successful when a 28 – *GHz* gyrotron and an 18 – *GHz* klystron were used at the same time [66], and that is an evidence that fixed-frequency generators are not adequate for TFH. The performances obtained by two-fixed-frequency klystron were also not so good as the ones obtained by a fixed-frequency generator plus a TWT.

As an example, for tin production, a current on the order of 2 to 3 *eμA* was measured for the charge state 29+ when a power of 1.4 *kW* from the 18 – *GHz* klystron and a power of 1.0 *kW* from the 14.5 *GHz* klystron were used. By keeping unchanged all the parameters and the 18 *GHz* klystron power, a TWT operating in the range of 8 to 18 *GHz* replaced the 14.5 *GHz* klystron. Even if the maximum power from the TWT was only 200 *W*, it permitted one to obtain more than 3 *eμA* of the charge state 29+ with a frequency of 17.600 *GHz*. Another interesting operating point was observed at the frequency around 17.000 *GHz*, which was less critical than any other setup and permitted one to obtain the state 30+ with the

²The CSD represent the distribution of the charge states in terms of the extracted current for each one. Usually for magnetic confinement plasmas it has a characteristic peaked profile, whereas for non-confined plasmas, like for MDIS, it has a "decay-like" trend.

same beam current, only by means of minor changes of the gas pressure. It must be underlined that the vacuum inside the plasma chamber improved by decreasing the total power from 2.4 to 1.6 *kW*, which permitted one to move the CSD to higher charge states and to have an excellent reproducibility, even after weeks. In order to have a comparison of these results with the performances in SFH mode, the Sn^{30+} current obtained with the same power from the 18 *GHz* klystron was only 0.7 *eμA*. All these data are explained by the description of the microwave-plasma interaction given here and they demonstrate the advantage of TFH operations, when the chosen frequency has the right electromagnetic field distribution in the plasma chamber. Moreover the closer frequencies are the best choice for SERSE. Not only for the couples 18 and 17.6 *GHz* or 18 and 17.0 *GHz* but also for operation around 14 *GHz*; it was observed that the optimum was obtained for 14.5 *GHz* klystron and TWTA at 14.3 *GHz*. This observation is also coherent with the unsuccessful test with a 28 *GHz* gyrotron and an 18 *GHz* klystron.

4.5 Magnetic field regime transition and "ultra-hot" electrons

Up to now the effect of the frequency on the sources performances has been investigated. Anyway, as mentioned elsewhere, a strong limitation to the trend towards higher charge states and currents for the 3rd generation ECRIS comes from the strong production of very hot electrons with energies on the order of *MeV* or higher. The hot electrons generate X-rays that heat the cryostat surrounding the superconducting magnet of the high field ECR ion source.

The first evidence of the production of *MeV* electrons came from the experiments carried out in 2000 with the SERSE source [66], devoted to the study of the scaling laws validity up to 28 *GHz*. Additional studies with the VENUS source at LBNL were carried out in 2004 and 2006 [7], [93]. The production of hot electrons inside the plasma, generating X-rays that are partially absorbed in the cold mass of the superconducting magnet of the high field ECR ion source, may heat the cryostat above the safety limits. For this reason, the exploitation of superconducting magnets ECRIS is subjected to the limitation of the heat load. Since 2004, the VENUS ECR ion source has been the most powerful tool to produce intense beams of highly charged heavy ions and quite soon the same phenomenon observed in [66] appeared. From the experimental results it seems that the production of such energetic

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electrons is strictly connected with the magnetic field profile. Hence a study was carried out on VENUS by varying B_{inj} (injection maximum field), B_{ext} (extraction maximum), B_{min} (minimum field on axis) and B_{hex} (the radial field) [40]. For each case, the charge state distributions (CSD) and the X-ray spectra were recorded and correlated with the source parameters. The collected data have been analyzed by taking into account that from measured bremsstrahlung spectra neither the plasma electron energy distributions nor the plasma densities can be directly determined, even though the spectra are related to both [12]. However, the amount of high energy X-rays is related to the amount of high energy electrons. Since X-rays up to MeV energy has been detected, it means that the heating mechanism can accelerate electrons up to these extremely high energies. A complete description of the experimental set-up and of the results of X-rays measurements carried out with the VENUS source is available in [40] and [12]. Hereinafter some of the results will be presented.

4.5.1 Experimental evidences

The presence of MeV electrons in ECRIS puts in evidence that the knowledge of ECR heating mechanism is incomplete, because the so-called adiabatic limit was exceeded by a substantial percentage of electrons. It will be shown hereinafter that the generation of these electrons is related to the magnetic field gradient, while the other source parameters, including the microwave power, affect the amount of X-rays of any energy, but not the distribution function.

According to the Stochastic Heating Theory the maximum value for the kinetic energy achievable for electrons after many crossing of the resonance is determined by the so-called absolute stochastic barrier W_b . It scales with the electric field at the resonance [37] and for the values of electric fields typical for VENUS ($|\vec{E}| \sim 3 \div 5 \cdot 10^5 \frac{V}{m}$ [12]) it is:

$$W_b \sim 250 \text{ keV} \quad (4.5.1)$$

Hence in the following we will define "hot" the electrons between 30 and 250 keV , and "ultra-hot" the ones above 250 keV . The analysis of the experimental results will be carried out by studying separately the dynamics of each population.

For all the experiments reported oxygen was used as mixing gas and the plasma was heated using 28 GHz only. About 3–5 kW of microwave power was used and the gas flow was kept constant while the magnetic field was changed.

4.5 Magnetic field regime transition and "ultra-hot" electrons

The influence of the B_{min} value on the formation of highly charged ions is shown in fig. 4.5.1.

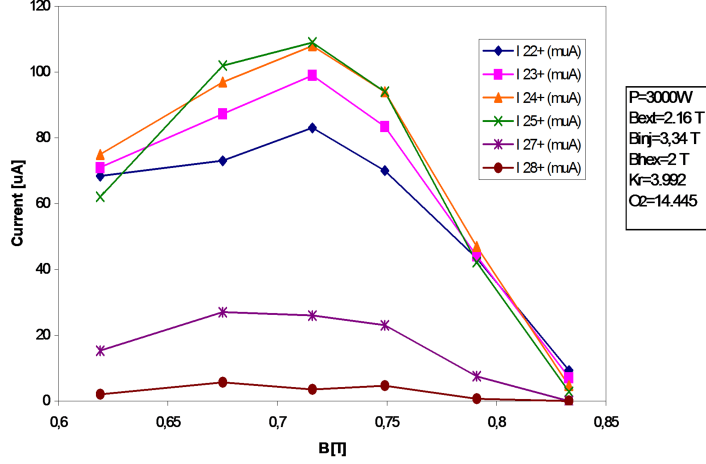


Figure 4.5.1: Trend of the Kr Charge States with respect to the B_{min} field.

It can be seen that there is a quite narrow range that gives higher current for $Q = 25+$ (but the same behaviour is observed for any charge state $Q \geq 22+$), i.e. $0.65 < B_{min} < 0.75 T$. Note that the current of each charge state is peaked at $B_{min} = 0.717 T$. As the total amount of power injected in the plasma chamber during the test for B_{min} scaling remains the same, we expect that this phenomenon is only related to the magnetic field gradient³. In addition we can observe in fig. 4.5.2 that the trend for the low energy electron population ($W < 100 keV$) is similar to that one of the high charge states. On the other way, the total amount of the X-rays at higher energies ($W > 200 keV$) is constant up to $0.75 T$ and increases about 20% for B_{min} increase to $0.85 T$ (i.e. the mirror ratio decreases). In particular, the amount of electrons above $800 keV$ increases of about one order of magnitude passing from $B_{min} = 0.72$ to $B_{min} = 0.85 T$. In the same interval the current of highly charged ion decreased abruptly. We know that the ionization cross sections are the highest for electron energies much lower than $100 keV$. Then the decrease of the number of low energy electrons ($W < 100 keV$) and the corresponding increase of high energy

³The magnetic gradient is strictly connected with the mirror ratio: the higher is the mirror ratio, the higher will be the gradient in the resonance zone, as the distance between the magnets remains the same.

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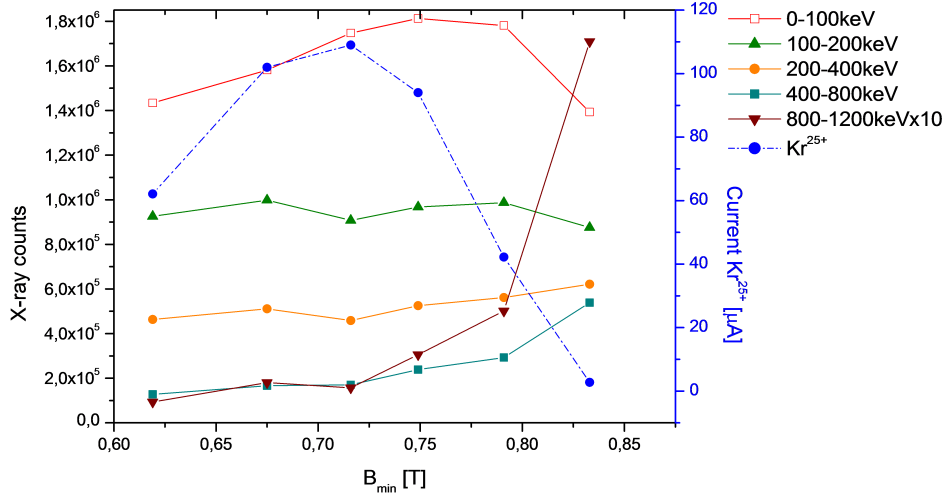


Figure 4.5.2: Axially emitted X-rays integrated for different energy ranges and Kr25+ current vs the B_{min} field.

ones is detrimental for high performances. From analysis of fig. 4.5.2 comes out that the "ultra-hot" electrons are produced preferentially in the case of high B_{min} , then in the case of low mirror ratio.

Similarly the changes of the B_{ext} field have been considered for the same power (3 kW) [12], showing that they may probably regulate the transition of the electrons among the two principal electron populations (warm and hot). The changes of CSD measured for variable B_{min} and B_{ext} presented similar characteristics. The amount of highly charged ions decreased with the increase of the ultra-hot electron production that was strongly influenced by the magnetic field gradient. In particular, if the gradient decreases (i.e. B_{min} increases or B_{ext} decreases), the amount of counts in the X-ray spectra high energy tail increases as a consequence of the enhancement of the hot and ultra-hot electron production. Figure 4.5.3 features the strong changes in X-ray spectra occurring as the B_{ext} decreases. The hot electron tail grows up considerably for the lower values of B_{ext} .

Two different regions can be individuated in these spectra: $400 < W_k < 800$ keV and $W_k > 1$ MeV, where W_k is the electron kinetic energy. In these two regions the spectra present a quite linear trend with respect to the energy in the logarithmic scale. An effective electron temperature can be calculated and it is inversely proportional to the EEDF slope. This

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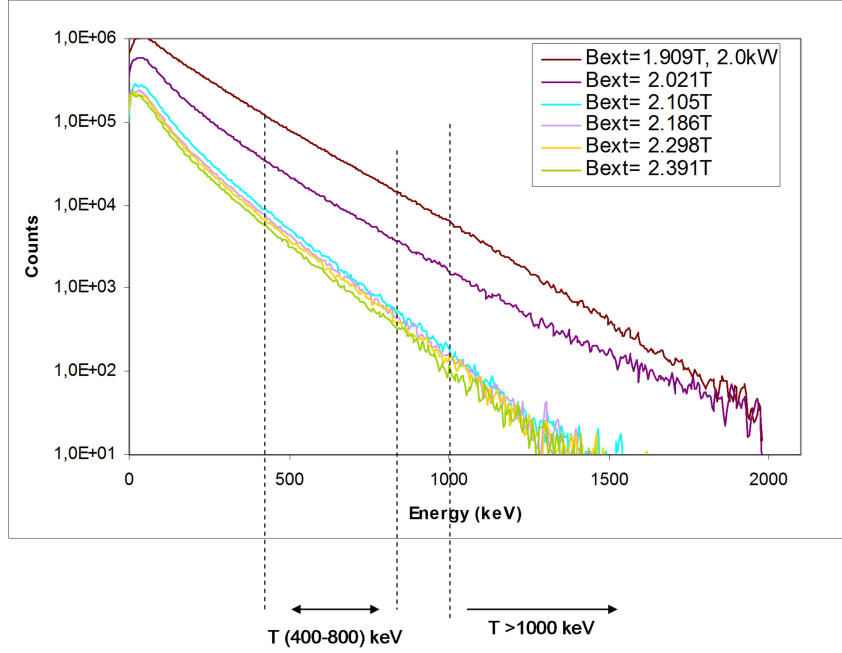


Figure 4.5.3: Axially emitted X-ray spectra for different values of the B_{ext} field.

"spectral temperature", that is only a signature of the high energy electron component distribution function, gives some interesting information about the heating mechanisms. Note in figure that the regions evidenced have two different temperatures that are reported in figure 4.5.4.

A marked "jump" occurs when B_{ext} passes from 2.105 to 2.021 T. Especially the temperature of the highest energy component grows up to about 245 keV, with an increment of more than 100 keV. Also the temperature of lower energy component increase, as well as the maximum energy. When B_{ext} decreases below 2 T the temperature of the 400 – 800 keV component remains almost unchanged, whereas the other one strongly decreases. Note that this behavior may be connected to the excessive decrease of the extraction field, which leads to the violation of the High-B mode requirements. For excessively low magnetic fields the confinement worsening decreases the electron lifetime and high energies cannot be reached anymore, except for a low amount of particles.

No direct information about the temperature of the warm electron component are available from the experiment. Anyway, according to the dis-

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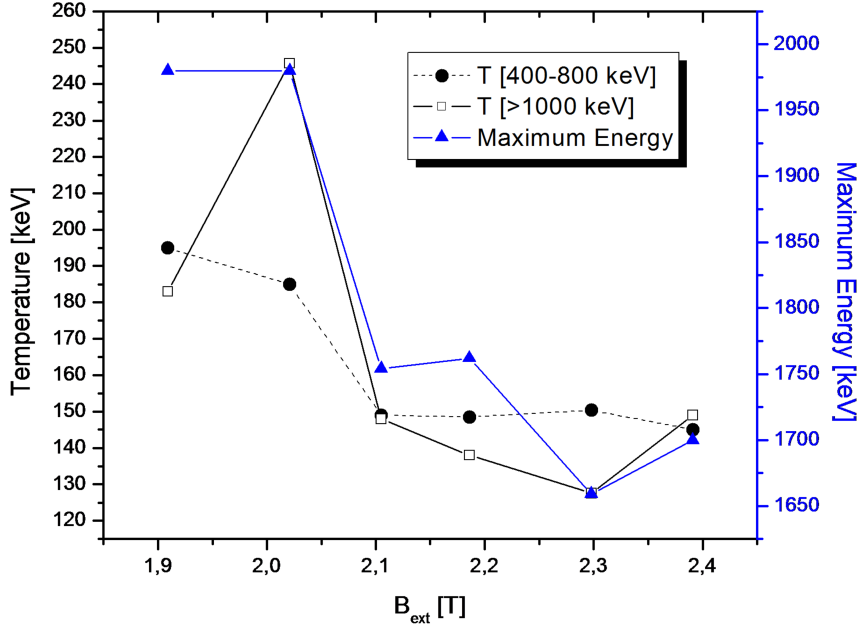


Figure 4.5.4: Spectral temperature calculated from the slope of the spectra shown in figure 4.5.3. The temperatures of the two considered populations have been plotted, together with the maximum energy obtained by the point where the spectrum intercepts the x axis.

cussion of section 1.2 it is possible to guess some indirect estimations of the electron temperature by analyzing the CSD. The warm electrons temperature corresponds approximately to $3 \div 5I_z$, where I_z is the ionization potential of the charge state on which the CSD is peaked. Without any quantitative consideration about T_e^w , we can investigate its trend observing how the CSD peaks moves from one magnetic configuration to another. When $0.67 < B_{min} < 0.75$ ($B_{min} = 0.717 T$ in our case) the CSD is peaked on 25+, while for higher or lower B_{min} values the CSD is peaked at slightly lower CS (Charge State). There is probably a fixed value of B_{min} for which the warm electron population reaches a maximum temperature. Above the optimum value the highest charge states are depressed. A similar behavior has been observed also for B_{ext} . Therefore it is clear from the experimental data that the heating mechanisms which lead to higher hot component temperatures subtract energy to the warm electrons. This leads to the decrease

of $\langle q \rangle$ and then to the worsening of the source performances.

It was observed that axial X-Rays spectra do not depend significantly on the B_{inj} variation, probably because the values of B_{inj} used for ECRIS are so high (more than three times the resonance field) that a variation of 10 to 20% does not change the field profile appreciably. The hot electron temperature and percentage seem to remain about the same for different values of B_{inj} . Only a slight decrease of the total amount of axial X-rays was observed (about 10%) for a B_{inj} decrease of 15% [12]. In terms of the CSD variation it has been observed a less pronounced peaked trend if compared with B_{min} and B_{ext} scalings.

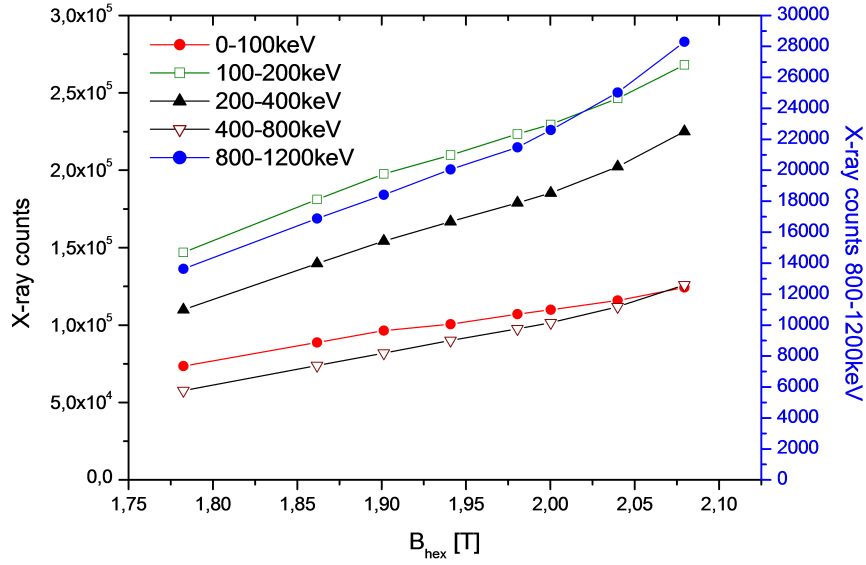


Figure 4.5.5: Axially emitted X-rays integrated for different energy ranges with respect to the hexapolar field.

In fig. 4.5.5 the X-ray measurements on-axis are shown for different hexapolar field values. The total amount of X-rays rises by a factor two for a change of B_{hex} of 15% only. Note that in this case the number of X-rays of all the different populations increases as the mirror ratio $\frac{B_{hex}}{B_{min}}$ increases. In this case the trend is practically inverse with respect to the other scalings. We may assume that the increase of the X-ray counts as B_{hex} increases can

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be really linked to the enhancement of the electron density: in fact the total flux of radial X-rays (determined by analyzing the dose of X-rays detected transversally with respect to the chamber axis) remains constant for different values of B_{hex} over a large range. The change in the spectra acquired along the axis is not much different for the different population, thus the variation of the absolute value indicates that B_{hex} determines mostly the electron density but not the energy distribution. The hypothesis of a substantial increase of the electron density with B_{hex} is coherent with the shift of CSD towards higher charge states. The current for Kr^{27+} increases from 9 to 26 μA when the hexapolar field is increased from 1.86 to 2.13 T and P_{RF} is kept constant (3 kW), while Kr^{25+} increases from 53 to 109 μA . We may conclude that the hexapolar field features a more "classical" behavior, as the CSD improves without any peaked trend like in case of the other fields scalings. This is probably due to the fact that the hexapolar field operates near the High-B mode requirements boundaries. Hence the improvement of the electron confinement due to the suppression of any MHD instability plays the major role to determine the plasma dynamics.

The experimental results pointed out that the microwave power P_{RF} has a limited influence on the electron energy distribution while its effect is remarkable on the total amount of warm and hot electrons [40]. In fact the total number of the detected X-rays increased about linearly with P_{RF} , as shown in figure 4.5.6.

The hot electron temperature depends weakly on the increasing power. A larger number of warm electrons is evident. In fact an increase of highly charged ions production with P_{RF} occurs because of a higher number of ionizing collisions, as shown in figure 4.5.7. The experimental results point out that the microwave power has a limited influence on both the electron temperatures (warm and hot) while its effect seems to be remarkable on the X-Ray rate, i.e. finally on the amount of heated electrons. In fact the fig. 4.5.7 points out an evident increase in the production of the higher charge states, but the CSD remains peaked to the same charge state.

It is known that the mean charge state is related to the electron density by the relation:

$$\langle q \rangle = n_e \tau_i \quad (4.5.2)$$

so we can conclude that the microwave power strongly affects the rate of the electron heating, creating a great amount of warm and hot electrons with a remarkable increase of n_e ; At higher powers the amount of hot electrons considerably increases but the warm and hot electron temperature grow in a less marked way. The increased amount of warm electrons, that are

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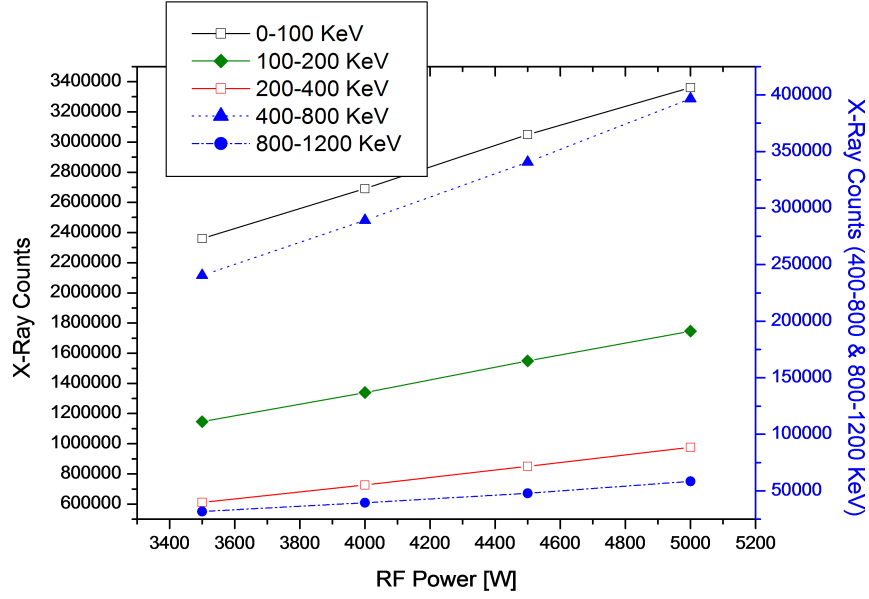


Figure 4.5.6: Axially emitted X-Rays for different energy ranges with respect to the microwave power.

created very quickly during the ECR-H process, can significantly reduce the ionization time so that it becomes much lower than the recombination time due to the charge exchange. In this way we should explain the increase of the HCI shown in fig. 4.5.7 (e.g. Kr^{27+} is $35 \mu A$ at $3.5 kW$ and $50 \mu A$ at $5 kW$).

Other experiments demonstrate a saturation of some sources parameters with the power. It seems, in our case, that a saturation for the electron temperature occurs already for a microwave power $P_{RF} = 3 kW$. Only by acting on the magnetic field gradient (as said above) we can further increase the temperature. It is remarkable that these phenomena can be anticipated by simulating the electron motion in a ECR plasma by means of a single particle approach (see section 4.8). It is possible to show that the mean energy of the electrons has a saturation value probably due to the dependence of the ECR zone crossing time on the electron energy. When the electron energy increases the crossing time decreases and the heating processes becomes slower, if we keep in mind that the acceleration time strongly depend on the crossing time.

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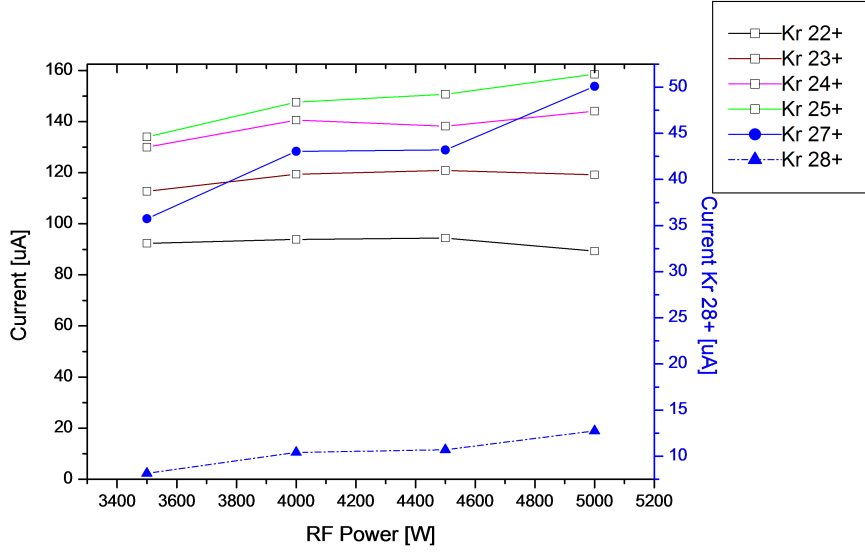


Figure 4.5.7: Trend of different *Kr* charge states at different microwave powers.

Another proof that the "ultra-hot" electron temperature is principally determined by the magnetic field structure comes from fig. 4.5.8. This figure features that for the higher B_{min} field ($B_{min} = 0.8\text{ T}$ and $P_{RF} = 3\text{ kW}$) the very hot electrons ($W > 400\text{ keV}$) are much more than ones created at $P = 5\text{ kW}$ and $B_{min} = 0.717\text{ T}$. Hence a high B_{min} , with a low $\vec{\nabla}B$, removes a conspicuous number of electrons from the warm population and shift them in the hot population.

On the contrary, high microwave power levels remove electrons from the cold population and shift them towards the warm population, by strongly increasing the density of electrons candidates for the ionization processes.

If we increase the plasma chamber pressure, by increasing the amount of the injected *Kr*, we observe (fig. 4.5.9, red curve) a strong charge exchange which decreases the amount of HCl; then the CSD peak shifts towards lower CS and the higher CS are no longer produced. The results shown in figure have been obtained with an RF power of 5000 W . Even with this high level of power the time needed for the step by step ionization exceeds the charge exchange time. No influence on the electron temperature and on the EDF have been observed, as confirmed by figure 4.5.10. However, a weak decrease in the production of the "ultra-hot" electrons is evident from the figure.

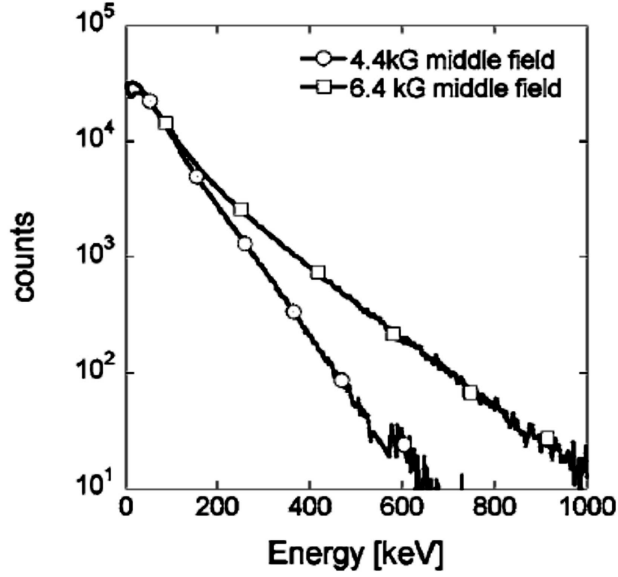


Figure 4.5.8: The axial x-ray spectra for a B_{min} of 0.44 T (8.8 kW power) in comparison to a B_{min} field of 0.67 T (6 kW power).

The lower energy ranges seem to be unchanged by the pressure variation; this is a sign that the heating process has not been strongly influenced by the increased pressure, i.e. the low and medium energy electrons do not lose more energy because of an increased number of collisions in an appreciable manner. Only the particles at very high energy are influenced by the pressure. Their cross section for e-i and e-n collisions (see section 1.2) is very low and they can be considered collisionless for very long times. However, when the pressure increases, the mean time between a collision and the following one becomes shorter than the time required to reach high energies.

In conclusion the pressure has a moderate impact over the EEDF, but it is the major constraint on the charge states distribution.

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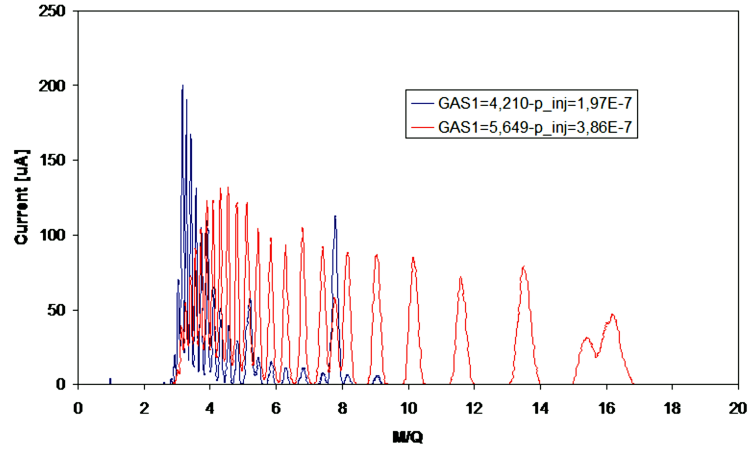


Figure 4.5.9: Change of the CSD for two different values of the background pressure.

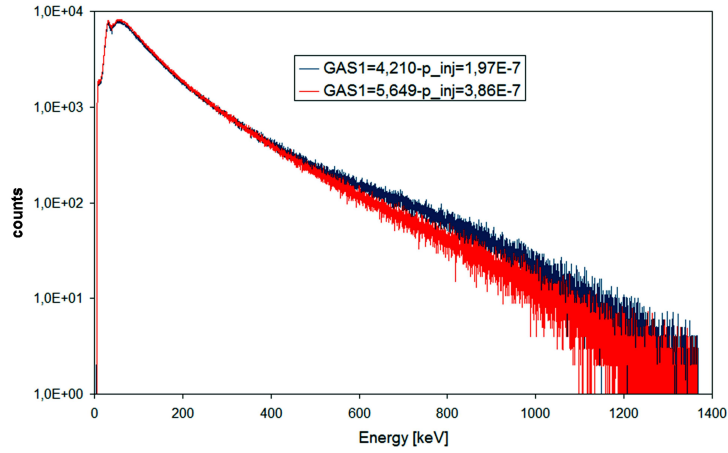


Figure 4.5.10: X-ray spectra for two different values of pressure ($1.97 \cdot 10^{-7}$ and $3.86 \cdot 10^{-7}$ mbar) at 5000 W.

4.6 Theoretical investigation: the single particle (SP) approach

The several results obtained more or less recently with different techniques of plasma heating (frequency tuning, two frequency heating, the use of broad-band generators, etc.) remained without a clear explanation for a long time. In fact the processes involved the plasma-waves interaction are outside from the ECRIS standard model, which considers mainly the global aspects of the ion sources. Hence our group started in 2005 a series of detailed studies concerning the Electron Cyclotron Resonance heating, and some preliminary results were presented already in the degree thesis at the end of 2005 [25].

4.6.1 Single particle and collective aspects of ECRIS plasmas

The single particle approach, formalized by Lieberman and Lichtenberg and reported in section 1.11, is justified if the electric field associated to the thermal fluctuations E_{th} is negligible as compared to the incident electromagnetic wave amplitude E_w , i.e. the condition $E_w \gg E_{th}$ has to be satisfied. E_{th} (expressed in $\frac{V}{cm}$) can be evaluated as it follows [17]:

$$E_{th} = 10^{-3} (n_e T_e)^{\frac{1}{2}} \quad (4.6.1)$$

where n_e is the electron density (in cm^{-3}) and T_e is the electron temperature (in eV). For example, typical values of density and temperature for the SERSE source are: $n_e = 3 \text{ to } 4 \cdot 10^{12} \text{ cm}^{-3}$, $T_e = 5 \text{ to } 10 \text{ keV}$. Then the single particle approach is valid as long as $E_w \gg 10^5 \frac{V}{m}$. It is important to point out that this model has permitted to evaluate the upper limit of the electron energy that can be achieved during the multiple passages across the ECR surface, the so-called "adiabatic barrier". Lieberman and Lichtenberg proposed that the electron, which is escaping from the resonance surface, forgets his phase relationship with respect to the microwave field before it reaches again the resonance zone in another part of the surface. This process goes on until the adiabatic invariant of the electron motion begin to play its role: the electron acceleration stops because its phase relationship cannot be further "randomized". Although many experimental results confirm the assumptions of this model, many others demonstrate that some phenomena, which occur in ECRIS plasma, remain out of this scheme:

1. The existence of several electron populations at different temperature requires several additional hypothesis to be explained, and it does not

4. MICROWAVE COUPLING TO ECRIS PLASMAS

comes out so easily from the model;

2. It does not explain the spatial fluctuations of the electron temperature in ECRIS plasmas (e.g. hot corona and warm plasma core) [17];
3. The existence of "ultra-hot" electrons is not coherent with this model if some additional phenomena are not taken into account. Many electrons having energies higher than the adiabatic barrier and have been observed in SERSE and in VENUS ion sources [66], [12].

Evidently the exhaustive comprehension of the ECRIS plasma dynamics has to take into account even collective aspects of plasmas and the non-linear phenomena. During the plasma start-up the ECRIS plasma density is quite low so that the single particle approximation remains valid. It will be still valid until the plasma density reaches the typical level of the steady state. At this moment the electrons feel not only the incident electromagnetic field, but rather a compound field, which is composed by a superposition of the microwave field and of the electrostatic fields due to the plasma fluctuations or waves. This superposition makes possible an exchange of energy between the electromagnetic and electrostatic plasma waves [17]. This mechanism of energy exchange was proposed by Golovanivsky as the one able to heat the plasma core and to explain the formation of hot electron layers in ECRIS plasmas. More generally, it can explain the occurrence of ultra-hot electrons acceleration within the layer where Bernstein or Langmuir waves are excited [94]. In this work we limit any consideration to a single particle framework, according to the developments of the stochastic heating theory provided by Canobbio [39]. The thickness of the ECR zone, according to the relativistic effects involved in the electron acceleration, is a key parameter and it has been a guideline in the interpretation of simulation results.

4.6.2 Assumptions for the simulation of the wave-particles interaction

As discussed in section 1.4 the study of the wave propagation in plasmas is complicated when a magnetostatic field is applied to the plasma and the most exhaustive theory of plasma-wave interaction is the Vlasov one. This picture is further complicated by the mode conversion (electromagnetic waves can be converted in electrostatic waves and vice versa) as discussed before. During the last three years a simplified approach has been followed at INFN-LNS in order to describe the electromagnetic field propagation inside the plasma chamber of the SERSE ion source. The main basis of this

approach has been the calculation of the resonant modes of the cavity where the plasma is created [90]. Then the plasma chamber has been considered as a resonant structure, according to the experimental results presented in [85] and especially in section 4.3. This is an innovative approach to the study of the plasma formation in ECR ion source, as it considers the possibility that stationary waves can be formed in the plasma chamber during the plasma creation. The calculations have shown that in the case of isotropic plasmas (i.e. plasma without magnetostatic field), the resonance frequency of each mode is simply shifted to higher values, the shift being proportional to the plasma density. However no analytic calculations are possible in the case of magnetized plasmas, if the magnetic field is not spatially constant. Unfortunately neither the microwave simulators (as the CST Microwave Studio or the HFSS codes [95], [96]) nor other simulations tools are able to give any prediction about the electromagnetic field distribution in the case of magnetized plasmas, and then we will consider modes patterns corresponding to empty cavities, by looking only to relative comparisons between the plasma dynamics in case of different modes structures.

