Investigations on the electron dynamics in the Neutral Loop Discharge

Dissertation
zur Erlangung des Grades eines
Doktors der Naturwissenschaften
in der
Fakultät für Physik und Astronomie
der Ruhr-Universität Bochum

von
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aus Iasi

Bochum 2009
Dissertation eingereicht am


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Chapter 1
Introduction

Plasma discharges in their multiple facets of generation [1, 2] are widely used in industrial applications and manufacture process. Material surface properties modification, etching and deposition of thin films and coatings [3] used in various applications such as photovoltaic solar cells [4], micro and nano-technologies [5], surface modification of biomaterials, sterilization [6, 7], to name here a few, are involving in the production process plasma discharges. There are several ways for generating plasma discharges. Yet with increase of the wafer size, better control of the plasma uniformity over a large area is necessary. Moreover to achieve the desired industrial process requirements the plasma discharges must have a sufficient high electron density, low electron temperature and operate at relatively low pressures, usually bellow 1 Pascal in order not to damage the treated material surface. In this work, special consideration is given to radio-frequency (RF) driven plasma discharges, and in particular on the investigation of a novel plasma discharge type, the Neutral Loop Discharge [25].

The Neutral Loop Discharge (NLD) uses a special magnetic field configuration, with a neutral loop (NL) region where the magnetic field vanishes. The NLD discharge allows operation at pressures considerably lower than the conventional ICP discharges, due to an efficient collisionless electron heating mechanism. One of the open questions on the fundamentals of the Neutral Loop Discharge is how this electron heating mechanism really works. A single picture slab model developed previously by Yoshida and Uchida stated that the chaotic movement of the electrons in the Neutral Loop region can be the source of its effective heating process through a collisionless randomization of their phase-space trajectories [26, 28]. However experimental investigations to verify this theory are lacking.

The aim of this work is to investigate the NLD electron heating mechanism. For this a novel planar inductively coupled magnetic NLD was design and built to facilitate both
fundamental and application oriented investigations. The versatility of the experimental setup allows operation of several plasma discharges. ICP, NLD and a planar type Helicon discharge can be operated and investigated. Since the plasma chamber remains the same, this provides a more reliable comparison between the respective plasma parameters of each discharge type.

Particular to the NLD, special interest was paid to provide diagnostic access to the Neutral Loop (NL) region, as well as to the application region of the discharge, located close to a possible substrate plane. From previous investigations it is believed that the NL region of the discharge is the main plasma production region. However this work proved different.

An improvement on the discharge characteristics can be done only by understanding how the discharge operates. A wide spectrum of diagnostic methods have been applied, however non-intrusive such as optical diagnostics where preferred. Thomson scattering, LIF together with Radio Frequency Modulated Optical Emission Spectroscopy (RF-MOS) and Phase Resolved Optical Emission Spectroscopy (PROES) to name here a few have been applied to study the plasma parameters in a large operating conditions.

1.2 Overview of the thesis

Chapter two gives a short introduction on RF plasma discharges and in particular to Inductively Coupled Plasma discharges (ICP) and the Neutral Loop Discharge (NLD). A discussion on the collisionless stochastic electron heating in the ICP is presented. This short comparison between the two discharges is aimed for a better understanding of the following chapters. The NLD heating mechanism, developed previously, together with the basic formalism on the Yoshida slab model is described.

Chapter three presents in detail the experimental setup together with the diagnostics used in this work. It starts with an overview of the industrial NLD reactor and introduces the new design concept applied in this work. Each diagnostic method has a short summary on its characteristics, followed by the implementation on the setup and examples of measured data.

In chapter four a thorough plasma source characterization is done. Spatially resolved Langmuir probe measurements in the plasma production chamber and at a
possible application plane, both in the NLD and ICP have been made. Thomson scattering, Radio Frequency Modulated Optical Spectroscopy (RF-MOS), and Optical Emission Spectroscopy (OES) gave a deeper insight on the electron characteristics of these discharges. These measurements proved for the first time the presence of a diamagnetic drift in the discharge.

In chapter five an investigation by means of Thomson scattering the electron drift is done. An additional electron drift component in the direction of the scattered vector is leading to a shift of the distribution function. Phase resolved Thomson scattering measurements have been performed for the first time in an ICP and NLD discharge. Oscillations of the RF induced electric field penetrating into the plasma are leading to oscillations of the electron velocity distribution function. By changing the phase with respected to the RF signal this oscillations could be measured. From this the local current density can be inferred, combining the oscillation amplitude with the density. A novel spectroscopic diagnostic technique, the RF – Modulated Optical Spectroscopy, can also provide a direct access to electron oscillation velocity.

Chapter six looks into the wave phenomena in the Neutral Loop Discharge. Preliminary measurements proved the existence of a wave traveling parallel to the magnetic field. For the given plasma conditions and taking into consideration the wave dispersion relation, Whistler waves can propagate within these boundaries. The flat antenna actually behaves as a wave launcher. Moreover a wave – particle interaction is a powerful collisionless electron heating mechanism in the low pressure regime. For the Neutral Loop Discharge this is found to be the dominant heating mechanism.

This work ends with a summary of the issues addressed and suggestions for further study on the wave heating and propagation in the Neutral Loop Discharge.
2.1 Inductively Coupled Plasma (ICP) discharge.

Although they are known for almost 150 years, Hittorf in 1884 being the first one to obtain this discharge type, radio-frequency plasma discharges had been studied with great interest only in the last 50 years. The gas breakdown and plasma generation process is achieved by inducing an electric field through a planar antenna or a coil winded around the discharge vessel as shown in figure 1.1. The research has been focused mostly on understanding the dominant power coupling mechanism, i.e. when the discharge operates in either capacitively (E-mode), inductively (H-mode) or in a so-called Hybrid mode. However self-igniting the plasma in the H-mode alone is not possible [8, 9]. Thus the E-mode will always precede the H-mode. Increasing the applied RF power a sudden “mode-transition or E-H jump” occurs. This transition effect and the plasma behavior have been widely investigated in [9, 10, 11, 12, 13] and a thoroughly review on this subject is beyond the subject of this work.

An advantage of the ICP reactors is that the antenna is not directly exposed to the plasma, but separated by a dielectric, usually made from quartz. This together with the simplicity of the design and an efficient plasma production made the Inductively Coupled Plasma (ICP) discharge a viable alternative for material processing, becoming widely used in nowadays industrial discharges. However in respect with other discharges, the efficient power coupling in the ICP is somehow restricted by its natural density limits. The RF frequency limits in which this discharges operates for material processing are bounded in the low region by the critical sheath frequency and in the upper region by the reactor size. A schematic view of two commonly used ICP reactor types is shown in Figure 2.1 and a more detailed in-depth is given in [14].
When an RF electrical field is applied to the antenna, a time dependent \( z \)-directed magnetic field \( B(z, t) \) is created accordingly. In cylindrical coordinated this is expressed using Faraday’s law:

\[
\nabla \times E = -\frac{\partial B}{\partial t} - \frac{1}{r} \frac{\partial}{\partial \theta} (r E_\theta) - \frac{\partial E_r}{\partial r}
\]

(2.1)

The \( B(z, t) \) time variation will induce an azimuthal electric field (\( \theta \) direction), encircling \( B(z, t) \), thus creating rf currents. Using Maxwell equation one can relate the electric field and the current density \( J \):

\[
\nabla \times B = \mu_0 J + \mu_0 \varepsilon_0 \frac{\partial E}{\partial t}
\]

(2.2)

where \( \varepsilon_0 \frac{\partial E}{\partial t} \) represents the displacement current. In an ICP this term can be neglected since \( \omega_c \ll \omega_p \).

Thus equation 2.2 becomes:

\[
\nabla \times B = \mu_0 J
\]

(2.3)
Rewriting and combining equations 2.1 and 1.3 and expressing the current density as a function of the electric field by Ohm’s law $J_\theta = \sigma_p E_\theta$:

$$\frac{1}{r} \frac{\partial}{\partial r} \left( \frac{\partial}{\partial r} \left( r E_\theta \right) \right) + \frac{\partial^2 E_\theta}{\partial z^2} - \frac{E_\theta}{r^2} = i \omega_0 \mu_0 \sigma_p E_\theta$$ (2.4)

where $\mu_0$ represents the permeability of the free space and $\sigma_p$ is the plasma conductivity defined as:

$$\sigma_p = \frac{e^2 n_e}{m (\nu_{ce} + i \omega_f)}$$ (2.5)

with $\nu_{ce}$ representing the electron collision frequency for the momentum transfer. A plot of $\nu_{ce}$ as a function of pressure taken from [15] is shown in Figure 2.2. The plasma conductivity is an important parameter in efficiently tuning the power coupling into the discharge. However the electric field penetration into the plasma is limited to a scale of comparable length to the skin depth $\delta$.

![Figure 2.2: Electron collision frequency as a function of pressure and driven RF frequency taken from [15].](image-url)
Therefore the electrons are accelerated and gain most of their energy from the induced electric field within the skin depth region. For a better understanding of radio frequency plasma discharges a short description on the interaction between electromagnetic waves and plasma is given. We assume here that a plasma discharge is already ignited and we focus only onto RF waves impinging on the plasma bulk. Based on the applied frequency of the incident wave, one can distinguish two cases:

1) Collective plasma – wave interaction when \( \omega > \omega_{pe} \)

where \( \omega \) is the incident wave frequency and \( \omega_{pe} \) is the plasma frequency defined as:

\[
\omega_{pe} = \sqrt{\frac{e^2 n_e}{\varepsilon_0 m_e}}
\]  \hspace{1cm} (2.6)

In this case the plasma behaves like vacuum with \( \varepsilon_0 \); the wave “sees” the entire plasma and interacts collectively with it. The wave will experience a lossless propagation through the plasma.

2) Localized plasma – wave interaction when \( \omega < \omega_{pe} \)

In this case no wave propagation occurs through the plasma bulk. Electrons will sample localized heating regions. Energy transfer from the wave can occur within a finite region which is usually referred to as Landau damping. The spatial decay of a TEM wave due to wave – particle interaction is investigated later in this work in Chapter 6.

In an ICP the antenna can be seen as a wave – launcher, generating transverse electromagnetic waves (TEM) propagating in the \( z \) - direction containing both electric \( E \) and magnetic \( B \) field components.

Electrons in RF discharges acquire energy from the oscillating induced electric field. The electron heating is the core process of sustaining the plasma, as electrons are loosing their energy through elastic and inelastic collisions with neutrals and charged particles. There are two mechanisms for the electron heating in plasmas, depending on the operational pressure of the discharge.

In the high – intermediate pressure regime (20 – 1 Pa) Ohmic heating is the dominant mechanism. Through collisions the phase coherence of the electron motion is destroyed. The randomization of the electron motion is the key mechanism of the energy transfer to heavy particles. In an RF discharge the induced electric field is oscillating
periodically and makes no net gain on the electron energy. The Ohmic heating is however efficient only when the collision frequency is high. As mentioned above the energy transfer from the induced electric field to the electrons is made within the skin depth region. Assuming an exponentially decaying spatial profile of the electric field in the $z$-direction:

$$E(z) = E_0 e^{-\gamma z} e^{i\omega t}$$

with $\gamma = \frac{1}{\delta} - ik$ being the complex wave vector and equation 2.4 one obtains for the classical skin depth [16]:

$$\delta = \frac{\delta_0}{\cos \frac{\epsilon}{2}}$$

with $\epsilon = \arccos \left( \frac{1}{1 + \frac{\nu}{\omega_{pf}}} \right)^2$ and $\delta_0 = \frac{c}{\omega_{pe}} \left( 1 + \frac{\nu^2}{\omega_{pf}^2} \right)^{1/4}$

For the limiting case of plasmas with large or small collision frequencies, $\nu >> \omega_{pf}$ and $\nu << \omega_{pf}$ respectively the skin depth becomes:

- $\nu >> \omega_{pf}$ (collisional limit): $\delta = \frac{c}{\omega_{pe}} \left( \frac{2\nu}{\omega_{pf}} \right)^{1/2}$ - the energy is collisionally dissipated within the skin layer
- $\nu << \omega_{pf}$ (collisionless limit): $\delta = \frac{c}{\omega_{pe}}$ - no net energy transfer

However the ICP discharge operates as well bellow 1 Pa. This suggests that there is another mechanism accounted for an efficient power transfer. This so called non-collisional heating has been investigated experimentally e.g. by Godyak and Popov in [17]. Following it has been treated in [18, 19, 20, 21] and several models [22, 23] have been proposed. For a CCP discharge the stochastic electron heating is treated in [24] and the references.
This process can be understood also as a collisional process, but this time the electrons are “colliding” with the sheath in front of the electrode, much as in the Capacitively Coupled Plasma Discharges (CCP). For the ICP this “sheath” is the skin layer. The electrons can gain energy from the electric field if their transit time through the skin layer is less than the RF period. Most of the nowadays industrial discharges are operating in this pressure regime, and therefore understanding their heating mechanisms can improve the discharge characteristics. Godyak in [17] demonstrated that the stochastic heating is the dominant mechanism at this low pressure.

2.2 Neutral Loop Discharge (NLD)

Increase demand on technological applications of plasmas in particular for plasma etching, low pressures, high density plasma discharges become desired. Although ICP discharges have proven to be advantageous in many applications, yet in the low pressure regime their operation becomes unstable. Another critical aspect of the ICP’s is that they are not achieving the required process uniformity over large areas.

These industrial requirements motivate development of new types of plasma discharges to meet the processing needs. The Neutral Loop Discharge (NLD) introduced by Taijiro Uchida [25] in the mid nineties had shown a promising alternative to conventional discharges. Chapter 2 of this work, extensively analyze the NLD properties in a newly design reactor, here only a short overview on its fundamentals and proposed electron heating mechanism [26, 27, 28] is given. A general review on the NLD discharge is made in [29].

The use of a special inhomogeneous magnetic field configuration with a so called Neutral Loop (NL) region where the magnetic field vanishes makes the NLD special. Moreover the magnetic field strength needed to form the NL is much lower than in ECR or Helicon plasma.

The NLD plasma is generated by applying an RF electric field along the NL. The plasma characteristics show a clear dependence on the magnetic field strength, RF power and operating pressure [30] compared to an ICP operated in similar conditions. The effect of the magnetic field configuration not only lowers the operational pressure bellow 0.1
Pascal, it can also spatially control the position of the NL ring and thus allowing a better uniformity over large processing areas. Reference [25] points out to a private communication where a first preliminary investigation proved high sputter etching rates with a high uniformity. The high spatial controllability of the plasma and plasma parameters made the NLD a strong candidate for industrial processing. The ULVAC Company is nowadays producing industrial NLD reactors. Typical etching profiles obtained by the ULVAC NLD Etcher are shown in Figure 2.4.

![Etching profiles obtained in the ULVAC reactor](image)

Investigations carried out in the NLD have been predominantly focused on its application potential. High etching uniformity, reported in [31], has been achieved by temporal and spatial control of the NLD by applying a sinusoidal current on the magnetic coils. NLD very fine pattern SiO$_2$ etching and dry etching have been investigated in [32, 33, 34]. Successful application of the NLD on organic low-k materials etching is reported in [35]. Experimental characterization of the electron properties of the NLD plasma as a function of applied RF power, magnetic field strength, pressure, RF driving frequency
have been reported in [36, 37, 38, 39]. However only few fundamental investigations on
its underlying electron heating mechanism have been reported in [40, 41] based on a
relative simple single particle picture proposed by Yoshida in [26]. This proposed
electron heating mechanism is presented in the following. The model core is based on
solving the electron equation of motion in a strongly inhomogeneous magnetic field with
a vanishing field region, the NL, and an applied oscillating electric field. The assumption
on the magnetic field configuration is shown in figure 2.5.

\[
\begin{align*}
\frac{\text{d}^2 \mathbf{X}}{\text{d}t^2} &= -e \left[ \mathbf{E} + \mathbf{v} \times \mathbf{B} \right] \\
&= \frac{m}{e} \left[ \mathbf{E} + \mathbf{v} \times \mathbf{B} \right]
\end{align*}
\] (2.9)

Textbook solutions for the electron equation of motion [42, 43, 44] in an
inhomogeneous magnetic field assume that the magnetic moment \( \mu \) is an adiabatic
invariant of the basis Lorenz equations. Adiabatic invariance states that the particle
trajectory changes “slowly”, with respect to the spatial variation of the magnetic field. In
other words if the characteristic scale of the inhomogeneity of the magnetic field is larger
than the particle gyroradius in one period, then, the net action on the particle is
conserved. However in the vicinity of a neutral point this assumption does not hold
anymore, thus charged particles trajectories become “chaotic”. This disordered motion is
called “meandering motion”. Figure 2.6 shows the result of solving equation (2.9).
This is a typical “meandering” motion of an electron moving in an “NLD” inhomogeneous magnetic field and oscillating electric field. The magnetic and the induced electric field are described by:

\[ B = \begin{pmatrix} 0 \\ 0 \\ B_z \end{pmatrix} \quad \text{and} \quad E = \begin{pmatrix} E_x \\ E_y \\ 0 \end{pmatrix} \quad (2.10) \]

Taking into consideration equation 2.9 and using the normalized variables:

\[ \tilde{t} = \omega_c t \quad \tilde{x} = \frac{x}{L} \quad (2.11) \]

with \( \omega_c = eB/m \) being the cyclotron frequency in the magnetic field \( B \) the x and z components of the motion equation can be expressed as:

\[ \ddot{x} = \omega_c \dot{x} \quad (2.12) \]

\[ \ddot{z} = -\omega_c \dot{x} + E \]

Integrating (2.12) with respect of \( \tilde{t} \) one obtains:

\[ \tilde{x} = \frac{x^3}{2} + (F + C)x \quad (2.13) \]

The nonlinear dynamics of the electron is characterized in equation 2.13 by the parameter \( F \left( \frac{E}{E} \right) \). Based on these model assumptions in [45] a calculation for a linear magnetic quadrupol, shown in the inset of Figure 2.7, on the influence of the magnetic
field on the particle energy is made. The electron trajectory is shown in Figure 2.7 and the projection on the X and Y axis of the electron trajectory are shown in Figure 2.8.

Figure 2.7: Electron trajectory in the NLD magnetic field configuration. Inset: Symetric linear magnetic quadrupol.

Figure 2.8: Projection of the electron trajectory on the X and Y axis (from Figure 2.7) as a function of time.
As shown in Figure 2.7 the electrons are confined within the NL region. When they acquire enough energy they can escape from the magnetic trap. However the magnetic field configuration, allows the electrons to escape only along the separatrices. The electron energy gain from this meandering motion is plotted in Figure 2.9 as a function of residence time in the magnetic trap.

![Figure 2.9: Electron energy gain as a function of residence time in the NL.](image)

The typical operational pressures of the NLD are below 0.5 Pa. Thus a collisionless heating mechanism must be considered for sustaining the plasma. The inhomogeneous magnetic field configuration introduces a randomization in the electron trajectory space. This has the same effect on the phase – mixing of the electrons as will collisions do resulting in an energy transfer from the electric field. However, investigations made in this work and shown later on in Chapter 6 identified another heating mechanism for the NLD.
Chapter 3
Experimental setup description and applied diagnostic techniques

3.1. Introduction

Plasma reactors are widely used both in research and in industrial applications. With the increased demand on technological applications of plasma discharges, particular to plasma etching and deposition, high aspect ratios and uniformity over large areas are desirable. Operational pressure of such discharges plays an important role in achieving these goals. High pressure processes are increasing the ion scattering probability through collisions in front of the substrate region. This will result in a poor etching anisotropy. A comparative view between anisotropic and isotropic etching is shown in the figure 3.1. In [1] a detailed overview on the plasma etching is presented.

Figure 3.1: Isotropic and anisotropic etching profiles
3.2 Experimental setup.