4.6.3 Calculation of the electromagnetic field existing inside the cavity

The ECR plasma is created inside a stainless steel or an aluminium made resonant chamber. In the case of SERSE a cylindrical cavity with aluminium walls has been simulated: its radius is 6.5 cm and its length is 45 cm. The microwaves, traveling in WR62 waveguides, are injected through two different ports located in the injection flange (see fig. 4.6.1), which are designed to avoid the coupling between the two ports (one is rotated by 90° with respect to the other). The electromagnetic field existing inside the plasma chamber can be calculated according to the equations shown in [13]. The maximum amplitude A_n associated to the wave's fields can be calculated thanks to the quality factor of the plasma chamber Q , that can be defined as follows:

$$Q = 2\pi \frac{\langle W \rangle}{\langle W_P \rangle} = \omega \frac{\langle W \rangle}{\langle P_P \rangle} \quad (4.6.2)$$

where $\langle W \rangle$ is the total energy contained into the cavity and averaged over a wave period; $\langle W_P \rangle$ is the energy lost averaged over a wave period and $\langle P_P \rangle$ is the power lost into the resonator averaged over a period. The Q is usually very high and for a completely closed and in vacuum chamber it can be on the order of 10^4 . In addition, for definite geometries, it can be calculated analytically [13].

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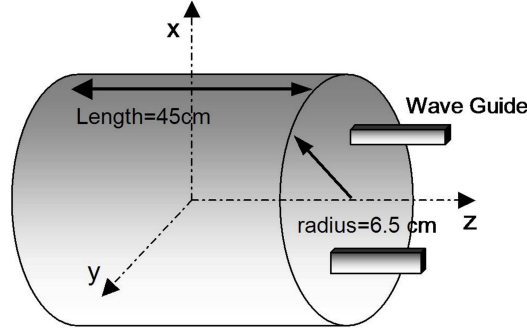


Figure 4.6.1: Sketch of the SERSE plasma chamber together with a schematic representation of the waveguides used for the microwave injection.

Values of the Q factor were determined for the cavity of SERSE, and the results are reported in [90]. Q values on the order of 10^5 have been determined in case of vacuum chamber and aluminum walls ⁴. Since we have:

$$W = \frac{PQ}{\omega} \quad (4.6.3)$$

where P is the total power injected into the cavity and W the total energy contained into the cavity and given by the equation [13]:

$$W = \frac{1}{2} \mu \int_V \vec{H} \cdot \vec{H}^* d^3r = \mu A_n^2 \frac{\pi l a^2}{8 h_{n\nu}^2} \left(\frac{r^2 \pi^2}{l^2} + h_{n\nu}^2 \right) \left(1 - \frac{n^2}{h_{n\nu}^2 a^2} \right) J_n^2(h_{n\nu} a) \quad (4.6.4)$$

we obtain: $A_n \simeq 2.62 \cdot 10^3 \frac{A}{m}$ if the total RF power injected into the cavity is 250 W; $A_n = 5 \cdot 10^3$ if $P_{RF} = 1000 \text{ Watt}$ ⁵.

Note that the treatment followed up to now does not take into account of losses due to the plasma and/or waveguide.

4.6.4 The simulation of the magnetic field structure

As regards the simulation of the magnetostatic field for the plasma confinement, an analytical approximation of the SERSE magnetic field has been

⁴Note that Q is a dimensionless parameter.

⁵The RF power values reported here are reasonable for many of ECRIS, especially for the SERSE source

appointed.

The central region of the axial magnetic field can be modeled by the following equation:

$$B_z = B_0 + B_1 z^2 \quad (4.6.5)$$

where z is the coordinate along the plasma chamber axis, B_0 and B_1 are two constants determined according to the minimum value of the magnetic field for B_0 and according to the intensity of the magnetic field gradient for the B_1 parameter. The equation 4.6.5 is practically the (1.6.1). The components along x , y and z of the simple mirror field can be approximated by the following expressions:

$$\begin{aligned} B_x &= -B_1 x z \\ B_y &= -B_1 y z \\ B_z &= B_0 + B_1 z^2 \end{aligned} \quad (4.6.6)$$

Once inferred the field equation along the z direction the other components have been derived by imposing the $\vec{\nabla} \cdot \vec{B} = 0$ condition. The minimum-B magnetic configuration has been simulated by adding the hexapolar field components. Summing up we have:

$$\begin{aligned} B_x &= -B_1 x z + 2Sxy \\ B_y &= -B_1 y z + S(x^2 - y^2) \\ B_z &= B_0 + B_1 z^2 \end{aligned} \quad (4.6.7)$$

where S is a constant which indicates the rapidity (then the gradient) of radial field increase. The equations (4.6.7) satisfy the $\vec{\nabla} \cdot \vec{B} = 0$; the parameters B_0 , B_1 and S can be determined according to the real field of the SERSE source and considering the plasma chamber dimensions. Considering that the maximum field at the injection is 2.7 T , and that the maximum radial field is 1.55 T we obtain:

$$B_0 = 0.4 \text{ [T]}; \quad B_1 = 25.55 \left[\frac{\text{T}}{\text{m}^2} \right]; \quad S = 310 \left[\frac{\text{T}}{\text{m}^2} \right] \quad (4.6.8)$$

The figures 4.6.2 and 4.6.3 feature a three dimensional view of the magnetic field, showing that it increases going from the center to the periphery of the plasma chamber, both radially and axially. In particular a three cusps structure is evident in proximity of the injection side of the plasma chamber (the same structure is present at the extraction end).

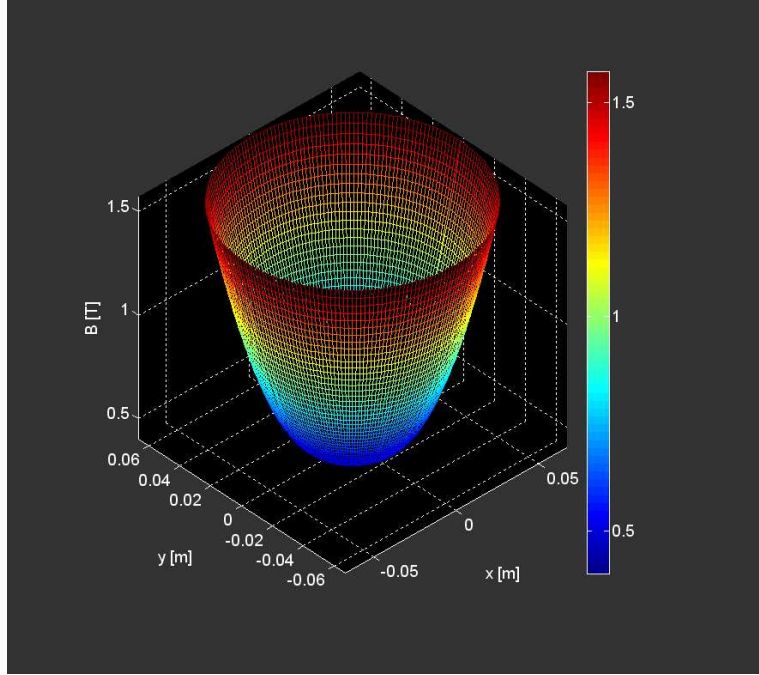


Figure 4.6.2: Trend of the radial magnetic field of the SERSE source on the transversal plane with respect to the plasma chamber axis. The hollow shape is evident (magnetic trap), which allows to stabilize the plasma confinement also in the radial direction.

The magnetic field plotted in fig. 4.6.4 is close to the real field. However the simulated field has been considered symmetric with respect to the midplane⁶.

4.6.5 Model assumptions summary

For our simulations we used the basic theory that describes electromagnetic fields inside cylindrical cavities, because the simulations show that the ECRH process is only slightly biased by the superposition of modes (the energy increase changes are less than 5%). However a more detailed field calculation is reported in [4]. In order to understand how different patterns of electric field in the plasma chamber can affect the electron heating, the evolution of the electron energy is followed through many crossing of the

⁶This simplification of the axial field was done mainly to reduce the calculation times.

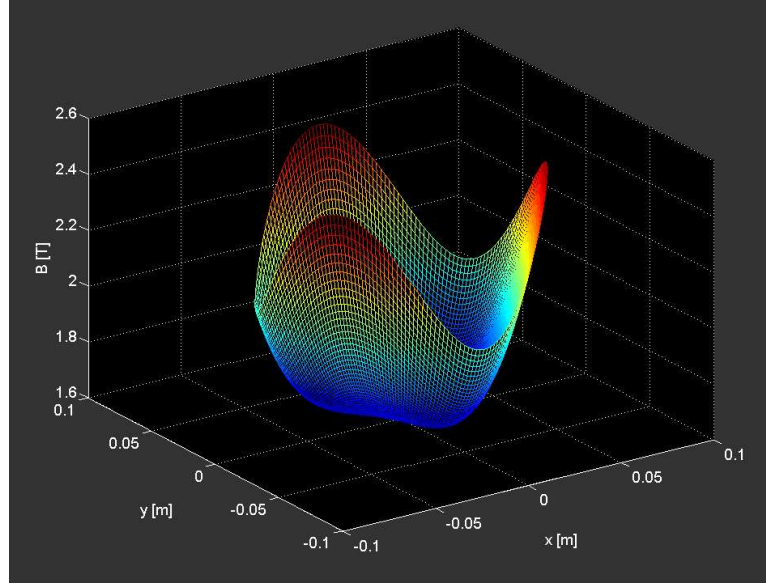


Figure 4.6.3: Trend of the radial magnetic field of the SERSE source on the transversal plane with respect to the plasma chamber axis and at the injection side ($z = -0.225$ cm). The three cusps structure at the plasma chamber external sides is evident.

resonance, in SFH or TFH mode. The model, based on the single particle approach, cannot give quantitative predictions on the final energy of the electrons subject to different electric field shape, because of its oversimplification; on the contrary we will obtain some comparative results, which are able to inform about the most effective mode to heat the plasma. According with the assumption done in section 1.10, the single particle approach is able to describe the plasma start-up as long as the electron density does not reach the cutoff. In addition, we said that the wave's electric field amplitude on the resonance surface has to be much higher than the plasma's self-generated electric fields ($E_w \gg 10^5 \frac{V}{cm}$). Calculations show that E_w in a SERSE-like ion source may be in the order of $10^5 \frac{V}{cm}$ and this means that the conditions of single particle approach are not strictly satisfied, even if we expect that the trend of the electron heating is reasonably determined. As ECR plasmas can be considered collisionless, then the collisions have been omitted in the equation of motion, provided that the simulation lasts a time so low to avoid the collisions effects ($t_s < 50 - 100$ nsec as the Spitzer

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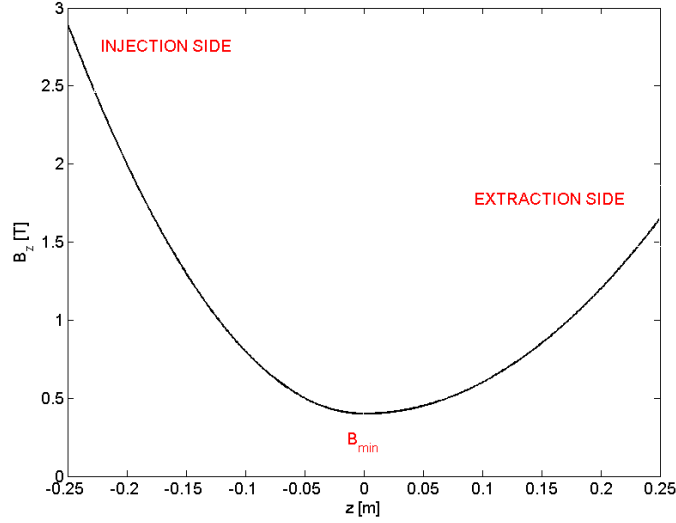


Figure 4.6.4: Trend of the axial magnetic field of the SERSE source.

and the ionizing collision frequencies lie in the range of 10^6 to 10^7 sec^{-1}). Hence we can simulate the energy absorption process of the electrons but not the following build-up of the different ion charge states, which is beyond the scope of the present work. We considered the modes that have been analyzed in [90] (in particular the modes closer to 14 and 18 GHz : the $TE_{4,4,23}$ around 14 GHz and the $TE_{4,6,22}$ around 18 GHz). The electrons are in thermal equilibrium at the starting time $t = 0$ (i.e. they follow a Maxwellian energy distribution), with mean temperature of 500 eV , i.e. we simulate the heating process from cold electrons (with $T_e \simeq \text{few eV}$) to warm electrons (T_e some keV). The electrons start from random positions located inside the resonance surface, as shown in fig. 4.6.5 (a) and (b).

4.6.6 Solution of the single particle equation of motion

Once defined the model, the simulations have been done by using MATLAB [24]. Basically the software solves the relativistic equation of motion:

$$m_0 \frac{d}{dt} \gamma(v) \vec{v} = q \vec{v} \times \vec{B}_S + q \vec{E}_{em} + q \vec{v} \times \mu \vec{H}_{em} \quad (4.6.9)$$

for the electrons, giving as output the information about the final en-

4.6 Theoretical investigation: the single particle (SP) approach

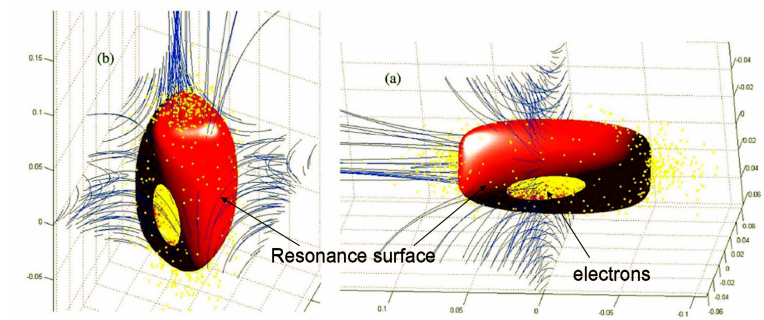


Figure 4.6.5: Location of the resonance surface (in red), of the plasma electrons (yellow) and of the magnetic field lines (in blue) as determined in our simulations. The two figures feature two different view of the system.

ergy and confined fraction of the electrons. In the equation (6) $\gamma(v)$ is the relativistic factor, which depends on the electron velocity v , $\vec{B}_S$ is the non-uniform magnetostatic field for the electron confinement, $\vec{E}_{em}$ and $\vec{H}_{em}$ are the electric and magnetic field associated to the electromagnetic waves and μ is the magnetic permittivity. Although the simulations cannot give yet any kind of information about the plasma evolution in a timescale comparable with a collisional regime, they give interesting information about the trend of the heating; this allows to evaluate how much time the electron takes to leave the cold population to the warm one, a crucial parameter that we may name ECR-heating rate v_h . This quantity can be defined as the energy acquired by the electron per nanosecond during the ECR process. The lower is the time needed to the electrons to reach high energies, the higher is the ionization efficiency, as the atoms can be ionized up to high charge states in shorter time and the recombination and charge-exchange effects are less important. This particular request is deemed to become more and more important for the increasing charge states; it was almost negligible in the past because charge states up to 20^+ were considered. Hence in the next pages we will analyze the simulation results in terms of v_h .

In order to solve numerically the equation (4.6.9), the three component along the cartesian coordinates must be determined. Actually the MATLAB solver of ordinary differential equations is able to solve only equations of first order by means of a Runge-Kutta algorithm. The three differential equations of second order obtained by projecting the (4.6.9) on the three cartesian coordinates can be reduced to a system of six equations of the first

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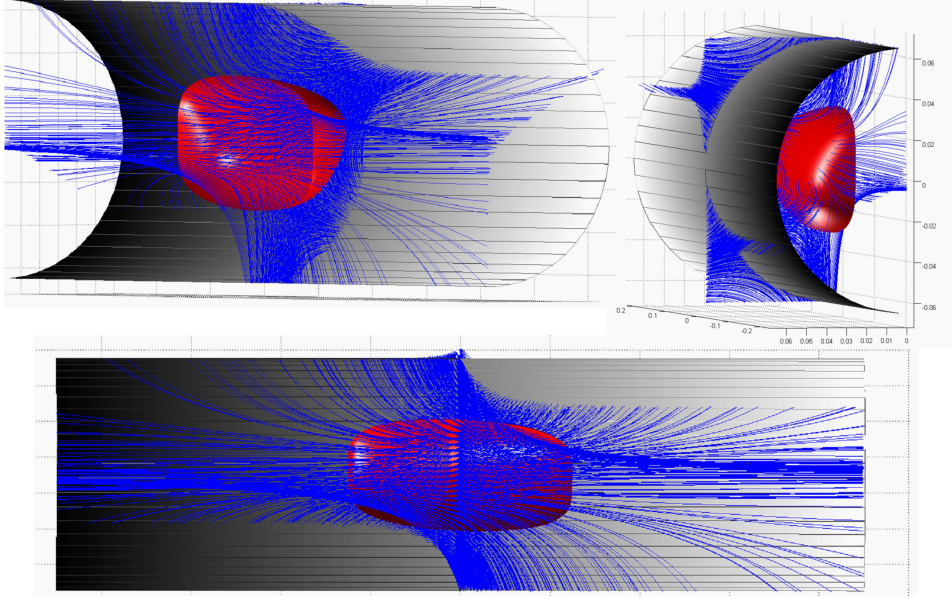


Figure 4.6.6: Sections of the plasma chamber of the SERSE ion source with the resonance surface (in red) and the magnetic field lines (in blue). The regions where the field lines intersect the plasma chamber are clearly shown, and they are of primary importance for the determination of the ion dynamics, as it will be shown in section 4.13.8.

order:

$$\begin{aligned}
 \dot{x} &= v_x \\
 \dot{y} &= v_y \\
 \dot{z} &= v_z \\
 \dot{v}_x &= F(v)[(v_y B_z - v_z B_y) + (v_y B_{em_z} - v_z B_{em_y}) + E_{em_x} - \\
 &\quad - \frac{1}{c^2} (E_{em_x} v_x + E_{em_y} v_y) v_x] \\
 \dot{v}_y &= F(v)[(v_z B_x - v_x B_z) + (v_z B_{em_x} - v_x B_{em_z}) + E_{em_y} - \\
 &\quad - \frac{1}{c^2} (E_{em_x} v_x + E_{em_y} v_y) v_y] \\
 \dot{v}_z &= F(v)[-B_x v_y + v_x B_y - B_{em_x} v_y + v_x B_{em_y} - \\
 &\quad - \frac{1}{c^2} (E_{em_x} v_x + E_{em_y} v_y) v_z]
 \end{aligned} \tag{4.6.10}$$

where:

$$F(v) = \left(\frac{q}{m} \sqrt{1 - \frac{|v|^2}{c^2}} \right) \quad (4.6.11)$$

depends on the relativistic factor γ . As the typical Larmor frequency of the electrons moving in the SERSE magnetic field is on the order of 10 *GHz* then a time step of 10^{-12} *sec* has been chosen to solve the equations (4.6.10). This time step permits to calculate about 100 points per single Larmor orbit.

4.7 Explanation of the frequency tuning effect

A series of simulations were carried out in order to verify the effect of different microwave frequencies on the electron dynamics. It has been demonstrated in section 4.2 that slight variations of the injection frequency strongly change the performance of the ECRIS both in terms of extracted currents and charge states. The experiment described in section 4.3 demonstrate that standing waves, i.e. resonant modes, persist also in case of dense plasmas filling the cavity. Then we would like to investigate, more than the global distribution of the field inside the chamber, the mode pattern over the resonance surface, in order to understand how the field structure affects the electron heating. The resonance surface can be calculated numerically once known the magnetic field structure, as it corresponds to that surface with constant field equals to $\frac{m\omega_R F}{q}$. The figure 4.7.1 features the resonance surface of the SERSE source calculated numerically by means of MATLAB⁷. Note that by varying the frequency of about 50 *MHz* the field distribution over the surface changes remarkably. In particular if one consider the zones of maxima the electric field is more than 10^2 higher in case of figure 4.7.1 (b) than in case of figure 4.7.1 (a).

The excitation of near modes, even with distances of few *MHz*, strongly changes not only the maximum field over the surface, but also the distribution of minima and maxima zones. Simulations carried out considering different modes have demonstrated that some of them are able to heat the electrons up to energies much higher than others (even ten times higher). The figure 4.7.2 features the strong fluctuations of the energy gained by the electrons after 50 *nsec*, i.e. the v_h parameter defined before. The simulations have been done in proximity of 14 *GHz*. The maximum energy after 50 *nsec*

⁷The subroutine *isovolume* has been used to plot the resonance surface once known the global B field structure

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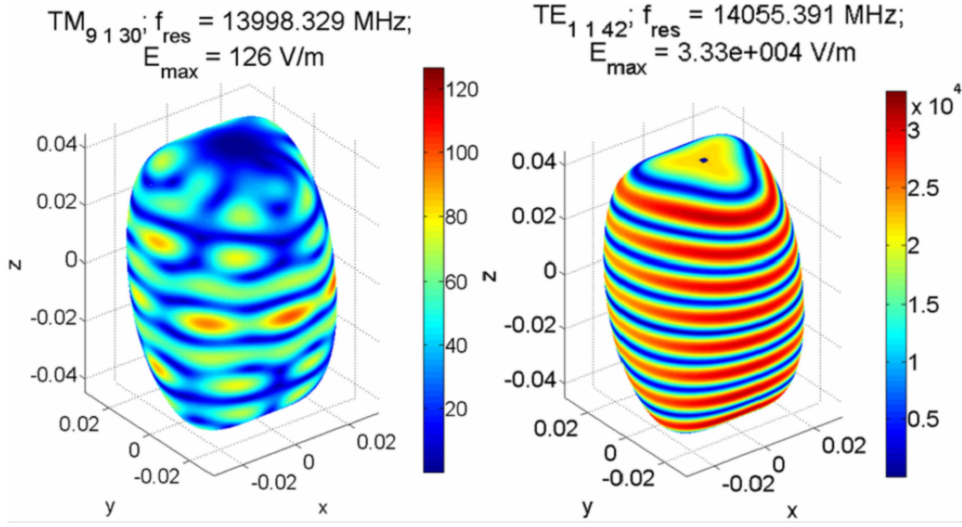


Figure 4.7.1: Simulated resonance surface of the SERSE source with the distribution of the electromagnetic field (false color representation).

has been reached when the $TE_{1,1,42}$ mode has been used for plasma heating, whereas the minimum energy has been reached for the mode $TE_{9,1,30}$. Note that these two extreme cases corresponds to the field distribution shown in figure 4.7.1. In our study the net power flowing into the chamber is kept constant, and therefore the observed effects are only due to the different field patterns.

In order to explain the simulation results shown in figure 4.7.2 we have to take into account that the great amount of energy transfer from waves to electrons occurs on the resonance surface because of the ECR condition. On the other side the electrons gyrate around the field lines with the Larmor frequency, with their guiding center moving along the magnetic lines. Then the structure of the magnetic field determines the the crossing points of the electron trajectories with the resonance surface. The figure 4.7.3 points out that only in particular zones of the resonance surface the energy transfer occurs, and the amount of the energy gained by the electrons strongly depends on the mode structure, as only in case of high field zones in proximity of the trajectory crossing points it is possible for the electron to gain very high energies.

Note that the higher is the heating rapidity, the higher is the replacement of the electrons lost from the confinement because of the Spitzer collisions.

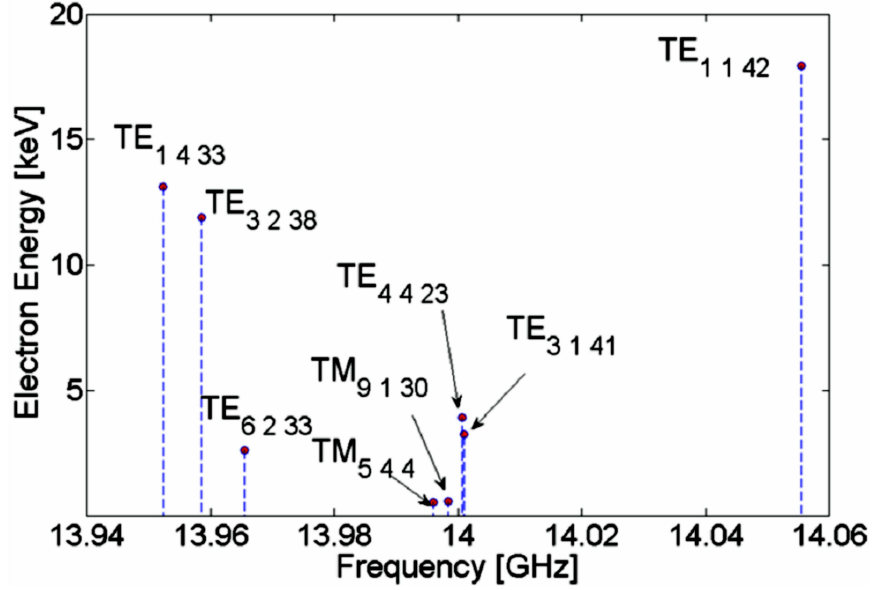


Figure 4.7.2: Simulation of electron heating rapidity for different modes.

This effect is of primary importance because it permits to obtain high densities of warm electrons. The electromagnetic field pattern over the ECR surface for several frequencies have been shown in figure 4.7.4. It is clear that passing from 14 to 18 GHz the pattern of the resonant modes over the ECR surface changes considerably; this means that other than the effects due to the scaling laws (see section 1.9) we have to take into account the differences in the field patterns when the frequency increases. In particular a further series simulations was carried out in order to investigate the electron heating for far frequencies with two different resonant modes. The $TE_{4,4,23}$ mode has been simulated with 300 W of power, and the results have been compared with the ones obtained with the $TE_{4,6,22}$ mode at 1000 W of power. It has been calculated that after 50 $nsec$ the mean electron energy in case of the $TE_{4,4,23}$ mode is of about 1.5 keV , whereas for the $TE_{4,6,22}$ mode the energy is 1.01 keV . Although the power for the second mode is three times higher than the first one, the energy is even lower. Obviously the simulations cannot predict the final density obtained using the mode near 18 GHz , as the increase of the pumping wave frequency should increase the cutoff frequency, according to the scaling laws. But the collective aspects are not taken into account in our model, then only the electron heating can

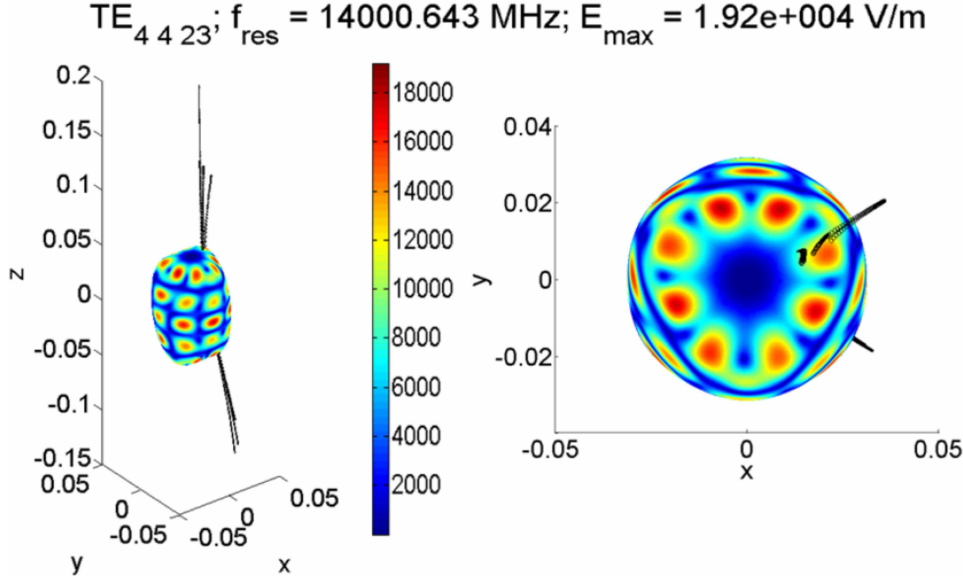


Figure 4.7.3: The electron trajectory and the simulated resonance surface with the electromagnetic field are shown. The crossing point between the electron trajectory and the resonance surface is of primary importance for the electron heating.

be investigated. This means that the increase of the density due to the cut-off can increase the source performances. This effect is usually observed in all the sources operating nowadays. Even if the frequency is increased, the benefits due to the scaling laws may be reduced by the strong fluctuations of the heating rapidity due to the changes in the mode structure.

The simulations demonstrate that the use of the right mode allows to strongly reduce the injected power. By using the $TE_{4,6,22}$ mode with 2000 W of input power a final energy of 1.26 keV has been obtained, that is still lower than the final energy obtained with the mode $TE_{4,4,23}$. Then the frequency tuning effect may allow to experimentally reduce the power required to obtain high performances, with great advantages in terms of reliability and stability of the sources operations.

4.8 Saturation of the heating rapidity with power

We have simulated the electron motion in an ECR plasma with a magnetic field structure similar to the VENUS and the SERSE. Several RF power

4.8 Saturation of the heating rapidity with power

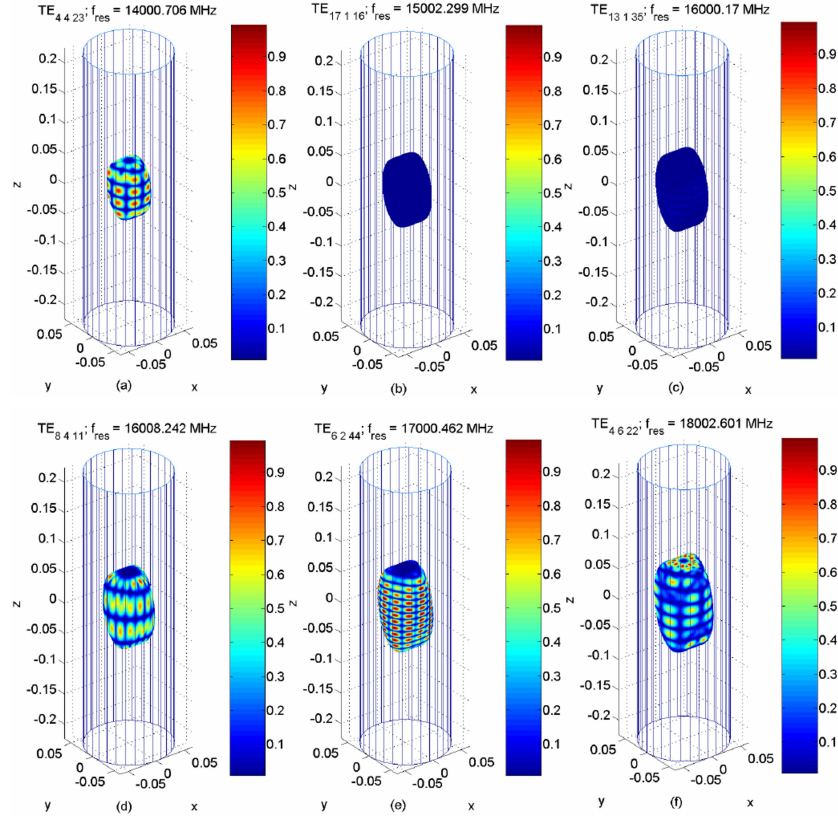


Figure 4.7.4: Six resonance surfaces with the electromagnetic field pattern for different frequencies. Note that for particular frequencies the amplitude of the electric field over the resonance surface is practically zero. Passing from 16 to 16.008 GHz the field pattern completely changes.

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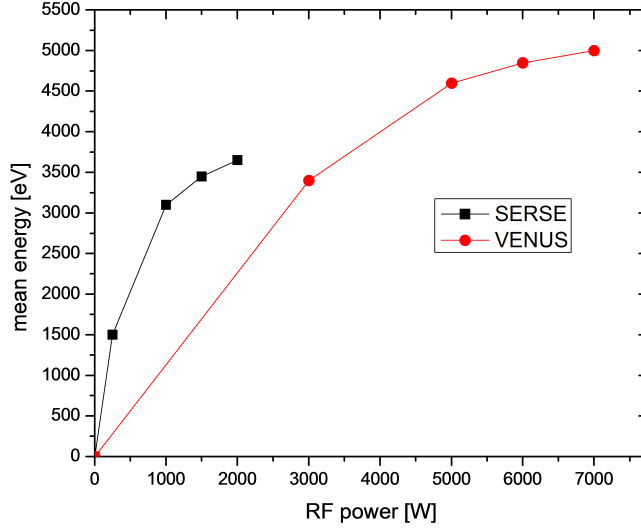


Figure 4.8.1: Mean electron energies (in eV) at different microwave power calculated by means of numerical simulations based on the single particle theory of ECR plasmas. Simulations carried out for SERSE and VENUS are shown.

were simulated, and the results have been analyzed in terms of the mean electron energy after 50 ns , i.e. considering again the heating rapidity. The values of the mean energies have been plotted with respect to the power in figure 4.8.1. Note that the two sources differ especially for the gradient steepness: for VENUS the magnetic field gradient is about two times than for SERSE in proximity of the resonance surface. A more detailed analysis about the gradient distribution over the resonance surface will be done in section 4.11. Here we would like to understand how the gradient affects the heating rapidity. In addition, the simulations give a reasonable explanation to the saturation of the ECRIS performances with the power, as observed in many experiments and as mentioned in section 4.5. It can be seen that an energy saturation occurs for VENUS above $P_{RF} = 5\text{ kW}$ (see fig. 4.8.1) in case of maximum longitudinal and transversal mirror ratio ($B_{inj} = 3.38\text{ T}$, $B_{min} = 3.38\text{ T}$, $B_{inj} = 2.36\text{ T}$, $B_{hex} = 3.38\text{ T}$). A larger microwave power does not change so much the electron heating, in terms of electron energy distribution, as observed experimentally.

As the simulation point out the rapidity of the electron acceleration from the cold to the warm population, that is of primary importance for the build up of the higher charge states, the saturation of the electron acceleration

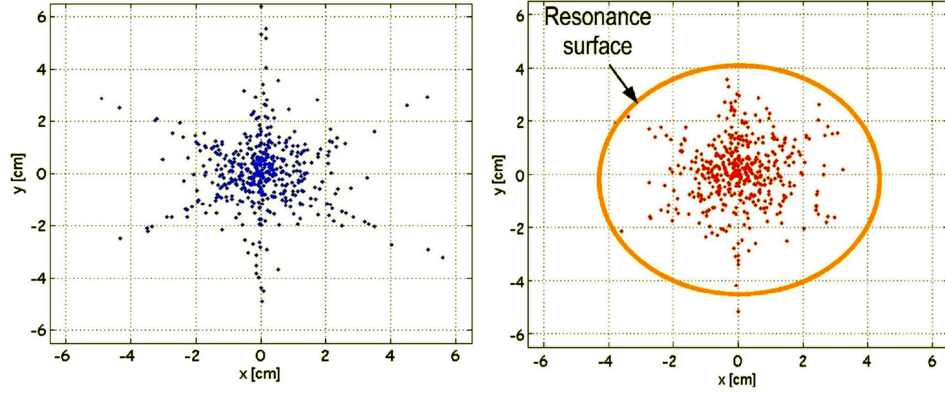


Figure 4.9.1: Projection along a transversal plane of the plasma electrons positions. With only the magnetostatic field (left) the electrons flow away from the plasma along the six loss cones. In presence of the electromagnetic field (right), the electrons are strongly plugged inside the resonance surface (a circle in the plane projection) and the confined fraction increases.

rapidity causes the saturation of performances observed in ECRIS. Note that for SERSE at 18 GHz the saturation trend is less marked and it occurs at lower RF power than for VENUS. In addition, the saturation energy results lower than the VENUS one. According to these results we may argue that the gradient steepness mainly influences the power values for which the saturation occurs (the higher gradient for VENUS corresponds to a higher power saturation value), while the higher frequency and the higher electric field on the resonance surface permit to the VENUS electrons to reach higher energies.

4.9 Two Frequency Heating simulations

Single particle simulations have demonstrated that the electromagnetic field not only provides the plasma heating but it also plays a role on the electron confinement. A considerable fraction of electrons which escape from the plasma if subjected only to the magnetostatic field, are plugged inside the ECR surface if they interact with the electromagnetic wave. This effect is able to recover those electrons which otherwise escape from the confinement as their velocity vector lies inside the loss cone. Figure 4.9.1 clearly shows this effect. The electron flows along the six loss cones of the confinement

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	(14 + 18)GHz (1000 + 300)W	18 GHz, 2000 W	18 GHz, 1000 W	14 GHz, 300 W
Recovered electrons [%]	38.2	35.2	19.2	12.9
Energy (keV)	1.9589	1.2612	1.0141	1.5509

Table 4.3: Percentage of recovered electrons (i.e., electrons that exit from the loss cone by means of the electromagnetic field) and the electron energy after 50 ns are shown for different heating configurations (SFH and TFH) and at different powers.

structure are strongly reduced by the wave, and the electrons are confined inside the resonance volume. This effect is of primary importance to stabilize the ECR plasmas and it also plays a crucial role in case of the so-called Two Frequency Heating - TFH. The advantages of the TFH are evident and they are hereby summarized:

- 1) We double the resonance zone width, so that electrons that do not gain energy during the first crossing can be heated by the second one. Figure 4.9.2 shows the locations of these two resonance regions inside the plasma chamber calculated for the initial setup of our simulations.
- 2) Electrons that have not been plugged by the first wave can be recovered by means of the second one, with a further increase of the plasma confinement.

A series of simulations has been carried out in order to verify the effect of the TFH. In the case in figure 4.9.2, the two resonance zones are far away. This allows the electrons to randomize their gyro-phase with respect to the electromagnetic wave. The numerical study has dealt with the case of two electromagnetic waves at different frequencies supplying the mode $TE_{4\ 6\ 22}$ (1000 W) and the $TE_{4\ 4\ 23}$ one (300W), respectively. It is important to notice in table 4.3 that, by means of the TFH, the final energy is almost two times higher than in the case of SFH at 18 GHz, whereas the amount of recovered electrons (i.e., equal to 38.2%) is two times the one obtained in single frequency operation (18 GHz; 1000 W).

It is the case to remember, according to the saturation curve shown elsewhere (figure 4.8.1), that in order to double the final energy in SFH, we had to triplicate the power, whereas in TFH, we increase the total power to about 30%. The two frequencies heat the plasma much better than the single one at 2000 W: The final energy is almost 50% higher, and even the amount of recovered electrons is slightly increased. Then, the effect of the second

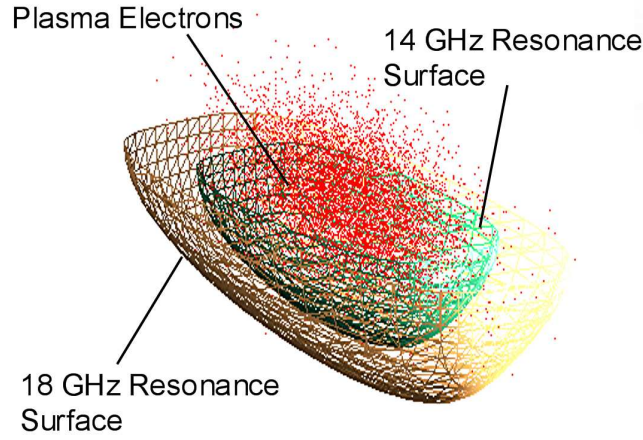


Figure 4.9.2: Section of the two resonance surfaces showing the position of the plasma electrons.

frequency is to increase the electron temperature and improve the electron confinement. This means that also the ion confinement time increases, and, finally, the extracted ion beam has higher charge states and currents. A more detailed analysis can be carried out by looking to figure 4.9.3. The cumulative probability is plotted on the vertical axis, with respect to the electron energy. A point on one of the plotted curves gives the number of electrons⁸ whose energy is below the corresponding energy value indicated by the abscissa. Let us define the cumulative probability for a given energy as CP , then the number of particles above a given E value will be, obviously, $1 - CP$. It gives information about the different number of electrons which populates different energy ranges. It is clear, from the figure, that the $TE_{4\ 4\ 23}$ mode is able to remove the greatest amount of electrons from the cold, initial low energy population. For this mode in SFH the number of the electrons above $2\ keV$ is higher than the $T_{4\ 6\ 22}$ at any power, being the $\sim 20\%$ for the former and the $\sim 15\%$ for the latter at $2000\ W$ and even more for $T_{4\ 6\ 22}$ at $1000\ W$. For higher energies, above $6\ keV$, the relative difference increases: for $TE_{4\ 4\ 23}$ $(1 - CP) \sim 13\%$, whereas for $TE_{4\ 6\ 22}$ at $2000\ W$ $(1 - CP) \sim 8\%$. The gap between the two modes drops approaching the maximum energy, probably because of the limited duration of the simulations. This is another evidence of the potentialities of

⁸Note that the vertical axis has been normalized to 1.

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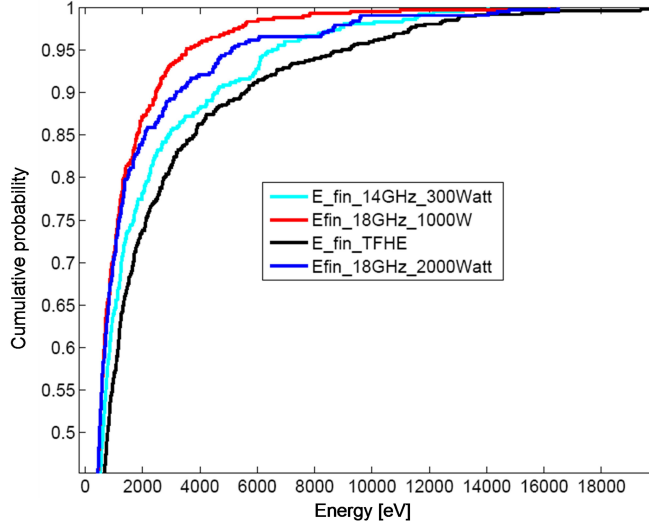


Figure 4.9.3: Cumulative probability plot of the electron energy distribution function for the different simulated modes and power. Once fixed an energy on the horizontal axis, the correspondent value on the vertical one gives the relative number of particles below that energy, or the $1 - CP$ relative number of electrons above that energy value, where CP is the cumulative probability.

the frequency tuning. The heating rapidity does not depend on the power directly, but on the electric field amplitude over the resonance surface. After 50 ns , by means of the TFH all the electrons have the time to acquire more energy from the wave. This capability to remove great amounts of electrons from the cold population is one of the main advantages of the TFH, as it permits to quickly populate the warm plasma component.

The TFH allows the increase of both the electron density and the ion confinement time, thus increasing $\langle q \rangle$; it optimizes the electromagnetic wave energy transfer to the electrons so that a higher quantity of heated electrons can be created before they leave the confinement region. In fact, the experimental results of TFH obtained with SERSE show an increase of the electron temperature and of the ion confinement (the CSD peak shifts toward higher charge states).

4.10 Two Close Frequency Heating

It is a common idea that in the TFH mode the two frequencies have to be well separated, so that the phase randomization of the electron gyromotion can occur. Then, we expect that, by injecting two waves with very close frequencies into the plasma chamber, in order to obtain two close resonance zones, a sort of *electron surfing* is obtained: the electrons that leave the first ECR zone may be further accelerated by the second one, depending on the phase shift between the two waves. In this case the process occurs without the phase randomization, and if the two waves have the right phase difference the energy transfer to the electron is maximized.

According to the Canobbio theory (see section 1.11.1), the phase continuously changes during the electron-wave interaction and the electron acceleration stops when the phase between the electromagnetic wave and the electron gyromotion is equal to $\pi/2$. Let us define τ_c as the crossing time of the ECR zone (the ECR surface becomes an ECR volume because of relativistic and Doppler effects); for $t \gg \tau_c$, the electron forgets his phase and the ECRH process goes on [36], [37]. However, for $t \geq \tau_c$, i.e., just few gyrop periods after the ECR zone, the phase is still fixed to the value $\pi/2$ that it reached at the end of the electron acceleration. The frequency of the second wave has to be chosen so that the corresponding ECR zone is located where the first one finishes; first, we calculated the thickness of the first ECR zone. The central region of the axial magnetic field can be modeled by means of the equation (4.6.5). From the ECR condition, we have the following:

$$B_{ECR} = \frac{m}{q} \omega_{RF} \quad (4.10.1)$$

where ω_{RF} is the pulsation of the RF field, m and q are the mass and the charge of the electron, respectively, and B_{ECR} is the magnetic field at the resonance. As m and q are two constants (by neglecting the relativistic variation of the electron mass, at least for electrons of some keV only), we can differentiate B_{ECR} with respect to ω_{RF} , obtaining

$$dB = \frac{m}{q} d\omega_{RF} \quad (4.10.2)$$

dB can be also obtained from (4.6.5)

$$dB = 2k_2 z dz \quad (4.10.3)$$

then, by substituting from (4.10.2) into (4.10.3), we have

$$\frac{m}{q} d\omega_{RF} = 2k_2 z dz \quad (4.10.4)$$

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by means of (4.10.4), we can determine dz

$$dz = \frac{d\omega_{RF}}{z} \frac{m}{2k_2q} \quad (4.10.5)$$

now, we can consider finite quantities instead of infinitesimal ones

$$\Delta z = \frac{m}{2k_2q} \frac{\Delta\omega_{RF}}{z} \quad (4.10.6)$$

where the thickness Δz of the resonance zone was calculated by Canobbio [39]:

$$\Delta z \simeq 1.4 \frac{c}{\omega} \left[\frac{\left(\frac{eE}{mc\omega}\right)^2}{\left(\frac{c}{\omega}\right) \left(\frac{\Delta B}{\Delta z}\right) \left(\frac{1}{B_{ECR}}\right)^3} \right]^{\frac{1}{5}}. \quad (4.10.7)$$

Finally, for $\Delta\omega_{RF}$, we obtain

$$\Delta\omega \simeq 1.4z \left(\frac{2k_2q}{m}\right) \left(\frac{c}{\omega}\right) \left[\frac{\left(\frac{eE}{mc\omega}\right)^2}{\left(\frac{c}{\omega}\right) \left(\frac{\Delta B}{\Delta z}\right) \left(\frac{1}{B_{ECR}}\right)^3} \right]^{\frac{1}{5}} \quad (4.10.8)$$

where z is the position at which the resonance occurs (being $z = 0$ the center of the plasma chamber), c is the speed of light, and E is the electric field in the resonance zone. If we replace the characteristic values of the aforementioned parameters for the case of the SERSE source ($z \sim 0.1$ m, $\Delta z \sim 1$ mm, and $k_2 \sim 25$ T/m²), we obtain the following:

$$\Delta\nu_{RF} \equiv \frac{\Delta\omega}{2\pi} \simeq 279 \text{ MHz} \quad (4.10.9)$$

Equation (4.10.9) gives the spacing between the frequencies of the two electromagnetic waves in order to have two contiguous resonance zones. By means of a simple theoretical approach, it is also possible to obtain an estimation of the phase shift between the two waves in order to improve the electron acceleration. It is known that the phase of a plane wave is as follows:

$$\phi = \omega t - kz + \theta \quad (4.10.10)$$

where k is the propagation constant, and θ is named the phase angle, determined by the initial conditions. As in the previous pages, we inferred that stationary waves can be considered even in the presence of the plasma inside the cavity; the phase of these waves can be written by only considering

the time dependence term, i.e., the argument of the sinusoidal term which determines the oscillations in time of the stationary wave. In addition, we have to consider the phases of the two waves with close frequency

$$\phi_1 = \omega_1 t_1 + \theta_1 \quad (4.10.11)$$

$$\phi = \omega_2 t_2 + \theta_2 \quad (4.10.12)$$

where t_1 is the time in which the electron reaches the first resonance zone and t_2 when it reaches the second one. As said before, the acceleration stops when the electron-wave phase difference becomes $\pi/2$; then, it means that the second wave should have a difference in phase of $\pi/2$ with respect to the first one in order to allow further electron acceleration. Hence, the condition $\phi_1 - \phi_2 = \pi/2$ has to be satisfied. The two frequencies must be very close, as shown in (4.10.9); then, we can approximate $\omega_1 \sim \omega_2 \sim \bar{\omega}$. By subtracting (4.10.11) to (4.10.12), we obtain

$$\phi_2 - \phi_1 = \bar{\omega}(t_2 - t_1) + \theta_2 - \theta_1 \quad (4.10.13)$$

$\Delta_t = (t_2 - t_1)$ is the time that an electron needs to travel from the first resonance surface to the second one. This means that somehow Δ_t depends on the electron temperature, as the time that the electrons spend to travel across the two ECR surfaces is smaller at higher than at lower temperatures. However, the electron energy, once the magnetic field profile and the two frequencies are fixed, depends only on the RF power. The higher is the power, the lower is the time needed for the electrons to go from the first ECR surface to the second one. Finally, we aim to evaluate the initial phase difference of the two waves $\Delta\theta = \theta_2 - \theta_1$, as it is the parameter that has to be set at the beginning of each simulation. Following (4.10.13), and remembering that $\phi_2 - \phi_1 = \pi/2$ in order to have the right phase condition on the second ECR surface, we obtain that

$$\Delta\theta \equiv \left(\frac{\pi}{2} - \text{const} \cdot \frac{\bar{\omega}}{P} \right) \quad (4.10.14)$$

i.e., the initial phase shift is a function of the power of the electromagnetic wave; in particular, once the frequency is fixed, the higher is the RF power, the higher must be the initial phase shift and vice versa. Let us show the results of some simulations of the two-close frequency heating (TCFH). First, the TE modes that exist in the SERSE plasma chamber in the range of $14 \pm 0.15 \text{ GHz}$ (i.e. in the range determined by (4.10.9)) have been calculated. Only the modes of the type $TE_{4 \ 4 \ r}$ have been chosen for the simulations in

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Mode	Power [W]	Confined Fraction	E final [keV]
$TE_{4\ 4\ 23}$	1000	0.904	$4.228 \cdot 10^3$
$TE_{4\ 4\ 23}$	500	0.886	$2.454 \cdot 10^3$

Table 4.4: The results of SFH with the $TE_{4\ 4\ 23}$ mode at two different powers.

order to avoid the possibility that any variation in the final electron energy was ascribed to different electromagnetic field patterns. Then, the modes $TE_{4\ 4\ 22}$ (13.827 GHz) and $TE_{4\ 4\ 24}$ (14.186 GHz) have been chosen for the first series of simulations. Furthermore, in order to prove the assumption in (4.10.9), other two modes of the type $TE_{4\ 4\ r}$ have been chosen but with frequencies higher than 14 ± 0.15 GHz . They are $TE_{4\ 4\ 25}$ (14.376 GHz) and $TE_{4\ 4\ 26}$ (14.572 GHz). We expect that, if the assumption (4.10.9) is true, the TCFH will not be effective by using these modes. Four series of simulations have been carried out.

- 1) Simulation of SFH with the $TE_{4\ 4\ 23}$ mode in order to evaluate the mean electron energy at 500 and 1000 W . This will be the energy reference for the comparison with the TCFH mechanism.
- 2) The second series has been devoted to simulations carried out with two waves carrying 500 W each ($TE_{4\ 4\ 23}$ mode at 14.001 GHz and $TE_{4\ 4\ 24}$ mode at 14.186 GHz) but with different initial phase shifts $\Delta\theta$. in order to evaluate the best phase shift condition. In addition, the test of (4.10.9), i.e., the determination of the ECR zone thickness, has been done.
- 3) The third series has been devoted to verify the equation (4.10.14), i.e., the phase shift dependence on the mean electron energy and then on the microwave power. Simulations with the same two modes have been carried out at 250 W each instead of 500 W each.
- 4) The last series of simulations has verified the effectiveness of the multiple close frequency heating, i.e., three waves with close frequencies have been simulated with a power of 333 W each. Even in this case, the condition of contiguous resonance regions has been satisfied by choosing the three modes $TE_{4\ 4\ 22}$ (13.827 GHz), $TE_{4\ 4\ 23}$ (14.001 GHz), and $TE_{4\ 4\ 24}$ (14.186 GHz) have been simulated at the same time.

$\Delta\theta$ [°]	Confined Fraction	E final [keV]	$\Delta E/E$ [%]
0	0.916	4.351	2.9
45	0.918	4.815	13.9
80	0.907	4.848	14.7
90	0.903	5.045	19.3
100	0.894	4.562	7.9
135	0.907	4.383	3.7
180	0.920	4.810	13.8
225	0.903	4.703	11.2
270	0.908	5.000	18.2
315	0.907	4.934	16.7

Table 4.5: Results of the TCFH are summarized ($TE_{4\ 4\ 23}$ and $TE_{4\ 4\ 24}$) modes; 500 W each) for different phase shifts of the two waves. In bold characters is the best result (the $\Delta E/E$ indicates the relative increase of the electron energy with respect to SFH at the same power).

Table 4.4 shows the results of the simulations in SFH, at 500 and 1000 W of power; these results will be used to make a comparison with the ones obtained in the TCFH at the same power.

Table 4.5 summarizes the results of the TCFH at 500 + 500 W. It can be noticed that, for $\Delta\theta = \pm 90^\circ$, there is a maximum increase of the electron energy, up to almost 20%. The confined fraction remains about the same in single or double frequency heating. Its fluctuations are less than the ones due to the inaccuracy of the numerical calculation. Two simulations have been done in order to compare the SFH and the TCFH when (4.10.9) is not satisfied, i.e., when the two resonance zones are well separated, according to the Canobbio calculation of the ECR thickness. The final energy for TCFH is even lower than for the case of SFH (4228 eV in SFH and 3263 eV for the TCFH, with the same total power). This means that the TCFH can be even detrimental if the right conditions are not satisfied. Other simulations have been done in order to verify how the TCFH depends on the microwave power.

In Table 4.6, the results obtained at 250 + 250 W are reported. Some considerations can be done about these data.

- 1) The maximum increase of the final energy is obtained for $\Delta\theta = 60^\circ$, i.e.,

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$\Delta\theta$ [°]	Confined Fraction	E final [keV]	$\Delta E/E$ [%]
45	0.897	2.940	19.7
60	0.907	3.121	27
75	0.903	3.044	24
90	0.894	2.939	19.7

Table 4.6: Results of the TCFH are summarized ($TE_{4\ 4\ 23}$ and $TE_{4\ 4\ 24}$ modes; 250 W each). In bold characters is the best result.

for a phase shift lower than the previous one, when we had 500+500 W. This is in agreement with the (4.10.14); as the total amount of power has been decreased, the phase shift, according to (4.10.14), decreases. Expression (4.10.14) cannot provide an exact estimation of the phase shift given by a fixed power, so that a variable phase shifter is needed for daily operations of the ECRIS, when the plasma parameters have to be adapted to the needs of the accelerators.

- 2) The relative increase of the final energy is equal to 27%, higher than the 20% obtained at lower power; the energy saturation plays probably an important role.

These results demonstrate that the TCFH is an effective method to quickly remove electrons from the cold population moving them towards the warm one.

These results suggest that other possible electron heating mechanisms are possible by taking advantage from the phenomenon that we named *electron surfing* due to the proximity of the resonance surfaces. This means that *multi-frequency heating* should be possible, and it should permit to maximize the electron heating velocity v_h by minimizing the total power needed to obtain high charge states and current. The multifrequency heating of the ECR plasmas is characterized by many contiguous ECR zones in order to exploit the electron surfing. The simulation results of the three-frequency heating show an increase of 30% of the final energy (three waves at 500 W each with respect to a single wave at 1500 W have been compared). It is important to underline that the energy reached from the electrons in the case of three close frequency heating is comparable with the one reached in SFH but with 2000 W of power instead of 1500 W. Figure 4.10.1 features that the cumulative probability curve for the Three-CFH is appreciably shifted

4.11 Considerations on the hot electrons based on the single particle approach

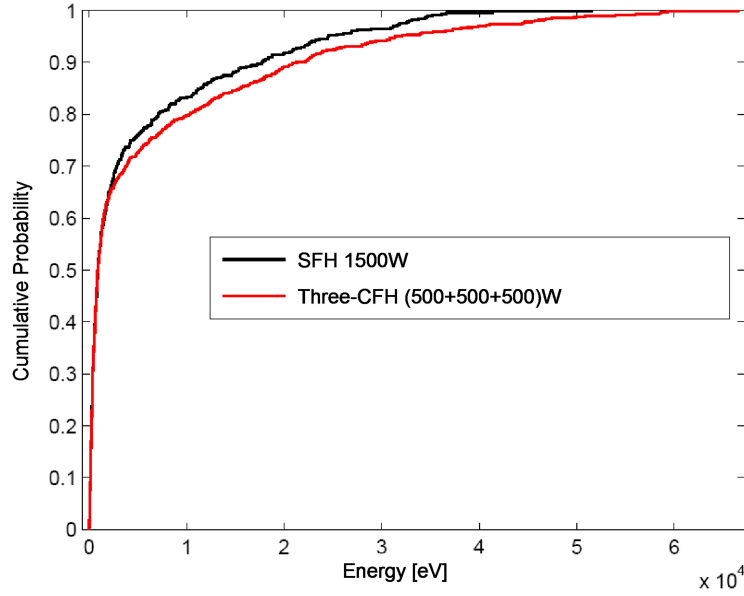


Figure 4.10.1: Cumulative probability comparison between SFH and Three-CFH.

towards higher energy both at intermediate and high energies. It is important to underline that not only the warm component is more populated in case of Three-CFH, but also the higher energy part of the curve is filled by high energy electrons more than in case of TCFH. It is remarkable that after 50 ns some electrons are able to reach almost 60 keV of energy, thus demonstrating all the effectiveness of the method. Anyway the shift of the curve tail towards high energy may be dangerous because after several μs even MeV electrons could be created.

4.11 Considerations on the hot electrons based on the single particle approach

In the previous sections it has been demonstrated that the electromagnetic field topology not only explains the phenomenon of the increase of source performances by means of the frequency tuning, but it also helps to explain the two frequency heating (TFH) and others [5]. Here we will focus our attention on the role played by the magneto-static field topology, especially as regards the distribution of the magnetic field gradient over the ECR

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surface. Such a distribution, analyzed on the basis of the Canobbio theory of the ECR heating in presence of non uniform magnetic fields, helps to better understand the production of high energy electrons in 3rd generation ECRIS. The ECRIS plasma need to be studied not only in terms of global parameters, but also in terms of local interactions among electrons, ions, magnetic field, and external waves.

The analysis of the experimental results reported in section 4.5 outlines several new items of 3rd generation ECRIS: the appearance of ultra-hot electrons, the non-saturating trend of the CSD even for very high mirror ratios, and finally a net dependence of source performances on relatively small variations of the magnetic field. Generation of ultra-hot electrons is a key point, because it is clear that the source performances deteriorate with the generation of such electrons. If the magnetic field configuration, the microwave pattern and power and other sources macroscopic parameters facilitate the electron heating up to high energies, then the source performances may decrease despite the operational conditions fulfill the usual ECRIS standard model. According to preliminary simulations, the appearance of high electron energies needs quite long electron lifetimes. In fact the ECR heating becomes slower as the energy increases, and the electrons probably require more than 1 *ms* to reach *MeV* energies. The measurements performed with SERSE [66] and recently with VENUS [40] show that the amount of ultra-hot electrons increases when the mirror ratio decreases; probably such variation of magnetic field does not affect the lifetime in an appreciable manner but it strongly affects the heating mechanism.

The ultra-hot electrons production strictly depends on the magnetic field profile, and in particular on the gradient. The Canobbio theory considers two different regimes for the energy transfer from an electromagnetic wave to the electrons moving in a magnetic field gradient. The first one is the "gentle gradient" regime, the second one the "strong gradient" regime. The maximum energy achieved in each passage through the ECR zone strongly depends on which of these regimes apply to our case. The transition value for the electric field from one regime to another is:

$$K_t = \left[\frac{\delta}{3.8} \right]^{\frac{3}{2}} \frac{m\omega c}{e} \quad (4.11.1)$$

with:

$$\delta = \frac{1}{B_{ECR}} \left(|\vec{\nabla} B| \right)_{ECR} \frac{c}{\omega} \quad (4.11.2)$$

where $\vec{\nabla} B$ is the magnetic field gradient, ω is the microwave pulsation, c is the speed of light, m and e are the electron mass and the electron charge

4.11 Considerations on the hot electrons based on the single particle approach

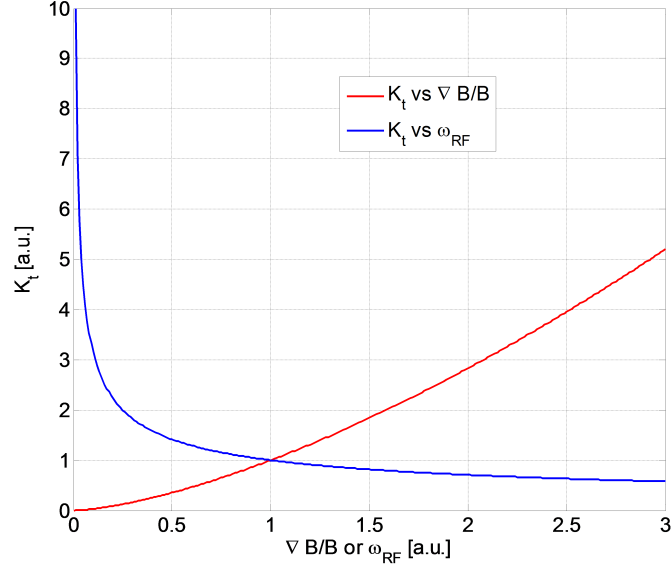


Figure 4.11.1: trend of the K_t parameter with respect to the magnetic gradient (red line) and to the RF frequency. For each curve the other parameter is fixed, then this plot features the dependence of K_t on gradient and frequency separately.

respectively [39]. Equation (4.11.2) relates the magnetic field gradient to the electric field amplitude at the ECR resonance (hence to the power). In addition it points out that the transition field depends on the microwave frequency. The possibility to operate in gentle or strong gradient regime depends not only on the magnetic field gradient, but also on the electric field at resonance and on the microwave frequency. If the microwave electric field amplitude E is higher than K_t ("gentle gradient"), the electrons can reach energies comparable with the value achieved in a constant magnetic field and the acceleration stops when the relativistic detuning occurs [10]. If E is comparable or lower than K_t ("strong gradient" regime), the maximum achievable energy is much lower than the previous one. The K_t dependence on the frequency and on the magnetic field gradient must be investigated more in detail. Figure 4.11.1 shows that if a reference value of K_t is fixed (in a.u. we can choose $K_t = 1$ as a reference value), by increasing the frequency (keeping constant the gradient) K_t decreases, whereas by fixing the frequency and by increasing the gradient K_t strongly increases (it depends

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on $\left(\frac{\vec{\nabla}B}{B}\right)^{\frac{3}{2}}$.

Calculations show that passing from 2nd to 3rd generation ECRIS the $\frac{\vec{\nabla}B}{B}$ parameter is almost doubled on a large area of the resonance surface (if the maximum mirror ratio is used), the frequency increases of a factor 1.5, while the RF power increases of a factor 3 (corresponding to a factor 1.7 for the electric field). Fig. 4.11.1 shows that by doubling the $\frac{\vec{\nabla}B}{B}$ parameter the electric field should be tripled in order to overcome K_t thus remaining in a gentle gradient regime (the corresponding increase of frequency slightly lowers this factor, but anyway the increase in electric field must be higher than 2). This means that the transition towards strong gradient regimes is facilitated and in some cases the source performances worse because of this effect. In case of strong gradient regime, if the mirror ratio is strongly increased, the energy gain may be considerably reduced and the mean plasma electron energy decreased. Hence the production of highly charged ions becomes more difficult.

The VENUS magnetic gradient near the resonance is about $12 \frac{T}{m}$ for $B_{min} = 0.673 T$. In this case K_t is about $1.15 \cdot 10^5 \frac{V}{m}$ and it scales with the magnetic field. For higher magnetic field gradients a higher P_{RF} is needed to overcome K_t and to maintain the "gentle gradient" regime in order to have a fast electron heating.

In order to better describe the behavior of the source for different longitudinal and radial profiles of the magnetic field, it is important to simulate the $\vec{\nabla}B$ distribution on the ECR surface as shown in figure 4.11.2. The calculation reproduces the field gradient distribution for a typical source condition, in particular in the case of the maximum field ratio used for the measurements shown in [12], [40].

Note that the previous "global" estimation of $\vec{\nabla}B$ is not able to describe the real situation, and it may lead to misleading conclusions. In particular the gradient is low in proximity of the extraction and injection side of the resonance surface (slightly higher for the injection side) and it is much higher on the lateral side. The field lines help to determine the electron trajectories, as their guiding center moves along them. Where the field lines cross the resonance surface, there the electrons absorb the energy from the wave. Jongen [97] and Koivisto [98] have demonstrated that the energy gain strongly depends on the gradient at the resonance, then it is clear that there are regions where the electrons gain a great amount of energy for single pass through resonance (because the gradient is not so pronounced), and others where the heating is modest because of the gradient steepness. Therefore

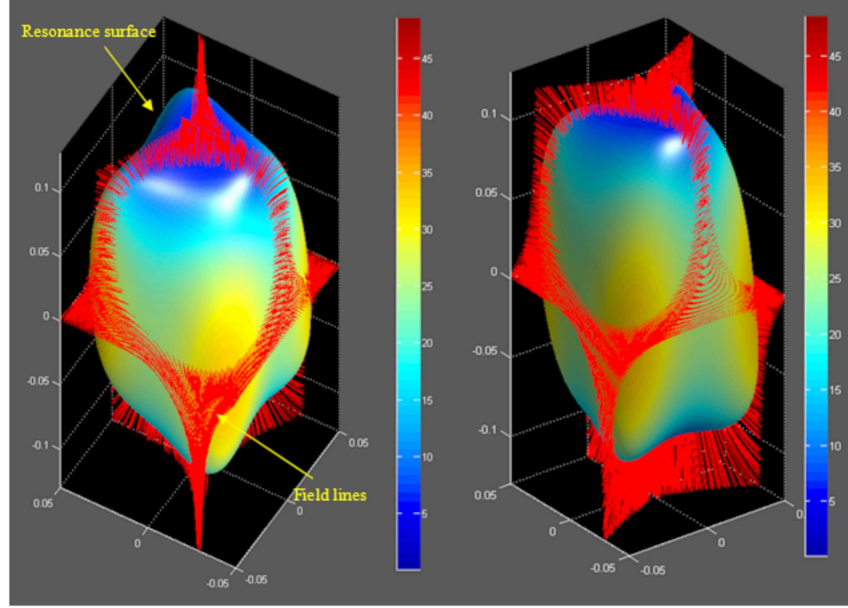


Figure 4.11.2: Plot of simulated $\vec{\nabla}B$ over the resonance surface for the VENUS source in case of the following magnetic configuration: $B_{inj} = 3.38 \text{ T}$, $B_{min} = 0.65 \text{ T}$, $B_{ext} = 2.36 \text{ T}$, $B_{hex} = 2 \text{ T}$.

only in limited regions of the plasma chamber the ultra-hot electrons can be originated, and the corresponding X-rays can be generated, mostly along the plasma chamber axis.

Let us study the variation of the $\vec{\nabla}B$ distribution for different configurations; in particular we will take two extreme cases (maximum and minimum mirror ratio, both longitudinally and radially) investigated during the experimental measurements. First, we study the changes of $\vec{\nabla}B$ over the ECR surface for variations of B_{min} . It is clear from fig. 4.11.3 that the increase of B_{min} strongly changes the gradient distribution both longitudinally and radially. In particular the gradient near the extraction hole passes from $6.1 \frac{T}{m}$ to $3.6 \frac{T}{m}$ in proximity of the plasma chamber axis. This means that an increase of the 25% in the B_{min} field produces a decrease of more than 40% in the gradient.

At the same time K_t changes from $4.3 \cdot 10^4 \frac{V}{m}$ to $1.9 \cdot 10^4 \frac{V}{m}$, i.e. the transitional electric field is more than halved. On the lateral side of the surface, the gradient decreases from about $35 \frac{T}{m}$ (low B_{min} case) to less

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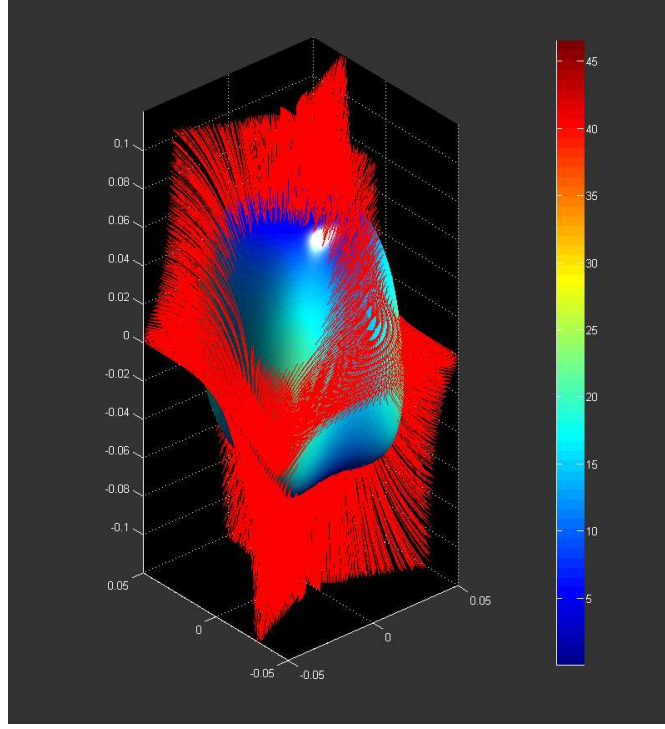


Figure 4.11.3: Plot of simulated $\vec{\nabla}B$ over the resonance surface for the VENUS source in case of the following magnetic configuration: $B_{inj} = 3.38 \text{ T}$, $B_{min} = 0.85 \text{ T}$, $B_{ext} = 2.36 \text{ T}$, $B_{hex} = 2 \text{ T}$.

than $25 \frac{T}{m}$ (high B_{min} case). For gradients around $30 \frac{T}{m}$, K_t is approximately equals to $5 \cdot 10^5 \frac{T}{m}$. Then the B_{min} variation influences globally the gradient distribution and finally the K_t parameter.

Calculating the wave electric field amplitude for VENUS according to the equation 4.6.4 (considering a $TE_{4,4,47}$ mode at 3000 W of power), we obtain that it may reach, on some localized areas of the ECR surface, values on the order of several $10^5 \frac{V}{m}$, and in other areas values below $10^4 \frac{V}{m}$. This means that the VENUS source probably operates around the transition region between the two different regimes and this can explain the trend of the CSD with B_{min} and B_{ext} (see fig. 4.5.1). For a value of B_{min} around 0.65 T the electric field at the resonance was not sufficient to guarantee the "gentle gradient" regime on a large area of the resonance surface (especially laterally, where the high hexapolar field produces gradients up to $35 \frac{T}{m}$). Then

the electrons are weakly accelerated in these regions. Hot and ultra-hot electrons are created in the near axis regions of the ECR surface, but their amount is considerably lower than in case of high B_{min} . Such a mechanism explains the decrease of VENUS performances with very high mirror ratio, despite of the improvement of the MHD stability according to the High-B mode principle. This behavior was not observed in the case of the SERSE source [73] because it was far from the transition region.