The investigations carried out in this work have been performed on a newly designed Neutral Loop Discharge experimental setup. A schematic view of the design used in other scientific investigations and industrial research applications is shown in Figure 3.2. Due to the special magnetic field configuration produced by the three coaxial electromagnetic coils the Neutral Loop is located in the middle coil plane. This makes the investigations in the NL region challenging. Moreover optical investigations such as Thomson scattering, Optical Emission Spectroscopy, Radio-Frequency Modulated Optical Emission Spectroscopy (RF - MOS) or Phase Resolved Optical Emission Spectroscopy are almost impossible to be carried out in this configuration.

![Figure 3.2: NLD Experimental setup as described in Reference xx](image)

The newly designed NLD reactor, follows the concept presented above, previous successfully applied in industrial applications. The NLD plasma discharge typically operates at pressured bellow 0.5 Pa with relatively high electron densities and sufficiently low electron temperatures. As outlined in Chapter 2 earlier investigations carried out in the NLD where mainly focused to improve its industrial applications. Only few
investigations where aimed on the fundamental understanding of the discharge sustainability at these low pressures and in particular on the electron heating mechanisms. In order to further improve and understand this discharge, knowledge of the heating mechanisms is of high interest. Therefore spatially and temporarily investigations in the NL plane of the Electron Energy Distribution Function, electron density and temperature are required. The novel NLD design shown in Figure 3.3 uses a planar type antenna design and a special configuration of the electromagnetic field coils.

Figure 3.3: The Bochum NLD experimental setup
The stainless steel setup consists of two main chambers. The plasma chamber located above the last coil, with a diameter of xx cm and the processing chamber located below the bottom coil with a diameter of xx cm. Both chambers allow radial investigations of the plasma properties from built-in view-ports.

Three electromagnetic coils are producing the same structure of the magnetic field as in the original NLD setup. In this configuration the respective coils are of different radii. This achieves the shift of the NL plane from the plane of the middle coil, and therefore an easier diagnostic access in the NL region. To attain the desired magnetic field configuration the currents in the top and bottom coils are flowing in the same direction whereas in the middle coil it flows in the opposite direction. This arrangement produces in the NL region a quadrupol magnetic field bent into a torus. The simulation of the magnetic field lines is shown in Figure 3.4.

![Figure 3.4: Simulated magnetic field lines](image)

The 0 point on the x-axis represents the chamber center. The z-axis runs from the quartz located at -24 mm up to the bottom of the discharge chamber. The 0th position on the z-axis is represented by the NL position. In the downstream region, the processing
chamber, the magnetic field is intersecting at a perpendicular angle a possible position of a wafer substrate. This will improve the surface etching quality, since ions, guided by the magnetic field are coming at a perpendicular angle. Calculations of the magnetic field for different coil radii and positions where performed to achieve an optimum position of the NL plane [46]. A simulated magnetic field strength and the position of the Neutral Loop, courtesy of Timo Gans for a particular coil configuration are shown in Figure 3.5. The gray circles are marking the position on the z-axis of the coils. The black line locates the plane of the middle coil. For this simulation the Neutral Loop is located at approximately 2 centimeters bellow the quartz (not shown in this plot) and at a radius of 6 centimeters.

Figure 3.5: Calculated magnetic field for coils powered at: top: 141 A, middle: 144 A, bottom: 117 A. The scale is in mT.

In the above plot the position and the structure of the Neutral Loop are clearly shown. The simulation computes as well the magnetic field lines structure. The arrows are pointing in the direction of the magnetic field vector. For later applications the direction in which the current is flowing in the three coils can be reversed. Figures 3.6 and 3.7 are showing the corresponding magnetic field for the 2 cases.
Figure 3.6: Magnetic field simulation in “normal” coupling of the DC current to the coils. The simulation is made with the coil configuration 141 – 114 – 117. Courtesy of Timo Gans

Figure 3.7: Magnetic field simulation in “reverse” coupling of the DC current to the coils. The simulation is made with the coil configuration 141 – 114 – 117. Courtesy of Timo Gans
In order to access the NL region and the downstream region the experimental setup is equipped with several lateral view-ports and a bottom flange as shown in Figure 3.3. The bottom flange can be easy retrofitted to a window from which the plasma can be visually monitored or bottom optical measurements can be performed.

The plasma is generated by applying an oscillating radio-frequency electric field along the Neutral Loop through a planar 4 fold 2 ½ turn ICP antenna. The antenna used in this work has a special design for minimizing parasitic capacitive coupling and to achieve an improved spatial uniform electric field along its diameter. This design is described extensively in [8]. A schematic view of the antenna and the calculated electric field from Kadetov’s Ph.D. work [8] in vacuum is shown in Figure 3.8.

![Figure 3.8: (Left) The calculated induced electric field by the antenna, (Right) The antenna design](image)

The antenna is operated at 13.56 MHz by a Dresler RF generator and separated from the plasma by a quartz dome with a width of 2 centimeters. The maximum outputted power is 2000 Watt. A matching network assures the optimal power transfer to the plasma by matching the variable plasma impedance to the impedance of the RF generator (50 Ohm).

The versatility of the experimental setup allows the plasma ignition in multiple modes, NLD with all the three coils powered, ICP without any external magnetic field and as a
planar Helicon type discharge with only the bottom coil powered. Photos of the plasma discharge are shown below for the NLD and ICP operational modes. The Helicon is widely treated in Chapter 6 of this work.

Figure 3.9: The NLD plasma discharge at coils powered at 141 – 144 – 117

Figure 3.10: The NLD plasma discharge at coils powered at 141 – 132 – 117
The photos presented in figures 3.9, 3.10, 3.11 and 3.12 were taken in an NLD discharge in Argon at 0.1 Pa and 1000 Watt applied RF power.

Figure 3.11: The NLD plasma discharge at coils powered at 141 – 126 – 117

Figure 3.12: The NLD plasma discharge at coils powered at 141 – 108 – 117
With the magnetic field turned off an ICP discharge can operate in the chamber down to a pressure of 0.3 Pa. For discharge stability and convenience reasons the lowest ICP pressure for measurements was 0.5 Pa.

Figure 3.13: The ICP plasma discharge in Argon at 0.5 Pa and 1000 Watt
3.3 Applied plasma diagnostic techniques

A wide range of diagnostic techniques are available to investigate plasma discharges. In this work both intrusive: Langmuir probe, B-dot probe and non-intrusive: Thomson scattering, Radio-Frequency Modulation Optical Spectroscopy (RF-MOS) where applied. Additionally a Hall probe system was used to map the magnetic field configuration and to prove the validity of the simulated filed configuration. A short description of each of the applied diagnostic techniques together with measured examples is given.

3.3.1 Mapping of the magnetic field with a Hall probe

The inhomogeneous magnetic field structure and strength was mapped with the help of a Hall probe [47]. This mapping is very important and in particular in the Neutral Loop region since there the calculated magnetic field strength comes from differences of big numbers. The results for two selected field configurations are shown in Figures 3.14 and 3.15.

![Figure 3.14: Measurement of the magnetic flux density and determined magnetic field lines with the coils powered at 141 – 144 – 117](image)
Figure 3.15: Measurement of the magnetic flux density and determined magnetic field lines with the coils powered at 141 – 132 – 117

The blue dot in the right hand side plots is marking the measured position of the NL. The R-axis has as 0th position the discharge chamber symmetry axis.
In the numerical calculation results shown in section 2.2 are not taken into account the finite dimensions of the coils. A small change in the position of the coils can produce a significant shift in the position of the null point. For this reason the coils on the experimental setup are fixed and separators are keeping the coils on the same position with respect to the wall.

The probe, due to its length and width, was able to access a region of 14 centimeters in the Z-direction and 12 centimeters in the radial, Y-direction. The probe measures directly the magnetic flux $B$, which has Tesla as a base unit. The magnetic field in vacuum is then given by:

$$H = \frac{B}{\mu_0}$$  \hspace{1cm} (3.1)

Commonly the term magnetic field is used rather than magnetic flux; therefore this will be used along this work too. Unless plasma currents are contributing to the magnetic field, there is no difference, up to a constant factor. The magnetic field lines can be easily
determined from the measured magnetic flux density considering a cylindrical coil configuration. For this, in the general case one can write:

\[ \vec{B} = \nabla \times \vec{A} = \frac{1}{r} \left[ \frac{\partial A_z}{\partial \phi} - \frac{\partial}{\partial z} (rA_\phi) \right] \cdot \vec{e}_r + \left[ \frac{\partial A_r}{\partial z} - \frac{\partial A_z}{\partial r} \right] \cdot \vec{e}_\phi + \frac{1}{r} \left[ \frac{\partial}{\partial r} (rA_\phi) - \frac{\partial A_z}{\partial \phi} \right] \cdot \vec{e}_z \]  

(3.2)

Since for the given configuration \( A_z = A_r = 0 \), equation (2.2.) can be rewritten to:

\[ r \vec{B} = \left[ \frac{\partial}{\partial z} (rA_\phi) \right] \vec{e}_r + \left[ \frac{\partial}{\partial r} (rA_\phi) \right] \vec{e}_\phi \leftrightarrow r \left( \vec{e}_\phi \times \vec{B} \right) = \nabla (rA_\phi) \]  

(3.3)

By integrating the equation 3.3 along any arbitrary path with \( \phi = \text{const} \) one determines the field lines.

The Hall probe system used is installed on a home-made computer controlled Z-Y stage. The error in the exact determination of the NL position is of about \( \pm 2 \text{mm} \) due to the step-motors moving accuracy. Several measurements were made in order to prove their reproducibility and accuracy. The probe is previously calibrated against a zero magnetic field head. The measured magnetic field strength is almost zero in the NL region up to a maximum of about 14 mTesla in the lower region of the chamber. These values agree well with the simulated ones. The location of the NL is determined within the expected error margins.

### 3.3.2 Langmuir probe

The Langmuir probe is one of the oldest and most used diagnostic methods for investigating plasma discharges [47, 48, 49, 50, 51, 52]. This method is relatively simple both in construction and in its application, and allows measurements of the electron temperature, density and respectively of the Electron Energy Distribution Function. The method itself can be shortly summarized as follows. A voltage biased wire is inserted into the plasma and the respective current drawn from the plasma through the probe is recorded. This will result in the (I-V) characteristic curve. This is used later in determination of the plasma parameters. Although its simplicity in nature, a reliable interpretation of the measured data becomes quite complicated and several theories have been developed. Moreover probe measurements in magnetically enhanced plasma discharges are subject of more careful interpretation due to the magnetic field influence.
on the current drawn by the probe. Extensive descriptions on the theories are given in () and here only the basic assumptions are described.

We introduce here a fundamental parameter for characterizing plasma discharges, the so called Debye length. Although the plasma is considered electrically quasi-neutral, however on small scales electrical interactions between charged particles exists. To maintain the quasi-neutrality each charged particle tends to screen out the electrical influence of its neighboring particles. The distance on which this charge separation occurs is named Debye length and is defined by:

\[
\lambda_D = \sqrt{\frac{\epsilon_0 k_B T_e}{n_e e^2}} = 743 \frac{T_e}{n_e} (\text{cm})
\]

(3.4)

Given that the NLD Discharge makes use of an additional magnetic field, special care was taken to investigate this effect on the probe measured data. As mentioned, the presence of the magnetic field can influence the current flux drawn by the probe. In cylindrical probe geometry Laframboise in [53] theoretically showed that this influence is related to the applied magnetic field strength and the probe radius and it can be quantified by using the following relation:

\[
\beta = \frac{r_{\text{probe}}}{r_{\text{Larmor}}}
\]

(3.5)

where

\[
r_{\text{Larmor}} = \frac{1}{eB} \sqrt{\frac{mn_e k T_e}{2}}
\]

(3.6)

Therefore the lower the magnetic field strength the smaller the influence on the current drawn by the probe. However this theory also shows that the magnetic field influence is at its minimum when the probe is inserted perpendicular to the magnetic field lines.

In this work at the measurement plane, the magnetic field strength varied between 0.1 mTesla up to 12 mTesla. For this strength the electron Larmor gyro-radius is in the order of \(10^{-6}\) meters, much smaller than the radius of the probe, resulting in a negligible influence on the measured data.
The probe used for the measurements is a commercial out-of-the-box Scientific System Smart-Probe. The probe is equipped with a computer controlled Linear Drive unit which allows the probe to move in a range of 300 mm in steps of 10 mm and a boxcar. For the given reactor configuration shown in Figure 2.3 this allows scanning of the entire NL region, when the NL diameter is smaller than 8 cm. Spatially resolved measurements where performed both in the plane of the NL and in the downstream region. Moreover given the magnetic field configuration the probe tip sees the magnetic field lines at a relative perpendicular angle, which in the view of [53] minimizes the influence of the magnetic field on the acquired data. The probe was inserted through a flange system at a distance of 24 mm bellow the quartz.

The probe tip is made of Tungsten and it has a diameter of 0.5 mm and a length of 50mm. Since measurements up to powers of 1000 Watt where performed at a relative close distance from the quartz, coating of the probe tip can occur. The Scientific System probe software allows automatic cleaning through measurement cycles of the probe tip. In radio-frequency generated plasma higher RF harmonics can disturb the probe measurements. To compensate for this, a passive built-in RF compensation located inside the ceramic tube close to the tip is used. The floating potential is measured by an external electrode attached to the probe ceramic tube. A schematic view of the Langmuir probe system is shown in Figure 3.16.

![Schematics of the Langmuir probe used for the measurements in this work.](image)

The voltage bias applied range was between -25V and 25V. The measurements where performed by averaging over 3 consecutively acquired data points in a sweep of 100 steps.
A typical I-V curve and the respective EEPF plot are shown in Figure 3.17 and 3.18 respectively for a NLD discharge operated at 0.1 Pa and 1000 Watt. The probe software allows determination of the plasma parameters either directly from the I-V characteristic or through the Electron Energy Probability Function (EEPF).

Figure 3.17: The Smart – Probe Langmuir system acquired I-V curve in NLD at 1000 Watt, 0.1 Pa. Coils powered at 141 – 144 – 117 A.

In the case of a Maxwellian distribution the representation in a semi-logarithmic scale of the EEPF as a function of energy is represented as a straight line.
Figure 3.18: The Smart – Probe Langmuir system acquired EEPF in NLD at 1000 Watt, 0.1 Pa.

The measured electron density and temperature shown in Figure 3.19 are calculated by integrating the measured EEPF. For most of the plasma discharges a departure from the Maxwellian distribution is present. For this case the so called Druyvesteyn distribution is used for a better approximation. Therefore the electron density is expressed as:

\[ n_e = \int_0^\infty g(\varepsilon)d\varepsilon \]  

(3.7)

and respectively for the electron mean energy:

\[ <\varepsilon> = \frac{\int_0^\infty \varepsilon \cdot g(\varepsilon)d\varepsilon}{n_e} \]  

(3.8)

The effective electron temperature is thus given by:

\[ T_{\text{eff}} = \frac{2}{3} <\varepsilon> \]  

(2.9)
A radial scan in the NL plane is shown in Figure 3.19. The plasma was operated in Argon at 0.1Pa and 1000Watt. The coil coils where powered at 141 – 137 – 117 A.

Figure 3.19: Radially resolved Langmuir probe measurements of the electron temperature and density in the NL plane.

The electron temperature and density values are together overlaid on the dimensions of the setup elements. The plot marks the position of the antenna, the quartz dome and the maximum induced electric field. A wider analysis on the plasma properties obtained by Langmuir probe measurements is presented in the forth chapter.
3.3.3 B-dot probe

The time varying RF magnetic field can be measured by an induction probe [54], from the induced current in a conducting coil shaped as a loop. The output voltage from the coil is thus proportional to the area, the number of turns in the coil and time characteristic of the magnetic field. The RF voltage induced into the probe is given by the Lenz law:

\[ V = -\frac{d\Phi}{dt} \]  

(3.10)

where \( \Phi \) is the RF magnetic flux through the loop. Integration of the B-Dot probe output voltage gives the time-averaged magnetic field values. At a given RF frequency the magnetic field amplitude directed perpendicular to the plane of the loop is given by:

\[ B = \frac{U_{\text{probe}}}{\omega N_{\text{loop}} A_{\text{loop}}} \]  

(3.11)

where \( \omega \) is the RF frequency, \( N_{\text{loop}} \) are the number of turns and \( A_{\text{loop}} \) is the area of the loop. From the B-dot probe measurements the electric field and the current density can be determined.

For the measurements made in this work the probe consisted from a two turn loop made of a stainless steel wire with a diameter of 10 mm mounted on a ceramic tube. The inside connectors are shielded by an inner shielding mounted inside the ceramic tube. Both extremities of the tube have been vacuum sealed with TorrSeal vacuum glue. A 50 \( \Omega \) load terminator is used at the oscilloscope to measure the probe signal. The probe can be rotated inside the plasma with any arbitrary angles, therefore being able to measure both \( E_p \) and \( E_r \) components of the induced electric field. A calibration of the probe against a known magnetic field from a Helmholtz coil was made to assure an accurate measurement and is shown in Figure 3.20. By inserting the probe into the plasma a variation in reflected power was observed up to of a maximum of 5 Watt at a maximum applied power of 1000 Watt. Again here the measurements are limited by the probe overheating.
Figure 3.20: B-dot magnetic field calibrated against Helmholtz coil.

Figure 3.21: B<sub>z</sub> and B<sub>θ</sub> for an Helicon discharge in Argon at 0.5 Pa, 1000 Watt. Bottom coil is powered with 59 Amperes.
Figure 3.22: $E_\phi$ component for an NLD discharge at 0.1 Pa, 1000 Watt. The small, medium and large coil configurations are respectively: 141 – 120 – 117 / 141 – 137 – 117 / 141 – 150 – 117 Amperes.
3.3.4 Thomson scattering

Non-invasive determination of electron energy and velocity distributions in low-temperature, non-thermal plasmas has always been a great challenge for diagnostics. Initially Thomson scattering was successfully applied in the fusion research experiments. The basic principles of Thomson scattering are well established in the literature [55, 56, 57, 58, 59]. In recent years Thomson scattering had proved to be a very versatile technique and it has been applied to low pressure as well as atmospheric pressure discharges [60, 61, 62, 63, 64, 65, 66].

Compared to other methods its advantages are clear:
- High temporal resolution down to sub-ns [67].
- Direct access to the Electron Velocity Distribution Function – easy access to the electron temperature.
- Electron density determination – intensity of the scattered spectra is in a direct proportional relation
- No direct interference with the investigated plasma discharge

The main physical obstacles are low electron densities combined with a small cross-section for Thomson scattering, the high amount of laser stray light coming from optics and reflection on apertures and walls, Rayleigh scattering at higher neutral densities, and background emission from the plasma. Moreover under certain conditions ionization from metastable states, multi-photon ionization from ground states or detachment of negative ions can be a problem. Major progress on the suppression of stray light and Rayleigh scattering was made by use of Triple Grating Spectrometers in combination with ICCD cameras.

The basics of this method lie in the interaction of the plasma electrons, i.e. scattering, with an electric field of a laser beam. In this interaction the electrons are “seeing” a shorter or a larger wavelength depending on their movement towards or away from the source. These accelerated electrons will start to emit themselves an electromagnetic radiation with a frequency higher or lower with respect to the laser one.
A schematic of a typical scattering process is shown in Figure 3.23.