In addition, the distribution of the electromagnetic field, related to a particular mode excited at a given frequency, should be compared with the magnetic gradient distribution, in order to give more quantitative details about the electron heating process. In particular, it can be shown that most of the excited cavity modes have a minimum corresponding to the plasma chamber axis, and maxima located around it [99]. This probably leads to a low amount of ultra-hot electrons produced along the source axis (as confirmed by X-rays spatially resolved measurements carried out from Biri [100]), whereas maxima located in the near axis region are crucial for ultra-hot creation because of the low magnetic field gradient. This is a very complicate task and we will limit the present discussion to qualitative estimations.

According with the results shown in figure 4.5.1 one can argue that for $B_{min} < 0.717 \text{ T}$ the source was operated in the strong gradient regime and the electron acceleration and the following electron heating have been slower, unless the electric field amplitude at the resonance be increased (if the "warm" electrons creation rate is not equal to the electron loss rate, the sources performances worsen). If B_{min} is further increased some unwanted phenomena begin to play. First, the plasma confinement worsens because of the decreasing mirror ratio. Second, and probably more important, the magnetic field gradient becomes lower in a large part of the ECR surface, then the electrons gain higher energy even for a single crossing of the resonance (above the adiabatic barrier, as demonstrated from preliminary estimations and PIC simulations [101]): in this way the great amount of the pumping wave energy is lost because of considerable energy absorption from high energy electrons, that are completely useless for ionization. Summing up we have:

- a low Mirror Ratio R cannot guarantee a sufficient confinement and the electron density, temperature and ion confinement time decrease; outgasing on the plasma chamber walls is produced and the charge exchange limits the higher charge states production;

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- in addition, low R leads to the production of a great number of ultra-hot electrons with the following decrease of the source performances;
- the increase of the Mirror Ratio (B_{min} decreases) improves the electron confinement and it reduces the production of high energy electrons: the "warm" component of EEDF and the electron density increase. An improvement of the source performances is obtained until the transition from the gentle gradient regime to the strong gradient occurs; then the heating rate of the electrons is limited, which negatively affects the ECRIS performance. Here strong and gentle gradient regimes must be considered in terms of predominant areas over the resonance surfaces where such regimes can be dominant.

When $B_{min} > 0.717 T$ the "gentle gradient" regime is active on a large part of the ECR surface so that the temperature depends also on P_{RF} . For increased magnetic field gradient and constant power the transition towards to the strong gradient regime occurs and a strong increase of the electric field is necessary at the same time to come back to the gentle gradient regime. Explanation of the source performances in terms of the topology on the resonance surface allows to understand also the dependence of the CSD trend and hot electrons production on the hexapolar field. Measurements show that increasing the radial mirror ratio the source performances increase both in terms of charge states and extracted currents. No strong increases in the production of high energy electrons has been noticed. These evidences confirm the validity of our modeling approach: variations in the hexapolar field change the magnetic field gradient on the lateral side of the ECR surface, but it still remains above $30 \frac{T}{m}$ in regions where the field lines cross the resonance, as demonstrated by figure 4.11.4.

As demonstrated before, values of $\vec{\nabla}B$ from 3.6 to $6.1 \frac{T}{m}$ correspond to K_t parameters comparable with the electric field amplitude of the pumping waves. Then the gradients on the lateral part of the resonance surface are much higher than the values at the injection or extraction side, also in case of the minimum hexapolar field used for the experiment (maximum 2.1 T, minimum 1.78 T). Hence the hexapolar field variation weakly affects the electron transition from one population to the other so the variation of the performances can be described by taking into account the usual ECRIS scaling laws and the High-B mode concept. The hexapole has a limited influence on the EEDF dynamics, but it affects the electron density. The poorer is the radial confinement, the lower is the electron density. Conversely the B_{min} and the B_{ext} variation strongly affects the plasma dynamics in

4.11 Considerations on the hot electrons based on the single particle approach

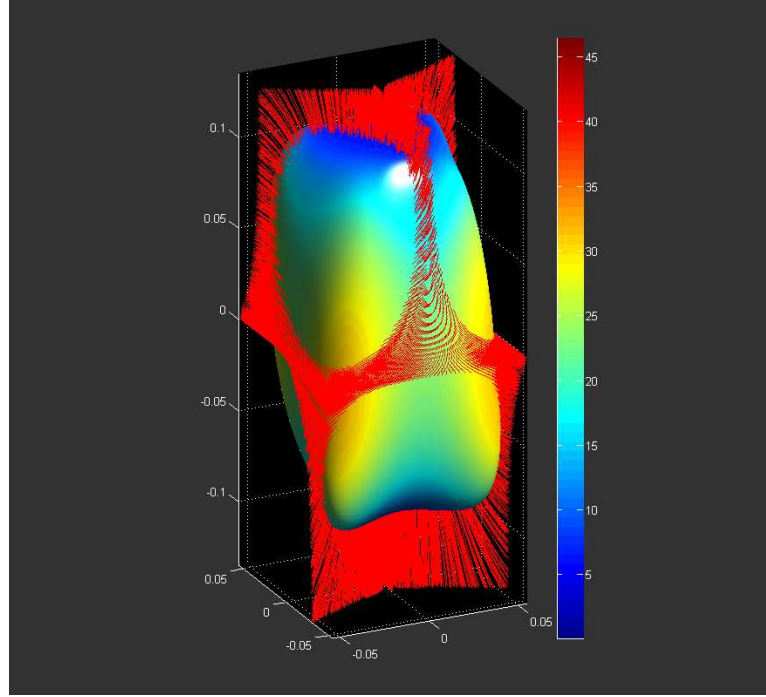


Figure 4.11.4: Plot of the simulated $\vec{\nabla}B$ over the resonance surface for the VENUS source in case of magnetic configuration: $B_{inj} = 3.38 \text{ T}$, $B_{min} = 0.65 \text{ T}$, $B_{ext} = 2.36 \text{ T}$, $B_{hex} = 1.78 \text{ T}$.

terms of wave-particle interaction, because of strong variations of the on the resonance surface. The dependence of the mean electron energy on the microwave frequency has been investigated in [98] and [40] featuring a greater amount of hot electrons produced with VENUS at 28 GHz with respect to 18 GHz operations. The same experimental evidence was found with SERSE during the source operation at 28 GHz , as shown in fig. 4.11.5.

In addition to the contribution of the wave phase velocity, as reported in [102], the strong dependence of the ultra-hot electron production on the microwave frequency may be connected with equation (4.11.2). In fact, from this equation it results that the electric field boundary K_t depends on the gradient conditions and on $\sqrt{\omega_{RF}}$ (see fig. 4.11.1). The magnetic field gradient over the resonance surface side near the extraction region is about $5 \frac{T}{m}$ at 28 GHz , and $3.2 \frac{T}{m}$ for 18 GHz . By using the equation (4.11.2) it can be noticed that the K_t parameter decreases from $4 \cdot 10^4 \frac{V}{m}$ (18 GHz) to

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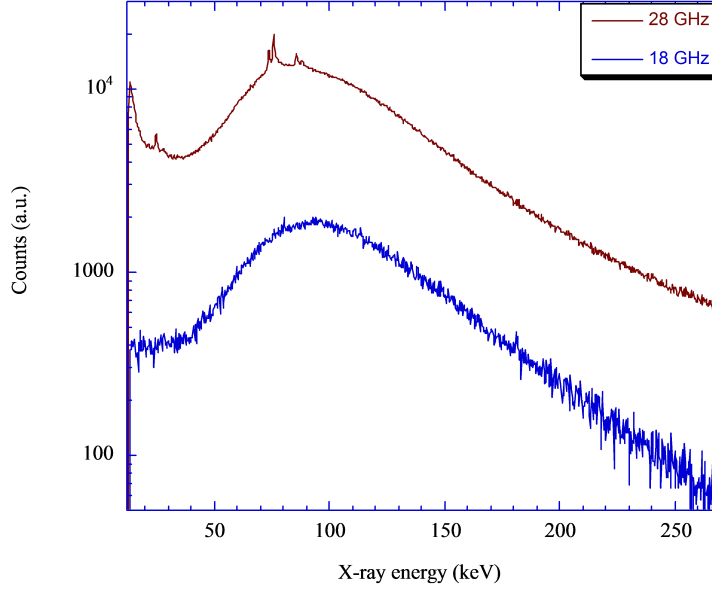


Figure 4.11.5: Comparison of X-ray spectra for SERSE, obtained at 18 and 28 GHz .

$3 \cdot 10^4 \frac{V}{m}$ (28 GHz). Again this values corresponds approximately to the estimated electric field values on the ECR surface, then transitions to "gentle gradient" regimes are easier in case of 28 GHz operations and the electrons are largely heated up to high energies. Probably the decrease of the X-ray spectrum slope reported in [40] at 28 GHz depends on RF diffusion and wave phase velocity [102], but the larger amount of hot electrons produced at 28 GHz may be explained by the regime transition and topology. These effects were evident for the SERSE source in case of plasma heating at 28 GHz [66].

Figure 4.11.1 features that in this last case the transitional electric field decreases because the frequency decreases and the magnetic field remains approximately the same, then more electrons can be accelerated up to high energies even for single resonance crossings.

The experiments have shown that the EEDF shape does not change too much with power, but the maximum energy reached by the electrons increases with it. On the other side a variation of the magnetic field profile can change the regime (gentle or strong) and the saturation curve shown in

figure 4.8.1. The simulations results are in agreement with the experimental data. The mechanisms discussed so far are able to explain most of results but they do not explain the generation of the ultra-hot electrons observed ($W > 400 - 500 \text{ keV}$) which behave differently than the warm ones, as they increase with B_{min} .

The "ultra-hot" electron temperature is mainly determined by the magnetic field structure, and in particular a lower gradient over the ECR surface favorites the production of high energy electrons. There is evidently a more complicated heating mechanism that is switched on by a lower mirror ratio.

4.12 Summary of simulations and comparison with the experiments

The comparison of the experimental data with the simulated ones has permitted to suppose that a completely different electromagnetic field pattern is produced if the frequency is slightly changed, but only some of the modes are able to heat the plasma more efficiently than others. Although the simulations are based on a oversimplified model and most of the collective aspects are not taken into account, they are able to give some relative estimations of the electron heating rapidity, and this permits to explain one of the most important effects of ECRIS operation, the frequency tuning. In fact the reported results demonstrate that the effective electric field over the resonance surface is the key parameter to determine the effectiveness of the electron acceleration. Slight variations of the input frequency considerably change the field pattern over the resonance surface, i.e. that zone where the wave energy is absorbed from the electrons and the great amount of ionizing collisions occur. Such variations of frequency not only determine the variations of the extracted current intensity, but it also affects the beam formation dynamics, as it has been experimentally observed during the experiment on the CAPRICE source (see section 4.2).

The situation described in section 4.7 is similar to the CAPRICE results discussed in 4.2, as the simulation results demonstrate that variation of frequency on the order of tens of MHz (over a pumping frequency of 14 GHz , like in in case of CAPRICE) strongly affect the electron heating, then also the source performances.

The results obtained with SERSE fed with TWT or Klystron microwave generators and shown in figure 4.2.1 can be also explained thanks to the simulation results. The TWT-based generators, in fact, are able to change the output frequency, then by using them it is possible to inject into the plasma

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chamber the best frequency in order to maximize the heating rapidity v_h . It is known that the beam formation strongly depends on the ion dynamics, then one can argue that the frequency tuning also affects the ion dynamics. As the ion dynamics is not included in our simulation model, then no conclusions can be done on the ion beam formation mechanism variation when the frequency changes. In section 4.13.8 we will investigate the influence of the field pattern variation over the resonance surface on the beam formation, but it will require the PIC simulations in order to investigate the ions response to the gradients of the electron densities when the field pattern changes.

The strong dependence of the ECR heating on the excited modes in the plasma chamber has been put in evidence by the simulations both in single-frequency heating and TFH. The simulations have also shown that, along with the usual effect of the electromagnetic waves on the electron acceleration, a strong ECR plug in effect occurs, which is at the basis of the increase of the electron confinement when the plasma is subjected to the combined effect of two electromagnetic waves. This explains why the best performances are obtained when the TWT amplifiers are used, because they permit one to make the power available at different frequencies. One would expect that the TFH gives the best results when the two frequencies are sufficiently far, so that the electron has a sufficient time to randomize its phase with respect to the second surface. However, in the case of SERSE, many interesting results were obtained with a separation between the two frequencies below 1 GHz. It is probable that in this case the tuning effect due to the use of a TWT plays the main role, and the use of the two frequencies optimizes the pattern of the electromagnetic field on the two resonance surfaces, this effect being stronger than the plug-in due to the double crossing of the ECR zones.

It is important to underline that, as a consequence of the power saturation trend, the TFH is a powerful method that is able to decrease the total amount of power needed to produce high charge states and currents. This is a great advantage for the stability and reliability of the ECR ion sources.

According to the results obtained with the VENUS ion source, the total amount of the X-Rays at high energies ($E > 100keV$) strongly depends on the values of the B_{min} and B_{ext} fields. For values of B_{min} above 0.85 T (the maximum field is about 3.3 T at the injection and about 2.1 T at the extraction) the current of highly charged ion abruptly decreases, while the amount of very hot electrons strongly increases. A similar transition in performances was observed also for the B_{ext} field variation. It looks

like that some magnetic field structure serves as good ion source, whereas others are more effective high energy X-ray sources, as hot electrons are effectively produced. Many efforts have been done in the last years in order to understand the creation of such electrons. In particular, the main goal is to find some methods that can reduce the amount of X-ray produced via bremsstrahlung, but on this purpose a complete comprehension of hot electrons creation mechanism is needed. Appearance of high energy electrons is detrimental not only for helium heating. Even if heating rapidity is quite low at low energies, it increases as energy increases, as it depends on gyro-radius. For low resonator quality, high energy electrons will reduce significantly the electromagnetic field amplitude because of their high ability of energy absorption, and this may lead to relaxation oscillations and unstable operations. Some related phenomena were described by Girard about the Quadromafios ECRIS [103] and other sources. These machines have 75 *cm* between mirror, a low mirror ratio and a low resonator volume. Any of these characteristics leads to bad performances.

The theoretical treatment reported in the previous pages has shed some light on the formation mechanism of ultra-hot electrons, but more detailed considerations will be done according to the results presented in the next part of the thesis.

4.13 A different theoretical approach to the plasma heating in ECRIS

The single particle approach used so far is not able to give detailed information about the EEDF evolution on long timescales. The numerical code used at INFN-LNS was able to simulate a relatively small number of electrons, and this negatively affects the investigation of the electron dynamics in the phase space in terms of the distribution function $f(\epsilon)$. Hence a completely different approach has been used to investigate the electron dynamics in ECR plasmas, based on the so called **PIC-Particle In Cell** simulations. The main advantage of this method is that the Maxwell equations are self-consistently solved in plasmas. This allow to evaluate not only how the electron population evolves under a given electromagnetic and magnetostatic field, but also how these fields react in turn. This will be of primary importance especially to determine the possible evolution of resonant modes inside the plasma chamber due to the presence of the plasma itself. In the next section the fundamental aspects of the code employed for the simulations are summarized.

4.13.1 Simulation description

The main goal of numerical experiments was to describe real phenomena that take place in ECRIS. Because reliable simulations are time expensive, as a first step, collisionless plasma in axially symmetric mirror trap in 2D R-Z approximation was looked for. No ionization was taken into account. The PIC code KARAT⁹ was used for the numerical solution of Maxwell and charged particles equations. Plasma chamber was modelled as a cylindrical volume (3.5 cm radius, 14 cm length) with ideal metallic walls, absorbing plasma particles. Pumping wave was launched through axially symmetric coaxial openings. The numerical accuracy was permanently controlled through the energy balance. Initially uniform plasma was chosen to occupy a volume larger than the volume surrounded by the resonance surfaces. Initial electron and ion temperatures were equal to few electron volts. Launched wave amplitude was able to provide maximum radial electric field in range of $1 - 2 \frac{kV}{cm}$, in order to obtain typical values for high power ECRIS. The dependence of heating on electric field amplitude was studied in the range $0.1 - 10 \frac{kV}{cm}$. A simple mirror configuration with mirror ratios in the range $1.7 - 4.5$ was chosen. The choice of a simple mirror configuration is

⁹The KARAT code has been developed by A. V. Tarakanov [33].

mandatory if one would to reduce the calculation times down to reasonable values (on the order of some hours). The simulation of a minimum-B structure implies the use of a 3D model which strongly enhance the complexity of the code. However, as we are interested to understand what are the influences of different mirror ratios and electric fields on the electron and ion dynamics, the Simple Mirror approximation is quite realistic. The choice of the plasma density, and then of the particles number, is of primary importance to ensure that the final results are reliable. As the higher is the simulated plasma density, the longer is the simulation time, a compromise has to be checked, in order to reduce the times of calculations without any appreciable reduction in the simulation accuracy. Then all the plasma parameters have been rescaled in order to ensure a quasi-real situation; for example, we know from the experiments that ions remain cold even after several μs of electron heating, because of the equations (1.2.8). This effect has a great influence on the discharge dynamics. The electron and ion density distribution evolve so that the charge neutralization occurs on distances higher than λ_D , but nevertheless some density layers can be formed on distances comparable or even smaller than the Debye length. In particular, we have to ensure that the relative relationships among the various plasma parameters are the same of a real experiment. Then the simulations were set so that the relations among all the main parameters like the standing wave nodes number, the particles number, the Debye length, etc. were fulfilled. Then for the less time expansive simulations a large volume plasma with $\omega_p/\omega_{RF} \sim 0.03$ was studied. However, some simulations have been performed with higher densities, rescaling also the other parameters, up to¹⁰ $\omega_p/\omega_{RF} \sim 0.1$.

In order to investigate the different dynamics of plasma located inside or outside the resonance surface, the simulations were divided into two groups: initially, an uniform plasma was chosen to occupy a volume larger than the inner resonance surface volume. The initial electron temperature was $5 - 7 eV$, whereas the ion temperature was $1 - 2 eV$. In most cases hydrogen ions were considered, in some cases heavy additive was used as "passive" impurity. This situation can be visualized by looking the figure 4.13.1. As stated before, the wave injection into the plasma chamber was simulated through a coaxial opening, and the initial electric field distribution can be calculated by the code, and it is shown in figure 4.13.2.

The PIC simulations are able to give us precious information not only

¹⁰The simulations at higher density were carried out mainly to investigate the effect of the plasma on the resonant mode pattern.

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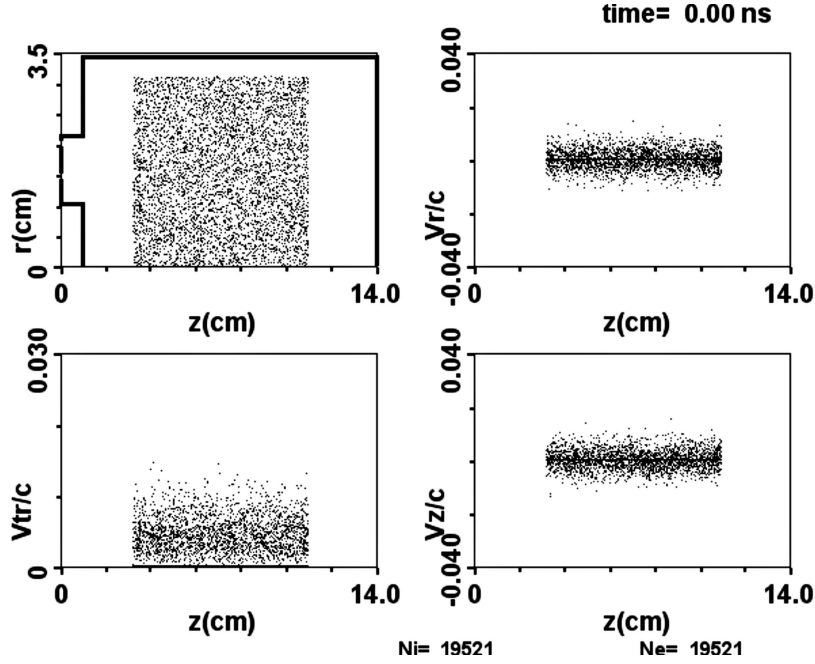


Figure 4.13.1: Plasma at the $t = 0$ instant. Clockwise from the top-left side: particle distribution inside the plasma chamber; velocity distribution along the radius, the z direction and in the transversal direction with respect to the plasma chamber axis according to the initial temperature of a maxwellian electron component.

about the electric field distribution influence on the particle dynamics, but also how in turn the particle dynamics affects the field distribution. The code gives an answer if whether (as proved by single particle simulations) or not the resonant modes affect the ECRIS performances.

To observe the long term dynamics of electrons and ions, the electric field distribution was chosen as smooth as possible. The possible effects of strongly spatially varying electric fields will be discussed in section 4.13.8.

Finally, it must be noticed that also in this case the plasma is considered collisionless. This also helps to reduce the simulation time, but on the other side it reduces the plasma evolution time that can be simulated. No ionization processes are taken into account, then no secondary electrons are provided by electron-ion collisions. Any consideration done in the following about the secondary electrons are only conceptual and they serve to make clear what occurs in the case of a real experiment.

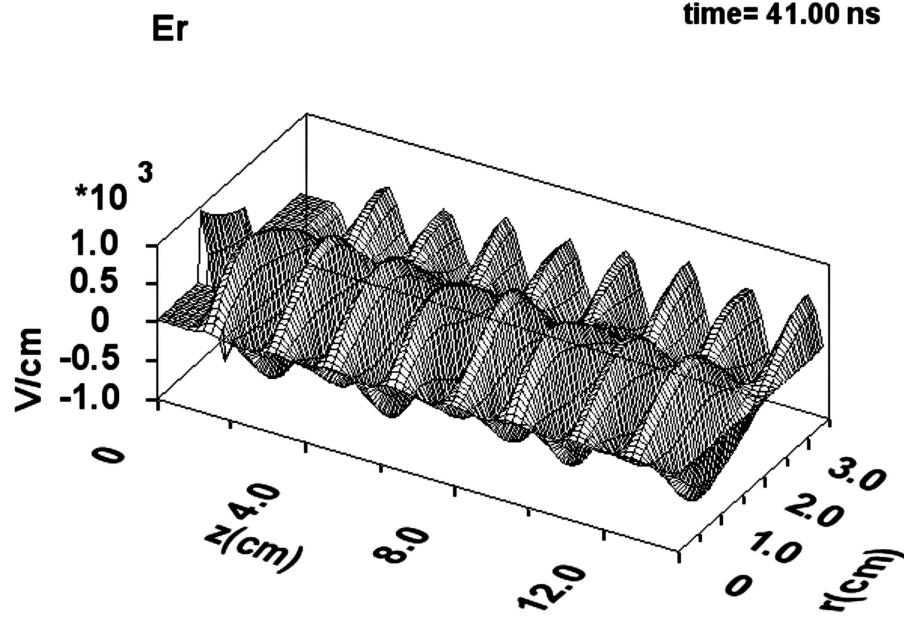


Figure 4.13.2: Radial component of the electric field at the beginning ($t = 0$) of the PIC simulation. The electric field distributes according to the resonator structure.

4.13.2 Basic results about the electron and ion dynamics

The first series of simulations analyzed were performed with a pumping power providing a maximum radial electric field variable between 1 and 2 kV/cm . A detailed analysis will be carried out in case of $\frac{B_{min}}{B_{ECR}} \sim 1.6$. In this case, the longest plasma confinement time was observed¹¹. A fast and effective acceleration of all the low energy electrons located near to the ECR occurs after a single pass through the thin region of resonance.

Just after 3 ns a considerable number of electrons overcame 1 keV of energy. A great amount of electrons are not accelerated because they had not

¹¹Note that this ratio between the minimum and the resonance field is the same of VENUS when the minimum amount of X-rays were observed (see section 4.5). However, in that case the source performances in terms of charge states and extracted currents were not the best.

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time to reach the resonance zone. One of the most important results which come from the simulations is the formation of a "rigid" electron density distribution structure. A quasi stationary state of plasma density distribution is formed on a timescale of several tens of bounce oscillations and its shape is shown in fig. 4.13.3.

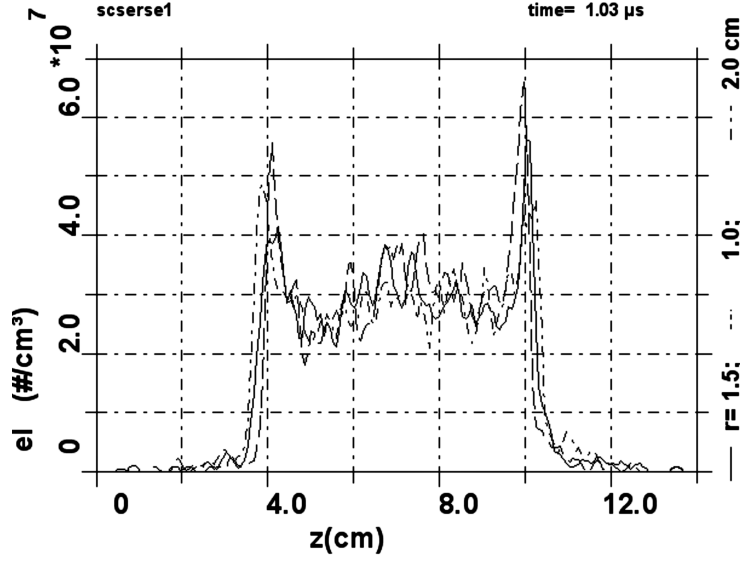


Figure 4.13.3: Spatial distribution of the electron density after 1 μ s.

The figure shows that the electrons are mainly confined inside the resonance volume, and the inner density is much higher than the outer one. This feature of the long term electron dynamics was predicted by Liebermann and Lichtenberg [36]-[37], but it was guessed to be a transitional state of the density distribution. On the contrary the simulations show that this structure is quasi-stationary and it is preserved for a long time also in presence of collisions.

If secondary, cold electrons were provided by the ionizing collisions, the simulations demonstrate that they are rapidly accelerated, thus becoming quite soon sloshing particles between the resonances. This mechanism generates the density structure observed in figure 4.13.3 and it can be theoretically explained. The electrons that reach the resonances and that are reflected inside them can be named **resonantly sloshing electrons**; on the contrary, those electrons reflected after the resonance crossing can be

named **non-resonantly sloshing electrons**. A detailed investigation of the electron dynamics and especially of the transition from the first group to the second one demands to consider the electron dynamics from an initial low energy state¹².

According to the theory proposed in [36]-[37], the impulse approximation permits to investigate the relationship between the perpendicular electron energy before and after the resonance crossing in the two limiting cases: resonantly sloshing (rs) and nonresonantly sloshing (ns) electrons. Mathematically, the difference between these two different cases is evaluated by means of the Γ parameter, that is:

$$\begin{aligned}\Gamma &= \left(\frac{2L\omega}{V_{\perp}} \right)^{\frac{2}{3}} = \left(\frac{2L}{r_e} \right)^{\frac{2}{3}} \\ L &= \frac{B}{dB/dl}\end{aligned}\tag{4.13.1}$$

where L is the characteristic length of nonuniformity of magnetic field B , ω is pumping frequency, $V_{\perp}$ is electron velocity perpendicular to the force line, l is the coordinate along the force line, r_e is the electron gyro radius. If the electron's perpendicular and parallel (with respect to force line) energies satisfy the inequality $\epsilon_{\perp} > \gamma\epsilon_{\parallel}$, γ being the relativistic factor, it will be called resonantly sloshing. In the opposite case the electron will be called non-resonantly sloshing. The perpendicular and longitudinal energies are evaluated at the resonance position. The values of the perpendicular energies after the resonance crossing for the two groups of electrons, ϵ_{ns} and ϵ_{rs} , are related with the perpendicular energy ϵ_i before the resonance crossing through the relations:

$$\begin{aligned}\epsilon_{ns} &= \epsilon_i + \Delta\epsilon_{ns} + 2 \cos(\phi_0 - \phi_p) \sqrt{\epsilon_i \cdot \Delta\epsilon_{ns}} \\ \Delta\epsilon_{ns} &= \epsilon_{osc} \frac{16L\omega}{\pi V_z} \\ \epsilon_{rs} &= \epsilon_i + \Delta\epsilon_{rs} + 2 \cos(\phi_0 - \phi_p) \sqrt{\epsilon_i \cdot \Delta\epsilon_{rs}} \\ \Delta\epsilon_{rs} &= \epsilon_{osc} 2^{\frac{4}{3}} \frac{L\omega}{\pi V_{\perp}}\end{aligned}\tag{4.13.2}$$

where $\epsilon_{osc} = \frac{mV_{osc}^2}{2} = \frac{e^2 E^2}{2m\omega^2}$ is a characteristic oscillatory energy, E is the amplitude of the wave electric field having the polarization in the same verse

¹²For simplicity we can consider as initial low energy electrons those ones created because of the ionization, or provided by electron donors or bias disks [63]. As low energy it is defined the energy significantly lower than the energy gained after a single passing through the resonance.

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of the electron gyromotion, V_z is the electron velocity parallel to the force line at the resonance, ϕ_0 and ϕ_p are the phases of the electron gyromotion with respect to the wave before and after the resonance crossing. These results were obtained by using the approximation that the electron motion in proximity of the resonance can be considered unperturbed by the wave field; otherwise, especially the low energy electrons are not free in proximity of the resonance, as their dynamics is strongly influenced by the wave. The effect of the wave on the particle dynamics can be estimated in case of resonances occurring not exactly at the bottom of the magnetic field lines¹³. A low energy electron that crosses the resonance coming from the high magnetic field side has a longitudinal energy on the order of the ponderomotive potential; being ϵ_{pp} this energy, we have:

$$\epsilon_{pp} = \epsilon_{osc} \frac{\omega^2}{4\omega(\omega_c - \omega)} \quad (4.13.3)$$

Near the resonance, this energy is on the order of:

$$\epsilon_{pp} = \epsilon_{osc} \frac{L}{4w} \quad (4.13.4)$$

where w is the width of the resonance zone defined in [36]-[37] as:

$$w = L \left(\frac{r_e}{2L} \right)^{\frac{2}{3}} \quad (4.13.5)$$

The value of ϵ_{pp} is not small, it exceeds typical energies of secondary electrons (i.e. electrons created as a consequence of the ionization processes). For example if $E = 1 \text{ kV/cm}$, $w = 1 \text{ mm}$, $L = 10 \text{ cm}$ and $\omega = 8 \cdot 10^{10} \text{ Hz}$, then $\epsilon_{pp} \sim 50 \text{ eV}$, with $\epsilon_{osc} \sim 2 \text{ eV}$. Then we can suppose that the great amount of the electrons, even the secondary ones, at the first resonance crossing have energies comparable with ϵ_{pp} . If we use the value calculated before and we insert it in the above equations (4.13.2) we obtain that an electron, after a single pass through the resonance, has an energy $\epsilon_{ns} \simeq 1.5 \text{ keV}$. Note that this value is in agreement with some electron energies determined in the simulations. Also in single particle simulations electrons with similar energies were observed. Electrons at these energy become resonantly sloshing and they begin to oscillate bouncing between the resonance regions. This applies in case of electrons coming to the resonance from a far initial position; other low energy electrons that appear (because of ionization) or

¹³For ECRIS the resonance position does not coincide with the magnetic field minimum value (i.e. the bottom of the magnetic field structure).

are reflected in the vicinity of the resonance gain a higher perpendicular energy ϵ_{rs} as they stay at the resonance for a longer time¹⁴. According to the equations (4.13.2) the energy gain for low energy electrons is practically phase independent, and it is constant for all the electrons (in fact for low initial energy the phase dependent term is small). Hence, the formation of a resonantly sloshing electrons population with energies in the range of several *keV* on a time scale of some bounce oscillations is a characteristic feature for modern ECRIS discharges. This is also confirmed by single particle simulations in the real ECRIS magnetic field structure: the figure 4.9.1 clearly shows that the electrons are mainly confined inside the resonance volume enclosed by the "egg-shaped" ECR surface. By analyzing more carefully the results summarized in section 1.11.2, it can be understood that according to [36]-[37], the dynamics of electrons with energies between the initial energy gain and the stochastic limit energy W_s is stochastic. Liebermann and Liechtemberg individuated two different regions in the phase space for the electron energy evolution. To some extent the heating is governed by stochastic processes and it is quite rapid. Once reached energies comparable with the stochastic boundary W_s , the heating goes on with a lower speed up to the absolute stochastic barrier W_b . After few tens of bounce oscillations the low energy electrons reach the stochastic boundary, then speed of "spreading" to higher energies reduces. The electrons diffuse till absolute adiabatic limit W_b with decreased speed. From the first equation of (4.13.2) it follows the scaling of the electron energy gain after a single pass: $\epsilon_{rs} \propto E^2$. According to section 1.11.2, the absolute stochastic barrier W_b scales as $E^{\frac{3}{8}}$. Hence for a high enough amplitude of the pumping wave, the steps in energy after a single resonance crossing become comparable with characteristic energy spread of the electrons. Then it may occur that the adiabatic barrier is overcome even after just one electron-wave interaction at the resonance. This mechanism is very important because it may help to explain a part of the experimental results reported in section 4.5. Some of the ultra-hot electrons may be created just after a single resonance crossing, thus overcoming suddenly the adiabatic barrier. Moreover, for electrons at very low energies the energy gain is higher. According to these results we can deduct the structure of the electron energy distribution function.

In fact, another significant effect that we observed in our simulations is the RF scattering of the electrons into the loss cone. The scattering due to RF diffusion determines electron lifetime. In the best cases it equals about 10^4 periods of bounce oscillations (from few μs till about $20 - 30 \mu s$

¹⁴The lower is the parallel velocity, the higher is the energy gain for the electrons.

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). It should be pointed out that the electrons are scattered into loss cone in the absence of collisions and only because of the RF scattering. Electron lifetime is not so long as traditionally thought. For typical parameters it is just one order of magnitude lower than that determined by Bohm value for cold plasma (see section 1.3.3). The non adiabatic effects leading to the μ invariance breakdown, discussed in section 1.6.1, are not so decisive to determine the mean electron confinement time, but nevertheless they determine the maximum electron lifetime (about 10^6 bounce oscillations in case of simulations parameters) for non RF-scattered electrons. This time is enough for the electron heating up to very high energies (hundreds of keV). It is obvious that the scattering is crucial for high energy features of EEDF, because it influences the number of electrons which can reach high energies, but it also influences the low energy part. If losses of low energy electrons exist, we can expect that low energy states will be occupied at smaller amount, and most of electrons will be located between the stochastic limit energy W_s and the absolute adiabatic limit W_b . The boundary of the plasma localized inside the resonance surface consist of high energy electrons only, and their energy lies between W_s and W_b . These energies for modern ECRIS correspond to few tens of keV . These electrons are very stable with respect to the action of static electric fields that develop at the boundary of resonantly sloshing electrons as a result of the ion reaction to the electron density drop. These static electric fields arise because of the ambipolar diffusion (see section 1.3.3). Hence, for ions born in the inner resonance part of the plasma the electrons' population is a sort of "rigid structure". Rather short electrons' mean lifetime explains the existence of problems related to the electron deficit in ECRIS.