![Figure 3.23: Typical 90° experimental Thomson scattering geometry](image)

In essence electrons are experiencing a double Doppler-shift effect, one with respect of the incoming light source and second with the respect of the observer. Depending on the particle velocity direction along the scattering vector $\vec{k}$ the two Doppler shifts can add or cancel each other:

$$k = k_s - k_i \quad (3.12)$$

where $k_s$ represents the vector along the observation direction and $k_i$ is the vector in the direction of the incident radiation. In terms of frequency shift this can be written as:

$$\omega_s - \omega_i = \vec{k} \cdot \vec{v} \quad (3.13)$$

The scattering vector $\vec{k}$ can be therefore expressed as:

$$|k| = 2|k_i| \sin \frac{\theta}{2} \quad (3.14)$$

Following from Figure 3.23 one can note that the maximum scattered light is achieved for electrons with velocities parallel to $\vec{k}$ and minimum when the electrons have velocities perpendicular to $\vec{k}$. Moreover the scattering intensity is influenced by the angle between the incident and the scattered radiation and has a maximum for $\theta = 90^\circ$. 

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3.3.4.1 The Salpeter approximation

The above relative simple approach to Thomson scattering, would not have been possible without the work of Salpeter [68]. An in-depth analysis of his work is shown in [55, 56, 57]. Here only a summary is given.

By solving the Vlasov equation for electron density fluctuations in plasmas, Salpeter in [68] made an important approximation in its calculations, namely the separation of the ion and the electron contributions to the scattered spectra. The result of the calculations made in () is the so-called dynamic form factor $S(k,\omega)$ which in end-effect determines the shape of the scattered spectrum and is expressed for a Maxwellian distribution function as:

$$S(k,\omega)d\omega = \frac{1}{1+\alpha^2 W(x_e)}\frac{\exp(-x_e^2)}{\pi^{1/2}}dx_e + Z\left(\frac{\alpha^2}{1+\alpha^2}\right)^2 \frac{1}{1+\beta W(x_i)}\frac{\exp(-x_i^2)}{\pi^{1/2}}dx_i \quad (3.15)$$

where:

$$x_e = \frac{\omega}{kv_e} \quad \text{and} \quad x_i = \frac{\omega}{kv_i} \quad (3.16)$$

$$\beta^2 = \frac{Z\alpha^2}{1+\alpha^2} \frac{T_e}{T_i} \quad (3.17)$$

$$W(x) = 1 - 2xe^{-x^2} \int_0^x e^{p^2}dp - i\pi^{1/2}xe^{-x^2} \text{is the plasma dispersion function}$$

Integrating equation 3.15 over $\omega$, where the first term represents the electron part and the second term the ion part leads to:

$$S_e(k) = \frac{1}{1+\alpha^2} \quad (3.18)$$

$$S_i(k) = \frac{Z\alpha^2}{(1+\alpha^2)^2} \left(\frac{1}{1+\beta^2}\right) \quad (3.19)$$

The $S(k)$ is expressed then as a summation of both terms:

$$S(k) = S_e(k) + S_i(k)$$

38
The above equations show that the electron part plays the major contribution to the scattered spectra when $\alpha << 1$ whereas in the opposite case the ions get the major contribution for $\alpha > 1$. This separation leads to two different scattering regimes treated next.

### 3.3.4.2 Scattering regimes

Following the previous analysis in laser Thomson scattering, one can distinguish between two main regimes described by the parameter $\alpha$, expressed as:

$$\alpha = \frac{1}{k\lambda_D} = \frac{1}{4\pi \cdot \sin(\theta/2) \lambda_D}$$

(3.20)

Additionally is introduced here the “sample length” distance on which collective effects are probed defined by $2\pi/k$. Sheffield in [57] showed that in general scattering takes place on electron density fluctuations of a $2\pi/k$ wavelength. Within a Debye length scale the electrons do not show any collective behavior. Figure 3.24 shows the difference between the two.

![Figure 3.24: Incoherent (right) and coherent (left) scattering regimes](image)

a) For $\alpha << 1$ Thomson scattering is said to be in the incoherent regime, scattering takes place on a sample length factor smaller than the Debye length. Therefore this will represent the velocity distribution of the free electrons in the plasma, which do not bare any phase relation between them. In this case for a Maxwellian distribution function the scattered light spectrum has a Gaussian shape. However if the distribution function is not
Maxwellian the scattered spectra will still reflect the velocity distribution function when the $\alpha < << 1$ condition is satisfied.

The scattered power in the solid angle as a function of frequency is given by:

$$\frac{dP_s}{d\omega_s} = P_i n_e L_{det} \frac{d\sigma_{\text{thom non}}}{d\Omega} \Delta\Omega \cdot S_k(\Delta\omega) d\omega_s$$

(3.21)

The $S_k$ function can be seen as the probability of a frequency shift of a scattered photon in the $d\omega$ range. Since a frequency shift occurs when the electron has a velocity component along $k$, this means that the spectral distribution function directly reflects the one-dimensional electron velocity distribution function.

$$S_k(\Delta\omega) d\omega = F_k(v_k) dv_k$$

(3.22)

For a Maxwellian distribution function $F_k(v_k)$ becomes:

$$F_k(v_k) = \frac{1}{v} \exp \left[ \left( -\frac{v_k}{v} \right)^2 \right] \text{ with } v = \left( \frac{2k_BT_e}{m_e} \right)^2$$

(3.23)

Rewriting equation 3.21 in terms of wavelength and taking into consideration equation 3.23 one can directly access the electron temperature as:

$$T_e = \frac{m_e c^2}{8k_B \sin^2(\theta/2) \left( \frac{\Delta\lambda_{1/e}}{\lambda_i} \right)^2}$$

(3.24)

where $\Delta\lambda_{1/e}$ is the HWHM of the scattered spectra.

Moreover by knowing the EVDF one can infer the EEDF simple by using the velocity to energy conversion: $E = \frac{1}{2} m_e v^2$. This can be done however only in the case of an isotropic distribution function.

b) For $\alpha > > 1$ Thomson scattering is said to be in the coherent regime. In this case scattering takes place on a sample length factor larger than the Debye length and collective behavior is shown in the scattered spectra.
3.3.4.3 The experimental setup for Thomson scattering measurements

Due to its relative small scattering cross-section, of the order of $10^{-25}$ cm$^2$ and the relative low electron densities in the RF plasma discharges, the laser source should be sufficiently strong enough to produce a detectable signal. However an increase in the laser power and intensity will produce a proportional increase in the stray light intensity. Moreover the laser should not “artificially” increase the electron temperature at the measuring location. Therefore much care must be taken to reject the stray light by using baffle systems and a proper choice of the used spectrograph and of the choice of the laser system.

Here we imply the usage of a Triple Grating Spectrograph (TGS) which by now is the state-of-the art system in stray light rejection. A typical TGS configuration is shown in Figure 3.25. This combined with an ICCD camera; Roper Scientific PI-MAX allows high spatial and temporal resolution measurement on the plasma characteristics [59, 64, 67].

![Figure 3.25: A typical Triple grating spectrograph schematics.](image)

The system used in this work for Thomson scattering measurements is shown in Figure 3.26.
Figure 3.26: The Thomson scattering experimental setup
A Continuum 9050 frequency doubled Nd:YAG laser with a wavelength of 532 nm, pulse length of FWHM = 8 ns and an energy of $E = 600 \text{ mJ}$ and a repetition rate of 50 Hz was used for the measurements presented in this work. The laser beam is guided by a pair of dielectric mirrors and focused by a plan-convex lens ($f = 1 \text{ meter}$) to a spot size of about 250 $\mu\text{m}$. The focus is 2 cm below the quartz surface and a radial distance of 6 cm from the chamber center, where the induced electric field is expected to have a maximum in vacuum. The polarization of the laser is perpendicular to the scattering plane. Stray light is reduced by Brewster windows and a set of apertures. Scattered photons from the focal region are imaged perpendicular to the laser beam by a pair of achromatic lenses with a focal length of $f = 500 \text{ mm}$ and a clear aperture of $\Phi = 98 \text{ mm}$ as a 1:1 image onto the entrance slit (height = 11 mm, $d = 0.32 \text{ mm}$) of a homemade triple grating spectrometer as shown in figure 2.26. In the imaging beam path a set of three mirrors rotates both the plane of polarization and the image of the scattering volume. The image of the scattering volume is then parallel to the entrance slit and the polarization perpendicular to the grooves of the gratings. A viewing dump located opposite to the observation window minimizes stray light detection. The laser beam dump is located about 2 m away in order to delay back-reflected light sufficiently with respect to the signal. The general design on the triple grating spectrometer is similar to other devices described in the literature [59, 64, 67]. Three holographic gratings with $n = 1800$ grooves $\text{mm}^{-1}$ are used. The first two gratings are operating substractively. A notch filter in the form of a mask with a central metal stripe having a width of $d = 2 \text{ mm}$ is set in between in order to block the stray light at the laser wavelength. The deflection angles of the gratings are $\delta = 25^\circ$, the corresponding angle of incidence and reflection are $\alpha = 16.87^\circ$ and $\beta = 41.87^\circ$ respectively. The dispersion in the exit plane is 1.21 mm $\text{nm}^{-1}$. For an entrance slit of diameter $d = 320 \mu\text{m}$ the corresponding resolution is $\Delta\lambda_{bp} = 0.34 \text{ nm}$. The intermediary slit is set to $d = 500 \mu\text{m}$. Imaging inside the spectrometer is made by six achromatic lenses with $f = 500 \text{ mm}$ and a clear aperture of 98 mm. The solid angle of the setup is $\Omega = 0.09 \text{ sr}$ and the spectrograph f-number is 5.1. Detection is achieved with an intensified CCD camera with a GaAs photocathode (Roper Scientific PI MAX II) for maximum quantum efficiency. The camera gate is set to 20 ns with the signal centered.
The effective size of the detector is 12 mm, allowing detection over a total spectral interval of $\Delta \lambda = 10.7\,nm$. The Thomson spectra are acquired over 90,000 laser shots (30 min integration time) by accumulation on the CCD chip. Typically only few photons are counted for a single laser shot. Background emission from the plasma is recorded over identical intervals with the laser off and subtracted subsequently. The laser is monitored with a fast photodiode for synchronization with the CCD camera gate and normalization for absolute density measurements. A homemade electronic synchronizes the laser pulses with the RF phase by locking the trigger on the monitor output of the RF generator.

The acquired Thomson spectrum for an NLD discharge operated in Argon at 0.5 Pa and 1000 Watt is shown in Figure 3.27.

Figure 3.27: The ICCD acquired Thomson scattering spectrum

Figure 3.27 clearly shows the center part of the spectrum (of about 100 pixels) being block by a mask with a width of 2 mm.
Apart from a straightforward determination of the electron temperature, Thomson scattering can be used for electron density measurements, the scattered spectrum intensity being proportional with the electron density. In practice, the calibration of the setup for absolute particle density measurements can be done by using either Rayleigh [55] or Raman scattering [59, 69, 70] on a known gas. Additionally, accurate knowledge of the scattering cross sections for both methods is required to calculate the correspondent calibration factors. In this work, the absolute particle density calibration method used is based on the one described by van de Sande in [59]. For Thomson scattering measurements, a Triple Grating Spectrograph (TGS) was used as shown in Figure 2.26. The spectrograph has a design an intermediary mask which blocks the Rayleigh signal. Therefore, usage of the Rayleigh signal to calibrate the absolute density measurements will require a change in the experimental setup configuration, consequently altering the system characteristic which is not desired. This limits the choice only to Raman scattering calibration. The scattering geometry is the same as for the Thomson scattering. The ICCD Raman spectrum is shown in Figure 3.28.

![The ICCD acquired Raman scattering spectrum](image.png)

Figure 3.28: The ICCD acquired Raman scattering spectrum
Nevertheless this method has its advantages since no modification of the setup is needed and the scattered Raman wavelengths are away from the region blocked by the mask. Moreover due to the high stray light rejection of the TGS an additional separation of the Raman spectrum from the stray light is not required. For this work the chamber was filled with Nitrogen at a pressure of 8000 Pa. The gas was considered to be at room temperature, 300 K.

Scattering is assumed to be incoherent and only rotational Raman scattering is considered. A high resolution Raman spectrum taken with the Triple Grating Spectrograph at an entrance slit of 50µm is shown in figure 3.29. This shows the Raman transitions observable in the wavelength span of the TGS. This spectrum is used later on calculation of the Raman scattering calibration factors.

![High resolution Raman spectrum](image)

Figure 3.29: High resolution Raman spectrum corresponding to Figure 2.28.

The absolute electron density [59] can be expressed from the Thomson scattering intensity with respect to the intensity of a rotational Raman spectrum by:

\[
n_e = n_{N_0} \cdot \frac{I_{\text{Thom,son}}}{I_{\text{Raman}}} \left( \frac{1}{d\sigma_{\text{Thom,son}}/d\Omega} \sum_{J} \frac{n_{J} \cdot d\sigma_{J,J'}}{d\Omega} \right) = n_{N_0} \cdot \frac{I_{\text{Thom,son}}}{I_{\text{Raman}}} \cdot \Gamma_{\text{Raman}} \tag{2.25}
\]
where $I_{\text{Thom-non}}$ and $I_{\text{Raman}}$ are the respective measured intensities, $\frac{d\sigma_{\text{Thom-non}}}{d\Omega}$ is the differential Thomson scattering cross section and $\Gamma_{\text{Raman}}$ represents the ratio of the total Raman scattering cross section and Thomson cross section.

To calculate $\Gamma_{\text{Raman}}$ the total Raman scattering cross section is required. This is expressed by the summation over all the observed Raman transitions, with their previously calculated cross-sections calibration factors. This calculation is taking into consideration the relative population of the different rotational states $J$. For the given conditions, at the room temperature $\Gamma_{\text{Raman}}$ is of about 567.
3.3.3.3 Radio Frequency Modulated Optical Emission Spectroscopy

In RF generated discharges the oscillating radio-frequency (RF) electric fields penetrating into plasma lead to a corresponding oscillation of the electron velocity distribution function. In parallel to modulating the anisotropic part of the velocity distribution function, the electric field also causes a modulation of the isotropic part or equivalently of the energy distribution function. In terms of a perturbation series this can be considered as a second order effect while the drift is a first order effect. Although the anisotropic drift does not act on the excitation rate of atoms, the modulation of the isotropic part leads directly to a corresponding modulation of the excitation rate. As a consequence, also the emission is modulated, although the finite lifetime of the excited state might cause a certain damping of the modulation in emission. The amplitude of this modulation can be directly measured from the phase resolved emission of the highly excited states. On this grounds a novel non – intrusive diagnostic technique, Radio Frequency Modulation Optical Emission Spectroscopy (RF - MOS), based on previous reports in [89, 90] has been developed. This method provides a direct determination of the electron oscillation velocity.

The basics of this method are described in the following. In the plasma bulk, the electron behavior can be describe by solving the Boltzmann for the given plasma conditions. For the general case the Boltzmann equation is expressed as:

\[
\frac{\partial f}{\partial t} + v \nabla f - \frac{e}{m} \left( E + v \times B \right) \frac{\partial f}{\partial v} = \frac{\partial f}{\partial t} |_{r} \tag{3.26}
\]

However the Boltzmann equation is a quite complex equation and it does not solve analytically quite easy. A method of solving the Boltzmann equations is the so called “two - term” approximation. In this approximation the velocity distribution function consists of an isotropic part \( f_0(v,t) \) and a smaller anisotropic contribution \( f_1(v,t) \). Furthermore the isotropic part can be further divided into a time independent and a time dependent part.

The following approximations and assumptions are made:

- no external magnetic field present \( \vec{B} = 0 \)
- \( \nabla f = 0 \)
- the electric field is modulated sinusoidally:

\[
\vec{E}(t) = \frac{E_0}{2} e^{i\omega t} \hat{e}_z + c.c. = \vec{E}_0 \cos(\omega t + \phi)
\]

(3.27)
- the distribution function is expressed as the contribution of all particle distribution functions

\[
f = \sum_{n=-\infty}^{\infty} f_n e^{i\omega t} + c.c.
\]

(3.28)
- \( f_{-n} = f_{n}^{*} \)

With these assumptions equation 3.26 becomes:

\[
in\omega f_n - \frac{e E_0}{2m} \frac{\partial}{\partial v_z} (f_{n-1} + f_{n+1}) = \frac{\partial f_n}{\partial t} |_{\epsilon}
\]

(3.29)

This equation can be treated for two cases: \( n = 0 \) and \( n \geq 1 \):

- \( n = 0 \) equation 3.29 becomes:

\[
- \frac{e E_0}{2m} \frac{\partial}{\partial v_z} 2 \text{Re}(f_1) = \frac{\partial f_0}{\partial t} |_{\epsilon}
\]

(3.30)

- \( n \geq 1 \) equation 3.29 becomes:

\[
in\omega f_n - \frac{\partial f_n}{\partial t} |_{\epsilon} = \frac{e E_0}{2m} \frac{\partial}{\partial v_z} (f_{n+1} + f_{n-1})
\]

(3.31)

Using the following normalization:

\[
\frac{v_i}{v} \rightarrow v_z, \quad \frac{\alpha}{t} \rightarrow \frac{e E_0}{2m \alpha \nu_{th}}, \quad \alpha \ll 1
\]

where \( \nu_{th} = \sqrt{\frac{2k_B T}{m}} \)

(3.32)

The assumption of \( \alpha \ll 1 \) is based on the fact that the high energetic states are excited by electrons in the tail of the distribution function where the modulation is particularly strong. We are interested only in the temporarily modulated emission from these states.

With 3.32 the equations 3.30 and 3.31 becomes:

- \( n = 0 \):

\[
- \frac{\partial f_0}{\partial t} |_{\epsilon} = \alpha \frac{\partial}{\partial v_z} 2 \text{Re}(f_1)
\]

(3.33)

- \( n \geq 1 \):

\[
in_{n} \frac{\partial f_n}{\partial t} |_{\epsilon} = \alpha \frac{\partial}{\partial v_z} (f_{n+1} + f_{n-1})
\]

(3.34)
However to solve further the equation set 3.33 and 3.34 several approximations are required.

In the first approximation we take into consideration that

\[ \lim_{n \to \infty} \left| \frac{f_{n+1}}{f_n} \right| = O(\alpha) \Rightarrow |f_{n+1}| \ll |f_n| \]

equations 3.33 and 3.34 are expressed as:

- \( n = 0: \quad -\frac{\partial f_0}{\partial t} \bigg|_{\alpha} = \alpha \frac{\partial}{\partial v_z} 2 \text{Re}(f_1) \) \hspace{1cm} (3.35)

- \( n \geq 1: \quad \inf_{\alpha} -\frac{\partial f_n}{\partial t} \bigg|_{\alpha} \approx \alpha \frac{\partial}{\partial v_z} (f_{n-1}) \) \hspace{1cm} (3.36)

2\text{nd}, since we are solving for the \( z \) direction it follows that \( v_z = v \cos \theta \).

Thus one can write:

\[ \frac{d}{dv_z} = \frac{1}{v} \frac{d}{d(\cos(\theta))} + \cos(\theta) \frac{d}{dv} ; \]

this is in end effect a change between the even and odd order of powers for \( \cos^k(v) \).

Therefore the isotropic part of \( f_i \) will be given by \( f_0 \).

The third approximation is done over the collision term for \( n \geq 1 \):

\[ \frac{\partial f_n}{\partial t} \bigg|_{\alpha} = \begin{cases} -vf_n, & \text{for } n \text{ odd} \\ 0, & \text{for } n \text{ even} \end{cases} \] \hspace{1cm} (3.37)

However this approximation influences only the odd terms of the sum. For the case \( n = 0 \) the above approximation is not valid. Now one can expand the distribution function up to the third term. The oscillatory part is given by \( f_2 \) and the respective \( f_0, f_1 \) and \( f_2 \) are given by:

- \( n = 0: \quad -\frac{\partial f_0}{\partial t} \bigg|_{\alpha} = \alpha \frac{\partial}{\partial v_z} 2 \text{Re}(f_1) = 2\alpha^2 \frac{\partial}{\partial v_z} \left( \frac{v}{1 + v^2} \frac{\partial}{\partial v_z} f_0 \right) \) \hspace{1cm} (3.38)

- \( n = 1: \quad (i + v)f_1 = \alpha \frac{\partial}{\partial v_z} f_0 \) \hspace{1cm} (3.39)

- \( n = 2: \quad 2if_2 = \alpha \frac{\partial}{\partial v_z} f_1 = \alpha^2 \frac{\partial}{\partial v_z} \left( \frac{1}{i + v} \frac{\partial}{\partial v_z} f_0 \right) \) \hspace{1cm} (3.40)

On the other hand the distribution function is said to be a sum of its isotropic and anisotropic parts, thus \( f_i = f_{i0} + f_{i1} \). Integrating the equations 3.38 – 3.40 over the solid angle for the \( n=1 \) and \( n=2 \) components one obtains:
\[ f_{\omega} = \frac{2\alpha}{\sqrt{1 + v^2}} \frac{\partial f_{00}}{\partial v} \cos \vartheta \sin(t + \varphi) \]  
(3.41)

\[ f_{2\omega} = -\frac{\alpha^2}{3v^2} \frac{\partial}{\partial v} \left( \frac{v^2}{\sqrt{1 + v^2}} \frac{\partial f_{00}}{\partial v} \cos(2t + \varphi) \right) \]  
(3.42)

Thus the oscillatory velocity is given by:

\[ \langle v_z \rangle = \tilde{u} = -\frac{2\alpha}{\sqrt{1 + v^2}} \sin(t + \varphi) \]  
(3.43)

Taking into consideration the previous normalization, the oscillatory velocity expressed in equation 3.43 becomes:

\[ \tilde{u} = -\frac{eE_0}{m\omega \left( 1 + \left( \frac{v}{\omega} \right)^2 \right)} \sin(\omega \tau + \varphi), \quad \varphi = \arctan \left( \frac{v}{\omega} \right) \]  
(3.44)

with the amplitude given by:

\[ |\tilde{u}| = \frac{eE_0}{m\omega \left( 1 + \left( \frac{v}{\omega} \right)^2 \right)} = \frac{2\alpha}{\sqrt{1 + v^2}} \]  
(3.45)

This approach is applied next in determination of the drift velocities from the optical modulation of the highly excited states. If the distribution function has an additional drift component \( \tilde{u} \) and with 3.32:

\[ f\left( (v - u)^2 \right) = f\left( v^2 \right) \left( 1 + \alpha^2 + 2 \cos \vartheta \right) = f\left( v^2 \right) + \frac{\partial f}{\partial \alpha} \bigg|_{\alpha=0} \alpha + \frac{1}{2} \frac{\partial^2 f}{\partial \alpha^2} \bigg|_{\alpha=0} \alpha^2 + \ldots \]  
(3.46)

For a Maxwellian distribution function, integrating 3.46 over the solid angle:

\[ \left\langle f\left( (v - u)^2 \right) \right\rangle_\Omega = f\left( v^2 \right) \left[ 1 + \left( \frac{2e}{3k_B T} - 1 \right) \alpha^2 \right] = f_0 + f_2 \]  
(3.47)

The result of equation 3.47 states that the distribution function with an additional drift component includes two parts, a modulated \( f_2 \) and a non-modulated \( f_0 \). The modulated emission from the plasma is recorded by an ICCD camera over several thousand RF cycles. The ICCD scans the 74ns RF cycle in steps of 2ns. The result is normalized over the RF period. Typical modulations are in range between 0.1% and 10%. The measured fluorescence signal is expressed as:
\[
\eta = \frac{I(t)}{\langle I \rangle_t} - 1 = \left( \frac{2E_{\text{excitation}}}{3kT_e} - 1 \right) \left( \frac{-\tilde{u}^2 \cos(2\omega t + 2\varphi - \phi_e) + \bar{u}\tilde{u} \sin(\omega t + \varphi - \phi_e)}{2\sqrt{1 + (2\omega \tau)^2}} \right) \frac{1}{\sqrt{1 + (\omega \tau)^2}} (3.48)
\]

Expanding equation 3.49 in a Fourier series one obtains for the first and the second coefficients:

\[
|A_1| = \left( \frac{2E_{\text{excitation}}}{3kT_e} - 1 \right) \left( \frac{2\bar{u}\tilde{u}}{v_{th} \sqrt{1 + (\omega \tau)^2}} \right) (3.50)
\]

\[
|A_2| = \left( \frac{2E_{\text{excitation}}}{3kT_e} - 1 \right) \left( \frac{-\tilde{u}^2}{2v_{th} \sqrt{1 + (2\omega \tau)^2}} \right) (5.51)
\]

Therefore the oscillatory and the drift velocity can be expresses as:

\[
\tilde{u} = \sqrt{\frac{2A_1 v_{th} \sqrt{1 + (2\omega \tau)^2}}{2E_{\text{excitation}} - 1}} (3.52)
\]

\[
\bar{u} = \frac{A_1 v_{th} \sqrt{1 + (\omega \tau)^2}}{2\bar{u} \frac{2E_{\text{excitation}}}{3kT_e} - 1} (3.53)
\]

Conversion software is used in this work to process the acquired images by a LaVision ICCD camera. The resulted output is the $A_1$ and the $A_2$ data. The raw plot is shown in Figure 3.30. A detailed investigation on the drift velocity as a function of plasma parameters is done in Chapter 5.
Figure 3.30: The raw first components of the Fourier transformation for a NLD discharge, 0.1 Pa in Argon at 1000 Watt.
The performed analysis over the raw data is shown in Figure 3.31 for the side and bottom view measurement of the drift velocity. The side view drift velocity plot is calibrated in the z-direction against the quartz position.

Figure 3.31: Side and Bottom view drift velocity in the NLD measured by RF-MOS
Chapter 4
Plasma source characterization

4.1 Introduction

As described in the previous chapter, several diagnostic techniques have been applied to investigate the plasma properties. This chapter discusses in details the measurements results. Most of the results presented are taken from previously published work [81, 82] from which this chapter is based. Langmuir probe measurements carried out in different gases, Argon, Xeon are shown in the first part of this chapter, both in the Neutral Loop Discharge and in an Inductively Coupled Discharge. In the NLD spatially resolved measurements were made in the Neutral Loop plane and in the application plane. A comparison between the two is presented. Moreover the differences between the NLD and ICP discharges are outlined.

This chapter follows with Thomson scattering measurements done in the same conditions as the Langmuir probe measurements. However due to the Thomson scattering experimental setup limitations, measurements could be performed only locally, i.e. at the NL positioned at 6 cm from the chamber center and 2 cm bellow the quartz. However the NL radius can be modified by changing the current in the middle coil. This is allowing measurements “inside” and “outside” the NL region. Again a comparison with an ICP discharge is presented.

In the Marie Currie fellowship collaboration framework, Optical Emission Spectroscopy investigations on the NLD discharge have been made together with Prof. Dr. Shivarova’s group. Here only a succinct overview is given.

The chapter ends by discussing partially results of a novel diagnostic technique, Radio-Frequency Modulated Emission Spectroscopy (RF-MOS). RF-MOS together with Thomson scattering confirmed existence of an electron drift in the NLD. Chapter 5 of this thesis goes more in-depth on this aspect.
4.2 Langmuir probe measurements

4.2.1. Langmuir probe measurements at the Neutral Loop position.

The discharge dynamics and therefore the electron heating are strongly dependent on the plasma parameters. The influence of pressure and induced RF power on the electron densities and temperatures is investigated by Langmuir probe. Spatially resolved measurements have been performed both in the NL plane and in the downstream region where a substrate can be placed as shown in Figure 3.3. The results were compared with measurements on an ICP discharge. In this experimental set-up an ICP discharge could be operated in stable conditions down to a pressure of 0.5 Pa.

The measurements for both the Neutral Loop Discharge and the ICP were taken at the position of the NL point. In Chapter 3 a detailed overview on the measurement and the simulation of the Neutral Loop position was made. The gradient of the magnetic field, around the null point of the NL in both axial and radial direction is shown in Figure 4.1.

![Figure 4.1: Calculation of the static magnetic field in axial and radial direction around the NL. The standard values of the coils, i.e.141 – 144 – 117 A was used. Courtesy of Timo Gans.](image-url)
The magnetic field gradient was chosen based on the NLD industrial design. Additional care was taken to ensure that the NL confinement region was not touching the quartz dome separating the antenna from the plasma. The standard coil configuration setup will have the Neutral Loop positioned at a radius of 6 centimeters with respect to the chamber center and approximately 2 centimeters below the quartz. An error of approximately ¼ of a centimeter in the vertical plane is considered for the NL location given by the resolution of the Hall probe measurements. The pressure was varied between 0.05 Pa and 10 Pa. Due to the extreme heating of the Langmuir probe, reliable measurements could be taken up to 1000 Watt of applied RF power. This chapter refers to Chapter 2.3.2 for a short introduction of Langmuir probe measurements in plasma discharges and in particular to probe measurements in magnetic fields.

The measurement of the electron density in an Argon discharge at 1000 Watt as function of pressure for both NLD and ICP is shown in the figure 4.2. The measurements in the NLD presented in the following were made in the standard coil configuration, i.e. the coils powered at 141 – 144 – 117 Amperes.

![Figure 4.2: Electron densities in ICP and NLD as a function of pressure in an argon discharge at 1000 Watt.](image)

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In the region below 1 Pascal the electron density in the Neutral Loop Discharge is higher than in the ICP. This shows that the efficacy of the magnetic field confinement becomes effective at low pressures. Nevertheless in both cases the electron densities are decreasing with pressure decrease. The discharge is sustained through a collisionless heating mechanism, which becomes efficient in this low pressure regime. At higher pressures the neutral gas particle density is increasing and through ionization the electron density increases. In this regime the electron heating is mainly done through collisions. Measurement of the electron temperature dependence with pressure for the same conditions is shown in Figure 4.3. This shows the comparison of the electron temperature between the Neutral Loop Discharge and ICP. The possibility of sustaining the discharge at pressures below 0.5 Pa is also clearly shown in this plot.

![Figure 4.3: Electron temperatures in ICP and NLD as a function of pressure in an argon discharge at 1000 Watt.](image)

The electron temperature is strongly increasing with pressure decrease. An interesting aspect is that the measured electron temperatures in both the NLD and ICP are
showing similar values at all pressures where the discharge operates. Since the measurements are taken in the center of the NL region, the temperature there is expected to be highest. In the previous chapter in figures 3.9 – 3.13 the photographs of both discharges taken from the bottom view where shown. In the NLD case, relative large spatial variations can be observed compared to a quite flat profile for the ICP. It can be concluded that, for the ICP a relative flat profile for the electron temperature is expected, whereas for the NLD can exhibit relatively strong spatial variations. Therefore the average electron temperature in the NLD can be expected to be lower than in the ICP, when the temperatures measured at maximum position are comparable. The same behavior of the electron temperature and density measured by Langmuir probe is observed as well in [83] although the measurement power was lower, 300 Watt.

The ratio of the electron temperature and density in the Neutral Loop Discharge to Inductively Coupled Discharge is shown in figure 4.4.

Figure 4.4: Ratio of the NLD to ICP mode of electron temperatures and densities, as a function of pressure, in an Argon plasma at 1000 Watt.
The measurements conditions are similar to the ones in figure 4.3. The electron temperature ratio is almost constant for all pressures, whereas the electron density is increasing with pressure decrease. In the high pressure regime the density in both NLD and ICP mode is almost similar, with a ratio of almost one. However at lower pressures the density in the ICP mode is decreasing more rapidly compared to the NLD mode. Bellow 0.5 Pa the ICP discharge becomes unstable and does not operate. This shows an efficient confinement of the NLD in the low pressure regime.

The Electron Energy Distribution Function (EEDF) measured at the NL location in both NLD and ICP discharges show a Maxwellian distribution function of similar temperatures with a slight high electron density in the NLD mode and shown in Figure 4.5.

![EEDF graph](Image)

Figure 4.5: Comparison between the EEDF in the NLD and ICP for Argon, 0.5 Pa, 1000 Watt
The experimental behavior of the discharge parameters as function of pressure and power can be described and understood by using a simple global model in terms of power and particle balance. A uniform density discharge model is assumed, based on the one described in [14]. In the low pressure regime, the electrons mean free path can be comparable with the discharge dimensions. Consequently electrons are responsible for the ionization, excitation and dissociation processes. The rate of these processes is strongly dependent on how many electrons have enough energy to complete the process. The rate constants for ionization, excitation and momentum transfer are expressed by $k_{\text{iz}}$, $k_{\text{ex}}$ and $k_{\text{el}}$ and are dependent on the electron temperature.

In this model the electron temperature $T_e$ is determined from the particle balance, where the absorbed power $P_{\text{abs}}$ is equating the particle losses. The density profile is assumed to be uniform in the NL region. Moreover in this region probe measurements confirm a Maxwellian electron energy distribution function as shown in figure 4.5 for an NLD Argon discharge.

The balance equation is expressed as:

$$K_{\text{ionization}} \cdot n_{\text{gas}} \cdot n_0 \cdot \pi R^2 l = \left(2\pi R^2 h_i n_0 + 2\pi Rlh_i n_0\right) \cdot u_{\text{Bohm}}$$

(4.1)

where in RHS

$$\left(2\pi R^2 h_i n_0 + 2\pi Rlh_i n_0\right) = A_{\text{eff}}$$

(4.2)

is the effective area for particle loss and $u_{\text{Bohm}}$ is the Bohm velocity. In the LHS, $K_{\text{ionization}}$ is the ionization rate, $n_{\text{gas}}$ is the neutral gas density, $n_0 = n_e = n_i$ are the electron and ion densities, $R$ is the radius of the cylindrical plasma volume and $l$ is the height of the plasma volume. The production and loss of particle terms are both proportional with the plasma density, therefore $n_0$ cancels out. Follows that equation 4.1 can be expressed simply as:

$$\frac{K_{\text{ionization}}(T_e)}{u_{\text{Bohm}}(T_e)} = \frac{1}{n_{\text{gas}} \cdot d_{\text{eff}}}$$

(4.3)

where

$$d_{\text{eff}} = \frac{1}{2} \frac{Rl}{Rh_i + lh_i}$$

(4.4)
is the effective plasma size for particle loss. The dependence of $K_{\text{ionization}}$ and $u_{\text{Bohm}}$ of the electron temperature $T_e$ is assumed to be known. Solving the equation 4.3 for $T_e$ for the given discharge parameters $A_{\text{eff}} = 0.015 \, m^3$, $h_l = h_v = 0.03 \, m$, $P_{\text{abs}} = 1 \, kWatt$, $\epsilon_i = 15.68 \, eV$ ionization energy for Argon, $\sigma_{\text{Argon}} = 7 \cdot 10^{-21} \, m^2$ ionization cross-section for Argon, $M_{\text{Argon}} = 40 \, amu$, $n_{\text{gas}} = 2 \cdot 10^{18} \, m^{-3}$ neutral gas density at the given pressure and a $\frac{V}{A} = 0.05 \, m$. Rearranging in a more convenient form $T_e$ is expressed

$$kT_e = \frac{\epsilon_i}{\ln \left( \frac{N \frac{V}{A} \sigma_{\text{Argon}}}{M_{\text{Argon}} m_e} \cdot \frac{1}{\eta} \right)}$$

and plotted in the figure bellow.

![Figure 4.6: Global model kTe dependence with pressure](image)

Equation 4.5 shows that $T_e$ is determined only by the particle conservation and is independent on the plasma density and therefore on the discharge input power.
The electron density follows from the balance equation from the total power absorbed to the total power lost and solving for $n_0$ yields:

$$n_0 = \frac{P_{\text{abs}}}{A_{\text{eff}} \cdot u_{\text{Bohm}} \cdot e \cdot \varepsilon_T} \quad (4.6)$$

Expressing the corresponding terms $u_{\text{Bohm}}$ and $\varepsilon_T$ the electron density becomes for a specified $P_{\text{abs}}$ and $T_e$:

$$n_e = \frac{P_{\text{abs}}}{\left(k\varepsilon_i + \frac{5}{2}kT_e\right) \cdot \eta \cdot \frac{A}{\sqrt{M_{\text{Ar}}}} kT_e^{1/2}} \quad (4.7)$$

and plotted in figure 4.7:

![Figure 4.7: Global model electron density dependence with pressure](image)

The electron density is determined by the total power balance in the discharge and is affected by the $T_e$ behavior.
A power variation in the NLD mode reveals that the electrons temperature is independent of power whereas the electron density scales almost linearly as shown in figure 4.8.

Figure 4.8: Electron densities and temperatures as a function of rf power in the NLD, Argon, 0.1 Pa.

The gray line shows a linear fit of the electron density through zero. This plot shows power measurements only down to 700 Watt. At lower powers the discharge shows a capacitive coupling contribution, thus the scaling becomes different. This explains the strong deviations, not shown in this plot, of measured densities at lower power from the straight line. Electron temperature is governed by the particle balance in the plasma and is thus independent of power. However the electron density depends on the power balance and thus scales linearly with power. The range of powers on which accurate measurements using the Langmuir probe is limited. At lower powers the power coupling into the discharge is not purely inductive having an additional capacitive coupling present, this making the measurements without RF compensation of the probe inaccurate.
At higher powers the Langmuir probe system can not be cooled properly and therefore it cannot sustain the high temperatures. The plasma ionization degree and “electron pressure” determined from the Langmuir probe data as a function of neutral gas pressure in an Argon discharge at 1000 Watt is shown in figure 4.9.

![Figure 4.9: Degree of ionization and electron pressure, as a function of pressure, for an NLD discharge in Argon at 1000 Watt.](image)

Both strongly increase with decreasing pressure. An extrapolation of the ionization degree and the electron pressure at lower pressures is also included in the plot. Operation at these lower pressures can be only achieved at higher applied RF powers, where the Langmuir probe can not be operated. It can be observed that at 10-2 Pa the degree of ionization reaches 1%. Assuming a linear electron density scaling with power and constant electron temperature with power the degree of ionization would reach up to 2%. These calculations are based on the assumption of the gas being at room temperature. However it can be expected that with higher gas temperatures the degree of ionization might reach even more than few per cent. This is relatively high for low temperature
plasmas and shows the high efficiency of the NLD at low pressures. This effect of the
electron pressure and gas temperature on the neutral gas density in the NLD is further
described in [84, 85] and shortly in a subsequent part of this chapter. This lower pressure
regime, below 0.1 Pa is actually the typical operational regime for applications. Another
interesting feature of the NLD is that in the NL the “electron pressure” can exceed the
neutral gas pressure. In this pressure region the dominant heating mechanism is
collisionless heating.

The discharge properties can be however controlled as well by changing the
proportionality factor of the applied magnetic field, i.e. by changing the steepness of the
gradient of the magnetic field. Measurements at the NL position for a discharge operated
in Argon at 1000 Watt are shown in figure 4.10. The proportionality factor represents
steps of 10% taken from the standard coil configuration (141 – 144 – 117 A). This
assures in the same time the position for the NL, i.e. at 2/3 of the antenna center.

Figure 4.10: Electron temperature and density for an NLD discharge at 1000 Watt and 0.1
Pa as a function of magnetic field strength.
Measurements have also been carried out in Krypton and show the same qualitative behavior for the different parameter variation. Measured densities are slightly higher than in Argon and the electron temperature slightly lower. This can be understood and expected since Krypton has lower ionization energy than Argon. Thus plasma can be sustained through a lower electron temperature and also increased ionization increases the electron density. Krypton measurements are shown in the spatial resolved measurements analyzed in the next subchapter.
4.2.2. Spatially resolved measurements in the NL plane and at the application plane.