A detailed analysis will be carried out in case of $\frac{B_{min}}{B_{res}} = 0.6$. Note that this is one of the typical condition in which the experiment reported in section 4.5 has been carried out. In that case the longest plasma confinement time was observed in the simulations. After few tens of bounce oscillations, the EEDF spreads till the stochastic limit W_s , corresponding to the energy at which any further spreading to higher energies becomes slow on the time scale considered, developing at speeds order of magnitude lower. On a timescale of several hundreds of bounce oscillations, the upper boundary of the EEDF stops at a value about 3.5 times larger than W_s . However, simulations on timescales of several thousands of bounce oscillations reveal that the heating goes on very slowly but electrons may reach few hundred of keV . In the simulations considered so far, $W_s \sim 18 - 20 \text{ keV}$, and it is reached approximately after 80 ns , while $W_b \sim 70 - 75 \text{ keV}$ is reached

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approximately after $2\mu s$. As a considerable number of electrons has been observed at energies well above W_b , this means that some additional stochastic mechanisms are involved in the plasma-wave interaction. In section 1.11 was underlined that an additional separatrix overlapping in the phase space may introduce other sources of stochasticity, thus permitting to overcome the adiabatic barrier. This barrier overcoming may explain the heating till very high energy in modern ECRIS like SERSE and VENUS.

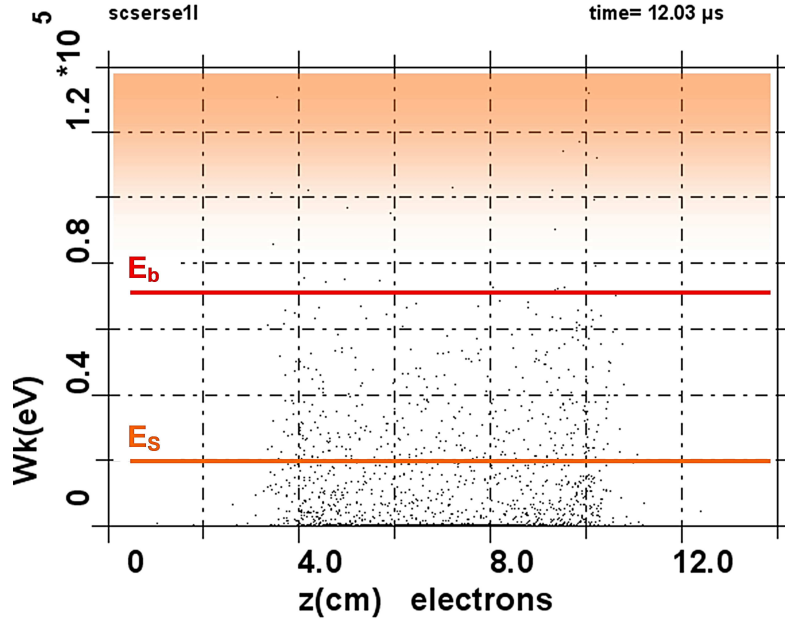


Figure 4.13.4: Electrons in the energy- z space. The heating above the absolute stochastic barrier is shown, together with the W_b and W_S energy boundaries. The figure refers to the electron energy distribution in space after $t = 12.03 \mu s$.

The figure 4.13.4 clearly features that few electrons have overcome the absolute stochastic barrier and they are above $100 keV$ of energy. After $25 \mu s$ electrons with a relativistic γ factor higher than 1.5 can be detected.

About the ion dynamics, it is clear from simulation that some of them are reflected in proximity of the ECR surface, others are accelerated away from the resonance. This means that the plasma potential (in particular the potential of the inner resonance plasma) is positive. The figure 4.13.5 shows the energy of the ions at different z positions; it is evident that some

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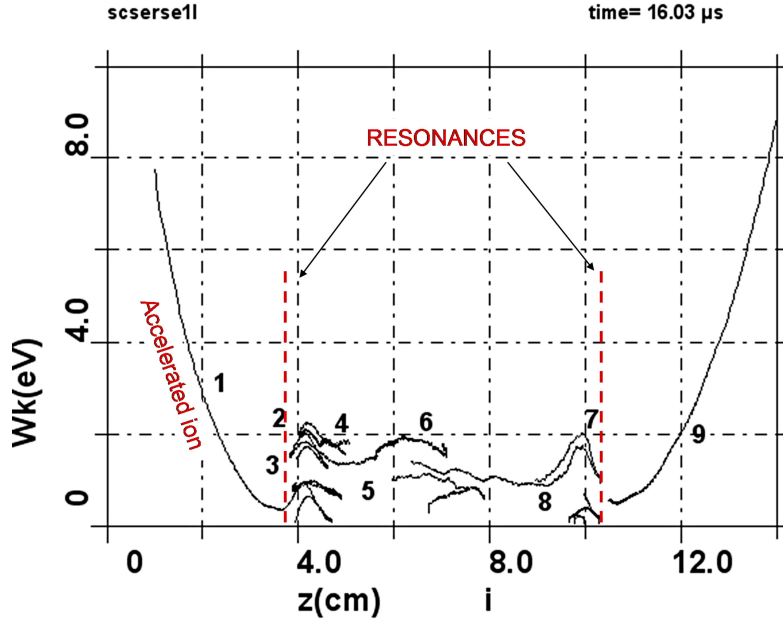


Figure 4.13.5: Energy of some proof ions with respect to the position along the z axis. The location of the two resonance zones is shown, and the ion acceleration outside the inner resonance plasma is evident.

ions are accelerated outside the resonance, because of the positive plasma potential which arises at the boundary of the inner resonance plasma.

Several simulations were made in order to verify whether the rigid and sharply bounded state of the plasma shown in figure 4.13.3 is an attractor for different initial conditions of the plasma distribution. In case of fully filled plasma chamber volume after $1\mu s$ (few hundred of bounce oscillations) the density of the outer plasma (which remains cold) was an order of magnitude lower than the inner density. Placing the plasma outside the resonance volume at the beginning, simulations demonstrate that it is partially trapped between the resonances and partially is lost at the chamber walls.

The figure 4.13.6 shows the initial location of the cold plasma. It is clearly displaced outside the volume enclosed by the two resonance surfaces. Quite soon the electrons crossing their closer resonance surface become resonantly sloshing electrons. The confinement time of the outer plasma is less than $1\mu s$, whereas all the electron passed through the resonance become resonantly sloshing and are well confined. After few μs a situation very similar

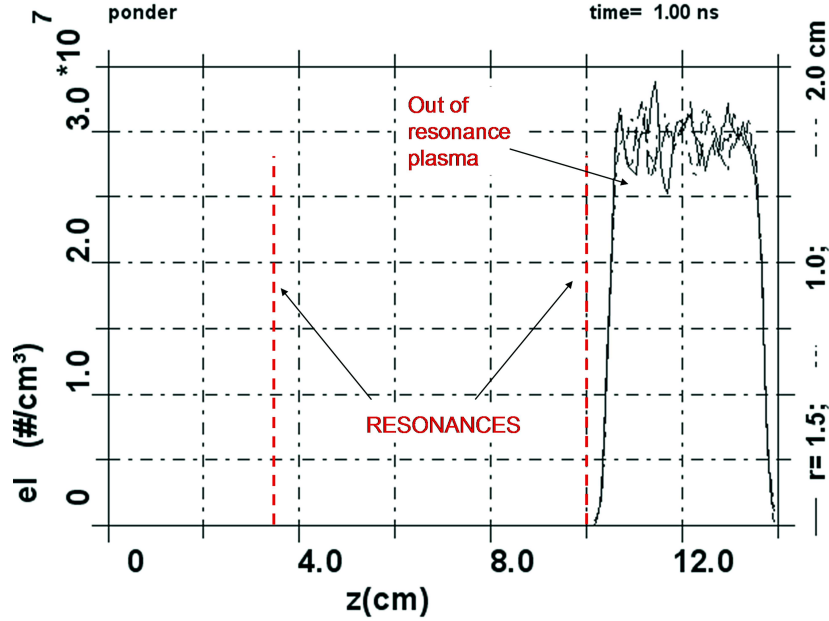


Figure 4.13.6: Initial plasma distribution inside the cavity: the cold plasma is located outside the inner resonance volume.

to the basic case of a plasma filling the entire cavity is established. This is confirmed by the figure 4.13.7. It is clear that the plasma has been plugged inside the inner resonance volume. The boundaries observed in figure 4.13.3 are completely restored and this affects again the ion dynamics.

4.13.3 Plasma effect on the electromagnetic field distribution

The main question concerning the results shown in section 4.7, for the explanation of the frequency tuning effect, is whether or not standing and structured waves exist in the plasma filled resonant cavity. Experiments shown in section 4.3 confirm that a resonant behavior persists also after the plasma ignition and in case of high densities.

By means of PIC simulations, because of their self-consistent nature, we can check for the effect of the plasma on the electromagnetic field. Preliminary results under the condition $\omega_p/\omega \sim 0.1$ for partly and fully filled resonators show that the plasma effect on the wave pattern is not strong.

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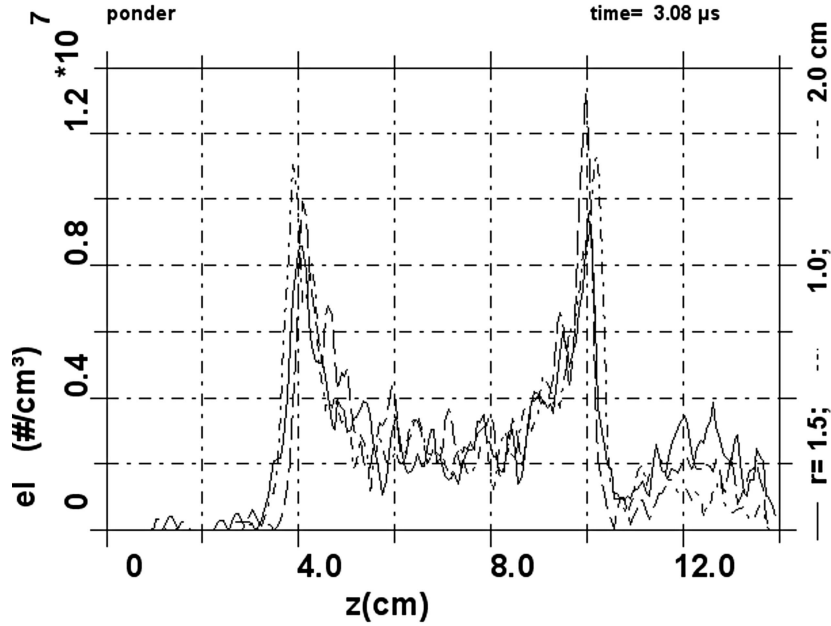


Figure 4.13.7: Plasma density distribution after about 8 μs . The outer plasma has been plugged inside the right resonance and reflected by the left surface.

The plasma only slightly redistributes the wave field, without any resonant effect. A nearly standing wave pattern can be easily maintained even in presence of the plasma; only slight variations in the amplitude and in the wavelength have been observed from simulations. This effect is based on three facts. First, inside the resonance surface the plasma density in typical ECRIS is lower with respect to the critical density; moreover outside the resonance surface the plasma density is significantly lower in quasi stationary state. Second, the system characterized by strong non-uniformities which do not affect coherently the waves propagation. Third, the resonantly sloshing electrons, that are formed under the influence of the wave, are stochastic and this means that they contribute very little to the coherent wave field. Hence, waves are well described, at least qualitatively, as a set of vacuum modes. For low density the simulation of the wave patterns are almost identical to vacuum modes. According to [104] at chamber pressures less than $5 \cdot 10^{-5}$ Torr, the frequency spectrum consists of a small number of peaks from almost undamped resonator modes. At the same time the plasma

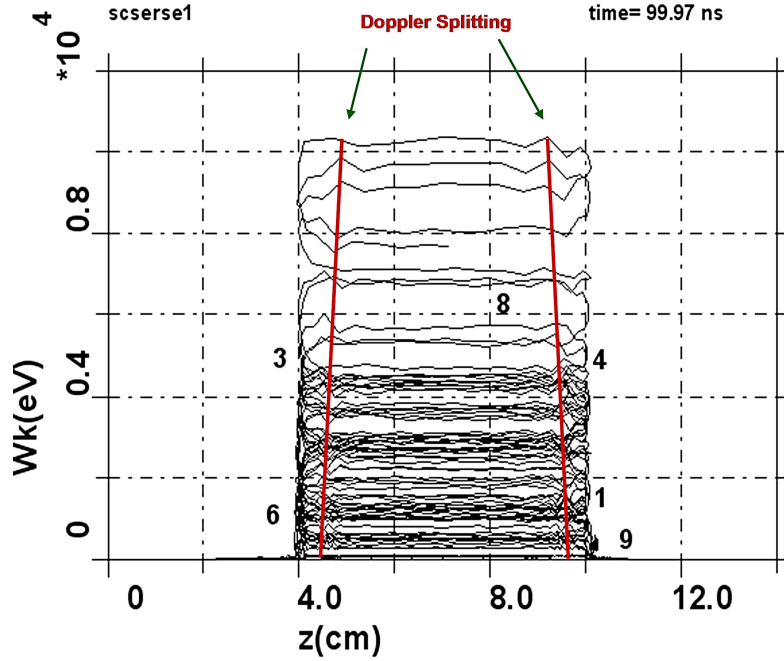


Figure 4.13.8: Evidence of a Doppler splitting of the resonance for the sloshing electrons.

remains quiescent. In [105] it is pointed out that high wave electric field strength is crucial for an effective interaction with an underdense plasma. The resonator provides such a high electric field through the excitation of a whispering mode gallery independently from the wave injection type. The property of the underdense plasma in a resonator enables to perform rather long-term simulations, lasting several tens of thousands of bounce oscillations, and this feature has permitted to provide an adequate modelization of the energy transfer to the electrons.

4.13.4 A more complex wave-plasma interaction: the evidence of the resonance Doppler splitting effect

The simulations have also revealed some energy changes at non-resonant positions, that can be seen in figure 4.13.8. These changes of energy in the inner part of the resonance volume are a signature of a Doppler splitting of the resonance. All the electrons undergo to another resonance because of

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the Doppler effect and of the standing wave nature of the excited resonator mode. For electrons with energy near 10 keV these energy variations occur, and for higher energy the Doppler shifted resonance is far away from the fundamental resonance. In figure 4.13.8 this trend is marked by the two red lines. The Doppler shifted resonance, for those electrons having the highest energy, is closer to the center of the magnetic trap. Hence the distance from the fundamental resonance grows with the energy, being this distance several times larger than the resonance width w defined in [36]-[37] as:

$$w = \alpha^{-1} \left(\frac{r_e}{2\alpha^{-1}} \right)^{\frac{2}{3}} \quad (4.13.6)$$

where r_e is the gyro-radius at the resonance, while $\alpha = \left(\frac{dB}{dl} \right) / B$ is the inverse of the magnetic field non-uniformity length. The mechanism which leads to the Doppler splitting of the resonance is quite complex as it involves also the resonant sloshing nature of the electrons and the standing wave properties of the electromagnetic field filling the resonant cavity. For simplicity, we can consider a standing wave existing inside a rectangular resonator¹⁵. For the wave electric field we can write:

$$\begin{aligned} E &= \frac{1}{2} E_{0-} [\cos(\Theta_+(t)) + \cos(\Theta_-(t))] \\ \Theta_+(t) &= \omega t + \vec{k} \vec{r}(t) \\ \Theta_-(t) &= \omega t - \vec{k} \vec{r}(t) \end{aligned} \quad (4.13.7)$$

where $\vec{r} \approx (const, const, z(t))$ is the time dependent radius vector of bouncing electrons unperturbed by the wave, E_{0-} is the amplitude of wave field with electron polarization, $\vec{k}$ is a wave vector defined by the particular excited mode. The equation (4.13.7) is the typical expression for a standing wave, that is formed by the superposition of two different waves traveling in the opposite direction¹⁶. The typical situation of a resonant sloshing electron bouncing inside the mirror trap in presence of these two waves is featured by the figure 4.13.9. When the electron moves from the center of the trap, towards the ECR zone (situation denoted by 1 in the figure), it feels one of the two waves traveling in its same direction, the other traveling oppositely. This means that in the region 1 it feels a blue shift of the frequency for the wave with $\vec{k}_2$, and a red shift for the wave with $\vec{k}_1$. The

¹⁵The treatment for the cylindrical cavities is similar but it needs a more complex mathematics.

¹⁶It is remarkable that a simple description like the one in (4.13.7) has been confirmed experimentally.

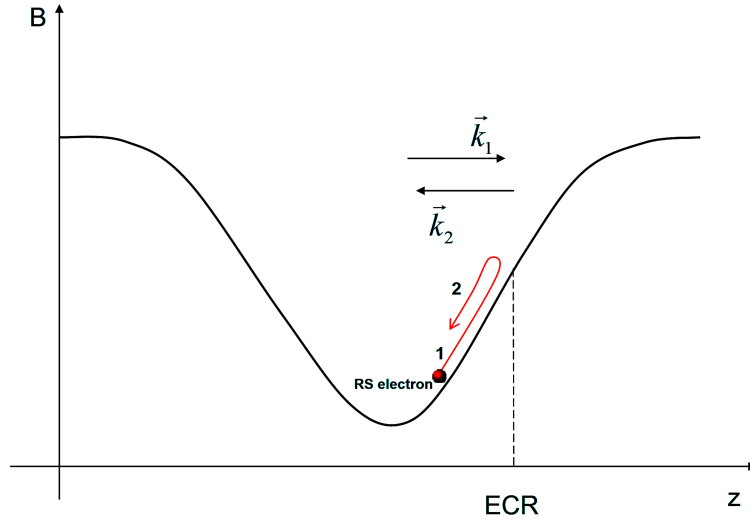


Figure 4.13.9: Resonant sloshing electron interacting with a standing wave inside the mirror trap.

amount of this frequency shift increases with the particle velocity (4.13.10).

It is clear from the figure that the electron is mirrored by the magnetic field before it reaches the blue shift resonance. This means that only the red shift resonance plays a role in the particle dynamics. Let us imagine an electron coming from the trap center with a parallel velocity $\vec{v}_{\parallel}$. It encounters the red shift resonance before the turning point, located in proximity of the unshifted ECR surface. During the interaction with the red shift wave, a small increase of its perpendicular velocity occurs. This leads to the reduction of the parallel velocity $v_{\parallel}$, with a mechanism similar to that occurring at the unshifted ECR, that is responsible of the resonant sloshing electrons creation. This sudden reduction of $v_{\parallel}$ reduces in turn the Doppler shifting. $v_{\parallel}$ continues to decrease during the particle approach to the unshifted ECR. Near this zone the $v_{\parallel} \sim 0$, then the dynamics of the electron is similar to what occurs in absence of the Doppler effect: it is accelerated at the ECR in the usual manner. Because of this mechanism the actual dynamics is not governed by a Doppler broadening of the resonance zone, but it is more similar to a **Doppler splitting which leads to an effective two fre-**

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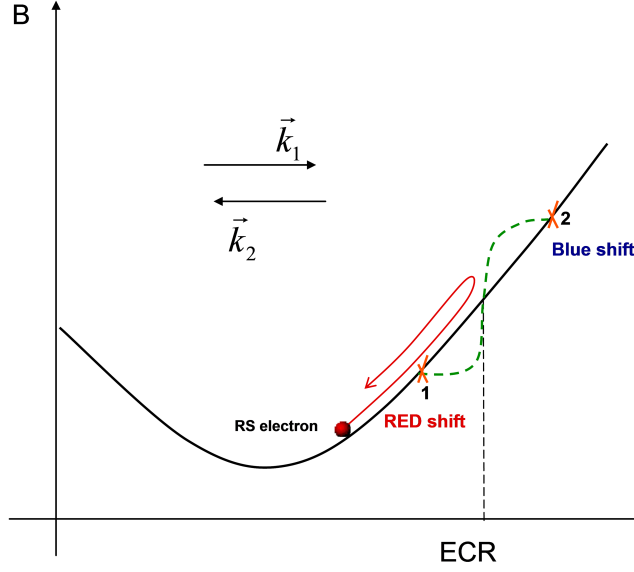


Figure 4.13.10: Doppler shifts of the resonance for a resonant sloshing (RS) electron.

quency heating¹⁷. When the electron turns at the unshifted ECR¹⁸, it will be accelerated along the force line because of the $-\mu\vec{\nabla}B$ force, and it will gain its parallel velocity again when it travels towards the trap center. This situation is denoted by 2 in the figure 4.13.10. Note that now the situation is perfectly symmetric with respect to the previous one (i.e. for the case 1). The electron will feel a red shifted frequency for the wave denoted by $\vec{k}_2$, and it will suffer again a second resonance. The amplitude of the electric field at the Doppler shifted resonance is 1/4 of the fundamental one. This because the electrons interact with just one of the two oppositely traveling waves¹⁹.

¹⁷This two frequency heating is quite different with respect to the mechanism described in section 4.10. In that case the frequency feeding the plasma were provided by two different generators, whereas in this case only one is the frequency, but it works, because of the Doppler effect, as two separated waves.

¹⁸Note that the ECR occurs at the same position as in the case of a Doppler unshifted frequency. This is due, as stated before, to the strong reduction of the electron parallel velocity in proximity of the turning point, located just in proximity of the resonance surface for resonant sloshing electrons.

¹⁹The standing wave interacting with the electron carries just half power of the incoming wave, as only the wave component with the same polarization of the electron motion can be absorbed at the resonance.

Note that not arbitrary splitting values are significant for long term dynamics. The value of the splitting should be such that two resonances overlap, that is the two resonance positions should be at a shorter distance than w . Otherwise, the electrons tend to turn within the inner resonance and the effect of the outer resonance is decreased [106]. The importance of the effective two frequency heating due to the Doppler effect as it constitutes one of the possible mechanism at the basis of the heating of very hot electrons.

Some additional words should be spent about the fine structure of the resonance splitting. For simple mirror configurations, with elongated force lines along the z axes, just one additional resonance is a good approximation. For ECRIS with minimum B configurations, a set of weaker resonances occurs at the periphery of the resonant surface, where the mutual orientation of the wave vectors and force lines is complex. Hence it is obvious that the complete set of resonances is strongly influenced by the conditions of the resonator modes excitation, i.e. by the number and spatial structure of the excited modes. Once again it may be concluded that the mode structure becomes crucial to optimize the performances of the ECRIS.

4.13.5 Heating above the absolute stochastic barrier

By looking to figure 4.13.4 it is clear that an above barrier heating occurs.

The dependence of heating on electric field amplitude was studied in the range $0.1 - 10 \text{ kV/cm}$. A simple mirror configuration with mirror ratios in the range $1.7 - 4.5$ was chosen. The upper energy limit that we observed reaches values of few hundreds of keV after $30 - 50 \mu\text{s}$ in typical cases. In cases of 7 kV/cm and low mirror ratio, the upper energy reaches the relativistic regime ($\gamma > 2$) after few μs . Again, the simulations confirmed that low mirror ratios encourage the formation of very hot electrons, as experimentally observed. In addition the wave amplitude, and then the resonator properties, influences the EEDF. As the wave amplitude at the resonance may be affected even by slight variations of the pumping frequency, because the resonator mode changes, then the EEDF may be strongly affected by some tunings of the pumping frequency.

The simulations were carried out for different positions of the resonance with respect to the center of the mirror trap. The case of resonance near the mirror bottom has some peculiarities. It was observed that the electrons overcome the adiabatic barrier if the wave amplitude E_0 or the L parameter are large enough. In case of electric field amplitude of few kV/cm the electrons overcome the adiabatic barrier only if the mirror ratio was lower than 2.5.

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Under the conditions for overcoming the adiabatic barrier, when the resonance is located near the bottom of the magnetic trap we obtain the highest EDF temperature. On the contrary (for that cases when only the fundamental resonance plays a role²⁰), the temperature and number of high energy electrons are smaller for a resonance close to the mirror. The Minimaños ECRIS [107] was an excellent example of trap of such kind. Its configuration leads to the smallest EDF temperature among all sources of such type. For nonrelativistic energies most of the electrons that overcome the adiabatic barrier remain resonantly sloshing; their reflection points approximately follow relativistic changes of the resonant magnetic field value [108]. These electrons are confined for a long time, and the energy evolution looks like a diffusion process. When γ further increases, the reflection point is located well before the ECR.

As example we may consider the case of a large amplitude wave, with $|\vec{E}| \sim 7 \text{ kV/cm}$. An electron having a relatively high γ is heated at the fundamental resonance. At this energy the mirror magnetic field value was 1.35 times larger than the non-relativistic resonant one, while the γ value of the electron, reflected near the mirror, is about 1.8. This means that the ECR can be obtained with a magnetic field $B_{ECR}^1 = \frac{\gamma m}{q} \omega_{RF} = \gamma B_{ECR}$, where B_{ECR} is the non-relativistic field at the resonance. Such an electron is reflected well before the corresponding relativistic resonant magnetic field. This effect is more evident for electrons heated at second harmonic (i.e. at $B = 1/2 B_{ECR}$): an electron reaches the fundamental resonance position at $\gamma = 2.6$. In case of heating at the fundamental resonance, the maximum obtained energy is one third higher than the one evaluated theoretically from the maximum fundamental resonant electric field value. The same ratio is valid for second harmonic heating. Electrons heated at the second harmonic appear at energies lower than W_b . At high energies they bounce between points, whose position depends on their energy. To observe the heating at second harmonic either large L values and large E_{0-} values are needed. The former because we need not so high mirror ratios, so that the electron can reach the resonance without the mirror reflection, the second one because we need high γ values to have the second harmonic surface inside the plasma chamber (i.e. the second harmonic B field must overcome the minimum trap field). This mechanism works better if no Doppler splitting of the resonance occurs.

²⁰i.e. when the Doppler splitting of the resonance does not play any role, as well as any second or higher cyclotron harmonics absorption.

Summing up, high energy tails of the EEDF depend on the force line length L , and on the position of the resonance with respect to the trap bottom and to the mirror. The most effective heating was observed in cases of resonance near bottom.

Next we investigate the importance of the two frequency heating. We are interested to study those electrons with high energy, as we would like to understand how the heating till extremely high energies occurs. Being ω_{refl} the nonrelativistic gyro frequency at the reflection point, according to the calculations reported in [101] we can state that $\omega_{refl} \leq \omega_{RF}$ in order to have an effective heating due to the resonance splitting. This result is in accordance with our simulations (for electrons with $\gamma > 2$, $\omega_{refl} \simeq 0.3\omega_{RF}$).

Liebermann and Lichtenberg developed a theory [106] for the two frequency heating which takes into account the changes of the perpendicular energy at the resonance which lead to changes in parallel energy at the second resonance. This leads to a four dimensional phase space dynamics. Then the separatrix overlapping discussed in section 1.11 is mainly due to a more complex structure of the phase space. It cannot be further represented as in section 1.11, and the two additional dimensions lead to a possible new way for the electrons to overcome the energy limitations imposed in a two dimensional dynamics. The presence of two resonances produces two important diffusion processes: a slow diffusion of the $v_{\perp}$ component above the W_b , and a new mechanism for diffusion in $v_{\parallel}$. These diffusion phenomena are similar to the diffusion caused by the interparticle collisions.

The phase space representation of the two frequency heating can be given by means of the so called *actions* [106]:

$$J \propto W_{\parallel} \quad (4.13.8)$$

$$I \propto \left(\frac{2W_{\perp}}{\rho} + W_{\parallel} \right) \quad (4.13.9)$$

$$\rho = 1 - \frac{1}{R_1} \quad (4.13.10)$$

where $W_{\perp}$ is the reduced kinetic energy perpendicularly to the force lines and measured at the first resonance, as well as $W_{\parallel}$, that is the reduced kinetic energy along the force lines. R_1 is the mirror ratio evaluated at the inner resonance. The electron dynamics can be qualitatively represented as in figure 4.13.11.

In a simplified model of energy independent increments after single pass through resonances (adequate for a weakly relativistic case) electrons with a positive perpendicular energy are above the line $I = J$. The further from

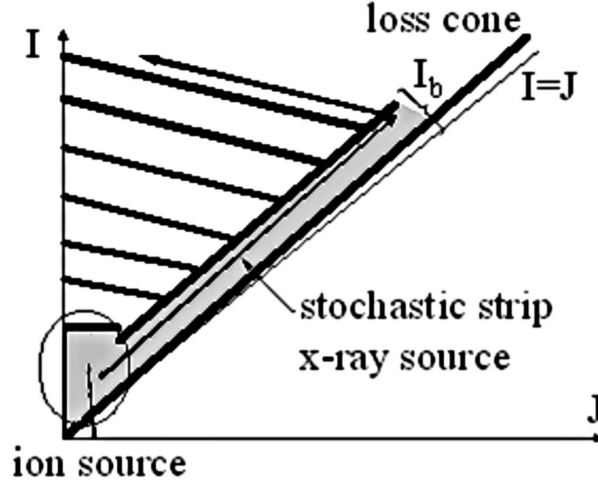


Figure 4.13.11: Description of the phase space in terms of the two actions I and J . The loss cone boundary is represented by the line near to $I = J$. The quasi-horizontal lines feature the regions where the Arnold diffusion may occur, leading to the formation of very high electron energies. The stochastic strip is mainly populated by high energy electrons especially in case of high mirror ratios and resonance near the bottom of the magnetic trap. For good performance ion sources the electrons must be located inside the circle at low I and J values.

this line an electron lies, the higher perpendicular energy it has. In the figure the general case is shown, where the loss cone line deviates from the line $I = J$ weakly. If the difference is neglected, electrons are absorbed at this line. In case of near mirror resonances, the loss cone strongly reduces the available volume for electrons, especially at high energies. In grey the so-called *stochastic sea* is represented, consisting of two parts: the one having a small longitudinal energy is the base for highly charged ion sources, as it contains many electrons with proper energy for the ionization. Because of the large initial perpendicular energy gain, the source of electrons is located at small values of J , i.e. inside the circle of figure 4.13.11. The other part is the strip that is a base for an x-ray source²¹. The electrons contained

²¹X-ray sources are able to produce a great amount of high energy electrons which generate X-rays via bremsstrahlung. The simulations demonstrate that the conditions to obtain good performance X-ray sources are different from that ideal ones for ion sources.

in this strip are characterized by a high parallel energy resulting from the rf scattering. The width of the strip corresponds to an energy near to W_b . Nearly horizontal thin strips contain phase resonances along which a slow Arnold diffusion occurs. The problem of high energy electrons is strictly connected with the electrons inside the stochastic strip. On a long timescale the thin strips are filled from the thick stochastic strip via Arnold diffusion. The electrons diffuse in a thin strip without losses, then after long enough time they become uniformly distributed along the thin strips; the value of f is equal to that at the stochastic strip edge. Hence, the values of f along the stochastic strip determine the high energy tail of the EDF. A more quantitative treatment can be carried out, and it is completely reported in [101]. Here we report the qualitative dependance of the temperature with respect to the fundamental source parameters: the force line length L , the pumping frequency ω and the electric field strength E . We obtain for the temperature:

$$T = \frac{2}{\pi} \frac{\delta\omega}{\omega_1} \frac{1}{\rho} W_b \quad (4.13.11)$$

where ω_1 is the inner frequency, $\delta\omega$ is the distance in frequency between the two resonances, and W_b is the absolute stochastic barrier. The formula (4.13.11) shows that for resonances far from the bottom of the magnetic trap the temperature is few times lower than W_b . Note that the frequency ω_1 differs from the fundamental resonance on a value of the order of $\delta\omega$. Then if the gyro-frequency at the center of the trap differs from the resonance one of a value less than $\delta\omega$ the energy changes differ from that considered in the model proposed by Liebermann and Lichtenberg. In that case the exact calculation of the temperature dependence on the above parameters must be calculated for the so called near bottom case.

Let us define with ν the characteristic length linked to the uncertainty in the determination of the inner resonance position due to the Doppler splitting. This width of possible locations of the inner resonance is due to the energy distribution of the electrons. Be $s_1 = \frac{l_1}{L}$ the dimensionless variable indicating the position of the first resonance along a given field line. Considering the case of our simulation, where the inner resonance is provided by the Doppler splitting of the fundamental one, and taking into account the case of *strong heating*, i.e. when the highest temperatures are observed, for resonances near the bottom of the mirror field the temperature has the weakest dependence on the electric field intensity, a moderate dependence on the pumping frequency, and the strongest dependence on the resonance position and on the force line length L . The main result coming from these

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considerations is that the electron dynamics must be described locally: the electron temperature differs in the different zones of the magnetic trap; the regions where the field lines are longer (in terms of L) are the favorite candidates to work as a privileged source of high energy electrons. The description in terms of the L parameter is detailed, but the gradient determination may help to understand whether or not a given source produce more or less high energy X-rays.

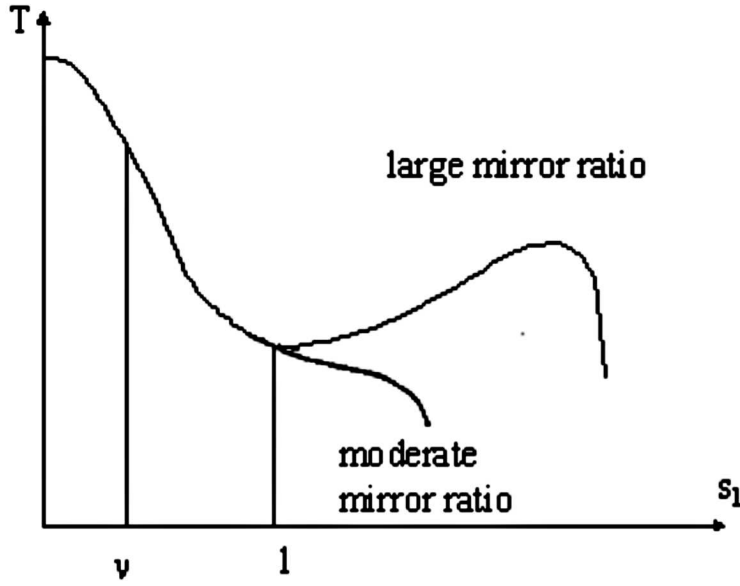


Figure 4.13.12: Trend of the electron temperature with respect to the location of the resonance position along the field lines.

The qualitative temperature dependence on the resonance position is shown in figure 4.13.12. In case of one-stage heating²² the temperature drops in case of heating at resonances near the mirrors. In this case the influence of the loss cone is strong. In fact, for small L and fundamental resonance near the mirror the electrons with energies larger than W_b were not observed in simulations. The temperature growths observed in the figure are due to two different reasons: for $s_1 > 1$ the temperature grows as W_b

²²For one-stage heating we mean the heating without any role played by the second or higher cyclotron harmonics.

increases²³; on the contrary, the strong increase of the temperature for $s_1 < 1$ is due to the strong influence of the two frequency heating. Practically, in this case the Doppler splitting provides a wide region filled by resonances near the bottom of the trap, each corresponding to a plasma electron. In this case a sort of heating with a quasi-flat magnetic field occurs, leading to an impetuous production of high energy electrons. Experimentally, this condition is detrimental because of the high production of X-rays at very high energy. Note that despite W_b decreases when the distance from the trap center decreases, the temperature increases. This model also provides a powerful method to estimate the electron temperature for ion sources using different rf frequencies. The higher is the frequency, the higher is the electron temperature. This result is in agreement with the experimental data obtained with SERSE and VENUS.

4.13.6 Discussion about the electron dynamics and comparison with the experimental results

The simulations point out that the electric field amplitude at the resonance is of primary importance, more than the RF power. There is not a direct relation between the electric field amplitude and the RF power, as the field strength also depends on the excited mode(s) pattern [4], [5].

In the case of the two frequency heating the heating rapidity increases, because of the raising of the adiabatic barrier limiting the heating from below. It can be optimized by setting opportunely the distance between the two resonances corresponding to the two applied frequencies (or the fundamental and splitted frequency in case of the single frequency with Doppler splitted resonance). Similar results about the rapidity of the heating have been obtained in section 4.10 through to the single particle simulations. The so called Two Closed Frequencies Heating Effect has been investigated as a tool to increase the number of warm electrons in short time. The PIC simulation and the two frequency heating theory confirm that this mechanism decreases the time needed for the electrons to pass from low energies to energies proper for the ionization in a short time. The single particle simulations also help to link this increase of heating rapidity with the relative phase of the two waves. Anyway no clear indications about the role played by the phase come from PIC simulations, and it will be investigated in the next months. However, as mentioned already in section 4.10, this mecha-

²³From the Liebermann and Lichtenberg theory it comes out that the absolute stochastic barrier W_b increases as the resonance position moves away from the trap center. Here s_1 represents the dimensionless distance from the first resonance point.