The spatial structure of the magnetic field is playing an important role on controlling both discharge uniformity as well etching uniformity. The discharge special magnetic field configuration, described previously in Chapter 2, can provide additional process control. The influence of the spatial profiles of electron properties is described in this subchapter under the variation of the position of the NL and of the magnetic field gradient [71]. Measurements where made in different NLD modes and at different discharge conditions. Radially resolved measurements on the electron temperature and density in both NLD and ICP discharge in Argon have been performed both in the NL plane and in the downstream (application) plane. The coils currents are kept at their typical values, top – middle – bottom: 141 – 144 – 117 with variations of the magnetic gradient strength of 10 % on each coil. The change of the profile homogeneity is shown in Figure 4.11 with the increase or decrease of the magnetic field strengths.

![Fig 4.11: Calculated radial magnetic field in the neutral loop plane at different magnetic field gradients.](image_url)
This is a numerical calculation of the gradient of the magnetic field in the radial direction along the plane of the null point at different coil current settings. The confinement region for electrons around the null point is 2Bc (where Bc is the cyclotron magnetic field) and the dashed line in the plot represents this region. The black line shows that for strong gradients the NL is a defined torus, where electrons are confined. At lower gradients though, the NL becomes larger and the electron confinement is not anymore so effective and electrons can escape and travel across the center of the discharge.

This will result in more diffuse plasma with a less defined structure and the plasma becomes more homogeneous and develops a disk like structure. By keeping the ratio of the coils current and varying the amplitude, allows to vary the gradient of the magnetic field around the Neutral Loop without changing its position. The steepness of the gradient can be characterized by a scale length L. This is defined to be the distance from the Neutral Loop to the Yoshida confinement point of 2Bc. All the scale lengths are expressed as multiples of the standard length. For the setup used in this work, the length is L= 14 mm. Multiples smaller than this correspond to a steep gradient, whereas large multiples to a more flat profile.

The electron temperature profile in both operational modes, ICP and NLD is similar. It peaks at the same position, illustrating the main production region in both modes is along the torus structure corresponding to the maximum induced electric field, at 2/3 of antenna radius. This is also the position of the NL in the used coil configuration.

The average electron temperature in both modes shows a similar behavior, with slightly higher values in the NLD mode. This is expected for the NLD due to its magnetic confinement, electrons gaining energy through a collisionless process. The spatial variation for the NLD is significantly stronger than for the ICP. Radially resolved Langmuir probe measurements of electron temperature and density in Argon plasma in both ICP and NLD mode at 1 Pa and 1000 Watt are shown in Figure 4.12. The measurements are in the axial plane of the neutral loop (NL).
Interesting spatial features can be observed for the electron density profiles. In the ICP mode the radial profile is relatively flat with a slight peak at the discharge center. In comparison the NLD exhibits a density profile has quite strong inhomogeneities and the electron density maximum is not located as one might expect at the NL, instead is positioned radially inward of the NL as shown in Figure 4.13. The explanation of this lies in the transport and diffusion of electrons within the discharge. In the ICP mode electrons produced along the torus of maximum induced electric field diffuse towards the center and therefore the radial distribution of electrons, peaking at the discharge center arises. However in the NLD mode the situation becomes rather complex due to its complicated magnetic field configuration. As in the ICP case, in the NLD the electrons tend to diffuse towards the center of the discharge, but the confinement provided by the magnetic field is blocking this diffusion process, and thus the density peaking radially inward with respect
to the electron temperature. The density in the NLD is much higher than in the ICP, whereas the temperatures are quite similar.

![Graph showing electron density and temperature vs. position](image)

**Figure 4.13:** Spatially resolved measurement of the electron density and temperature in an NLD discharge

An investigation of the magnetic field gradient around the NL and thus the confinement of the electrons around the NL point are also investigated. In the Yoshida model [17], the confinement of the electrons is within a region of a length of twice $B_c$. $B_c$ is defined as above, as the magnetic field strength which gives a cyclotron frequency equal to the RF frequency. In this definition is not implying that the electrons are actually gyrating at the cyclotron frequency within the NL region, rather their trajectories are expected to be chaotic. The radial profiles of the electron temperature and density for different magnetic field gradients in the NL plane are shown in Figure 4.14 and 4.15 at 1 Pa and Fig 4.16 and 4.17 for 0.5 Pa. Consequently the profiles in the ICP mode are shown in both plots.
Figure 4.14: Radial profiles of electron temperature in the plane of the NL for varying magnetic field gradients, Argon 1000 Watt, 1 Pa.

Figure 4.15: Radial profiles of electron density in the plane of the NL for varying magnetic field gradients, Argon 1000 Watt, 1 Pa.
Figure 4.16: Radial profiles of electron temperature in the plane of the NL for varying magnetic field gradients, Argon 1000 Watt, 0.5 Pa

Figure 4.17: Radial profiles of electron density in the plane of the NL for varying magnetic field gradients, Argon 1000 Watt, 0.5 Pa
The electron temperature in the NL itself remains almost constant for all gradients, whereas the average electron temperature decreases with increase of the gradient. These stronger variations around the NL for increased gradients show a stronger confinement of electrons in the NL region for stronger field strengths. ICP shows a similar temperature profile with less pronounced spatial variations.

One can observe from Figures 4.15 and 4.17 that the electron density is increasing with field strength. It demonstrates that the gradient of the field is not only influencing the efficiency of the discharge but also its structure. For increased gradients the torus shape of the NLD becomes clearly visible, whereas in the opposite case it becomes more diffuse with less defined density distributions.

Homogeneity in the plasma is an important factor for technological applications and thus the electron energy density is a key parameter. The energy density is the product of electron temperature and electron density and is plotted in Figures 4.18 and 4.19 for the same conditions as the above plots.

Figure 4.18: Radial profiles of the electron energy density in the plane of the NL for varying magnetic fields gradients, Argon 1000 Watt, 1 Pa.
Figure 4.19: Radial profiles of the electron energy density in the plane of the NL for varying magnetic fields gradients, Argon 1000 Watt, 0.5 Pa.

It can be observed that the profiles at gradients 3.5*L and 1.75*L have more uniform profiles than in the ICP mode. This is a possibility for improved uniform plasma treatment. Figure 4.11 illustrate how the plasma profile can turn from homogeneous to inhomogeneous with increasing field strengths.

The NLD can be operated in two different modes, with low gradients where in the production plane the plasma is very homogeneous, while on the other hand for large gradients the plasma is very inhomogeneous. But is this inhomogeneous mode which has the advantage of control over the conventional ICP discharges. The position and the gradient of the magnetic field can be controlled by modulating the coil currents during discharge operation to compensate for inhomogeneities and provide uniform processing at the substrate. Regulating the coil currents not only controls the gradient, and therefore the size of NL, but also it can control the position of the NL by scanning the torus radius. As just mentioned, homogeneity in the production region is one issue; however for application purposes plasma uniformity in the substrate treatment region is equally important. Therefore the transport between the production region and the application
Due to its magnetic field structure in the NLD, this can be rather complex [81]. In Figures 4.20 and 4.21 is plotted the electron temperature and density in the downstream region, for the same conditions as above measurements at 1 Pa. In figures 4.22 and 4.23 is plotted the same at 0.5 Pa. The electron temperature in the ICP mode shows a more homogeneous profile than in the NLD mode. This can be a consequence of the missing magnetic field structure. The average electron temperature in the NLD is lower in the NLD mode compared with the ICP. For the stronger gradients, where the NL torus is strongly defined, the electron temperature profile follows the same structure as in the NL plane, with a smaller radius. The electron density peaks at a inward position as in the NL plane.

Fig 4.20: Radial profiles of electron temperature in the application plane for varying magnetic fields gradients, Argon 1000 Watt, 1 Pa.
Fig 4.21: Radial profiles of electron density in the application plane for varying magnetic fields gradients, Argon 1000 Watt, 1 Pa.

Fig 4.22: Radial profiles of electron density in the application plane for varying magnetic fields gradients, Argon 1000 Watt, 0.5 Pa.
Fig 4.23: Radial profiles of electron density in the application plane for varying magnetic fields gradients, Argon 1000 Watt, 0.5 Pa.

Figures 4.21 – 4.23 are highlighting the strong coupling between the two regions and the transport of electrons from the NL region down to the processing region through the separatrix of the magnetic field without major losses.

Electron density profiles in the NLD mode, for low magnetic field gradients and showing a more homogeneous profile than in the ICP. The density is higher as well in the NLD than in the ICP. For surface treatment, the important parameter is the ion flux, since ion energy can be separately controlled by using an additional substrate bias. In Figures 4.24 and 4.25 is plotted the ion flux at a possible application region at 1 Pa and 0.5 Pa respectively. The ion flux is determined from the ion density and the Bohm velocity calculated out from the electron temperature as:

\[ d_{\text{eff}} = \frac{1}{2} \frac{Rl}{Rh_i + lh_i} \]
The ion flux uniformity is similar in the ICP mode as in the NLD mode, with low field gradients; however the fluxes are relatively higher in the NLD mode. Additionally the ratio of electron density in the NLD to the ICP mode is greater at the NL plane than in the application plane. This is due to the confinement of particles in the NL.
Even though electrons are confined in the production region of the NLD there is still a higher electron density, resulting in higher fluxes and thus allowing higher etch rates in the application region than the conventional ICP discharge. Decreasing the pressure at which the discharge operates reveals that both electron temperature and density become less uniform. The NLD discharge can be operated at significantly lower pressures than the ICP. Operation in this low pressure regime decreases ion scattering through collisions in the sheath, allowing perpendicular ion bombardment at the surface. Figure 4.26 shows the radial profile of the energy density in the NL plane and the ion flux at the application plane of the discharge for both a strong and weak field gradients.

For the above measurements the ICP can not be sustained at this pressure. This plot clearly illustrates the two operational modes of the NLD, the homogeneous mode using a weak field gradient and an inhomogeneous mode that allows better controllability.
Another feature of the NLD is the strong coupling between the NL region and the application region. Referring to Figure 3.3 one see how the separatrix structure of the magnetic field travels down through the plasma source to the substrate plane. Figure 4.22 shows the high density plasma streams down, along the separatrix, from the NL region to the substrate region and results in similar structures, with smaller diameters, in the application plane.

As discussed earlier the NLD offers additional control and the effect of varying the gradient of the field was discussed and investigated above. Now we will examine the influence of radially scanning the position of the null point. Figures 4.27 and 4.28 show plots of the electron density and temperature at the NL plane at 0.1 Pa and 1000 Watt.
Figure 4.27: Radial profiles of electron temperature in the NL plane for varying the neutral loop radius, Argon, 1000 Watt, 0.1Pa

Figure 4.28: Radial profiles of electron temperature in the NL plane for varying the neutral loop radius, Argon, 1000 Watt, 0.1Pa
One clearly observes as the radius of the NL is varied, the structure of the plasma is varied similarly. The electron temperature peaks always at the center of the NL and the density at a position radially inward of the NL. The energy density in the NL plane, plotted in Figure 4.29 shows for larger NL radii that the torus structure is more pronounced than at smaller radii, where electrons can diffuse across the center of the plasma.

![Figure 4.29: Radial profiles of energy density in the NL plane for varying the neutral loop radius, Argon, 1000 Watt, 0.1Pa](image)

Similar measurements were performed in the application plane where a possible wafer can be placed. The corresponding electron temperature and density are shown in figures 4.30 and 4.31. As expected the plasma streams down along the separatrices structures.
Figure 4.30: Radial profiles of electron temperature at a possible application plane for varying the neutral loop radius, Argon, 1000 Watt, 0.1Pa

Figure 4.31: Radial profiles of electron temperature at a possible application plane for varying the neutral loop radius, Argon, 1000 Watt, 0.1Pa
The ion flux at a possible application plane is plotted in Figure 4.32.

Figure 4.32: Radial profiles of ion flux at a possible application plane for varying the neutral loop radius, Argon, 0.1Pa

Another point to note is that regardless of the gradient or the diameter of the NL, the density decreases outside the region of the separatrix. This leads to inhomogeneity at the substrate surface. Therefore a design where the separatrix does not end at the substrate surface, such as a very large NL radius, would improve uniformity. The acquired plasma potential as a function of magnetic field strength is shown in figure 4.33. Although the electron temperature and density show strong inhomogeneous profiles, $V_p$ has a relative uniform one. This behavior is reproducible at all magnetic field strengths and at any NL diameter.
Figure 4.33: Radial profiles of the plasma potential for varying the magnetic field strength in the NL plane, Argon, 1000 Watt, 0.1Pa

Until now Langmuir probe measurements on the plasma parameters in Argon have been presented. The electron temperature and density where determined from the acquired EEDF as described in 3.3.2. Figures 4.34 and 4.35 shows the EEDF in an NLD discharge as a function of power at the NL for Argon any Krypton at 0.1 Pa. Both clearly show a Maxwellian distribution function. Moreover the spatial distribution of the electron density and temperature in Krypton show the same qualitative behavior for different parameter variation.
Figure 4.34: EEDF in an NLD, Argon discharge as a function of power at 0.1 Pa

Figure 4.35: EEDF in an NLD, Krypton discharge as a function of power at 0.1 Pa
Measured densities in Krypton in an NLD discharge are slightly higher and temperatures slightly lower that in Argon. This is due to the lower ionization energy of Krypton compared to Argon. Thus plasma can be sustained through a lower electron temperature and also increased ionization increases the electron density.

4.3 Neutral gas depletion in the NLD

Typical operation of the NLD discharge, in the low pressure regime can reach ionization degrees up to few percent [46, 71]. The high plasma densities at relatively low neutral gas pressures and the pronounced spatial structures of plasma parameters in the NLD can for low temperature plasmas, in localized neutral gas depletion in the region of the NL. In this region the electron pressure can exceed the neutral gas pressure, in particular at relatively low gas pressures, resulting in a depletion of the neutral gas density. Recent theoretical investigations on the depletion phenomena have been carried out in [74, 75], however less research has concentrated on depletion through high ionization rates [76, 77, 78, 79]. Moreover direct experimental measurements are lacking. In [72, 73] experimental investigations in an NLD discharge on the neutral gas depletion have been carried out. The depletion through both gas heating and high ionization rates is investigated. Langmuir probe and Thomson scattering measurements are used to determine the electron temperature and density to determine the effect of the depletion through high ionization rates.

Diode laser absorption spectroscopy on metastable argon atoms is used to measure metastable densities and gas temperatures [80]. Neutral gas depletion in the region of the NL is investigated by spatially resolved LIF measurements on argon metastables and TALIF experiments on ground state krypton atoms. Knowledge of the gas temperature is required for the determination of neutral particle densities [81]. Gas temperatures are deduced from the Doppler-profile of line integrated diode laser absorption measurements on metastable argon. The temperature linearly increases - quite significantly - with both power and pressure as shown in Figure 4.36.
Figure 4.36: Gas temperature and ratio of metastable density to ground state density, measured using diode laser absorption spectroscopy as a function of pressure, Argon, P = 1000 W

Line integrated absolute densities of metastable argon atoms are also measured using diode laser absorption spectroscopy. The ratio of metastable density to ground state density is determined by the electron temperature - as expected from a simple balance equation model with electron quenching being the dominant destruction mechanism [14]. At low pressures (< 0.1 Pa) the ratio reaches values up to a few times $10^{-3}$ as shown in Figure 4.32. Metastable atoms are, therefore, expected to play a very important role, in particular at low pressures.

Figure 4.37 shows the degree of ionization and electron pressure as a function of neutral gas pressure. These were determined using Langmuir probe measurements of electron temperatures and densities. As expected from global model predictions, both the degree of ionization and electron pressure strongly increase with decreasing pressure. The degree of ionization can reach up to several percent, even exceeding ten percent, which is relatively high for low temperature plasmas and shows the extremely high efficiency of
the NLD. A related and very interesting feature is that in the NL region the electron pressure can exceed the neutral gas pressure, especially at relatively low pressures.

Figure 4.37: Degree of ionization and electron pressure over neutral gas pressure as a function of pressure in an argon NLD at 1000W.

A detailed investigation on this particularity of the NLD is made in references [72 and 73].
4.4. Thomson scattering measurements

Thomson scattering measurements have been made for similar Langmuir probe measurements conditions previously shown. Since an ICCD camera is used for recording the scattered Thomson spectrum, it can be used either in an accumulation mode or in a special “photon counting” mode. An analysis between the two methods has been done to determine the best alternative. For a Raman spectrum the normalized result is shown in Figure 4.38.

![Figure 4.38: Accumulation vs. photo-counting mode for a Raman spectrum](image)

The special design of the TGS, shown in Chapter 2, has a high stray rejection factor of about $10^{-6}$. Thus the photon – counting technique will have an advantage only in improving the signal– to– noise ratio. However this slight improvement on the noise base level is shown only in the Raman spectrum. For a Thomson spectrum the difference between the two methods is not noticeable. Moreover the accumulation method makes the analysis on the electron density relative easier, since a recalibration of spectra’s is not required.
A typical Thomson spectra with respect to a pressure variation recorded with the Triple Grating Spectrograph in the NL region is shown in Figure 4.39. Thomson scattering allow measurements on the plasma parameters at operational pressures down to 0.01Pa. Due to the limit of the SoftShout automatic pressure valve control the pressure bellow 0.1 Pa can not be 100% trusted. However bellow 0.1 Pa the measurements with the Langmuir probe could not be performed. The Argon NLD discharge was operated at 1000 Watt with the coils in the SCC. In this configuration the NL radius is 6 cm. The center region of the spectra \((\pm 1.2\,\text{nm})\) is blocked by the masked. Previous Langmuir probe measurements in the NL region showed a Maxwellian distribution function, therefore one will expect in an incoherent scattering regime a Gaussian profile of the Thomson spectra.

![Thomson scattering as a function of pressure variation in the NLD](image)

Figure 4.39 Thomson scattering as a function of pressure variation in the NLD

The change in the distribution function from a Bi-Maxwellian to a Maxwellian from a low applied RF power, 150 Watt, to 1000 Watt is shown in Figure 4.40 for an ICP discharge operated at 2 Pa.
Figure 4.40: Bi-Maxwellian to Maxwellian distribution in an ICP discharge in Argon at 2 Pa and 1000 Watt

Two Energy Probability Function (EEPf) determined from the measured spectra are shown in Figure 4.41 for the two pressures. The wings are centered back with the
displacement given by the Gaussian fit. Both curves clearly show a Maxwellian
distribution with a temperature of 3.5 eV at 0.1 Pa and 2.2 eV at 1 Pa respectively. The
error in the temperature measurement is of maximum 5% at the lowest operational
pressure of 0.05 Pa.

Figure 4.41: Thomson scattering EEPF

An important aspect in controlling the discharge parameters is the possibility to
vary the applied magnetic field strength. This will result in a variation in size of the
confinement region around the null point. A reduction in steps of 10 % of the applied DC
voltage starting from the SCC is made. Gradients 0.5 and 0.3 will correspond respectively
to a coil configuration of (72 -74 -54) and (42 – 45 – 35). Simulations of the magnetic
field structure at these values confirm that the spatial position of the NL remains
constant.

Figure 4.42 and 4.43 show the electron temperature and density as a function of pressure
at different gradients and the respective comparison with an ICP discharge.
Figure 4.42: Electron temperature in NLD and ICP discharge as a function of pressure and magnetic field strength.

Figure 4.43: Electron density in NLD and ICP discharge as a function of pressure and magnetic field strength.
With the magnetic field turned off, a stable ICP discharge can be operated up to a lower pressure of 0.5 Pa. The Thomson scattering measurements shown in figures 4.42 and 4.43 are made in an NLD and ICP discharge ignited in Argon at a RF applied power of 1000 Watt. The measured EVDF reflected at all pressures and gradients in the NLD a Maxwellian distribution. For better plot visibility, only measurements at three gradients are shown. The behavior at other gradients is similar. The electron temperature in the Neutral Loop remains almost constant at all gradients, and strongly increases with pressure decrease, behavior shown as well by previous Langmuir probe measurements in the NL plane shown previously in this chapter.