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nism may lead to an excessive production of very hot electrons, and PIC simulations confirm that in case of long L , a great number of above-barrier electrons may appear that explains the high energy X-rays flux. The use of two close frequencies appears more as a powerful method to reduce the RF power injected into the plasma chamber to obtain a given CSD, thus increasing the reliability of the ECIRS, or, even better, to operate in conditions of very high mirror ratio. The heating rapidity due to the TCFH may help to avoid any peaked trend of the CSD, thus preserving the high performances without any strong production of high energy X-rays.

The ECRIS with minimum B magnetic structures have characteristic features connected with the strong non-uniformity of the magnetic field. The near axes force lines differ strongly from that at periphery, and the values of L differs several times, especially in case of high-B mode conditions. The structure of the field lines is quite complex, as shown in figure 4.6.6, and it is not available a detailed description of the plasma dynamics but only some general guidelines. Hot electrons at periphery (because of smaller L) are expected to be less significant, but much more sensitive to any variation of outer magnetic field, such as the extraction magnetic field B_{ext} . Small changes of these fields can increase the L value at the periphery. The effect of the B_{min} variation is even more pronounced, as it increases globally the value of L . This corresponds to the experimental data, showing a strong increase of the electron temperature for low mirror ratios, and especially for low B_{ext} of high B_{min} . It can be viewed also in terms of the global magnetic field gradient over the resonance surface, as shown in figure 4.11.3: by increasing B_{min} , the global distribution of the magnetic gradient over the resonance surface strongly changes. Note that L is practically inversely proportional to the gradient. In particular, considering the B_{min} and the B_{ECR} field for the equation of L in (4.13.1), $L \equiv \left[(dB/dl)^{-1} B \right]$, i.e. it is proportional to the inverse of the field gradient along the given field line. Then the treatment of the X-ray analysis presented in section 4.11 gives some preliminary information about the trend of the hot electrons production and their energy.

At first glance it seems that for mirror ratios like $R > 4$ the creation of ultra-hot electrons is practically suppressed. However, as in min-B ECRIS different regions of the plasma cavity have different mirror ratios, the role played by the magnetic field structure is fundamental. Therefore the simulations explain the results obtained with the VENUS source; for $B_{min} > 0.85 \text{ T}$ (the maximum field is about 3.3 T at the injection and about 2.1 T at the extraction) the current of highly charged ion abruptly decreases, while the

amount of very hot electrons strongly increases. A similar transition in performances was observed also for the B_{ext} field variation [12]. According to the formalism followed in the previous sections, we can say that the variations of s_1 caused by the variation of B_{min} were the most important. The smaller the dimensionless distance to the bottom s_1/L was, the higher the temperature was.

The position of the mirror points and of the resonances is crucial. By increasing B_{min} one reduces the distance between the resonances and also the distance from the bottom of the magnetic trap, thus leading to higher T_e . The B_{min} increase also increases remarkably the field lines lengths L , thus leading to higher temperatures.

Furthermore, from simulations it comes out that the second harmonic play a role in some cases. If even the second harmonic heating takes place (because of the relativistic variation of the $1/2B_{ECR}$ field), together with the Arnold diffusion, the temperature increases and a two-stage heating develops. This effect will be further investigated in future, as it may explain some experimental results like that shown in figure 4.5.4. The high energy component of the VENUS plasma features two different temperatures, and the highest one differs more from the lower when the extraction field B_{ext} is strongly reduced. May be this is a sign that for lower mirror ratios the second harmonic heating plays a decisive role to further increase the high energy tail of the EEDF.

To link the discussion of simulations described in section 4.11 with the results about the L dependence of the temperature, we can say that the discussion on the gradient was mainly based on a single particle passing through the resonance just one time. In that case we determined the possible energy one electron may gain by crossing the resonance in different conditions of magnetic field gradient and electric field strength. This approach give us some estimations of the energy acquired during a single interaction of the wave with the particle. More than a prediction about the final EDF, it gives us some information about the rapidity to overcome the first energy barrier W_S , and the number of electrons that are possible candidates to overcome W_b . It is clear that in case of gentle gradient regimes (i.e. long L) the number of electrons which quickly approach W_b is very high just after few crossing of the resonances. On the contrary for strong gradients, although a non negligible probability remains to overcome W_b , the number of electrons is small because the energy gain per single pass is relatively small. This reduces the number of very hot electrons.

The trend of the CSD observed in figure 4.5.1 was connected with the

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excessively strong gradient, which reduces the heating rapidity at intermediate energies needed for an effective ionization. This explanation is confirmed by the figure 4.13.12, which features the temperature drop for resonances distant from the trap center. Even if a not so marked drop was observed experimentally for the hot component temperature, may be one still stay far enough from the pronounced temperature drop. The heating rapidity, according to the single particle simulations, may be the key point. The strong gradient regime strongly limits the temperature of the warm population, the only one affecting the CSD.

A confirmation of the above considerations comes from the Quadromafios ECRIS [103]-[109], that has the worst performance from point of view of minimization of x-rays generation. This is not surprising, it has the largest L value ($\sim 1.5L_{VENUS}$), and the resonance was not far from the bottom. As a result Quadromafios typical temperature at frequency of 10 GHz equals the highest VENUS temperature detected at 18 GHz.

The PIC simulations also help to better understand how some additional tools used for ECRIS work (like bias disks and electron donors). They affects both the electron and ion dynamics (the latter as a consequence of the former). The application of biased disks for the enhancement of the highly charged ion currents is a common tool. In [64] a temperature increase under the influence of a biased disk was observed. It is likely connected with the reflection back to plasma of electrons from the stochastic strip at the loss cone boundary. The temperature growth with biased voltage is especially interesting. The higher is the voltage the more effectively the electrons are reflected. Negatively biased disk placed not far from mirror increase the lifetime of the moderate energy electrons having sufficient parallel energy.

4.13.7 The ion dynamics

The fundamental aspects of the ion dynamics have been discussed in section 4.13.2. It was clear that the inner resonance plasma produces a positive potential which accelerates the ions. By means of the equations reported in section 1.2 about the Coulomb long range collisions, it comes out that, for ECRIS typical parameters, the ion collision frequency in the inner resonance plasma is higher than the gyro-frequency. This means that the mean free path λ_{mfp}^i is smaller than the Larmor radius r_{ci} . This implies that ions of the inner resonance plasma are collisional and cannot be confined by the magnetic field. Outside the resonance the plasma density drops of about one order of magnitude and the ions become collisionless, being confined by the magnetic field. The ion dynamics changes according to the region

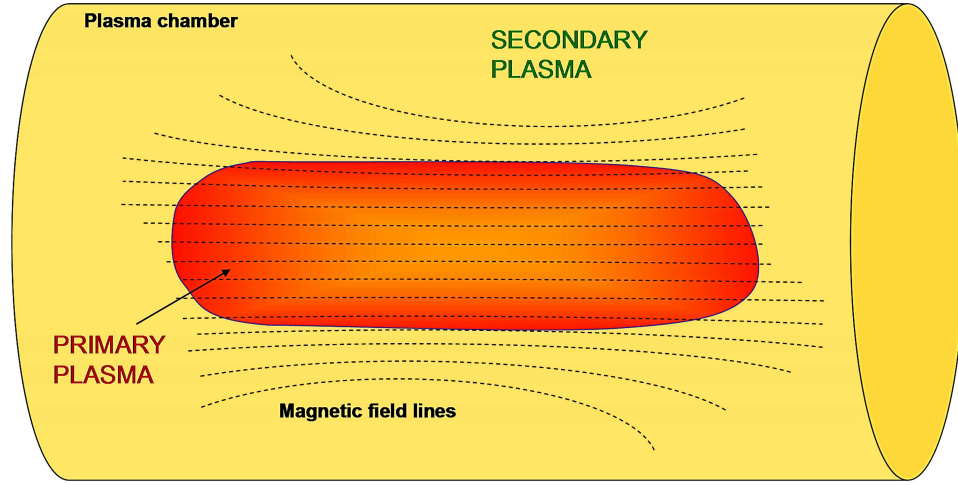


Figure 4.13.13: Long-term plasma distribution: distinction between the primary and secondary plasma for ECRIS.

where they are moving. The figure 4.13.13 shows the typical distribution for a plasma inside an ECRIS plasma chamber, in case of a minimum-B configuration. Although the simulations were performed in a simple mirror field, we can generalize, to some extent, the main results. Note that the plasma inside the resonance surface is the **primary plasma (PP)**, whereas outside the **secondary plasma (SP)** is located. The primary plasma is dense and hot, whereas in the secondary plasma the electrons and the ions are escaping from the primary plasma and consequently they have a lower density. The positive potential is localized on the PP-SP boundary. Note that although in figure 4.13.3 a two peak density distribution is shown, actually in the central part of the plasma chamber the density is quite flat and it coincides approximately with that one at the resonance. This because of the ambipolar diffusion which tends to equalizes any potential difference inside the plasma. Without the density flattening, low frequency ion oscillations may arise. The potential at the PP-SP boundary is quite stable because of the rigid structure of the resonant sloshing electrons. As the energy of the resonantly sloshing electrons lies mainly in the range of few tens of keV , the effect of the positive boundary potential on these electrons can be ignored. To investigate the ion dynamics inside the layer one can consider electron density distribution in layer as constant. Ions on their way through the layer need to match their density with fixed sharply dropping electron density, be-

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cause of the ambipolar diffusion. This is possible if ions are accelerated and reflected in part as well. The space scale on which the ions are reflected is on the order of the ion Debye length $\lambda_i \sim V_{Ti}/\omega_i$, where V_{Ti} is ion thermal velocity and ω_i is the ion plasma frequency. For real ECRIS, this distance is on the order of 10^{-2} cm. The other space scale is connected with density drop $L_d = ((dn_e/dz)/n_e)^{-1}$. The typical value of this length ranges from millimeter to a fraction of millimeter. On this scale ions are accelerated. The exact solution of ion beam problem under such conditions is complex. The accelerating layer is formed around the resonant surface, with the accelerating field directed perpendicular to the surface, as featured by the figure 4.13.14. Figure 4.13.15 features a possible ion trajectory in a minimum-B

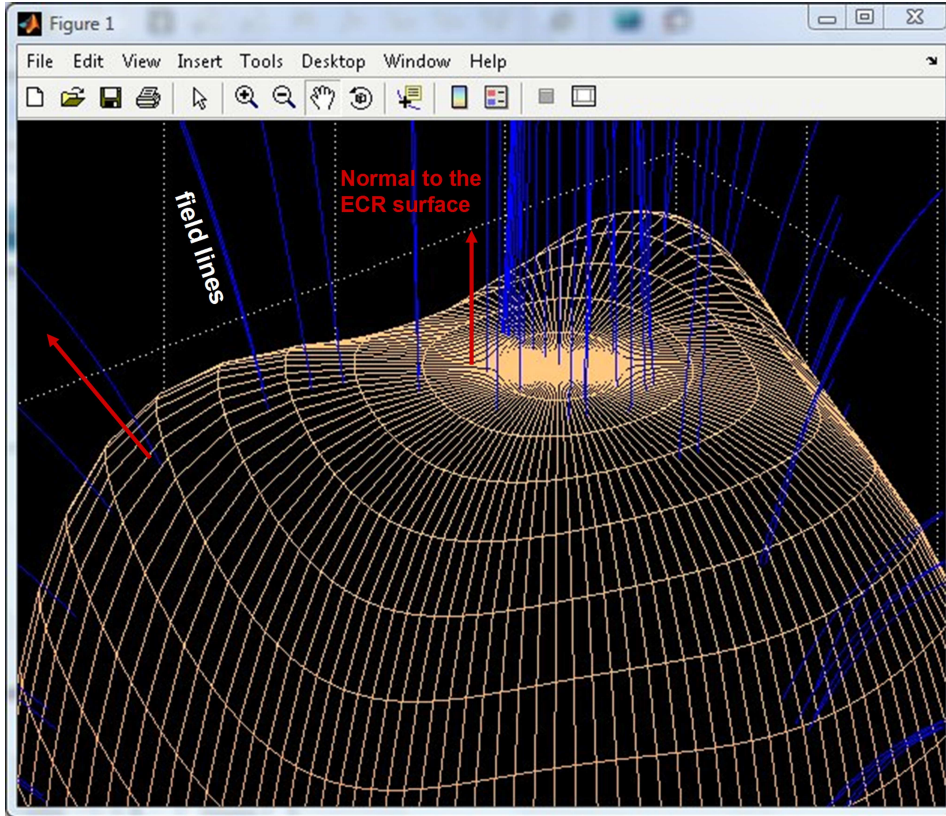


Figure 4.13.14: Detail about the ECR surface in a minimum-B source with the field lines (in blue) and the direction of the normal to the surface, which coincides with the electric field direction due to the plasma potential.

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structure. An example of ion trajectory is shown inside and outside the ECR surface. The inner trajectory is characterized by continuous scattering due to the collisions. Ions are not magnetized but nevertheless they are dynamically confined because of the electron confinement. The main question is that because of the quasineutrality the ions cannot escape in a greater amount with respect to the electrons: the electron and ion currents and densities must be about equal at any point inside the plasma chamber, and any change leads to the formation of the plasma potential.

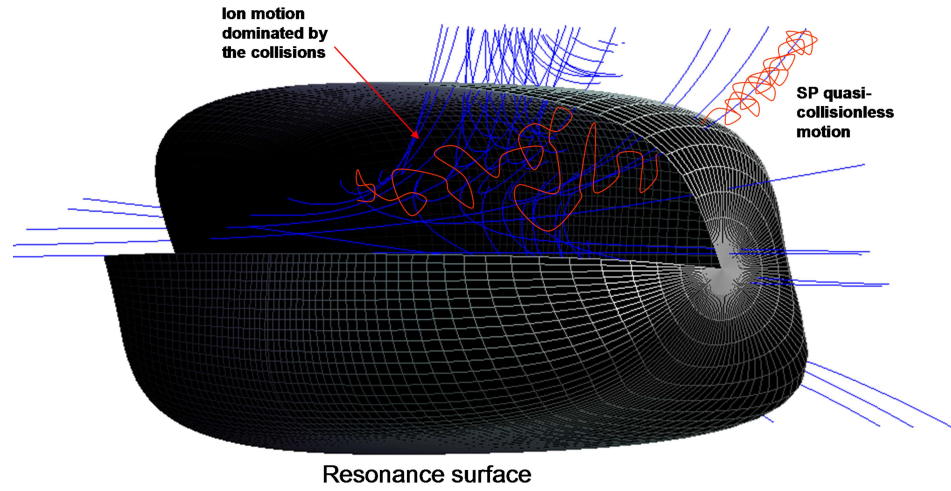


Figure 4.13.15: The resonance surface is shown with a part of the field lines: an example of inner and outer (with respect to the resonance volume) ion trajectories is given.

An ion passing through the accelerating layer will be accelerated perpendicularly to the resonant surface. For the near axis field lines (i.e. that ones on the top of the resonance surface in figure 4.13.14) the ions acquire a relatively high longitudinal velocity (with respect to the field line); as the perpendicular velocity is on the order of the thermal one (temperature is about 1 eV), and the longitudinal energy kick is of 10 eV at least (because of the primary plasma potential), the ions are accelerated in the near axis region, although they are magnetized outside the ECR surface as explained in figure 4.13.15. Finally they cannot be confined for a long time because their velocity vector is injected inside the field line's loss cone as a consequence

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of the PP boundary acceleration. Note in figure 4.13.14 that for field lines far from the axis their mutual orientation with respect to the accelerating electrostatic field changes in such a way that the ions suffer also an energy gain perpendicular to the force line. For field lines on the lateral part of the ECR surface the ions are mainly perpendicularly accelerated; this means that they are ejected from the loss cone and the magnetic confinement in the secondary plasma works better. Hence the large peripheral part of the resonant surface provides ions having their pitch angles outside the loss cone after the acceleration. Such ions are reflected before walls and return to the inner plasma. Hence the ion losses over the a large part of the resonance surface are extremely small. Remember that this description is based on the collisionless nature of the ions in the outer plasma, that is quite reasonable because of lower density. Strongly reduced ion output currents leads to higher accelerating potentials, thus creating ions having larger gyro-radii which increase the ion losses and create strong non-uniformities. Performances of ion sources will be bad in such a case and only the application of external electrons will reduce the accelerating potential, if the energy of injected electrons is not so high to make the probability of electrons trapping low. For simulations it comes out that electrons with a longitudinal energy equals to 1 keV gain about 0.5 keV of perpendicular energy after a single pass through the resonance. Such electrons reach the loss cone boundary and the trapping probability decreases. This effect was also obtained in single particle simulations featuring that the higher is the initial mean energy of the electrons, the lower is the number of trapped electrons. Another possibility to influence on accelerating potential is to reflect the escaping electrons. Most of these escaping electrons have an energy lower than 0.5 keV. Hence, the reflection by means of a biased disk will inject again most of them into the inner plasma, where they are effectively trapped. It looks like that the biased disk operates as an injector as reflector for the electrons. It may be concluded that *ions are accelerated outward at the resonant surface and perpendicularly to it. After the acceleration they are mirror trapped, except those accelerated at near axes region.*

4.13.8 Discussion about ion dynamics and comparison with the experimental results

The conclusions about the specific aspects of the ion dynamics were compared with the experimental observation. In many experiments metallic deposition on chamber walls has been observed, for example using SERSE as source of multiply-charged gold ions. Other than the typical triangular

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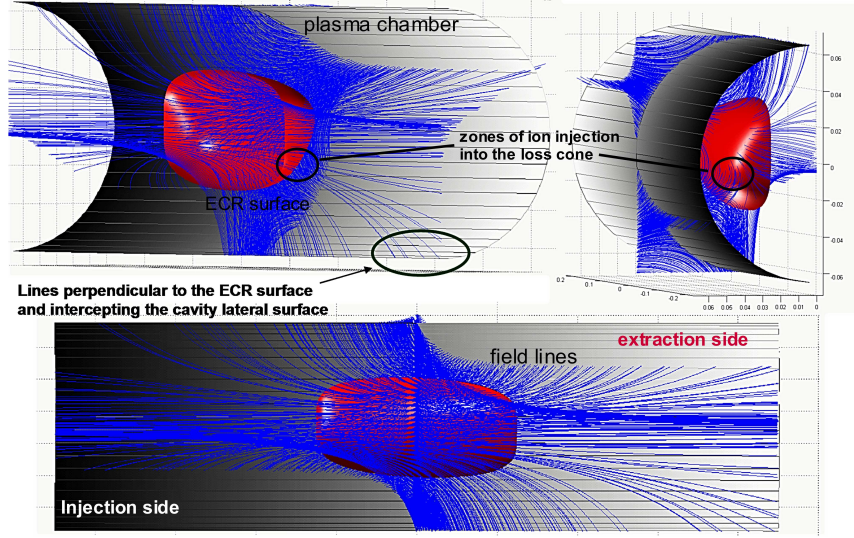


Figure 4.13.16: Structure of the magnetic field lines (in blue) together with a cutted view of the resonance surface (in red), showing which force lines intercept the ECR surface with an angle which lies in the loss cone and at the same time intercepting the plasma chamber walls.

stars at the extraction end, due to the structure of the magnetic field, three regions of gold deposition on the cylindrical surface were observed²⁴. These strips were the prolongation of the arms of the triangular stars. Other parts of walls had no visible deposit. By studying those force lines originated at the position of any strip, and intersecting the resonance surface at an angle lying in the loss cone, it is clear from figure 4.13.16 that some of them are good candidate for the ion acceleration towards the lateral cylindrical surface. Some of the field lines in proximity of the top side of the ECR surface are nearly perpendicular to the surface and also they intercept the plasma chamber laterally. The ion move along these lines and they cannot be confined by the magnetic field because of the small angle of their velocity vector with respect to the force line direction.

The metal deposition properties confirm the PP and SP model, with the accelerating electric field perpendicular to the ECR surface, that in addition is the boundary between the two different plasma regions.

The results obtained by means of the PIC simulations have permitted to

²⁴They had the form of three strips few centimeters long.

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link the ion dynamics to the electron dynamics well beyond the ambipolar diffusion.

The ion acceleration model from PP to SP works well when the PP-SP boundary is not perturbed by any electron density fluctuation. Anyway, we demonstrated that the wave electric field distribution over the resonance surface strongly affects the electron dynamics, and the electron density distribution shown 4.13.3 depends on the wave pattern on the resonance surface. This in turn affects the structure of the positive potential distributed over the ECR surface. At first approximation, considering a typical mode excited inside the plasma chamber, e.g. the $TE_{4,4,23}$, we can imagine that the potential structure follows the electric field structure over the resonance surface. In fact, where the wave electric field has a maximum, then the electrons will be easily resonant sloshing and the positive potential will also have a maximum. The contrary in case of small wave amplitude. This means that the situation appears like in figure 4.13.17.

Ions ejected from the inner resonance plasma through the lateral part of the ECR surface are reflected by the magnetic field and return towards the inner resonance plasma. There they are further ionized. Let us suppose that, because of the step by step ionization, the ions increase their charge state by one unit for each inner resonance plasma crossing. During this motion, the ions drift towards the near axis region²⁵.

Once arrived there, the electrostatic field direction over the PP-SP boundary (i.e. the ECR surface) is proper to inject the ions into the force line loss cone, then directly towards the extraction hole. If the motion was unperturbed by any other electric field directed along different directions, each ion is able to cross many times the primary plasma region thus achieving high charge states. In this case a considerable accumulation of highly charged ions occurs in proximity of the extraction hole. There the beam formation occurs, the ions follow the magnetic field lines and the electrons which are leaving the trap, hence the shape of the extracted beam is like the typical triangular star with a well focused core (see figure 4.13.18 1)).

The case of smooth electric field distribution over the resonance surface is the nearest to the unperturbed ion motion. On the contrary, the fine structure of the mode pattern generates electrostatic fields at the PP-SP boundary no longer directed along the normal to the resonance surface. This fields scatter the ions during their motion and an effect similar to the Spitzer collisions arises: ions can be injected into the loss cone with a greater frequency with respect to the unperturbed case. Then the ions are

²⁵For ion drifts in magnetic traps see simulations shown in [25].

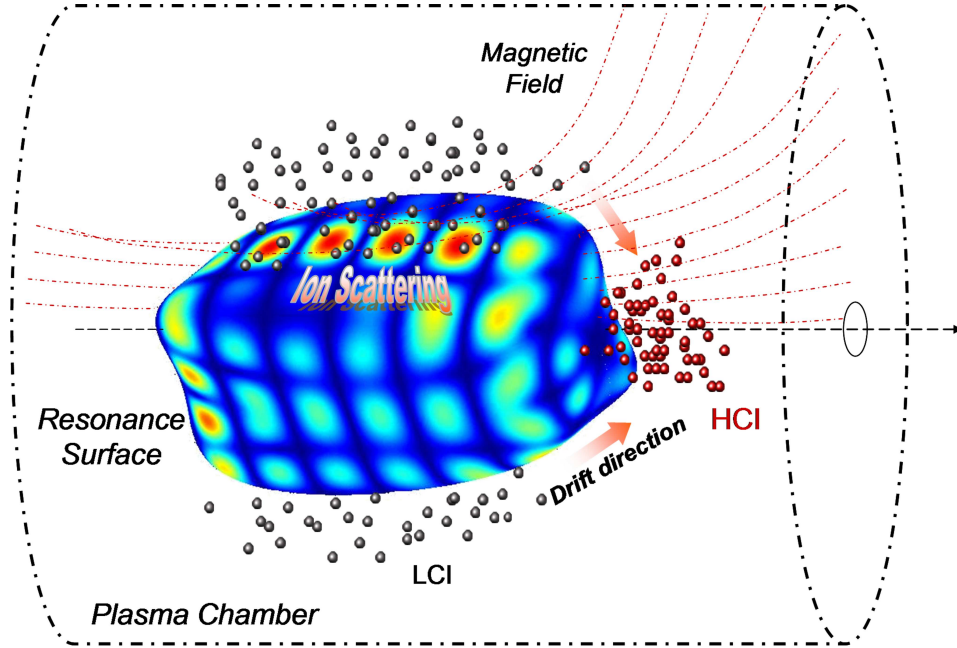


Figure 4.13.17: Description of the ion beam formation in ECRIS for multiply-charged ion production. The electromagnetic wave pattern of a particular cavity mode is plotted over the resonance surface. A simplified view of the field lines is also shown, together with a rough picture of the plasma chamber with the extraction hole. The figure also points out the direction of the ion drift from the periphery to the near axis region, showing the accumulation of highly charged ions in proximity of the extraction hole.

less charged and the source performances worse. The ions are injected into the loss cone well before they reach the near axis region and this leads to the beam broadening (see figure 4.13.18 2), thus worsening the beam brightness.

The beam brightness is a key parameter for ECRIS, then *the control of the beam formation by means of the frequency tuning appears as a powerful tool to optimize the source performances*. Comparing the figure 4.13.18 with figure 4.13.19 it is clear that the model described in this pages finds a surprisingly confirmation in the experimental data reported in section 4.2. There, the best performance in terms of extracted currents of the higher charge states was obtained when the beam was more focused. This is in accordance with the model of ion scattering over the ECR surface. In the experiment the frequency was varied of some tens of MHz , thus changing

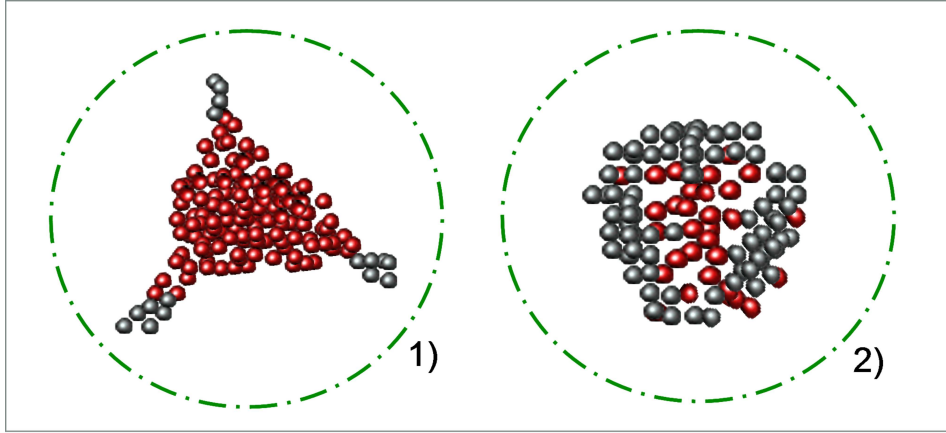


Figure 4.13.18: Beam shape for two extreme situation: 1) the beam is well focused (although it preserves the typical three-arms shape) and except the arms periphery is almost completely formed by highly charged ions (red circles); this situation occurs when the primary plasma potential is weakly perturbed by the wave structure. 2) The beam is broadened by the strong perturbation to the primary plasma potential due to the wave structure; in addition, it is mainly populated by ions at low charge state (grey circles).

the electromagnetic field distribution over the ECR surface. Well focused beams were obtained and high values of the extracted current for the higher charge states were alternated with lower currents and charge states, and broadened beams for close frequencies. This means that the field pattern over the PP boundary passes from quasi-uniform to strongly spatially variable structures.

Note finally that the ion scattering occurs non only in coordinate space, but also in velocity space. Then also the emittance may be strongly affected by the frequency tuning, and emittance measurements in ECRIS will be of primary importance together with detection of beam shape and spatially resolved OES.

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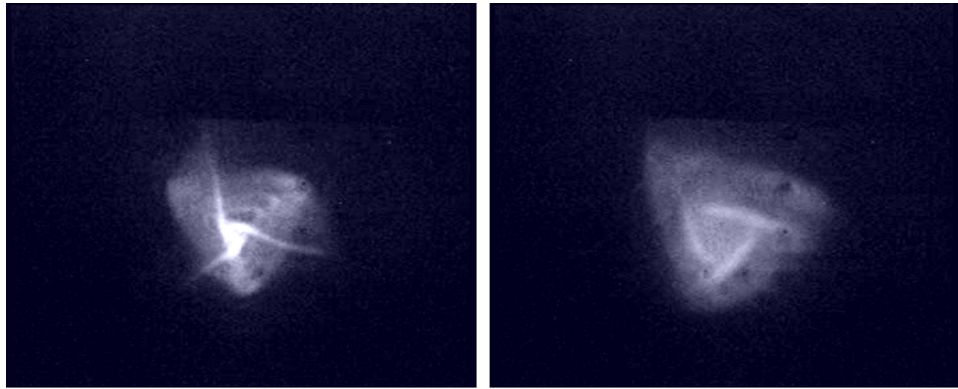


Figure 4.13.19: Beam shape determined experimentally at the extraction of the CAPRICE source [6] (see section 4.2) for two different frequencies (i.e. two different resonant modes).

4.14 Considerations about possible turbulent heating in VENUS

In section 1.12 some considerations about an alternative mechanism of plasma heating are outlined. It consists essentially in the turbulent heating via quasi-particle formation (the cavitons) with their following collapse, leading to the energy transfer to the electrons through the Landau damping. It is reported that such heating processes are usually triggered by plasma instabilities. We tried to evaluate the maximum energy that an electron of a VENUS-like plasma may reach if such turbulent heating takes place. The formation of cavitons should be triggered, in the case of the experiment shown in section 4.5 by instabilities linked principally to the magnetic field. For lower mirror ratios, and because of the high RF power level, plasma instability may transfer a lot of the incoming wave energy into quasi-particles and then to the plasma electrons via Langmuir Caviton collapse, according to the mechanism described in section 1.12. But another source of instability should be considered, not connected with the plasma fluidodynamics but more with phase space instabilities. In section 1.2 we discussed about the electrons runaway. This kinetic instability is due to the formation of electrons whose velocity is so high that collisions are practically absent. This means that they cannot be thermalized by the bulk plasma, and the EDF differs appreciably from a Maxwell-Boltzmann distribution. The formation of such a high energy component is certainly enhanced by the lowering of the mirror ratio, which makes the gradient near the resonance less pronounced. Hence runaway electrons may be created with larger rate in case of lower mirror ratios and further increased by high power levels. According to the equations of section 1.12, we can estimate the density fluctuations and the electron temperature at the beginning of the cavitons formation. We assume that electric field amplitudes around 10^6 V/m can exist in some regions of the plasma chamber for high Q modes, as confirmed by our simulations and also by R. Geller, who in his book [10] says that electric field up to 10^6 V/m are "realistic"; considering that the LCC may last about 50 *nsec*, electron temperatures in the range of 4 – 6 keV have been considered, as suggested by the single particle simulations²⁶. Then a realistic value of $\frac{\delta n_e}{n_e}$ lies in the range 0.005 – 0.01. For VENUS the caviton diameter is $\Delta \simeq 7$ mm, then the energy stored in the caviton's volume is about $1 - 7 \cdot 10^{-6}$ J. The caviton's diameter is much larger than the Debye length λ_D . During the LCC, the

²⁶In figure 4.8.1 it is shown the mean electron energy after 50 ns which to some extent corresponds to the electron temperature at that time.

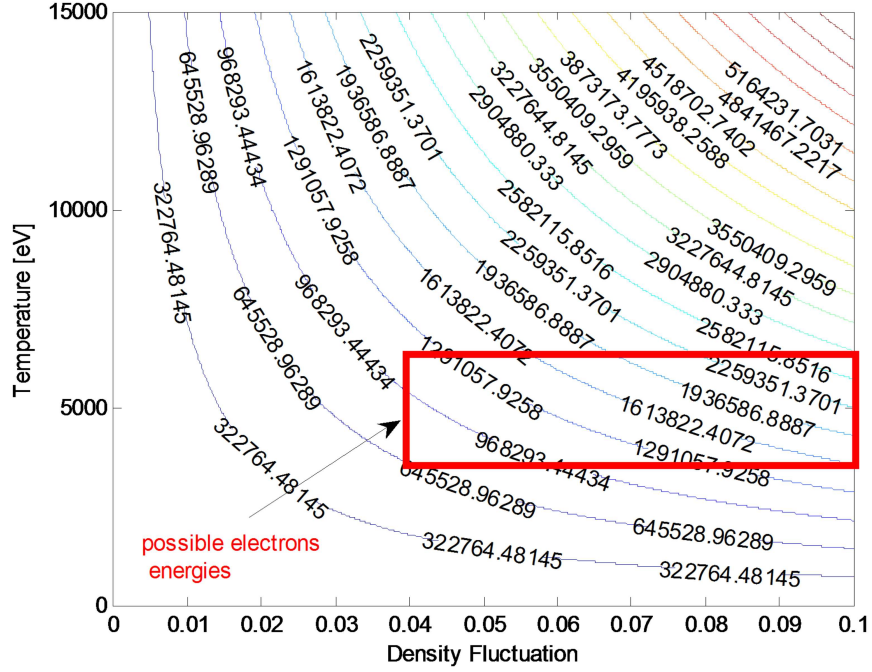


Figure 4.14.1: Electron energies for different T_e and density fluctuations as a consequence of LCC.

plasmons are compressed in a smaller volume but the electrons gain energy, by means of the Landau damping, only when Δ becomes comparable to λ_D . For the values of δn_e and T_e above calculated, some electrons can gain energy from 1 to 2.3 MeV. This values have been calculated by means of a MATLAB code which solves the equations of section 1.12. A more detailed description of the energies that can be achieved through the LCC is shown in figure 4.14.1. This figure clearly shows that the higher are the electron temperature and the density fluctuations at the beginning of the LCC, the higher will be the energy achieved by the electrons. The calculated energy has been found by dividing the total caviton stored energy among the electrons filling a Debye sphere. The model does not predict any limit about maximum achievable energies and it can explain the strong increase of the ultra-hot electron production when the mirror ratio decreases, and the slight increase with the power.

Chapter 5

An alternative plasma heating mechanism

5.1 The experiment at WEGA facility

As conclusion of the experimental activities described in PhD this thesis an experiment has been performed on the WEGA stellarator operating at the Max Planck Institute for Plasma Physics of Greifswald, Germany, in order to investigate the mechanism concerning the electromagnetic mode conversion leading to an electrostatic wave (the Bernstein Wave - BW) propagating and absorbed in an overdense plasma. The experiment principally aimed to the study of the emitted X-rays in different conditions of pressure and injected RF power. The scope was also to understand whether or not the emission of the X-rays was linked to the plasma density, and if the BW absorption suppressed the production of high energy electrons. The properties of the BW have been largely discussed in section 1.5, where it was pointed out the absence of any density cutoff for these waves. Hence, these waves are very interesting for nuclear fusion devices. Numerous experiments were performed with ECRH heating in toroidal devices. The obtained plasma parameters vary for different installations, but one common limit has been observed: the density of the obtained plasma is limited by the cutoff values. The cutoff density is different for different modes propagating in the plasma except for the BW, for which the cutoff does not exist at all.

For ECR heating in toroidal devices Ordinary mode (O-mode) or Extraordinary mode (X-mode) are externally launched to the plasma from high or low field side of the torus. Both modes have a cutoff limit for propagation. However, Slow Extraordinary mode (SX-mode) waves may be converted to

into a BW mode if they reach the Upper Hybrid Resonance (UHR) layer. O-mode waves launched from low field side may be coupled to the SX-mode in the region of the O-mode cutoff layer, and then the SX-mode is coupled to the B-mode. This process called OXB mode conversion, was described in 1973 by Preinhaelter and Kopecký [110]. In their work, it was shown that the OXB mode conversion might be optimized in terms of O-mode insertion angle with respect to the external magnetic field direction. Under conditions for the optimization of the mode conversion, the O mode is completely converted into a slow X mode, which then completely converts to BWs in the UHR layer.