The electron density however shows a stronger variation with changing the gradient profile. This clearly proves the role of the magnetic field confinement on the discharge characteristics. Moreover a direct comparison with an ICP discharge shows that the electron density is higher in the NLD, particularly in the low pressure regime at all gradients. Similar measurements where performed at higher powers, up to 2000 Watt. For the standard Gradient 1 configuration a power variation is shown in figure 4.44 and 4.45.

Figure 4.44: Electron density in the NLD at 0.1 Pa as a function of applied RF power
Figure 4.45: Electron temperature in the NLD at 0.1 Pa as a function of applied RF power

Applied RF power variation in the NLD at all magnetic field gradients, show the same tendency, both for electron temperature and density. A relatively linear increase is observed.

Due to the magnetic coils arrangement, the NL radius can be controlled by changing the middle coil applied current. This feature made the NLD discharge attractive for the material processing industry, since higher uniformity can be achieved at the substrate. The influence on the electron density and temperature by varying the NL radius is shown in figures 4.46 and 4.47. The NLD discharge was operated at 0.1 Pa and 1000 Watt in Argon. The coil currents are respectively (141 – 120 - 117) for a NL radius of 5.5 cm (small), (141 – 136 - 117) for a NL radius of 7.5 cm (medium), and (141 – 150 - 117) for a NL radius of 9 cm (large).
Figure 4.46: Electron temperature in NLD as a function of mid – coil applied DC current

Figure 4.47: Electron density in NLD as a function of mid – coil applied DC current
The electron temperature shows a relative constant profile at all gradients and NL radii. The electron density shows a strong dependency both with the magnetic field gradient, increasing with the magnetic field strength, and with the NL size, increasing with decrease in the NL size.

However since no spatial – resolved measurements are possible, an alternative is variation of the middle coil current, such as the scattering volume is set at some other position in the plasma. A variation from 137 – 150 Ampere has been simulated to maintain the NL position at the same plane with the laser beam. A ± 6 Ampere from the “center” position, i.e. 144 A will result in a variation of the NLD radius with ± 4 cm. The result for the electron temperature and pressure is shown in Figure 4.48. The bigger the applied current, the smaller the NLD radius becomes.

![Figure 4.48: Electron temperature and density in NLD as a function of mid – coil applied DC current.](image)

The electron temperature stays almost constant at all NLD radii. A similar behavior can be observed as well in the electron density. It was always assumed that the most efficient power coupling is at the NL. However in this case the NL position does not
correspond anymore with the position of the maximum induced electric field. Thus other effects should be taken into consideration for this behavior. Some aspects will be discussed in chapter 5 of this work.

4.5. Comparison between Thomson scattering and a novel Optical Emission Spectroscopy (OES) CRM model.

As shown above and in [59 – 67, 83] Thomson scattering (TS) has proved to be a very versatile techniques in low pressure discharges for determination of electron densities temperatures and within certain energy limits also the electron energy distribution function (EEDF). However techniques based on optical emission are clearly non-invasive, require only moderate spectroscopic equipment, are easy to implement, and measurements are usually fast. The aim is to measure electron temperatures and densities by applying OES and TS in parallel in a low pressure, high density discharge in argon [83].

A novel Optical Emission Spectroscopy (OES) [85] technique has been applied to investigate the NLD discharge. However, the underlying collisional–radiative models (CRMs) are often very complex and the validity of assumptions made in the model, for example, Maxwellian EEDFs, is often not clear. Further, additional input parameters such as the gas temperature are often required. Electron temperatures and densities obtained by both diagnostics are compared. The validity of assuming Maxwellian EEDF in the CRM results is carefully checked by TS.

The radiation emitted from the NL plane is collected tangentially to the loop by an optical system formed by a lens (focused at the NL point which is located at a radial distance of 6 cm from the symmetry axis in the centre) and an optical waveguide that images the plasma emission onto a spectrometer. Argon spectra are detected using an Ocean Optics HR4000 spectrometer. The spectral response of the optical and detecting system is accounted by using an Ulbricht sphere for calibration. A certain degree of line integration has to be expected along the detection cone of the optical system due to the tangential direction of observations. However, the region of homogeneity (ne, Te) around the NL point can be estimated from our previous work [8] as $\rho \approx 1.6$ cm, where $\rho$ is a
torus radius around the NL. This gives a total homogeneous length along the line of integration of the detection cone of about 10 cm, with the centered point at the focus. Therefore, line integration occurs predominantly along a homogeneous region and the disturbance caused to the emission spectra is assumed to be negligible.

Spectra of the NL emission where recorded over a wide wavelength range (400 – 900) nm and cover the whole experimental conditions explored in the study (applied rf power $P = 1$ kW, 1.5 kW, and 2 kW at a frequency of $f = 13.56$ MHz, argon gas pressure $p = 0.05$ Pa, 0.1 Pa, 0.5 Pa, 1 Pa, and 5 Pa). Spectra emitted from the plasma torus at the lowest ($p = 0.05$ Pa) and at the highest ($p = 5$ Pa) pressures, applied power $P = 1$ kW, and in the wavelength range (400 – 900) nm are shown in figure 4.49. The spectral range (400 – 690) nm, at two pressures, is enlarged in separate figures in the frame of the main graphs to give an idea of spectral line intensities. The presented experimental data of argon spectra are corrected for the spectrometer response and normalized to the most intense line in the spectra at 811.5 nm. This example indicates the sensitivity of the emission spectra to changes of the discharge operating conditions.

Figure 4.49: Emission spectra of NLD at $P = 1$ kW and gas pressure values $p = 0.05$ Pa in (a) and $p = 5$ Pa in (b).
The emission spectroscopic diagnostic method combines measurements of certain argon line-intensity with collision-radiative modelling of excitation kinetics at low pressures. The CRM employed for NLD plasma diagnostics was proposed initially for Ar-He gas mixtures [86], subsequently it was further developed for OES of non-equilibrium low pressure argon inductively coupled plasmas (ICPs) [87], and it is further extended here. The model includes kinetic processes determining the population densities of the first four \((1s_{5,2} - \text{Paschen notation})\) and next ten \((2p_{10,1} - \text{Paschen notation})\) excited levels, belonging to the \(3p^54s\) and \(3p^54p\) Ar level configurations, respectively. A more detailed view on this model is made in [84].

The results for Te and ne obtained by OES are compared with those obtained by TS measurements performed under identical experimental conditions. The dependences of electron density and electron temperature on pressure \((p = 0.05 \text{ Pa}, 0.1 \text{ Pa}, 0.5 \text{ Pa}, 1 \text{ Pa and } 5 \text{ Pa})\) at \(P = 1 \text{ kW}\) are presented in figures 4.50 (a) and (b). The open triangles correspond to TS measurements.
The OES data are presented with full circle and full triangle for the case when the perfect gas law or the law of partial pressure, respectively, is applied in the modelling. The agreement between the OES method and TS is excellent throughout, except for the lowest pressure of $p = 0.05$ Pa where the electron temperature determined by OES is about 1 eV lower than the value obtained by TS. We attribute this to the transition of the diffusion regime, assumed in the CRM, to the free fall loss of the metastables. The critical pressure $p = 0.1$ Pa agrees quite well with the estimate derived in section 3. It should be noted that agreement between OES and TS can only be achieved if the depletion of the neutral gas density by electron pressure is included. Neglecting this leads to a substantial overestimation of the electron density and a small deviation in the temperature. Electron densities and temperatures as a function of applied rf power and gas pressure are presented in figure 4.51.
Figure 4.51: Electron density in (a) and temperature in (b) dependence on the applied power.
A straight line fit on each $n_e(p)$ dependence through zero (figure 4.51 (a)) and a straight line fit on each $T_e(p)$ dependence (figure 4.51 (b)) are given. Variation of the applied RF power leads to an accordingly linear variation of the electron density. In contrast, $T_e$ is very little affected through changes of the applied power. Both results are in agreement with TS measurements and also with the expected typical behavior resulting from global energy and particle balance, respectively [14].
Chapter 5

Thomson scattering and RF MOS - Phase resolved measurement of anisotropic electron velocity distribution function and charged particle drifts in weakly magnetized plasmas

5.1. Introduction

Oscillating radio – frequency (RF) fields penetrating into a plasma are leading to a corresponding oscillation of the electron velocity distribution function. This oscillation can be measured by Thomson scattering in an ICP and an NLD discharge. From this the local current density can be inferred by combining the oscillation amplitude with the plasma density, also inferred by Thomson scattering. First time measurement of these oscillations by Thomson scattering are reported in [88] from which a part of this chapter is based. Moreover Thomson scattering measurements showed a shifted Maxwellian distribution function. This shift can be interpreted as a drift component in the direction of the scattering vector. The shift shows a high dependence with pressure. In the NLD this drift is identified to be the diamagnetic drift. This drift could be identified as well in the RF-MOS measurements.

Chapter 2 gave an introduction on generating and sustaining a plasma discharge with or without an externally applied magnetic field. Is not rarely that magnetic fields are applied to enhance the discharge characteristics, mostly being focused to lower the pressure regime in which the discharge operates. In this context a short introduction of the basic formalism on the charge particle dynamics in uniform and non-uniform magnetic fields is given. In the first part of this chapter fundamental concepts on the charge particle kinetics in a uniform magnetic field are introduced. However the applied magnetic fields to plasma discharges are typically inhomogeneous, thus having field
gradients perpendicular and parallel to $\mathbf{B}$. This will have as a consequence on the charge particle motion a “drift” component across the magnetic field. For a uniform electric and magnetic field the drift is the well known electric drift or $\mathbf{E} \times \mathbf{B}$ drift. In an inhomogeneous magnetic field, several drifts are identified, namely gradient-$\mathbf{B}$ drift, curvature drift, polarization drift. A wider description on each is given in [52, 53, 55].

Plasmas can be treated either kinetically, therefore by using a statistical approach or as fluids. A special attention is given in this work to the $\nabla \mathbf{B}$ drift and to the diamagnetic drift. The experimental setup was modified, the bottom flange being replaced with a window to allow RFMOS bottom measurements. For consistency sake the experimental setup configuration is shown in Figure 5.1.

Figure 5.1: Experimental setup schematic. The scattering vectors in relation to the induced electric field are included.
5.2. Motion of a charge particle in a uniform magnetic field

In a spatially uniform magnetic field, the equation of motion of a charged particle can be written as:

\[
m \frac{d \vec{v}}{dt} = q \left( \vec{v} \times \vec{B} \right)
\]

(5.1)

where q represents the particle charge. Since no electric field is present the only force acting on the particle is the Lorenz force. To resolve this equation one can conveniently separate the particle velocity into its components along and perpendicular to the magnetic field lines. Thus velocity \( \vec{v} \) becomes:

\[
\vec{v} = \vec{v}_\parallel + \vec{v}_\perp
\]

(5.2)

To obtain the vectorial expressions for the parallel and the perpendicular components the BAC – CAB rule is used. For a vectorial cross product this rule is written as:

\[
\vec{A} \times \left( \vec{B} \times \vec{C} \right) = \vec{B} \left( \vec{A} \cdot \vec{C} \right) - \vec{C} \left( \vec{A} \cdot \vec{B} \right)
\]

(5.3)

Applying equation (5.3) for a magnetic field \( \vec{B} \) and an arbitrary vector \( \vec{C} \) one can write:

\[
\vec{B} \times \left( \vec{B} \times \vec{C} \right) = \vec{B} \left( \vec{B} \cdot \vec{C} \right) - \vec{C} \left( \vec{B} \cdot \vec{B} \right) = \vec{B} \left( \vec{B} \cdot \vec{C} \right) - \vec{C} B^2
\]

(5.5)

Rewriting equation (5.5) in a more convenient way:

\[
\vec{B} \times \left( \vec{B} \times \vec{C} \right) = -B^2 \left[ \vec{C} - \frac{1}{B^2} \vec{B} \left( \vec{B} \cdot \vec{C} \right) \right]
\]

(5.5)

In the RHS of the equation (5.5) the dot product \( \vec{B} \cdot \vec{C} \) actually is expressing an arbitrary component of C in the direction of B. Thus one can define this to be \( BC_1 \). With this in mind the LHS will represent the perpendicular component of C. Therefore based on equation (5.2) one can define:

\[
\vec{C}_1 = \frac{\vec{B} \left( \vec{B} \cdot \vec{C} \right)}{B^2} = \frac{\vec{B}}{B} C_1
\]

(5.6)
Taking into consideration equations (5.6 and 5.7) the velocity components parallel and perpendicular to the magnetic field can be expressed as:

\[ v_p = \frac{B}{B^2} B \cdot v = B v_i \]  \hspace{1cm} (5.8)

\[ v_\perp = -\frac{B}{B^2} B \times (B \times v) = -B \times (b \times v) \]  \hspace{1cm} (5.9)

where \( b = \frac{B}{B} \) represents the unit vector along the magnetic field B. Basically from the above two equations, getting the parallel component of a vector with respect another vector, one has to take the dot product of the two, whereas for the perpendicular component the cross product.

Substituting \( v \) with (5.2) in the equation of motion:

\[ \frac{d}{dt} v = \frac{q}{m} (v \times B) \]  \hspace{1cm} (5.10)

For the parallel velocity component equation (5.10) becomes:

\[ \frac{dv_i}{dt} = 0 \iff v_i = v_{i0} = \text{constant} \]  \hspace{1cm} (5.11)

This shows that the particle velocity parallel to the magnetic field remains constant. In other words the magnetic field does not have any influence on the velocity component of the particle along its direction. The particle does not gain or loose energy. Assuming that the magnetic field is taken in the \( \hat{z} \) direction, \( B = B\hat{z} \) and integrating equation (5.11), the particle position along the magnetic field is:

\[ v_i = \frac{dz}{dt} \Rightarrow z = z_0 + v_{i0} t \]  \hspace{1cm} (5.12)

representing a uniform translation along magnetic field.

The velocity component perpendicular to the magnetic field is derived by taking the cross product from equation (5.9) with \( v \) given by (5.2):
\[ -b \times \left( b \times \left[ \frac{d v_\perp}{dt} + \frac{d v_\parallel}{dt} = \frac{q}{m} (v \times B) \right] \right) \Leftrightarrow \frac{dv_\perp}{dt} = \frac{q}{m} (v_\perp \times B) \]  

(5.13)

Integrating (5.13) with respect to \( dt \) and replacing \( v_\perp \) in the RHS with \( \frac{d r_\perp}{dt} \):

\[ v_\perp = \frac{q}{m} \left( r_\perp - r_\perp^{(0)} \right) \times B \]  

(5.15)

This equation can be further simplified by introducing the cyclotron frequency vector defined by:

\[ \vec{\omega}_c = -\frac{q \vec{B}}{m} \]  

(5.15)

with its magnitude given by: \( \omega_c = \frac{|qB|}{m} \). For an electron this is known as the electron cyclotron frequency. The sign of the \( \omega_c \) actually represents the direction in which the particle rotates with respect to the magnetic field.

With (5.15) equation (5.15) becomes:

\[ v_\perp = \vec{\omega}_c \times \left( r_\perp - r_\perp^{(0)} \right) \]  

(5.16)

Finally the particle trajectory is obtained by applying a cross product on equation (5.16) with \( \vec{\omega}_c \) and taking into consideration the BAC – CAB rule (5.3):

\[ r_\perp = r_\perp^{(0)} + \frac{v_\perp \times r_\perp^{(0)}}{\omega_c^2} = r_\perp^{(0)} + \rho \]  

(5.17)

with the gyro-radius \( \rho = \left| \vec{\rho} \right| = \frac{v_\perp}{\omega_c} \).

A sketch of the particle trajectory with the respective vectors is shown in Figure 5.2. The gyro-radius is one of the fundamental quantities in magnetized plasmas. This represents the particle helix radius as it travels along the magnetic field. It is often called “Larmor radius”.

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Previously a vectorial calculation was made. In Cartesian coordinates this translates into the particle position in terms of its guiding center for the x, y and z:

\[
\begin{align*}
  x &= x_0 + \frac{v_y}{\omega_{ce}} \sin(\omega_{ce} t + \phi) \\
  y &= y_0 - \frac{v_y}{\omega_{ce}} \cos(\omega_{ce} t + \phi) \\
  z &= z_0 + v_z t
\end{align*}
\]

which correspond to a helical path.

### 5.3. Motion of a charged particle in an inhomogeneous electric and magnetic field.

Real plasma systems are however using spatially and temporally inhomogeneous electric and magnetic fields. In this case the equation of motion becomes:

\[
m \frac{d \vec{v}}{dt} = q \left( \vec{E} + \vec{v} \times \vec{B} \right)
\]

(5.18)
In an inhomogeneous magnetic field the equation (5.18) becomes relatively hard to solve. However if the particle gyro – radius $\rho$ is much less than the magnetic field scale length inhomogeneity: $\rho / L \ll 1$ the motion equation can be solved in the small gyro-radius approximation with the following assumptions:

- The velocity of the particle in an inhomogeneous magnetic field is given by the sum of its parallel and perpendicular components:

$$ v = v_\perp + \delta v $$

(5.19)

where $\delta v = v_{\text{drift}} + v_\parallel$

- Given the $\rho / L \ll 1$ the magnetic field can be express in a Taylor expansion around the guiding center:

$$ B = B_0(x_0) + \left[ \rho \nabla B \right] $$

(5.20)

where the term $\left[ \rho \nabla B \right]$ contains the inhomogeneity of the magnetic field (divergence, convergence, shear).

Inserting 5.19 and 5.20 in 5.18 and taking into consideration the magnetic field configuration, equation 5.18 becomes:

$$ m \left( \frac{d}{dt} v_{\text{drift}} + \frac{d}{dt} v_\parallel \right) = q \left[ E_0 + v_{\text{drift}} \times B_0 + (v_\perp \times \rho \nabla B) - m v^2 \nabla \cdot (\nabla B) \right] $$

(5.21)

For getting the drift component one solves equation (5.21) for the gyro-averaged drift on the guiding center and takes only the components perpendicular to $B$.

The physical understanding of the $\nabla B$ drift follows from the fact that the curvature radius of the gyro – orbit is small when the particle finds itself in a stronger magnetic field and large when the magnetic field is weak. Thus net particle drift perpendicular to $B$ and $\nabla B$ occurs. The $\nabla B$ drift depends on the particle electric charge. This will give rise to a net current, leading to a charge separation in the plasma.
However plasmas are a collection of individual charged particles and can be either treated from a kinetically point of view or it can be considered as acting as a special kind of fluid. These two approaches are respectively kinetic and fluid approach.

- **Kinetic model for particle drifts in external magnetic fields.**

Assuming that the plasma has a Maxwellian distribution function in the kinetic description we solve the Boltzmann equation:

\[
\frac{\partial f}{\partial t} + v \cdot \nabla f - \frac{e}{m} \left( \vec{E} + v \times \vec{B} \right) \cdot \frac{\partial f}{\partial v} = \frac{\partial f}{\partial t}
\]

(5.22)

For simplification the following assumptions are made: no RF electric field and a collisionless case. Equation 5.22 becomes:

\[
v_y \frac{\partial}{\partial y} f - \omega_y g(y) \left( v_y \frac{\partial}{\partial v_y} - v_x \frac{\partial}{\partial v_x} \right) f = 0
\]

(5.23)

where \( \frac{e B}{m} = \omega_y g(y) \).