The region outside the UHR layer is evanescent for the EBW. Thus, the EBW can not be directly launched outside of the plasma but it must be excited inside the region bounded by the UHR layer. Thus, mode conversion from X to electrostatic Bernstein modes (EBW) is required to generate the EBW. There are two main techniques to generate EB waves in the center of the overdense plasma. The first is to launch an X-mode from the high field side of the given magnetic structure. The X-mode converts to the B-mode on the UHR layer, and then the B-mode propagates into the overdense plasma region.

The second way is via the OXB mode conversion process. In this case, the HF power is launched from the low field side of the torus in the O-mode. The O-mode converts to the SX-mode, the SX-mode to B-mode, and the B-mode propagates inside the plasma and heats the electrons. As affirmed in section 1.5, the Bernstein waves are absorbed at the cyclotron harmonics, and then they travel inside the plasma until they encounter the ECR or the higher order harmonics.

For the first time, OXB conversion was shown experimentally in 1987 in Japan Institute of Plasma Physics Nagoya University on the TPL device [111]-[112]. In that case the heating wave was launched obliquely via a horn antenna. The conversion was proved by two-dimensional high frequency (HF) probes measurements. The OXB plasma heating and current drive with BW in an overdense plasma was demonstrated also in the Stellarator WEGA operating at Greifswald [113].

5.1.1 Density limitations due to the heating at the plasma periphery

The main limitation of the WEGA Stellarator consists in the overdense plasma cutoff. Without any mode conversion, in fact, the wave can not reach the cyclotron resonance region in the center of the plasma in O-mode

5. AN ALTERNATIVE PLASMA HEATING MECHANISM

as well as in the X-mode. The wave meets cutoff barrier at the edge of the plasma where electron density reaches its critical value n_{cutoff} . It is assumed that the plasma is heated mostly by the multiple reflections of the O-mode between the cutoff layer and the vacuum chamber wall. Such heating mechanism results in the relatively low plasma density just slightly higher than the cutoff density. This mechanism of plasma heating provides a hollow shape of the electron density. In fact, the great part of the injected RF power is absorbed in the plasma periphery. Because of transport issues, a considerable fraction of energy escapes from the inner plasma. This means that in the plasma core only a reduced fraction of energy arrives, thus leading to a hollow temperature and density profile. In addition, the higher is the fraction of rf power absorbed at the plasma edge, the higher will be the edge density, hence the lower is the fraction of power able to reach the plasma core, because of the cutoff. We can find here similar problems as the ones of the ECR ion source plasma heating. The PIC simulations reported in section 4.13 demonstrate that the plasma can be divided into primary and secondary plasma, the former having the highest density and temperature. The figure 4.13.3 features a marked hollow shape of the electron density. Even if it has been demonstrated that the real distribution should be flattened, in order to avoid any low frequency ion oscillations, some experimental results [114] demonstrated that the plasma density has a slightly hollow shape. The figure 5.1.1 shows the experimental results obtained by means of the optical spectroscopy for a plasma heated by the ECR inside a cylindrical camera. The spectrometer was moved vertically on the injection flange of the plasma chamber. Especially for helium the figure features a slightly hollow profile for the detected signal, that is somehow proportional to the plasma density.

The results shown in figure 5.1.1 can be explained only by the model of the ECR zone shown in figure 5.1.2.

The model proposed in section 4.13 is in good agreement with such measurements. These results demonstrate that the peripheral heating of the plasma limits considerably the wave-plasma coupling, and then the maximum density that can be reached in the plasma core.

5.1.2 Overcoming the cutoff density: experimental evidences about the BW excitation in MDIS

Some preliminary experiments [21] performed on WEGA demonstrated that the cutoff density can be overcome without any induced mode conversion inside the plasma. The magnetic field was fixed, at the beginning of the experiment, to a value close to the ECR one. After the plasma ignition, the

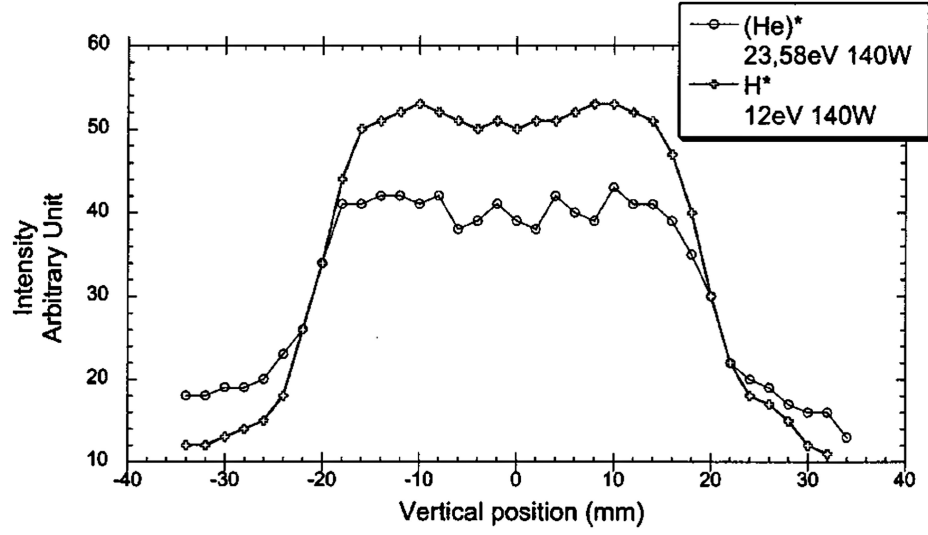


Figure 5.1.1: Profile of the intensity of the line at different vertical positions of the spectrometer and for two different ion species, helium and hydrogen. [114].

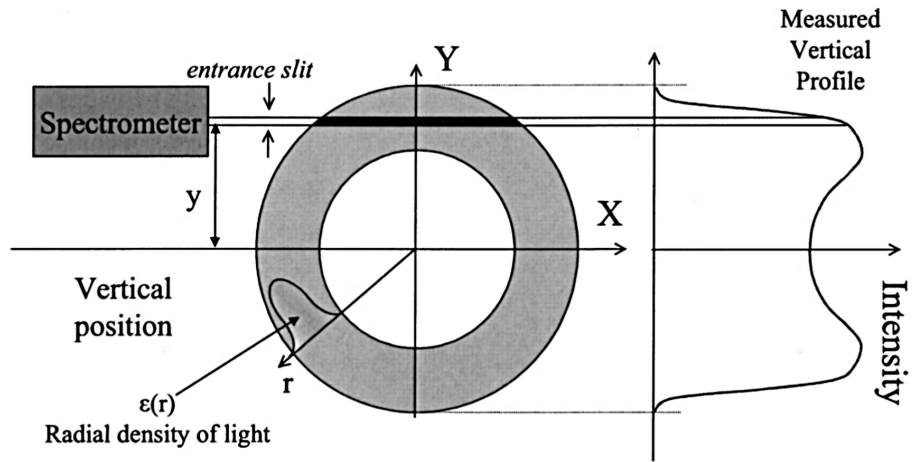


Figure 5.1.2: Longitudinal cut of the hollow heating zone. The spectrometer at the y vertical position measured the intensity of the atomic lines emitted from a thin slice of plasma. [114].

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magnetic field was decreased down to $0.5 B_{ECR}$. In this way the density overcome the cutoff value, and a maximum was observed for $B = 0.65 B_{ECR}$. Similar results were obtained in [115] and in [10]. In these cases the increase of the electron density was attributed to the second cyclotron harmonic absorption. This mechanism is still based on a single particle approach and does not take into account any collective motion of the plasma particle, neither the possible formation of electrostatic waves.

However, the experimental results cannot exclude the presence of electrostatic BW absorbed at the cyclotron harmonics.

Other evidences about the overdense plasma formation in under resonance field come from an experiment carried out in our laboratory [68] with the plasma reactor. A systematic measurement of the electron density was performed along the entire plasma chamber and for different values of gas pressure and microwave power. The figure 5.1.3 shows some results obtained during this experiment at 2.45 GHz .

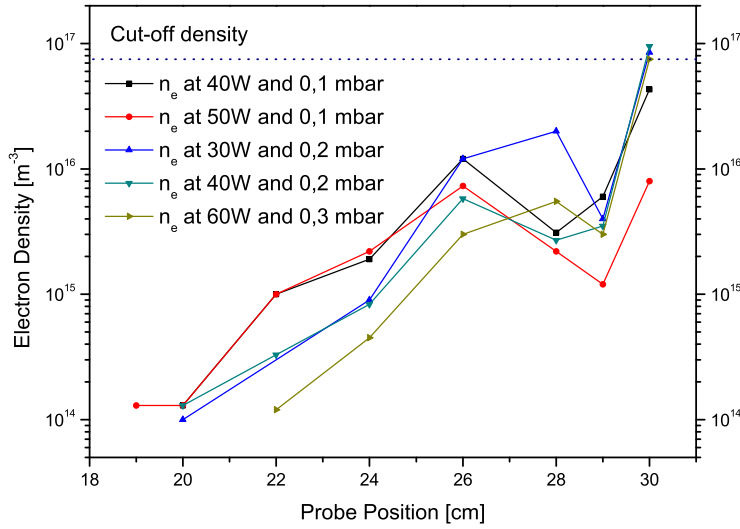


Figure 5.1.3: Trend of the electron density measured along the plasma chamber axis by means of a Langmuir probe for different values of power and pressure.

The density along the plasma chamber axis depends on the background

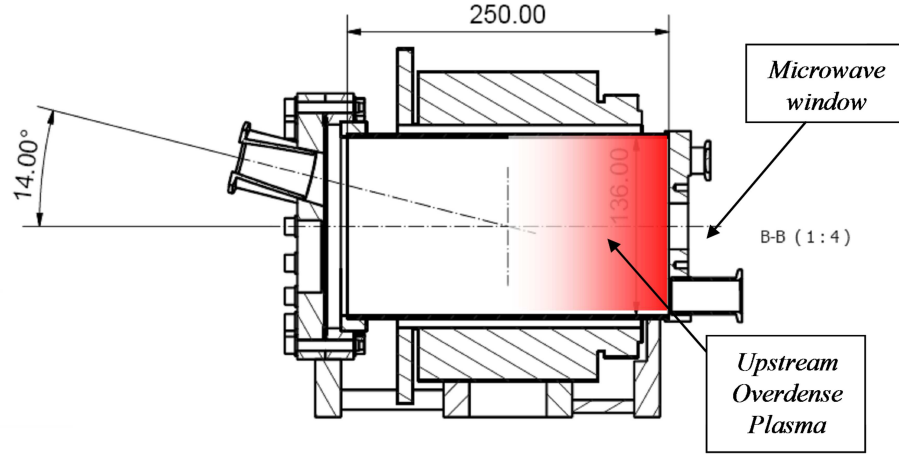


Figure 5.1.4: Design of the plasma reactor with the region of "upstream overdense plasma". Once created, it limits the propagation of the incoming electromagnetic wave.

pressure and on the microwave power. In three cases a density slightly higher than n_{cutoff} was measured¹. The magnetic field could not be varied because permanent magnets are used, than any possible alternative mechanism of wave-plasma interaction depending on the magnetic field can be investigated only by looking to the plasma properties in different positions. The higher was the density at 30 cm, which corresponds to a near microwave window position, the lower was the downstream density. The figure 5.1.4 helps to understand where the overdense plasma is located. From the sketched picture it is clear that the overdense plasma is located in the near window region, and it prevents the wave to propagate downstream. Hence the plasma distribution becomes strongly inhomogeneous inside the plasma chamber. This effect demonstrates once again that the peripheral heating is detrimental for the maximization of the plasma density. The mechanism leading to the creation of this overdense plasma in a region where the magnetic field is much lower than the resonance one is evidently an off-resonance mechanism, but a simple description does not applies to the results shown here. The single particle absorption and some collisional absorption processes, may explain why this effect is maximized at lower pressures and

¹Note that the density measurements with Langmuir probes are not so precise (typically less than 30%). Therefore from figure there is only a small evidence of the overdense plasma creation.

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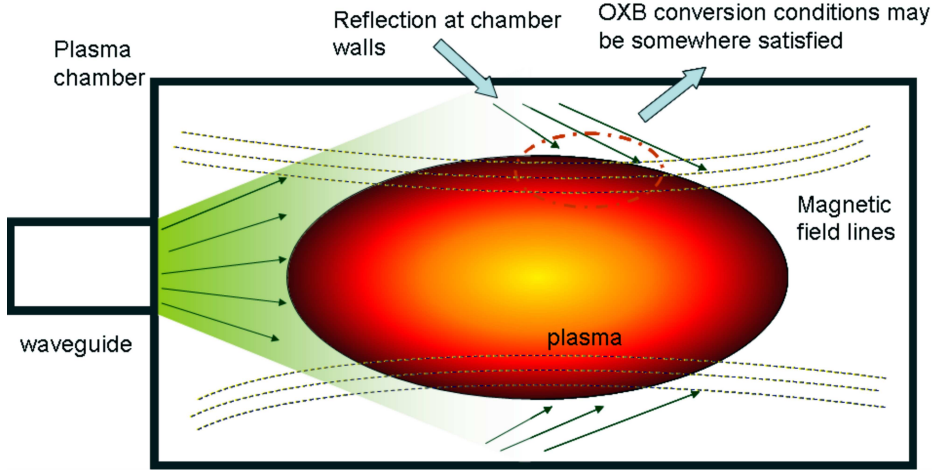


Figure 5.1.5: The large wavefront emitted by the 2.45 GHz waveguide and the possible reflections by the chamber walls are shown, together with the location of the plasma and the magnetic field lines in a MDIS (see section 2.2).

higher RF powers. However, an alternative approach can be supposed, similarly to what observed for WEGA. In the plasma reactor the microwaves are provided by means of a waveguide located longitudinally with respect to the chamber axis, as shown in figure 5.1.4. This means that the injected wave is mainly an O or R wave when it propagates inside the plasma. However, because of the complex structure of the magnetic field lines, and because of possible reflections at the chamber walls, the modes may convert each other and also X modes can be generated somewhere inside the cavity, through an $O - X$ conversion. At that time, the X waves can be directly converted into BW at the UHR layer. However, if the injection angle of the O mode is not the best one, the BW creation process has a low efficiency. The variation of the power and of the background pressure changes the plasma properties and in turn it affects the conversion efficiency, leading to the variation of the density trends shown in figure 5.1.3.

The fixed magnetic field is a limitation, as the decrease of the magnetic field would help to increase the plasma density up to overdense values.

The ECR is an effective method to heat the electrons, as it is a powerful channel to transfer energy from the wave to the electrons. Anyway, it works so well that electrons resonantly accelerated gain higher and higher energies thus limiting the number of lower energy particles that can be accelerated. It

was also observed that in conditions for which the downstream propagation of the wave occurs, the plasma is more homogeneous and the density and the temperature are quite flat. This means that the two processes, the ECR and the possible BW excitation, are alternative each other: when the ECR plays the major role, the great part of the energy is coupled with the plasma at the resonance and no power is available for the mode conversion and the BW excitation. Otherwise, if the optimal conditions for the BW creation are satisfied, the wave is mostly converted in BW, then absorbed at the second cyclotron harmonic.

If we would like to increase the plasma density by suppressing the number of high energy electrons generated by the ECR, we should be able to use the OXB or the X-BW direct conversion in the plasma trap. Even if the technology advanced significantly since the time of pioneers, Golovanivsky and Dougar-Jabon, who used the BW excitation in 2.45 GHz plasma sources, this objective is still a challenge.

5.1.3 Experiment on WEGA

The experiment on WEGA aimed to investigate the production of high energy electrons under extreme conditions of the Stellarator plasma, i.e. for maximum values of RF power and minimum pressure, in order to have a pressure and a power density comparable to the ECRIS ones. Our scope is to understand if high density plasmas can be obtained without the production of high energy electrons. This means that the methods applied on WEGA may be precious for future experiments on ECRIS, especially as regards the tailoring of the EEDF aiming to the maximization of the warm electrons. Typical values of plasma parameters like density and temperature are usually determined with the diagnostic shown in section 3.5, and they are the following:

- Electron density: $1 \cdot 10^{17} \text{ m}^{-3} < n_e < 5 \cdot 10^{18} \text{ m}^{-3}$;
- Plasma potential: $15 \text{ V} < V_p < 60 \text{ V}$;
- Temperature of slow electrons: $3 \text{ eV} < T_{es} < 12 \text{ eV}$;
- Temperature of fast electrons: $T_{ef} \sim 400 \text{ eV}$;
- Fraction of fast electrons: $f_n < 5\%$.

The values reported above point out that two typical plasma populations exist in WEGA. The higher energy one is named *suprathermal*, and it will

5. AN ALTERNATIVE PLASMA HEATING MECHANISM

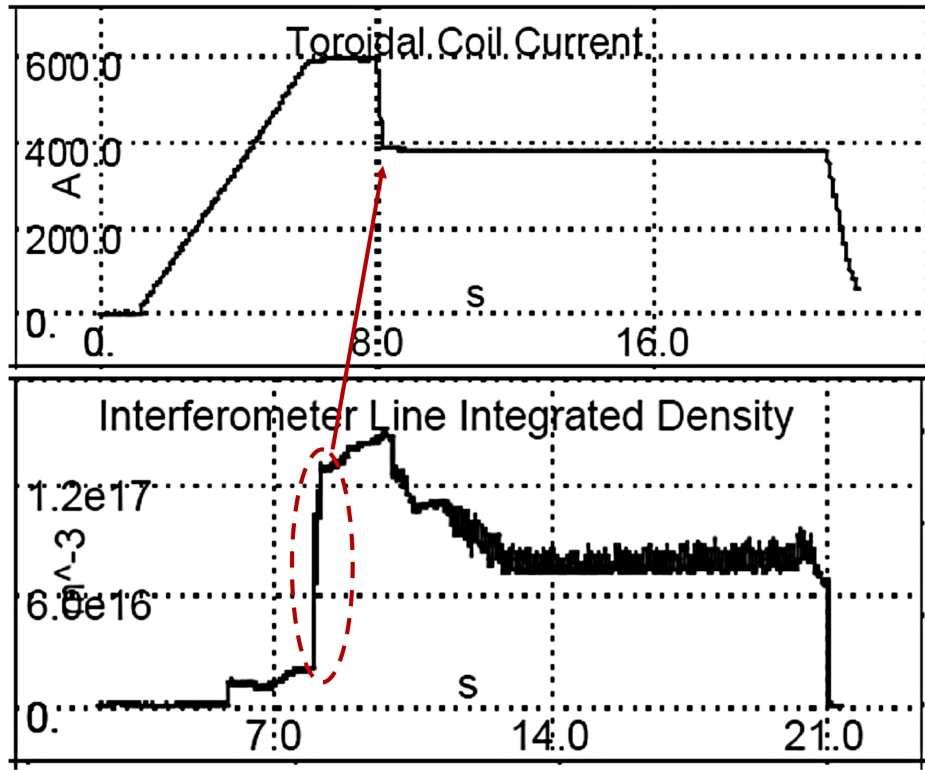


Figure 5.1.6: Trend of the electron density with respect to the time for a single shot compared with the coils current.

be important to understand if some electrons belonging to this population goes till higher energies when the extreme conditions described above are satisfied.

We started the experiment by looking to the variation of the electron density in typical conditions of background pressure and injected power. The power has been maintained constant during all the experimental acquisition; its maximum possible value was about 18 kW; no more power was possible to inject into the plasma chamber as in the cases of lower pressures the reflected power increases dangerously. Microwave shots lasting about 40 seconds were used, and for each shot the pressure and the magnetic field were varied. In figure 5.1.6 the trend of the plasma density is shown; it is plotted versus the time during a single shot and in comparison with the current flowing in the toroidal coils. The maximum current value corresponds to the ECR field.

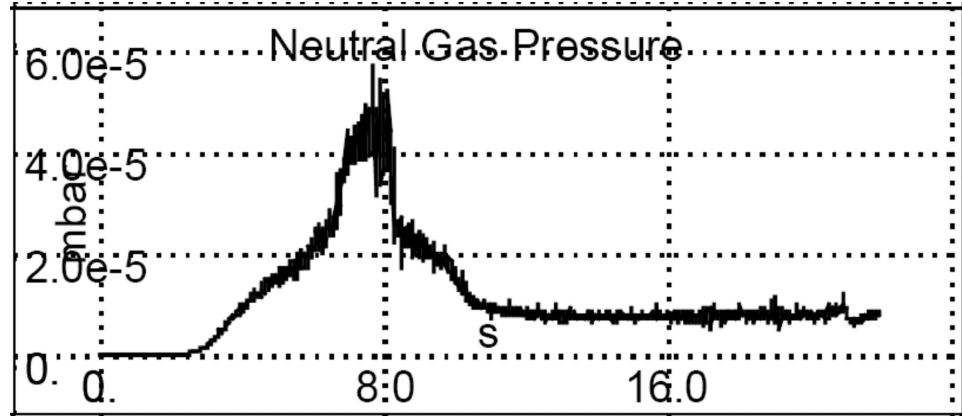


Figure 5.1.7: Trend of the pressure versus the time for the shot #25.

Note that when we decrease the magnetic field the density abruptly increases as expected. Lowering the magnetic field below the ECR value, the density increases above the overdense limit. In this case the process is optimized by the OXB conversion driven by properly shaped antennas and microwave mirrors (see [116] for more details). Anyway the jump of the density occurs also without any optimization of the OXB process. The density value shown in figure 5.1.6 has been obtained through the interferometer measurements. It gives an integrated density along the last closed magnetic field surface. A maximum density of about $7.5 \cdot 10^{17} \text{ m}^{-3}$, i.e. about one order of magnitude higher than the cutoff one was obtained and even higher densities were obtained for higher background pressures.

Under the conditions of figure 5.1.6, the soft X-ray detector did not give any signal. This is a signature of the absence of high energy electrons in case of a high density plasma. After the calibration with a chromium-copper alloy used as anode for the internal X-ray source, no signals around and above 800 – 900 eV were detected. Then we decided to decrease the inlet gas flow, thus decreasing the background pressure.

The figures 5.1.7 and 5.1.8 feature the pressure trend versus the time for two different shots, labeled 25 and 27. Note that at the beginning the pressure is quite high; the plasma ignition occurs to relatively high values of the pressure, in order to avoid any excessive power reflection due to a strong decoupling of the incoming wave. Then, in order to investigate the extremely low pressure limit, we decrease the inlet gas flow. However, in case of figure 5.1.8 it is clear that the maximum value of the pressure is

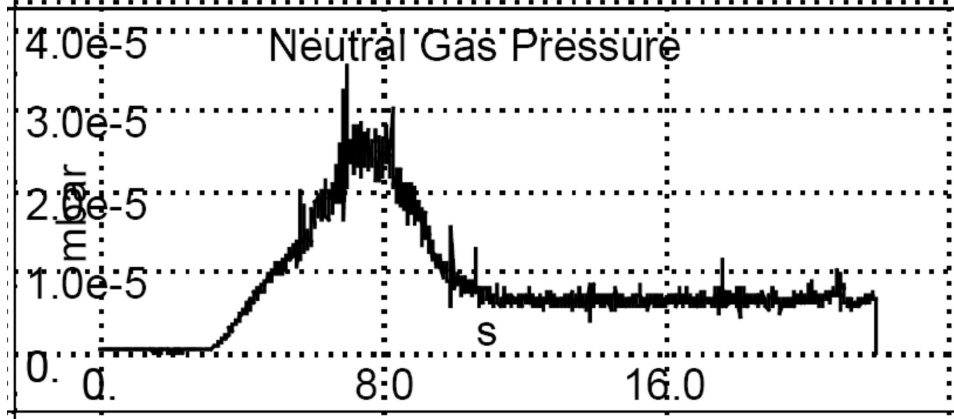


Figure 5.1.8: Trend of the pressure versus the time for the shot #27. The inlet gas flow has been decreased, then the pressure is lower on average.

lower than for the shot 25, as well as the stabilized low value after 10 – 12 sec. What happened about the plasma density is shown in figure 5.1.9. Note that in correspondence of the pressure decrease the density decreases as well; in addition, n_e begins to strongly oscillate. This is a sign that a very unstable plasma state is established. In this case, some high energy X-rays were detected. Their energy lies in the range 0.9 – 2.5 keV, but a reduced number was detected with an energy around 6.5 – 7 keV. The maximum energy is much smaller than the ECRIS one, because of the higher pressure and lower power density. The presence of these X-rays when the density has such pronounced oscillations suggests that they are produced only in case of low densities, otherwise they should be detected also in case of non oscillating density. The strong instable plasma creation may be interpreted as a transition between an ECR-heated state (with low density and high electron energy) and an OXB-heated state characterized by high density and low electron energy. Hence, as expected, the low densities facilitate the heating up to high energies. On the contrary, when a great amount of electrons is generated through ionizing collisions the wave energy is distributed among many particles, and this leads to a reduction of the single particle energy. By keeping constant all the parameters except for the magnetic field, which was set at the ECR value for all the shot duration, and for the RF power, which was decreased down to 8 kW, the plasma density was again measured, and it is shown in figure 5.1.10. Note that the density remains well below the cutoff value at all the times, even when the pressure

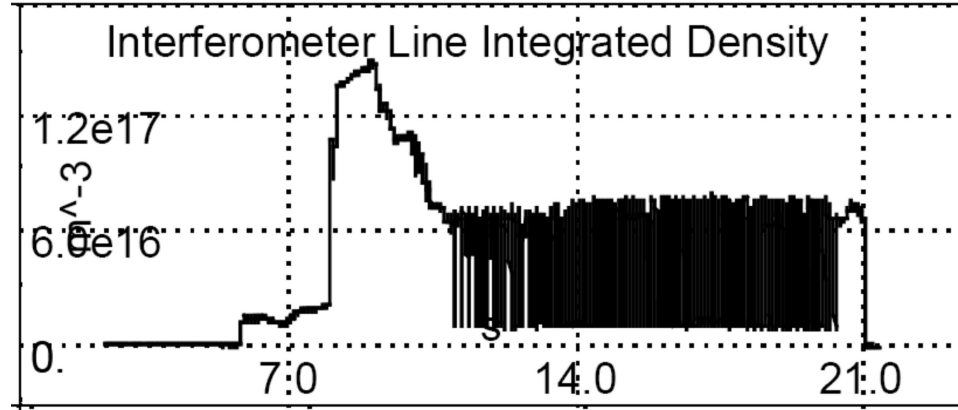


Figure 5.1.9: Trend of the electron density versus the time for the shot 27. Strong fluctuations are evident, in correspondence with the strong decrease of the background pressure.

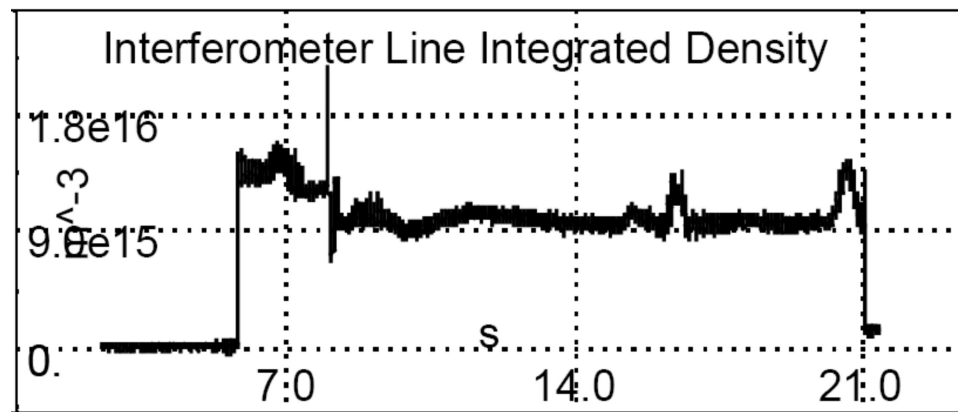


Figure 5.1.10: Trend of the electron density versus the time for the shot 69. The density remains well below the cutoff value for very low background pressures.

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is decreased². This is the confirmation that when the ECR is active inside the plasma the overdense plasma state is not obtained and no X-rays were detected. Measurements with the same set-up but with higher power levels were not possible because of the very high reflected power especially for the lower values of pressure.

5.1.4 Modeling of the BW launching in ECRIS

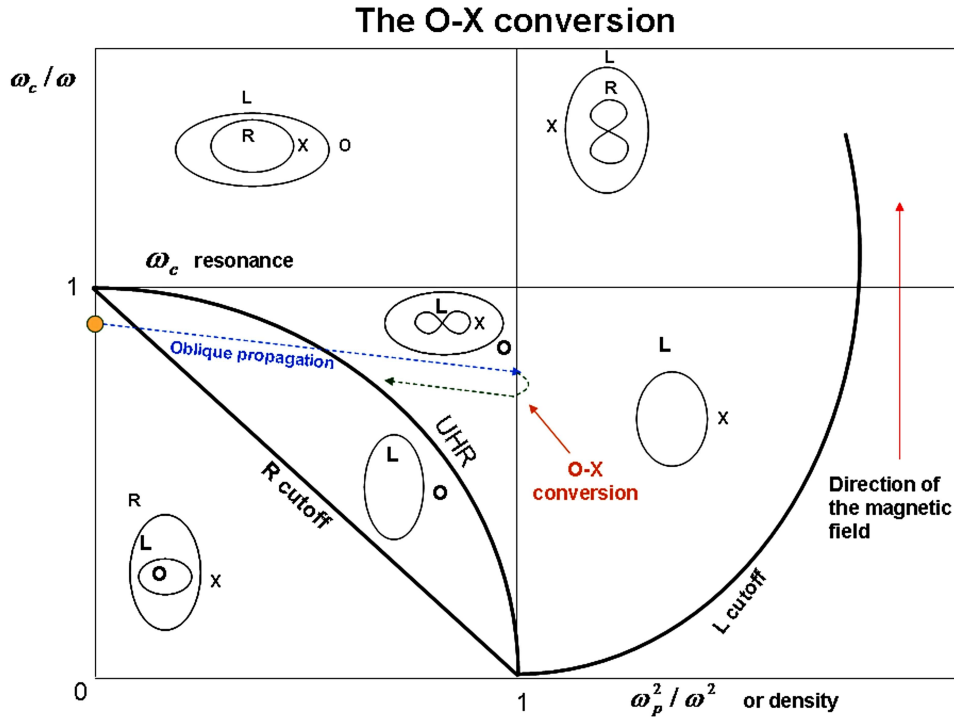


Figure 5.1.11: The CMA diagram showing the path of a obliquely launched O-mode with the following conversion in a X mode at the O cutoff.

The experimental results reported above permitted to imagine an experiment with modern ECRIS about the generation of BW via O-X-B or direct X-B conversion. The design of a BW heating for ECRIS requires completely new approach: as mentioned above, the ECR surface must be placed outside from the plasma chamber, and a proper displacement of the O and X

²The density trend is similar to the one shown in figure 5.1.8.

cutoffs, other than the UHR layer, must be chosen to generate BWs. The cutoffs and resonances relations are reported in section 1.4. There was also pointed out that the CMA diagram is a powerful method to trace the path of a wave propagating inside the plasma for different values of density and magnetic field. The figure 5.1.11 shows the path of an obliquely launched O mode³, traveling from a higher to a lower magnetic field towards the O cutoff. If the injection angle is proper, the O mode converts in a X mode at the O cutoff layer. The X mode comes back to the UHR layer, where it is converted in a pure electrostatic wave, the Bernstein wave, according to the mechanism described in section 1.5. The region between the UHR (which works as an X-cutoff) and the R cutoff is forbidden for the X wave (as shown in figure 5.1.12, then the X mode becomes evanescent there. Note that the starting point of the O mode is well below the ECR field. Once generated, the BWs can propagate in the shadowed region shown in figure 1.5.3. As the maximum field is below the ECR, the BW propagating towards higher B fields cannot be absorbed. Therefore we can consider only the BW which propagate towards the plasma core. Supposing that in the plasma core the density is well above the cutoff, and the magnetic field has a very slow decrease near the center of the plasma chamber, then the BWs travel towards the plasma core and they are absorbed at the second cyclotron harmonic, provided that the condition $B_{min} < 1/2B_{ECR}$ holds. This situation is shown in figure 5.1.13 on the CMA diagram.

For minimum-B structures the constant field surfaces are closed (see sections 1.6.4 and 4.6.3). In addition, according to the PIC simulations and to the MHD stability principles, we can say that the constant field surfaces are also constant density surfaces. To determine the possible path in the real space (i.e. inside the plasma chamber of an ECRIS) of the O, X, and B modes, we can suppose that a given density profile has been established inside the plasma chamber; after we determine the waves' paths, and we may suppose that the density profile is achieved self-consistently in the transition to the OXB process.

We determine the locations of the cutoffs and resonances with the profile shown in figure 5.1.14. If the displacement of these layers is proper the assumed density profile can be achieved self-consistently during the plasma ignition. The cutoffs and resonances layers, for the given density profile and for a different magnetic field configurations, have been simulated by means of a MATLAB numerical code. This code solves the equations reported in

³Hereafter the statement oblique launch will indicate an oblique propagation direction of the O mode with respect to the magnetic field.

The X mode cut-off and the Upper Hybrid Resonance

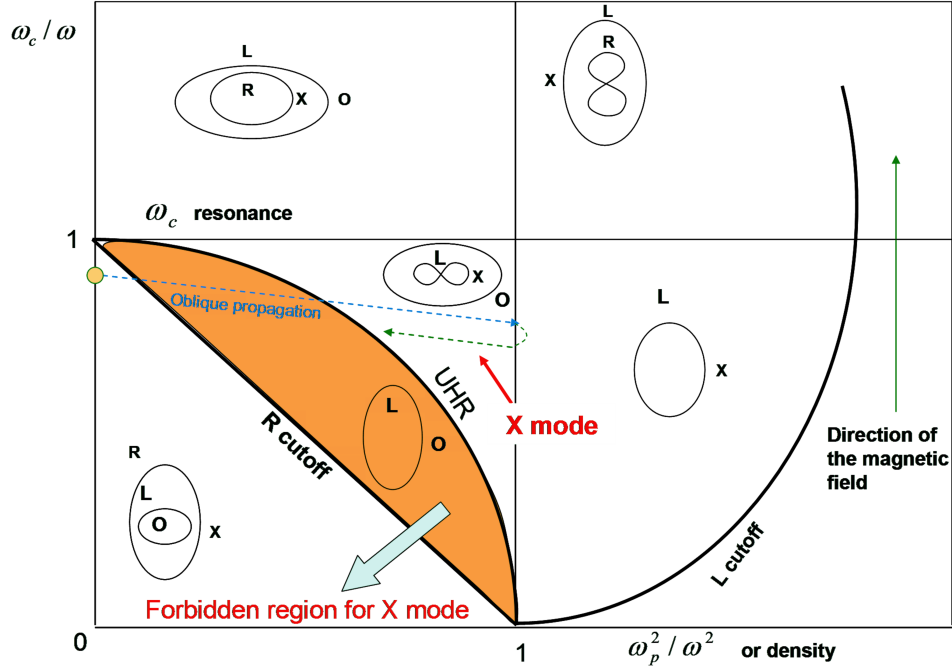


Figure 5.1.12: Forbidden region for the X-mode, that is converted into a B-mode at the UHR layer.

table 1.1 for different magnetic and density profiles, with the magnetic field given by the equations (4.6.7).