The solutions for equation 5.23 have the form:

\[
f = f \left[ C_1 \left( v_x^2 + v_y^2 \right), C_2 \left( v_x + \omega_y \int_0^y g(y') dy' \right) \right]
\]

(5.25)

Denoting \( G(y) = \int_0^y g(y') dy' \) and considering a shifted Maxwellian distribution function with an additional drift component of the form \( v_x + u \). With

\[
C_1 = \frac{m}{2kT_e}, \quad C_2 = \frac{m}{2kT_e}, \quad 2u = \frac{1}{v^2_{thermal}}
\]

\( f \) can be written as:

\[
f = C_3 \exp \left[ - \frac{(v_x + u)^2 + v_y^2}{v^2_{thermal}} \right] \exp \left[ - \frac{2u \omega_y}{v^2_{thermal}} G(y) \right]
\]

(5.25)

where the function \( G(y) = \int_0^y g(y') dy' \) represents the profile of the magnetic field.

Using Uchida’s slab model:

\[ g(y) = \frac{y}{L} \quad G = \frac{y^2}{2L} \]

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the density profile function can be expressed as:

$$h(y) = \exp\left(-\frac{1}{2}\left(\frac{y}{s}\right)^2\right)$$  \hspace{1cm} (5.26)

From the RHS of equation 5.25 combined with 5.26, results for the drift velocity:

$$u = \frac{kT_e}{eB_s} \left(\frac{L}{s}\right)$$  \hspace{1cm} (5.27)

where \(s\) represents the spatial gradient of the magnetic field. However this is a special solution for the case of a shifted Maxwellian. Conversely only a diamagnetic drift will not be sufficient to obtain a shifted Maxwellian. Furthermore there can be no drift without a density gradient.

- **Fluid dynamic model for particle drifts in external magnetic fields.**

  In the fluid dynamic one begins with the standard form of the momentum balance equation:

  The following assumptions are made:

  - no static electric field present, shown by Langmuir probe measurements at different gradients in Figure XX
  - the ambipolar diffusion is assumed to play no important role
  - anisotropic pressure
  - collisions are included

  With the above assumptions equation 5.28 becomes:

  $$-\frac{e}{m} (\vec{v} \times \vec{B}) - \frac{kT_e}{m} \frac{\nabla n}{n} - v_m \vec{v} = 0$$ \hspace{1cm} (5.29)

  Taking only the solution for the perpendicular velocity component on the magnetic field and considering no density gradient in the x direction \(\nabla n_x = 0\) equation 5.29 can be solved for x and y:

  $$-\omega_c v_y - v_m v_x = 0 \Rightarrow v_y = -\frac{v_m}{\omega_c}$$  \hspace{1cm} (5.30)

  $$\omega_c v_x - \frac{kT_e}{m} \frac{\nabla n}{n} - v_m v_y = 0$$ \hspace{1cm} (5.31)
Inserting 5.30 in 5.31:
\[
\omega_x v_x - \frac{kT_e}{m} \frac{\nabla_y n}{n} - \frac{v_m^2}{\omega_e} v_x = 0 \iff \left( \omega_x + \frac{v_m^2}{\omega_e} \right) v_x = \frac{kT_e}{m} \frac{\nabla_y n}{n}
\]  (5.32)

Thus
\[
v_x = \frac{kT_e}{m \omega_e} \frac{\nabla_y n}{n} \left( \frac{1}{1 + \left( \frac{v_m}{\omega_e} \right)^2} \right)
\]  (5.33)

Using as in the kinetic approach Uchida’s slab model, and taking into consideration relation 5.26:
\[
\frac{\nabla_y n}{n} = -\frac{y}{s^2}
\]  (5.35)

With 5.35 the velocity in the x direction can be now written:
\[
v_x = -\frac{kT_e}{eB_0s} \left( \frac{L}{s} \right)
\]  (5.35)

which is the same result as in the kinetic approach.

The minus sign means than the drift is defined in a negative direction. \( v_x \pm u \).

This effect of the density gradient is sketched in Figure 5.3, showing the electron orbits. Clearly there is a greater current to the right than to the left. These currents are the effect of the diamagnetic drift and they act to reduce the magnetic field inside the plasma.

Figure 5.3: Electron orbits in a magnetic field with a density gradient.
5.4. Thomson scattering and RFMOS measurement of charged particle drifts velocities in an NLD discharge.

Previous Thomson scattering measurements in the NLD have been reported in [92, 93, 95]. In the experimental setup used in this work Thomson scattering measurements showed a shifted Maxwellian distribution function. This shift can be interpreted as a drift component in the direction of the scattering vector. A plot showing the measured EEDF in an NLD operated at 1000 Watt in Argon at two different pressures 0.1 Pa and 1 Pa is shown in Figure 5.5. In chapter 3 oscillatory velocities can be quantified by using RFMOS; however the first component of the modulation is proportional with the drift velocity and the induced electric field. A detailed investigation by Thomson scattering and RFMOS on the drift behavior at different plasma conditions is shown in the following.

![Figure 5.5: Shifted EEDF in NLD in Argon at 1000 Watt, 0.1 Pa and 1 Pa. The coils are powered at 141 – 144 – 117 A.](image)

The shift difference from the “central” position, the middle of the ICCD camera chip (512 pixels), is clearly shown in Figure 5.5.
RFMOS measurements were made both from a side–view port and from the chamber bottom. The respective values at the same pressures as the above Thomson scattering measurements are shown in Figures 5.5 – 5.9.

Figure 5.5: RFMOS drift velocity in an Argon NLD discharge, in the 141 – 144- 117 A coil configuration, 0.1 Pa, 1000 Watt side–view. The calculated magnetic field for the same coil configuration is overlaid.
Figure 5.6: RFMOS drift velocity in an Argon NLD discharge, 0.1 Pa, 1000 Watt bottom view. Magnetic coils powered at 141 – 144 – 117 A.

Figure 5.7: RFMOS drift velocity in an Argon NLD discharge, 1 Pa, 1000 Watt side view. Magnetic coils powered at 141 – 144 – 117 A with the magnetic field overlaid.
Both diagnostic techniques are identifying an electron drift. In the Thomson scattering case, figure 5.5 dotted curves represent the Gaussian functions fitted to the data points. An interesting feature is observed by overlaying spectra taken at different pressures. The $\omega_c$ representing the central part of the Gaussian fit is not at the same position, rather it shifts towards the “red” or “blue” zone of the spectra at different discharge conditions. At pressures shown in Figure 5.5 the $\omega_c$ shift between them is of about 5 pixels. Since Thomson scattering accesses directly the electron velocity distribution function, a shift will correspond to an additional superimposed velocity component, and suggesting an anisotropic behavior of the distribution function. To quantify this shift a reference position has been taken. For this the mask was removed from the beam path and the scattered light from a wire placed inside the chamber located at the NL position is imaged onto the entrance slit of the spectrograph. At the exit plane i.e. on the CCD detector, the
respective spectrum is recorded. This spectrum is then fitted with a Gaussian function and the $w_e$ position determined. Finally $w_e$ is made to correspond with the central pixel position of the iCCD camera. Therefore the error in the pixel shift determination will be mainly due to uncertainty in setting this position. In the measurements presented in this work an error of 1 pixel is assumed. This will correspond in an error in the drift velocity of about $\pm 1.18 \cdot 10^4 \text{ m/s}$.

The electron drift dependence as a function of pressure and power for an NLD discharge are shown in figures 5.9 and 5.10 respectively. The drift velocity strongly increases with pressure decrease and has a linear dependence on the power variation and is good agreement with the above calculations. A higher drift velocity will be expected at lower pressures.

The RFMOS measurements taken from the side and from the bottom of the chamber are showing the same behavior, however since a line integration of the signal is used; a different value for the drift is obtained.

![Graph](image)

Figure 5.9: Drift velocity in an Argon NLD discharge at 1000 Watt as a function of pressure. The coils are powered at 141 – 144 – 117 A.
Figure 5.10: Drift velocity in an Argon NLD discharge at 0.1 Pa as a function of power. The coils are powered at 141 – 144 – 117 A.

Figure 5.11: Drift velocity as a function of magnetic field gradient in an Argon NLD discharge, 0.1 Pa, 1000 Watt
By changing the magnetic field gradient in steps of 10 % a relative linear change in the drift velocity is observed. At lower magnetic field strengths the drift velocity becomes smaller. This is in good agreement as well with the equation 5.35.

The average drift velocity is of about half the thermal velocity of a particle. In a linear quadrupol magnetic field configuration four symmetric regions where the maximum drift velocity is expected. However the NL structure has symmetry only for the upper and the lower regions. Changing the magnetic field configuration, the sign of the drift changes as well and plotted in Figure 5.12. A relative symmetry can be identified. The relative small different values between the two can be due to the effect of the earth magnetic field. In one configuration adds and in the other subtracts from the “total B field” seen by the electron. However for the RFMOS side measurements only one component of the drift velocity is accessible. The difference to the Thomson scattering measurements can be clearly seen.

Figure 5.12: Drift velocity as a function of pressure and magnetic field direction for an NLD discharge at 1000 Watt in Argon. The coils are powered at 141 – 144 – 117 A.
5.5. Phase resolved measurement of anisotropic velocity distribution.

5.5.1. Thomson scattering

As shown in figure 5.1 an intensified CCD camera (La Vision) with an S20 photocathode is placed in front of a window at the bottom of the chamber for emission measurements. The 2ns gate of the camera is also synchronized with the RF generator. For these measurements 5% of Neon was added to Argon feed gas in order to allow detection at $\lambda = 585.2 \text{nm}$ after collisional excitation ($e_{ex} = 19 eV$) of the short living ($\tau = 14.5 ns$) $2p_1$ state. A tunable filter (VariSpec) with a spectral width of about $\Delta \lambda = 10nm$ is set in front of the camera lens ($f = 50 \text{ mm}$) for spectral selection. The lens is focused to the surface of the antenna window with a fast aperture ($f = 1.8$) in order to minimize the field of depth.

Measurements of the electron velocity distribution function in Argon ICP discharge at a pressure of $p = 0.5 \text{ Pa}$ and an RF power of 1000 Watt is shown in figure 5.13 for two different phases ($t = 20ns$ and $t = 60ns$, relative timescale) differing approximately by $\pi$. The solid curves are Gaussian fits to the data points.

![Thomson spectra in an ICP discharge at p = 0.5 Pa at two different phases](image.png)

Figure 5.13: Thomson spectra in an ICP discharge at $p = 0.5 \text{ Pa}$ at two different phases
There are no data points at the center ±1.2 nm since there the THS notch filter is blocking the light passing the spectrometer. The displacement of the two spectra to shorter and longer wavelengths respectively is clearly shown. The re-centered velocity distribution function plotted on a logarithmic scale as a function of energy, the so called EEPF, is shown in figure 5.15.

![Figure 5.15: EEPF obtained from the same data](image)

The displacement used for re-centering is taken from the Gaussian fit. Further each data point is an average over twenty neighboring points from the original spectrum. Since the energy scale is non-linear with respect to the frequency or velocity scale, the distance between adjacent averaged data points varies between 0.2 eV at low energies to 0.5 at high energies. Both curves clearly show the same Maxwellian distribution function with a temperature of $T_e = 2.5 \pm 0.05 \text{eV}$. The dynamic range is approximately two orders of magnitude and the maximum energy, limited by the size of the detector, is about $\epsilon_\text{max} = 14 \text{eV}$, comparable to the ionization energy of Argon. Slight deviations from
a perfect Maxwellian at high energies are within the noise limits. However, stronger deviations at higher energies due to inelastic collisions or stochastic heating effects cannot be excluded in general. Measurements at other phases behave the same but are not included in the above plot for better visibility. The absolute electron density is \( n_e = 2.8 \times 10^{10} \text{ cm}^{-3} \).

The phase resolved displacement for the above conditions, pressure of \( p = 1 \text{ Pa} \) and RF power of 1000 Watt is shown in figure 5.15.

![Phase resolved oscillatory velocity](image)

Figure 5.15: Phase resolved oscillatory velocity obtained from the displacement of the Thomson spectra at \( p = 0.5 \text{ Pa} \) and \( p = 2 \text{ Pa} \).

The electron density at \( p = 2 \text{ Pa} \) is \( n_e = 1.6 \times 10^{11} \text{ cm}^{-3} \) and the temperature \( T_e = 1.8 \text{ eV} \). The solid lines are or a sine functions oscillating at the RF frequency of \( f = 13.56 \text{ MHz} \) with the amplitude, the phase and an offset being the fit parameters. The oscillation amplitudes at \( p = 0.5 \text{ Pa} \) and \( p = 2 \text{ Pa} \) are \( v = (13.9 \pm 0.5) \times 10^4 \text{ m/s} \) and \( v = (2.6 \pm 0.2) \times 10^4 \text{ m/s} \), respectively. The corresponding current density amplitudes are almost identical: \( j = 62 \text{ mA cm}^{-2} \) at \( p = 0.5 \text{ Pa} \) and
$j = 67 \text{ mA/cm}^2$ at $p = 2$ Pa. These measurements clearly demonstrate that the collective oscillation of electrons in an RF electric field can be measured by Thomson scattering with high accuracy. The lower velocity amplitude at the higher pressure is probably caused by an increased damping of the induced electric field due to the higher plasma density and the higher electron-neutral collision rate. For both pressures the collision rate is lower than the RF frequency and calculation of either the anomalous skin depth from Ishimaru's formula [91] or the collisionless skin depth yield values ranging from 1 cm to 3 cm. In any case, the field can be expected to drop over a distance of the same order as the distance of the laser scattering point from the quartz.

Similar measurements have been performed in an NLD discharge operated at 0.5 Pa and a RF applied power of 1000 Watt. The resulted electron velocity distribution functions at two phases (differing approximately by $\pi$) are shown in Figure 5.16. The displacement between the two phases is present as well, compared with the one in Figure 5.13.
However due to the strong presence of the drift in the Neutral Loop Discharge the oscillatory velocity does not pass through a 0th position. The respective phase resolved oscillatory velocity in shown in Figure 5.17.

Figure 5.17: Phase resolved oscillatory velocity in the NLD obtained from the displacement of the Thomson spectra at p = 0.5 Pa and 1000 Watt. The coils are powered at 141 – 144 – 117 Amperes.
5.5.2 RF MOS

The oscillatory velocity amplitude and its phase can also be measured by emission spectroscopy (RF-MOS) as outlined above. The basic concept is based on the fact that a velocity distribution function $f$ that is slightly displaced by a velocity $u$ can be expanded, where the dimensionless smallness parameter by which successive orders scale is

$$\alpha = \frac{u}{v_{th}} \ll 1,$$

with $v_{th} = \sqrt{\frac{2kT_e}{m}}$ being the thermal velocity or an equivalent velocity for non-Maxwellian distributions. The expansion has to be averaged over the full solid angle $\Omega$ since the excitation integral for atoms is an isotropic quantity. As demonstrated in Chapter 2 (2.3.3.3) the first order and all other odd order terms vanish and the second order term gives the dominant contribution depending on $u$:

$$\left\langle f\left(\left(\overrightarrow{v} - \overrightarrow{u}\right)^2\right)\right\rangle_{\Omega} = f(v^2) + \left\langle \frac{1}{2} \left( e_u \nabla_v \right)^2 \cdot f\left(\left(\overrightarrow{v} - \overrightarrow{u}\right)^2\right) \right\rangle_{\Omega} u^2 = f(v^2) \left[ 1 + \left( \frac{2e}{3kT_e} - 1 \right) \alpha^2 \right]$$

In the last step we have assumed $f$ to be Maxwellian, as is the case here. In ICP discharges $u$ can be composed of a potential constant drift and a harmonic term at the RF frequency $\omega$ and with a phase $\phi$. The square of $u$ then gives a constant, a term oscillating at twice the RF frequency, and a further term oscillating at the RF frequency itself. The latter term appears only if the potential constant drift is non-zero and pointing into the same direction as the oscillatory velocity:

$$u^2 = \left[ \overrightarrow{u} + \overrightarrow{\tilde{u}} \sin(\alpha \chi + \phi) \right]^2 = \left( \overrightarrow{u}^2 + \frac{1}{2} \overrightarrow{\tilde{u}}^2 \right) - \frac{1}{2} \overrightarrow{\tilde{u}} \cdot \cos(2\alpha \chi + 2\phi) + 2\overrightarrow{\tilde{u}} \cdot \overrightarrow{\tilde{u}} \sin(\alpha \chi + \phi)$$

If the measured temporally resolved fluorescence signal $I(t)$ from a state excited by electron collisions is normalized to the period $T$ averaged value, the resulting quantity $\eta$ is independent of the electron density, the atomic density, and transmission and sensitivity properties of the detection system. In general $\eta$ have two frequency components:
Here the additional phases are $\phi_n = \arctan(n\omega\tau)$, with $n = 1, 2$. In the above formula, damping due to the finite lifetime $\tau$ of the excited state is included by the square root terms. Further, in carrying out the excitation integral over the modulated part, as a simplifying assumption the energy factor $\varepsilon$ in equation 1 is replaced by the excitation energy $\varepsilon_{ex}$. More detailed numeric calculations actually show that this is a good approximation if $\varepsilon_{ex}$ is much larger than the thermal electron energy. From a Fourier analysis of $\eta$ then the oscillatory velocity, its phase, and also the potential drift can be determined. Atomic states used for RF-MOS should neither be excited from intermediate metastable states nor by cascading transitions from higher states and should have short lifetimes. This is the case for the Ne transition used in this work [16-17]. The result of such an analysis giving the oscillatory velocity and the phase is shown in Figure 5.18 and Figure 5.19 respectively for the case of $p = 0.5$ Pa.

![Figure 5.18: Oscillatory velocity amplitude obtained from emission spectroscopy (RF-MOS) at a pressure of $p = 0.5$ Pa](image)
Figure 5.19: Relative phase obtained from emission spectroscopy (RF-MOS) at a pressure of \( p = 0.5 \) Pa

In the analysis, the measured electron temperature of \( T_e = 2.5 \text{eV} \) was used. The velocity amplitude obtained at the same radius as the Thomson scattering is \( v = 12.4 \times 10^4 \text{ms}^{-1} \) which is very close to the value determined above. The phase is basically flat with a change of less than 1 rad over the entire cross section. In addition, the first harmonic component shows a small drift of about \( v = 0.5 \times 10^4 \text{ms}^{-1} \), a value close to the detection limit. Also in Thomson scattering, a drift of up to a few times \( 10^4 \text{ms}^{-1} \) and increasing with power was observed but again the value is too close to the detection limit to allow definitive conclusions presently. It should be noted that the emission measurement actually performs certain line integration along the axis vertical to the quartz window, although the optical field of depth was minimized by a fast aperture. In this respect the good coincidence with the value obtained from Thomson scattering seems a little lucky.
Comparison can be made with the induced electric field in vacuum as calculated by Maxwell's equations which yield after a few easy manipulations for the incremental electric field normalized to the current through the antenna coils:

\[
\frac{d E(r)}{I_{ant}} = \frac{\mu_0 \omega d \vec{r}}{4\pi r} 
\]

Here \( d \vec{r} \) is a vector element pointing along the coils. The total field is found by integrating numerically along all coil paths of the antenna geometry shown in Fig. 5.1. Very good agreement of the radial profile is found as shown in Fig. 5.20 and in more detail in Fig. 5.22.

Figure 5.20: Calculated induced azimuthal electric field in vacuum normalized to the current in the antenna horizontal cross section at an axial distance from the quartz of 2 cm. The square indicates the region where emission measurements were performed.

The reason for the good agreement despite the line integration problem becomes obvious by Fig. 5.21. This shows that the radial profile actually changes very little over the first few cm from the quartz which makes it insensitive to line integration. The
agreement with the vacuum profile is also consistent with the common assumption that
the dissipation of the energy of the RF field changes predominantly the axial and not the
radial profile [18].