Figure 5.1.15 shows the localization in the real space of the cutoffs and the resonances. The plasma chamber dimensions and the features of the magnetic field structures are similar to that of SERSE. The location of the different layers is shown in figure: the ECR surface is outside the plasma chamber, thus avoiding any single particle effective absorption of wave energy. The UHR layer is bounded outward by the X cutoff surface, and inward by the O cutoff. This displacement ensures that the O mode, coming from regions outside the plasma edge, converts in the X mode at the O cutoff for a particular injection angle. The detailed tracing of the wave path inside the plasma is shown in figure 5.1.16. According to the path on the CMA diagram shown in figure 5.1.13, the X waves travels towards the UHR layer, i.e. in regions of lower density and higher magnetic field. This means that,

The BW generation thanks to X-BW conversion at UHR

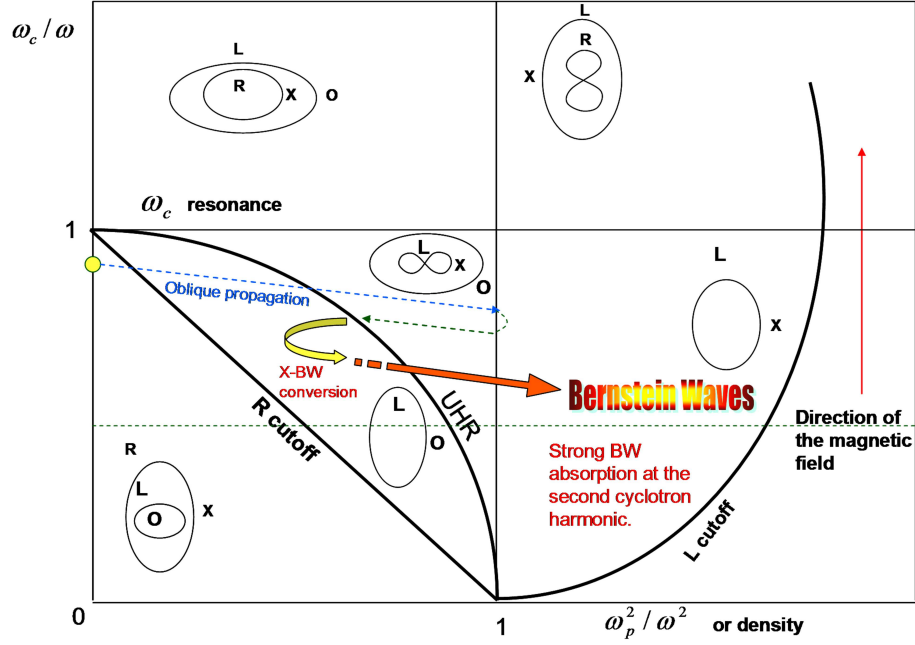


Figure 5.1.13: Description of the complete O-X-B conversion mechanism, with the trajectory followed by the BWs, effectively absorbed at the first cyclotron harmonic.

in real space, it will come back from the O cutoff layer to the UHR surface (see again figure 5.1.16). There, the X mode converts in a B mode; as shown in figure 5.1.13, the BW will move towards the lower magnetic field region without any density cutoff. It will be absorbed only when it encounters one of the cyclotron harmonic resonances. In figure 5.1.15 the first cyclotron harmonic has been located near the center of the plasma. This means that the BW are centrally absorbed and the central heating ensures, because of transport reasons, an effective heating of a great amount of particles, up to densities well above the cutoff one. The required profile for the magnetic field is: $B_{min} \simeq 0.32 \text{ T}$, $B_{inj} \sim B_{ext} \simeq 0.63 \text{ T}$ and $B_{hex} \simeq 0.6 \text{ T}$. The other required condition to optimize the OXB is that the magnetic field lines be almost parallel to the constant density surfaces, especially in the case of the O cutoff layer. This condition is not so easy to be satisfy in minimum-B fields; anyway, some regions approximatively satisfy this condition. Something about the mutual orientation of the field lines with respect to the O

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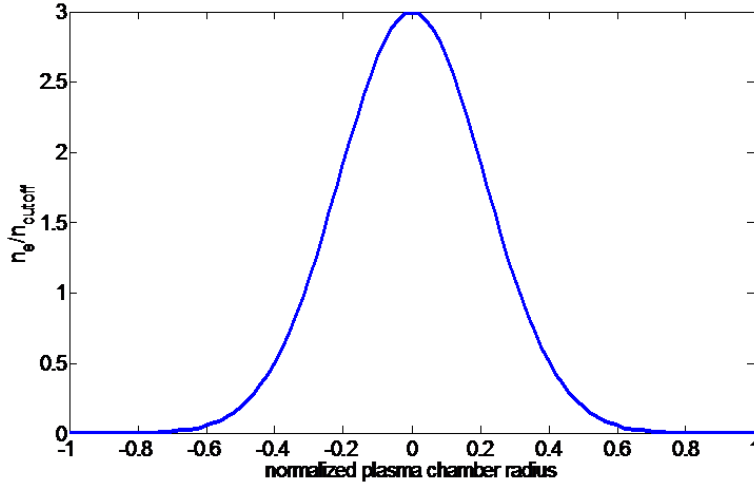


Figure 5.1.14: Assumed profile of the electron density inside an ECRIS with respect to the chamber normalized radius.

cutoff surface is shown in figures 5.1.17 and 5.1.18.

Note that in the lateral part of the O cutoff the field lines curvature is less pronounced, and for some centimeters the quasi-parallel condition among the force lines and the O surface is satisfied.

Hence the main question is how to bring the O wave just in the required region, as described in figure 5.1.17. The microwave injection in ECRIS is usually provided by waveguides located in the injection flange. Hence we need a direct injection of the wave towards the O cutoff layer, with a proper angle and in an adequate region of the cutoff surface as shown in figure 5.1.17. This can be obtained through two different injection schemes. The first one is shown in figure 5.1.19. A single cut antenna could be used, with the cut oriented so that the emitted wave directs towards the O cutoff layer. However in this way the great amount of the power reaches the cutoff surface with a non optimal angle, as the microwave beam is not focused. Moreover the space to locate the waveguide between the chamber walls and the O cutoff surface is very small. An improved scheme can be designed, requiring two metallic parabolic mirrors for the microwaves, placed as in figure 5.1.20. Now the microwaves are brought towards the O cutoff layer through the two parabolic (focusing) mirrors. The mirrors displacement should be adapted to launch the microwaves towards the O cutoff surface region. More detailed simulations about the microwave path inside the chamber will be

5.1 The experiment at WEGA facility

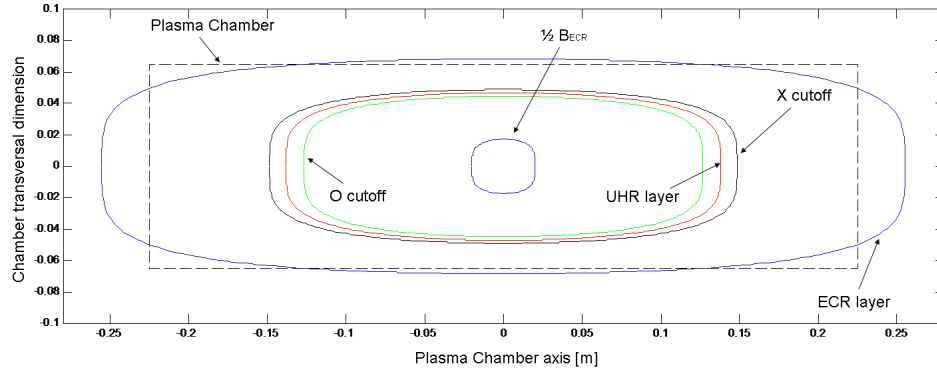


Figure 5.1.15: Real cutoffs and resonance localization for the different modes propagating inside the plasma of a typical ECRIS. The second cyclotron harmonic layer is also shown.

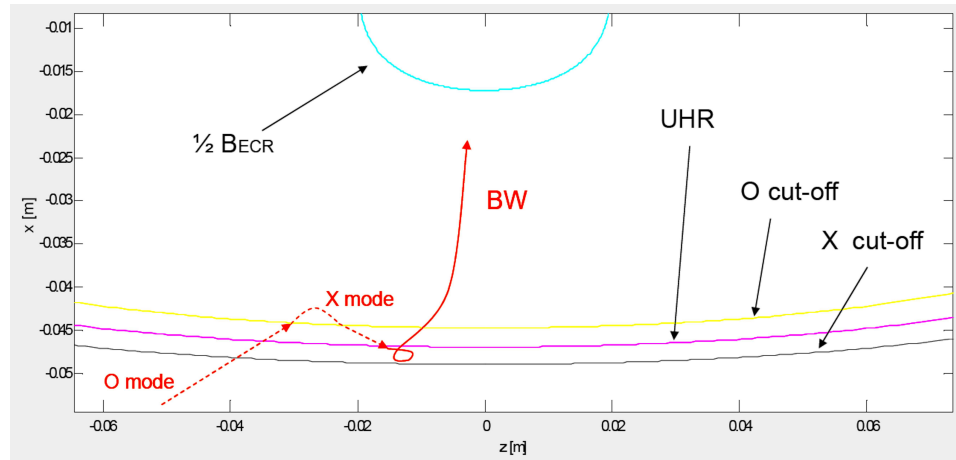


Figure 5.1.16: Tracing of the different mode propagation paths with the relative conversion and absorption layers in a SERSE like ECRIS system.

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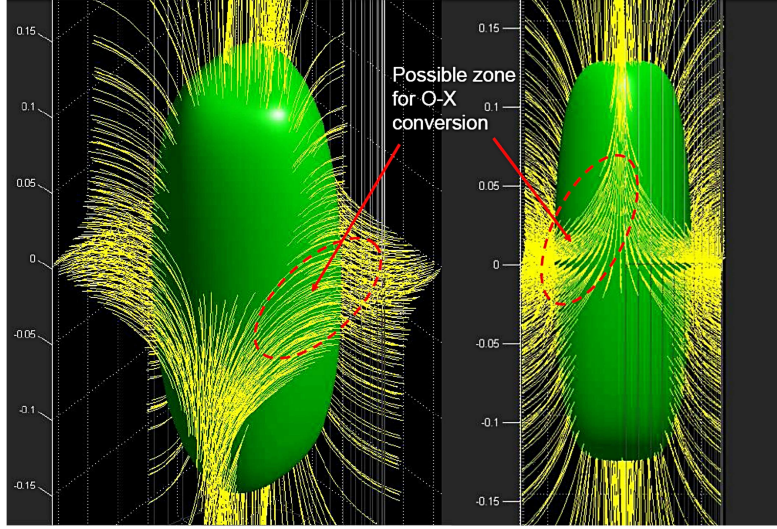


Figure 5.1.17: Three dimensional views of the O cutoff surface (in green) together with the field lines (in yellow). The region of possible optimized mode conversion is indicated.

necessary, in order to design an experiment based on such a launching system. Ray tracing of the electromagnetic waves traveling inside the plasma is also required. For a further optimization of the O-X-B conversion, the wave should be left hand elliptically polarized, with the main axis of polarization along the magnetic field direction.

5.1.5 Advantages and unknown phenomena of the OXB conversion in ECRIS

The main unknown phenomena about the OXB in ECRIS are the plasma startup and the MHD stability. The assumed density profile could be achieved self-consistently immediately after the plasma ignition, but at the beginning the various cutoffs and resonance layers (except for the ECR and the first cyclotron harmonic, which do not depend on the density) are located differently with respect to the ones shown in the previous pages. This may limit the heating till overdense plasmas and a sort of two stage heating, initially by means of the ECR, may be helpful. In cases of variable magnets it should be possible to increase the mirror and the hexapolar field so that

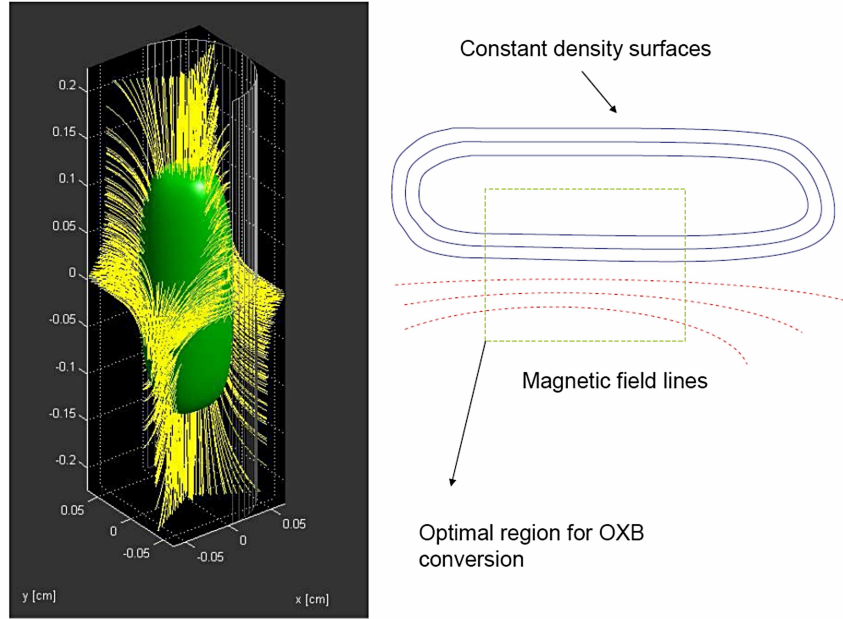


Figure 5.1.18: 3D view of the O cutoff layer together with a two dimensional representation of the constant density layers. The optimal region for the O-X conversion is emphasized.

the ECR surface exist inside the plasma. Once reached a density near to the cutoff one in the plasma core, the magnetic field should be decreased below the ECR value everywhere, so that the OXB conversion takes place and the peaked density profile is created.

The main advantages of the OXB conversion is obviously represented by the strong increase of the electron density and the experiment on WEGA here reported confirms that the higher is the density, the lower is the number of high energy electrons. However, to obtain highly charged ions temperatures on the order of keV or tens of keV are needed. Temperatures order of magnitude lower were observed on WEGA, because of the lower power density with respect to the ECRIS one. Rescaling the power density, we obtain that in a SERSE-like ECRIS it is almost one hundred times the WEGA's one. Although no well known scalings of the temperature with the power are available, it is reasonable to suppose that temperatures on the order of few keV in ECRIS are obtainable in case of BW heating.

The application of the innovative concept here outlined implies the revi-

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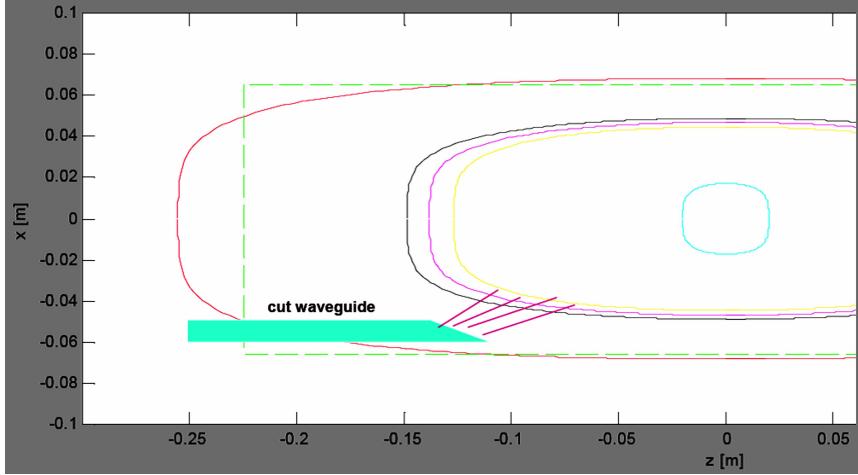


Figure 5.1.19: Scheme of a possible injection of the microwave in order to have the O-X conversion at the O cutoff layer. The dimensions of the SERSE source are considered as an example.

sion of all the guidelines contained in the "ECRIS standard model". Many problems may arise by the point of view of the plasma confinement, as the mirror ratios must be strongly reduced. Again, this problem can be treated by the two usual different points of view: the single particle and the macroscopic (fluid) dynamics. The main question concerns the MHD equilibrium and stability, which may result seriously damaged. We will need to experimentally determine a new "stability island". One possible technique to avoid any plasma instability may be to increase the periphery magnetic field once achieved the overdense plasma state, that is possible because of the high versatility of the magnetic system in ECRIS; however the effects of the ECR layer coming again into the plasma cannot be easily predicted and must be checked during the experimental operations.

Another key point is the possible formation of HEL-Hot Electron Layers observed by Golovaniwsky and co-workers in an experiment based on the BW heating of 2.45 GHz plasmas [42]. They observed a strong absorption of BW at the cyclotron harmonics, leading to very high energy electrons, above 700 keV . However, their microwave launching scheme was quite different with respect to the OXB injection. They launched an X wave perpendicularly with respect to the plasma chamber axis, and the magnetic field was more similar to that used in MDIS instead of ECRIS: it decreased

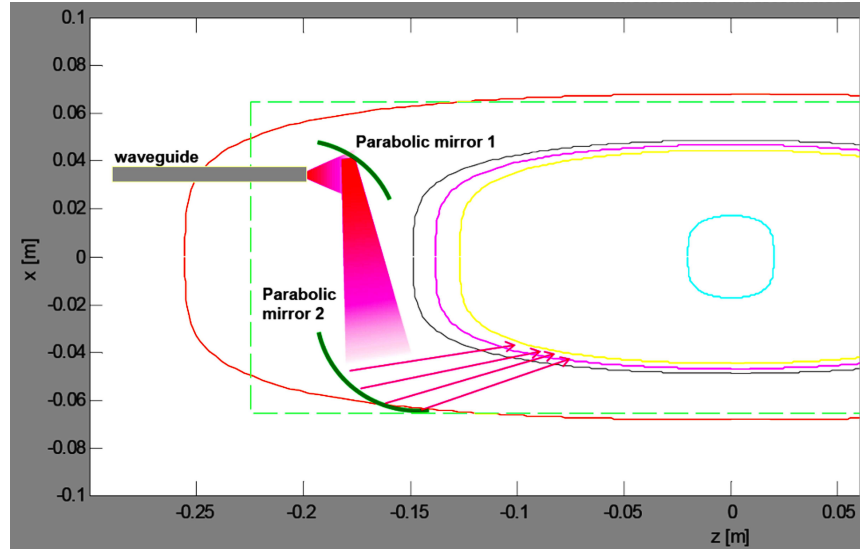


Figure 5.1.20: Improved design of the microwave injection for O-X-B conversion. Two parabolic mirrors are used to bring a focused microwave beam in the optimal region of the O cutoff layer.

from the center to the periphery of the plasma chamber. BW were generated by means of the direct X-B conversion and they were reflected by the UHR layer, moving exactly in the opposite direction with respect to the incoming X-mode, i.e. towards the periphery of the plasma chamber. In this way the power is deposited in the plasma halo⁴ leading again to the formation of an hollow shape of the electron density and, because of transport, to a non efficiently heating of the plasma core. A great amount of energy is deposited in small layers, correspondingly to the cyclotron harmonics, involving few particles which can reach very high energies. With the scheme proposed in figure 5.1.20 the HELs formation should be avoided, thus depositing the RF power centrally.

⁴The cyclotron harmonics were localized in the periphery of the plasma chamber

Chapter 6

Conclusion and Perspectives

6.1 The ECRIS R&D: "brute force" or "ad hoc" energy transfer from microwaves to electrons

The theoretical and the experimental results here reported demonstrate that the plasma dynamics must be further investigated to overcome the limits of the modern ECRIS. Single particle and PIC simulations contribute both to better understand the basic plasma physics of the ion sources, and in particular they permitted to explain several experimental results and unsolved questions. First of all, the frequency tuning and the two frequency heating are now made clear and the possibility of the multi-heating of the plasma is open. From the point of view of experimentalists, the VENUS measurements on the ultra-hot electrons and the CAPRICE beam shape characterizations have been the most important results here analyzed. Especially for VENUS, it has been demonstrated that the future developments of ECRIS will be affected by the presence of some detrimental phenomena: the reduction of high energy X-rays implies a limitation of the source performances due the violation of the scaling laws at very high magnetic gradients. The existence of ultra-hot electrons can be explained by means of the the PIC simulations, still remaining inside a "single particle" description. However, the turbulent heating evidently play a role and only further experimental results will be able to evaluate the relative weight.

The PIC simulations, because of their self-consistent nature, should be able to put in evidence any possible formation of plasmons and cavitons (i.e. the "cavitonization process"), but nothing has been observed so far. On the other side, although the PIC simulations explain, by means of the Doppler splitting and of the Arnold diffusion, the existence of electrons above the

6.1 The ECRIS R&D: "brute force" or "ad hoc" energy transfer from microwaves to electrons

W_b barrier, electrons up to 2 MeV have not been observed; this is probably connected to the limited time of the simulations. Future simulations would be able to simulate up to 1 msec thus giving some insights about the electron heating till extremely high energies. The turbulent heating hypothesis remains just an alternative but it is to be checked experimentally.

Another important results is that the single particle and the PIC approach have given comparable results, and in many cases an information coming from one approach was made more clear and complete from the other one. Although the PIC approach is certainly the most effective to describe the plasma dynamics of ECR ion sources, the single particle simulations, being based on the real ECRIS magnetic field, have given some additional information confirming the PIC results.

As written in the introduction, the synergy between the theoretical investigation and the experimental measurements has opened a new road for the future study of the ion sources. *Instead of the **brute force**, i.e. higher power and higher magnetic field as a consequence of the increase of the microwave frequencies, we expect that an appropriate transfer of the electromagnetic energy to the electrons boosts the ECRIS performances, through the combination of the techniques investigated in this thesis.*

On this purpose, our results about the TCFH are coherent with the ones obtained in the experiment carried out at JYFL, Finland, with a WNG-White Noise Generator [83]. They feature the effectiveness of the WNG particularly for the production of HCIs. The improvement seems related to the creation of a large resonance volume with quasi-uniform accelerating electric field, as the broadband spectrum excites a large number of contiguous resonant modes inside the plasma chamber. The uniformity of the global electric field (that is given by the superposition of several fields, each related to a particular mode) avoids the formation of fluctuating plasma densities and temperature gradients on the ECR volume edge and this improves the ion dynamics, as demonstrated in section 4.13.8. No significant effects on the electron dynamics seem to be responsible of the increase of the extracted current for the ions at higher charge states. This is confirmed by the X-ray spectra measured at JYFL with and without WNG [83]: there are no appreciable differences in energy distribution of X-rays, which confirms that the electron heating is weakly influenced by the use of WNG, at least as it regards the higher energy component detected with the X-rays. Then, in an ECRIS that operates with a microwave broadband spectrum, a great production of high energy X-rays is avoided, whereas in the case of large ECR volume obtained by means of flat magnetic field at the resonance

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[117]-[118], a notable increase of such X-rays occurs. This experimental evidence is encouraging also for TCFH. The optimization of the ion dynamics, which appears as the main advantage by using the WNG, may be obtained also by using the TCFH and by means of frequency tuning. In addition, if compared with WNG heating, the TCFH may permit one to perfectly "tailor" the electromagnetic field and then to minimize the power needed to reach high performances. It should avoid the strong reflection of input power, which was observed when using WNG, due to the mismatch of some modes that are not well coupled with the plasma chamber.

The experimental measurements reported in section 4.5 about the production of ultra-hot electrons show that their amount strongly increases as the magnetic field gradient (or the mirror ratio) decreases. These data also demonstrate that for very high-B gradients the CSD shifts towards lower charge states. Our simulations (see section 4.11) suggest that this effect is due to the strong gradient, which limits the amount of energy that an electron can gain on each crossing of the ECR zone. Hence, it may be convenient to operate at lower B gradients to get a better heating rate, but on the other side, the increase of the B gradient permits the decrease of the amount of ultra-hot electrons; a compromise is found experimentally. A faster energy transfer from the wave to the electrons due to the TCFH should be able to generate a great amount of warm electrons even in the presence of strong B gradient, but limiting the amount of ultra-hot electrons. This means that the next generation ECR ion sources should be able to operate in multi-frequency mode.

An almost ideal setup may be based on the microwave injection from four microwave generators, two klystrons, and two TWTs. It should be better to use four TWTs because of their ability to frequency tuning, but it is not strictly necessary. The two klystrons should operate at medium-high powers (1 to 2 kW), whereas even two low power TWT generators may be sufficient. By fixing two distant frequencies, we can associate each one to another close frequency. In this way, we will associate the benefits of the TFH to those of the TCFH, together with the optimization of the frequency tuning provided by the use of the two TWTs.

As the simulations put in evidence the crucial role of the frequency not only in terms of scaling laws, reached through large variation of the frequency, but also in terms of slight variation of the output frequency of the microwave generators, this means that a detailed study of new broadband microwave generators may allow to strongly increase the ECRIS performances without any expensive improvement of the superconducting magnets

technology and/or RF generator output power.

Then the improvement of the third generation ECRIS will require a wide band amplifier working with microwave frequencies around 28 GHz and able to provide several kW of output power. This device does not exist yet on the market, but the development and the realization of a microwave generator named Gyro-TWT (Traveling Wave Tube) working at high frequency and high power with broad bandwidth is in the plans of different laboratories interested to the huge advantages in terms of beam availability¹.

According with the experimental and numerical results presented in this thesis we can imagine to design future ECR ion sources by applying not only the variable frequency generators, but also a different magnetic field design to remove those regions where the field lines length L is high. The position of the ECR surface is also important, as the further it is from the minimum B value, the lower T_e is. The multi-frequency heating should be used, as suggested above, to avoid any performance worsening due to the excessive increase of the magnetic gradient. The existing tools like bias disk and electron donors should be used to optimize the sources performances.

In addition, a research programme has been defined at MPI - Greifswald in order to define the new trap design to fully exploit the OXB conversion.

6.2 Multi-frequency heating test preparation

The simulations reported in this thesis permit to qualitatively predict the plasma dynamics in case of multi-frequency heating. An experimental system able to provide a multi-frequency heating of plasmas, with the possibility of the several waves phase relationship control has never been designed in the past.

A similar set-up is shown in figure 6.2.1 according to the characteristics of the SERSE source. The injection flange of such ion source presents two $WR62$ waveguide apertures. This feature allows to feed the plasma chamber with two or more discrete frequencies. A Traveling Wave Tube Amplifier (TWT) and a Klystron may be connected to these inputs through the two $WR62$ microwave lines. Two high insulation waveguide reflectometers will be inserted in such microwave lines in order to monitor, by means of power sensors, the reflected power from each input and the transmitted power

¹Gyro-Traveling Wave Tubes (TWT) are microwave amplifiers able to provide several kW of output power. The wave amplification is achieved by means of a resonant interaction between an electron beam and the electromagnetic wave at the harmonic of the electron cyclotron frequency.

6. CONCLUSION AND PERSPECTIVES

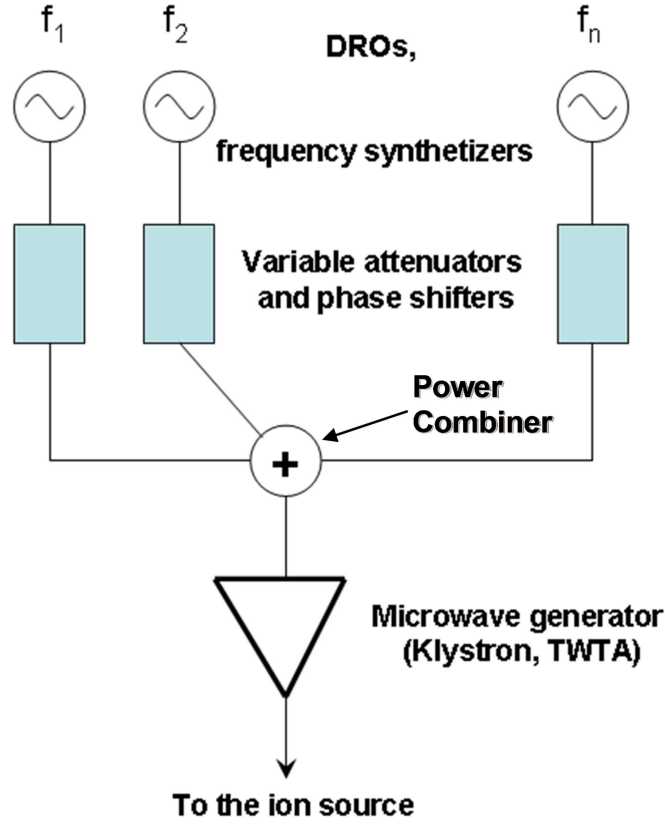


Figure 6.2.1: Design for an experiment aiming to the multi-frequency heating of ECRIS plasmas with the wave's phase control.

between the two ports. The forward power sent to the plasma chamber from the two lines will be decoupled with directional couplers and measured with two power sensors. The microwave frequencies, produced by the signal generators, will be optimized in order to achieve the maximum extracted current measured with a Faraday cup; the reflected power will be measured to correlate the beam output and the beam ripple to the microwave coupling to the plasma. A preliminary measurement will be done by sweeping the TWT over the whole frequency range in order to get a guideline for the next experiments with combination of TWT and Klystron including the relative phase shift. The registered charge state distribution will be used as a reference. The following tests will be made by varying the relative phase

between the two chosen microwave frequencies produced from the signal generators and amplified through the TWTA. A spectrum analyzer will detect any noise or distortion introduced by the two waves amplification. Then a third microwave frequency produced by the Klystron will be introduced. The microwave power will be systematically distributed between the three frequencies taking care in protecting the instruments from damaging power exceeds.

The measurements in terms of beam emittance will be essential in order to confirm the simulations related to the ion dynamics presented in section 4.13.8. The results will be compared with the simulations and the results of preliminary experiments carried out in the past years in particular at GSI [6]. The new scheme of multi-frequency heating is expected to increase the current of highly charged ions either for the existing ECR ion sources (working in the range $14 - 18 \text{ GHz}$) and for the next generation of ECRIS ($28 - 37 \text{ GHz}$). Similar tests will be carried out at JYFL and at other facilities according to the techniques here proposed.

6.3 Considerations on plasma diagnostics

During the last years, the main limitation in the understanding of ECRIS plasma physics has been related to the plasma diagnostics. Up to now the information about the changes of plasma characteristics at different operational conditions have been mainly provided by the charge state distribution; however this is only an indirect sign of plasma properties. Additional diagnostics tools are needed in order to get a complete characterization of plasma dynamics for different conditions: variations in power, magnetic field, electron supplies, microwave frequency can be studied by means of a combination of different diagnostic methods. Methods based on the detection of the plasma emitted radiation, because of their non-intrusive nature, are certainly an advantage, however the recorded information always corresponds to an integration by one direction over the whole plasma volume. Then the use of "spatially resolved" techniques will be of primary importance for a complete characterization of ECRIS plasmas especially for next generation sources, where the high energy content will make unavailable the invasive methods like probes or similar instruments. Several diagnostics techniques can be applied on ECRIS, and they will be summarized in the following.

OES-Optical Emission Spectroscopy. As electrostatic probes are not available for ECRIS plasmas (high temperature and densities may destroy the probe tips), the optical emission spectroscopy is a powerful tool to

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characterize the plasma in terms of electron density and electron temperature.

Spatially Resolved OES. The possibility to evaluate the distribution of cold, warm and hot electrons inside the plasma chamber may allow to test some theoretical models:

- according to the ECR-heating properties, a great amount of high energy electrons should be created in a slab located around the ECR surface. This assumption is confirmed by PIC simulations performed at low density (see section 4.13) which point out that high density and sufficiently energetic plasma is formed in a small volume surrounding the ECR surface. In this volume the greatest amount of ionization processes occurs and the ion dynamics is also affected by this strongly inhomogeneous plasma.
- An alternative scenario is suggested by Golovanisky [17], who stated that the plasma core heating occurs provided that an electromagnetic-electrostatic wave conversion occurs in proximity of ECR surface. Parametric decay should be observed in this case, then a plasma emission of electromagnetic waves at frequencies below the injected one.

The spatially resolved OES may help to distinguish between the two above models. Some confirmations about the validity of the first one come from the experimental results shown in section 5.1.1 and obtained by means of a spatially resolved optical spectroscopy method applied on a 2.45 GHz source.

However we expect that a more complete scenario takes into account both the previous mechanisms. A detailed investigation by means of the optical spectroscopy may allow to determine the weight of each one.

In section 4.2 it has been reported that the experiments demonstrate that the shape of the extracted beam strongly depends on the microwave frequency. Then the beam formation strongly depends on the physics of the following steady state discharge, and the diagnostic tools able to detect the beam shape after the extraction, and emittance measurements, are essential.

Another technique should be employed to detect density and energy content fluctuations inside the plasma. **The X-ray spectroscopy** allows to get information about the high energy component of plasma electrons. This technique is useful especially for the detection of "ultra-hot" electrons, i.e. those electrons which energy overcomes the stochastic barrier imposed by the single particle model of ECR-heating. At high power, as the one that

will be used in next generation ECRIS, the creation of high energy electrons will be of primary importance.

A more detailed diagnostic tool based on X-rays is the **X-ray pin-hole camera imaging**. Together with spatially resolved OES it permits to detect the spatial distribution not only of energetic electrons, but also of moderate and highly charged ions.

As the heating till high energies is a time-expensive mechanism, which requires several hundreds of μs to reach MeV energies, we need a **X-ray spectra characterization either for ECRIS operating in pulsed mode**.

The future experiments on the OXB conversion at MPI - Greifswald and those ones designed to check the possibility of plasma heating driven by turbulence will require the **measurements of plasma emitted electromagnetic waves**. This experiment allow to understand the mechanisms involved in wave-plasma interaction, by permitting to check this phenomenon for different plasma parameters. In particular, in some experiments performed during the last ten years a worsening of ECRIS performances for high power has been observed. This phenomenon may be related to some decays of the incident electromagnetic wave that can reduce the total amount of energy absorbed by the plasma. When it will be possible to verify the OXB conversion efficiency in ECRIS, electric field probes will be necessary, in order to investigate whether and where the conversion takes place. These probes have been designed for WEGA [21] and a slightly different design may be prepared.

Some of the above techniques are already used in ECRIS, especially as regards the X-ray diagnostics. Many others have been largely used in other plasma based devices, and especially in plasmas for fusion applications.

In our laboratory in the last years we have focused our attention on the characterization of the plasma chamber scattering parameters and more recently on the detection of standing waves in dense plasmas (see section 4.3). These methods are based on the Spectrum and Network Analyzers. Additional measurements of electric field distribution inside the plasma chamber will be performed also on high performance ECRIS. In particular, some systematic investigations about the electric field distribution under different condition of density and temperature (i.e. of magnetic field and vacuum pressure) are mandatory. There are many signs which indicate that the changes in the heating process may in turn affect the wave propagation inside the plasma chamber.

6.4 A future warm and overdense dream machine

By looking to the results of ECR ion sources in the last thirty years, it is evident that such a rush cannot be continued "ad libitum". For highly charged ions, the increase of available beam current accounted to one order of magnitude per decade. Moreover, the technological limits of the superconducting magnets and of microwave generators, as well as the increasing cost, suggest that the next steps towards the 4th generation ECRIS ($f > 56 \text{ GHz}$) will be slow, in spite of some recent proposal [C. Lyneis ECRIS08].

The alternative roadmap that has been proposed here may help to continue the "beam current rush", provided that some clues will be confirmed by the experiments proposed in this chapter, as it seems reasonable. A further step in this direction is offered by the design of a "warm overdense dream machine". The results obtained with WEGA about the OXB conversion, have opened new frontiers for the design of ion sources based not only on the ECR-heating, but also on the Bernstein waves heating. In collaboration with the Max Planck Institute, a preliminary test will be carried out in the second half of 2009 about the applicability of the OXB conversion on ECRIS-like devices. On the basis of the obtained results, a "dream machine" may be designed, able to operate with overdense plasmas without any limitation due to the production of high energy electrons. In addition it is clear that such a machine should operate with relatively low magnetic fields, if compared with those ones used nowadays for ECRIS of 3rd generation. The extraordinary advantages coming from the completely different heating mechanism should overcome the disadvantages due to the worsening of the confinement even if the RF power may be larger than 10 kW. Finally, the dream to perfectly tailor the EEDF (avoiding the formation of excessively high energy electrons, with a great amount of particles in the warm population), seems to be not so far and it may be the key for further increases in the next ten years. Of course, only the experiments will give the right answers to these questions.

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