![Graph showing electric field distribution](image)

**Figure 5.21:** Vertical cross section through the radial centre See figure 5.20.

If one assumes electrons to be accelerated only by the local field, the measured
velocity amplitude can be converted to an electric field amplitude of

\[ E = \frac{\tilde{n} m \omega}{e} = 0.67 \text{Vcm}^{-1} \text{ at } p = 0.5 \text{ Pa}. \]

First B-dot probe measurements give a value slightly smaller by about 15%. However, at present we are not sufficiently confident about the disturbance caused to the plasma by our probe and more detailed investigations are necessary. Especially, at the lower pressure one would expect stochastic non-local heating effects with a non-monotonous decay of the electric field strength [19-21].
Figure 5.22: Comparison of the oscillatory velocity profile obtained by RF-MOS (diagonal of figure 5(a)) with the calculated electric field profile (figure 5) at an axial distance from the antenna window of 2 cm. Left scale (●): measured oscillatory velocity amplitude, right scale (——): calculated induced azimuthal electric field in vacuum normalized to the current in the antenna.
Chapter 6

Wave phenomena in the Neutral Loop Discharge

6.1. Introduction

Magnetically enhanced low temperature plasma discharges, such as Electron Cyclotron Resonance plasma discharge (ECR) or Helicon discharge, are well known for their wave related phenomena’s. Particularly for these discharges the wave – particle interaction becomes an efficient collisionless electron heating mechanism, sustaining the plasma at very low pressures. However for this heating mechanism to be efficient the wave energy density must be higher than the electron energy density. Moreover the electrons should move close to, or at wave phase velocity for the resonant energy transfer to occur. In an opposite case the electrons will transfer their energy to the wave and thus not sustaining any longer the plasma.

The Neutral Loop plasma discharge uses as well a magnetic field for efficient plasma generation. However, until this date, investigations on the wave phenomena in an NLD discharge have been made in [96] from which this chapter is based. Previous studies on the effects of the magnetic field on inductively coupled plasma immersed in a weak magnetic field have been reported in [96, 97, 98]. Azimuthally symmetric pseudosurface and helicon wave propagation in inductively coupled plasma at low magnetic field was investigated in [99]. For plasma excited at 13.66 MHz, the characteristic frequency is between \( \omega_{ci} \ll \omega_{RF} \ll \omega_{ce} \ll \omega_{pe} \) where \( \omega_{ci} \), \( \omega_{ce} \), \( \omega_{pe} \) are respectively the ion cyclotron frequencies, electron cyclotron frequency and plasma frequency. The wave equation is solved in the following to obtain the dispersion relation for a TEM wave propagating parallel and perpendicular to the magnetic field. This gives actually the possible wave types which can propagate in the given conditions through the plasma.
Applying an electric field onto a medium will give rise to a charge separation. If this electric field has a time variation itself then the charge separation will also change in time. Therefore particle currents and displacement currents are produced. The displacement current and the particle current ads to act as the source of the magnetic field in the medium:

$$\nabla \times \mathbf{H} = \mathbf{J} + \varepsilon_0 \frac{\partial \mathbf{E}}{\partial t}$$  \hspace{1cm} (6.1)

Taking into consideration that $$\mathbf{B} = \mu_0 \mathbf{H}$$, $$\mu_0 \varepsilon_0 = \frac{1}{c^2}$$, $$\mathbf{J} = \sigma \mathbf{E}$$ equation 6.1 becomes:

$$\nabla \times \mathbf{B} = -\frac{i \omega}{c^2} \left( \mathbf{I} + \frac{i}{\varepsilon_0 \omega \sigma} \right) \mathbf{E}$$ \hspace{1cm} (6.2)

where \( \mathbf{I} \) represents the identity matrix. From equation 6.2 it follows that the dielectric tensor of the medium is defined by:

$$\varepsilon = \mathbf{I} + \frac{i}{\varepsilon_0 \omega \sigma}$$ \hspace{1cm} (6.3)

Equation 6.2 can be now written:

$$\nabla \times \mathbf{B} = -\frac{i \omega}{c^2} \varepsilon \mathbf{E}$$ \hspace{1cm} (6.4)

The wave equation is obtained from 6.4 by operating with $$\frac{\partial}{\partial t}$$. To eliminate $$\frac{\partial \mathbf{B}}{\partial t}$$ the Farady law is used: $$\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t}$$

$$\nabla \cdot \mathbf{E} - \nabla \cdot \left( \nabla \cdot \mathbf{E} \right) + \frac{\omega^2}{c^2} \varepsilon \mathbf{E} = 0$$ \hspace{1cm} (6.6)

Taking into consideration the simplest case, time dependent electric field is applied $$E = E \exp(-i \omega t)$$ and no magnetic field, the electron will respond with the velocity:

$$v = \frac{e}{i \omega m} E$$ \hspace{1cm} (6.6)

The current density associated with the corresponding electron density is
\[ J = -n_0 e v = -\frac{n_0 e^2}{i \omega m} E \]  

(6.7)

For \( E \approx \exp(i \mathbf{k} \cdot \mathbf{r} - i \omega t) \), and combining equations 6.6 and 6.7

\[ \left( \omega^2 - \omega_{pe}^2 - c^2 k^2 \right) E + c^2 \mathbf{k} \cdot (\mathbf{k} \cdot E) = 0 \]

(6.8)

Considering the case of transverse modes, \( \mathbf{k} \cdot \mathbf{E} = 0 \) and \( E \) not zero equation 6.8 becomes:

\[ \left( \omega^2 - \omega_{pe}^2 - c^2 k^2 \right) E = 0 \iff \omega^2 = \omega_{pe}^2 + c^2 k^2 \]

(6.9)

Equation 6.9 represents the dispersion relation. Figures 6.1 and 6.2 show the relation between \( k \) and \( \omega \) for waves that are propagating parallel and perpendicular to the magnetic field. In the parallel case the wave dispersion relation is given for the R and L polarized waves.

\[
\frac{c^2 k^2}{\omega^2} = 1 - \frac{\omega_p^2}{\omega^2} \frac{\omega}{\omega_c} - \text{L wave} \tag{6.10}
\]

\[
\frac{c^2 k^2}{\omega^2} = 1 - \frac{\omega_p^2}{\omega^2} \frac{\omega}{\omega_c} - \text{R wave} \tag{6.11}
\]

Yoshida et. all in [26] referring to the NLD, mention that no collective mode propagates perpendicular to the magnetic field at the operating RF frequency. This is correct as shown in Figure 6.1. However helicon (whistler wave, right hand polarized) waves can propagate along the magnetic field lines, shown in Figure 6.2. Helicon waves and helicon plasma discharges are widely treated in [100, 101, 102, 103, 104]. These waves are basically low-frequency bounded whistler waves propagating in an axially magnetized plasma column. If the wave vector \( \mathbf{k} \) would be at an angle \( \theta \) to \( \mathbf{B} \), the dispersion relation for the helicon will be given by [108]:

\[ k = \frac{\omega}{k_c} \frac{e n \mu_i}{B} \]

(6.12)

However, helicon sources are typically using particular antenna structures [106] to excite either \( m = 0, m = \pm 1 \) mode. Given the above helicon dispersion relation, the system is determinate, thus to get a stable solution one either fixes the wave number or
the electron density. A previous report on a helicon plasma source exited by a flat spiral coil is given by Stevenson et all in [106].

Figure 6.1: Wave frequencies propagating perpendicular to the magnetic field (Taken from reference [107])

Figure 6.2: Figure 6.1: Wave frequencies propagating parallel to the magnetic field (Taken from reference [107])
6.2 Wave propagation in the NLD

Floating probe measurements were performed in the NLD. The probe has a tip diameter of 0.6 mm and length of 10 mm connected to a rigid and temperature resistant RF cable. It can be freely moved inside the chamber for spatially resolved measurements from the side, or it can be mounted on a rotating wheel and mounted from the bottom view port.

![Diagram of X-Point Medium NLD]

Figure 6.3: Measurement configuration for the floating probe in NLD

Two dimensional profiles of the electric field amplitude and phase were measured by moving the probe vertically at various positions as shown in Figure 6.3. To generated the desired magnetic field configuration the coils were powered respectively 141 A – 144 A – 117 A. Figures 6.4 and 6.6 are showing the results for the phase and relative amplitude of the oscillating field at the RF frequency. Peaks in the amplitude occur when the probe crosses the separatrices as shown in Figure 6.4.
Figure 6.4: Measured wave amplitude for the above configuration

Figure 6.6: Measured wave phase for the above configuration
For the 7.7 position the probe intersects almost perpendicular the magnetic field lines. There is no visible change in the wave phase. The same behavior is shown by the measurements at the position 5.5. However the probe goes almost parallel with the magnetic field for few centimeters and this is seen by a phase change. The measurements done at the position 3.3 are scanning almost parallel the magnetic field line. The wave presence is now becoming clear and the phase velocity can be measured.

The above measurements exhibit a standing wave pattern and suggest the existence of a wave propagating along the magnetic field lines. Moreover side – view Phase Resolved Optical Emission Spectroscopy measurements shows also an interesting feature of the NLD discharge. A superimposed plot of the magnetic field lines over the $1\omega$ component is shown in Figure 6.6.

Figure 6.6: Calculated magnetic field lines superimposed on the Phase Resolved Optical Emission Spectroscopy measured $1\omega$ component in an NLD discharge. The coils are powered at 141 – 144 – 117 A., 0.1 Pa and 1000 Watt.

The modulation at the fundamental harmonic does not vanish, compared to the second harmonic which almost disappears. Within this boundary conditions Whistler
waves can be directly generated by the electric field coming from the antenna. However under normal conditions Whistler waves can not be observed with the Phase Resolved Optical Emission Spectroscopy since the energy of the electrons is not temporarily modulated. However if a drift occurs the symmetry is broke and electric field parallel to the drift leads to a modulation at the first harmonic.

Now if one plots again the magnetic field structure and with the above supposition, a clear separation of three distinct regions can be done as shown in Figure 6.7.

![Figure 6.7: Regions of the wave propagating in the NLD](image)

The three regions are separated by the points where the separatrices are intersecting the quartz. At small and large radii the magnetic field intercepts the quartz at almost normal to the surface and the poloidally induced electric field efficiently excites R waves. In the intermediate region (2) the magnetic field lines are parallel to the surface and the electric field enters into an evanescent X wave which penetrates into the NL region and is dumped there.
In frame of this picture the Whistler waves are traveling along the magnetic field lines without much of a dumping and their amplitude structure is shown in Figure 34. The probe at the position 33 is almost parallel to the separatrix between $z=100$ and $z=260$. At this position the magnetic field strength is of about $\omega_i = 1.6 \cdot 10^6 \text{s}^{-1}$ and the plasma density $\omega_p = 3.6 \cdot 10^6 \text{s}^{-1}$ are almost constant. Moreover the condition $\omega << \omega_i << \omega_p$ for a Whistler wave is satisfied and therefore one can estimate the phase velocity as:

$$v_{ph} = c \sqrt{\frac{\omega \cdot \omega_i}{\omega_p}} = 3.1 \cdot 10^6 \text{m/s}$$

This is in perfect agreement with the experimentally measured one of $3.4 \cdot 10^6 \text{m/s}$.

$$I_\omega \propto \vec{E}_{RF} \cdot \vec{u}_D$$

$$I_{2\omega} \propto E_{RF}^2$$

Figure 6.8: $1\omega$ and $2\omega$ components of Phase Resolved Optical Emission Spectroscopy in the NLD.

Bounded Whistler waves are Helicon waves which are obeying the R wave dispersion relation.

Helicon discharges are usually generated by certain antennas wound around a cylindrical tube and with a static magnetic field oriented along the tube axis [1].

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However, the induced azimuthal electric field from a flat coil antenna can easily be decomposed as local superpositions of L and R waves. With a static magnetic field perpendicular to the RF coil, L waves become evanescent and vanish after distances of the order of typically a few cm. The R wave propagates along the field lines and eventually turns into a Helicon wave due to the boundary conditions. The wave is guided by the magnetic field lines and the produced plasma stays confined [2]. Here we demonstrate Helicon operation in a setup originally designed for investigation of the Neutral Loop Discharge. The setup is described in [3]. A reduced schematic is shown in Figure 2. An additional RF antenna with 12 cm diameter and powered at 27 MHz is placed at the bottom, separated from vacuum by a window. The experimental setup is shown in Figure 6.9. The photos are taken from a lateral view port as shown. An optical calibration was done to match the magnetic field to the plasma.

Figure 6.9: Experimental setup for the helicon-type discharge.
In Figure 6.10 only the bottom antenna is powered with 1 kW. The discharge operates in Argon at 0.1 Pa. A magnetic field with an average strength of 10 mT is obtained with the three coils powered at 141A (top), 144A (middle) and 117A (bottom). Clearly the bright emission follows the magnetic field lines. Spectral measurements show that the blue light originates from Ar ions, although the near IR emission is more intense.

![Figure 6.10: Helicon-mode plasma generated with a flat antenna at the bottom; magnetic field lines and a length scale are superimposed on the right.](image)

The magnetic field configuration is superimposed and scaled to the taken photo. One can clearly see that the plasma is “bounded” by the magnetic field structure. However this is not the optimum configuration, but it proves the existence of Helicon Waves in the NL discharge.

In Fig. 6.11 the top antenna (20 cm diameter) is powered with 1 kW at 13.66 MHz in Argon at 0.1 Pa. The static magnetic field is provided by the bottom coil alone. In this
case the bottom window is replaced by a metal flange. Helicon waves require standing waves in radial and axial direction. Then the longitudinal wave vector must fulfil the condition $k_l L = m \pi$ with $m$ integer and $L = 0.63 \text{ m}$ the length along the field lines between antenna and the bottom. The dispersion relation for Helicon waves can be approximated by the whistler wave relation if the transversal wave vector is much smaller than the longitudinal wave vector. In this case one finds that $k_l^2 \sim \omega_p^2/\omega_c \sim P/I$.

![Figure 6.11: Helicon mode with the top antenna powered.](image)

In the latter step it was assumed that the plasma density scales linearly with the applied RF power $P$ and the electron cyclotron frequency linearly with the current $I$ in the static magnetic field coil. This gives with a proportionality factor $b$ and $m = m_0 + \Delta m$:

$$P = b(m_0 + \Delta m)^2 I$$
In fact it is found that matching with zero reflected power is only possible for certain fixed power to current ratios, as shown in Figure 6.13 for a fixed value of $m$. Various modes can be observed but for a given power matching is only possible within a certain current range. The slopes of the various modes are shown in Fig. 6.12. From a fit on derives $b = 0.013$ W/A and $m_0 = 9$.

![Figure 6.12: Slope of matching curves at different longitudinal modes versus relative mode numbers.](image)
Figure 6.13: Matching curve for $m = 10$

Maximum Landau damping should occur for $k_l \frac{v_{th}}{\omega} = 3^{1/2}$. For a thermal electron velocity of $v_{th} = 10^6$ m/s ($k T_e = 3$ eV) one calculates $m = 8$, which is in very good agreement with the measurement. For smaller $k_l$, the damping falls off quite drastically while it decreases more moderately for larger $k_l$. The $m_0 = 9$ mode corresponds to a wavelength of $\lambda_l = 12$ cm. For the transversal mode with $m_t = 1$, a much larger wavelength of about $\lambda_t = 60$ cm is calculated.

This is consistent with the above assumption on the ratio of the wave vectors. An estimate on the plasma density can be made by using the whistler dispersion relation for a given value of $k$. For a mean cyclotron frequency of about $\omega_c = 2 \times 10^9$ s$^{-1}$ at $I = 100$ A and $m = 9$ one finds a plasma frequency of $\omega_p = 4.4 \times 10^{10}$ s$^{-1}$, corresponding to an electron density of $n_e = 2 \times 10^{12}$ cm$^{-3}$. This density seems to be very reasonable and is about four times the value found in the NLD mode at the same pressure and RF power.
Chapter 7
Conclusions and further work

A new design of an RF Inductively Coupled Neutral Loop Discharge has been developed and built. A special configuration of the magnetic field coils was used to push the Neutral Loop (NL) from the middle coil plane allowing easy diagnostic access, both to the NL plane and to a possible remote processing region. This plasma chamber configuration, described widely in Chapter 2, allows operation of several discharges. With the magnetic field turned off, the discharge can be operated as a common ICP. Powering all the three coils, the operation in NLD mode is achieved. With only the bottom coil powered a Helicon discharge is operated. Thus a direct comparison between the plasma parameters of these discharges is straightforward made. However this versatility of the chamber allowed proving a possible different electron heating mechanism for the Neutral Loop Discharge. The results of this work can be summarized as follows:

- Detailed characterization of the plasma source was carried out using Langmuir probe measurements at different operation conditions. The measured EEDF at the NL exhibited a Maxwellian character at all conditions. This is of particular interest since the NLD can operate down to pressured of 0.05 Pa. Spatial resolved Langmuir measurements show a coupling between the NL region and a remote processing region. These regions are connected through the separatrices of the magnetic field. A feature of the NLD is that the electron temperature peaks in the center of the NL, whereas the maximum of the electron density is shifted radially inwards. This behavior is shown also in the previous investigation of an industrial reactor from which the design used in this work is based. A similar behavior of the electron temperature and density is apparent in the remote application region. Another factor is the influence of the magnetic
field strength on the inhomogeneity of the magnetic field. At lower
gradients, exhibits higher plasma uniformity; however the density is lower
due to a loss in the confinement strength.

- Thomson scattering measurements at the NL region confirm the same
  behavior of the plasma parameters. A Triple Grating Spectrograph
together with an ICCD camera was used to record the scattered spectra.
The obtained EEDFs exhibited a clear Maxwellian distribution up to
powers of 2000 Watt and pressures down to 0.05 Pa. However the EEDFs
exhibited a clear shift from a base defined center position. This shift can
be interpreted as a drift component in the direction of the scattering vector
and shows a high dependence with pressure. In the NLD this drift is
identified to be the diamagnetic drift. These are the first reported
measurements of drifts in low temperature plasma discharges with
Thomson scattering.

- The measured plasma parameters in the NLD with Thomson scattering
  have been compared with the results from an Optical Emission
  Spectroscopy (OES) novel method. This technique is based on the
  measurements of certain argon line ratios combined with a CRM model
developed for argon plasma kinetics at low pressures. An astonishing
coincidence of the plasma parameters was found. However this result is
achieved only by taking into consideration depletion of the neutral gas
density. In the NL region, the electron pressure can exceed the neutral gas
pressure, and in particular at relatively low gas pressures, resulting in a
depletion of the neutral gas density.

- Oscillating radio-frequency (RF) electric fields penetrating into plasma
  lead to a corresponding oscillation of the electron velocity distribution
  function. First time measurement of this oscillation by Thomson scattering
  in low-pressure ICP and NLD plasma where made. However due to the
  presence of the drift in the NLD these oscillations are substantially lower.

- A novel diagnostic technique, Radio Frequency Modulation Spectroscopy
  (RFMOS) was applied to measure the electron oscillation velocity and its
phase. The results obtained with RFMOS agreed well with the phase resolved Thomson scattering.

- For plasma excited at 13.66 MHz, the characteristic frequency is between $\omega_{ci} << \omega_{ke} << \omega_p$ where $\omega_{ci}, \omega_{ke}, \omega_p$ are respectively the ion cyclotron frequencies, electron cyclotron frequency and plasma frequency. In this regime, Whistler waves can propagate along the magnetic field lines. This thesis reports first time measurements of these waves propagating in the NLD.

- Wave propagation in the NLD leaded to two further developments
  - an electron heating mechanism for the NLD, through a wave – particle interaction
  - a planar type Helicon plasma Discharge

Further work is required to investigate the proposed electron heating mechanism in the NLD. Development and investigations on the planar type Helicon discharge can promote this discharge configuration as a suitable alternative to plasma processing due to its high versatility in controlling the plasma density and temperature.
Publications


References


Acknowledgements

First of all I would like to sincerely thank to my supervisor and mentor Prof. Uwe Czarnetzki for his support during this project both scientifically and financially. Prof. Czarnetzki contributed extensively on the theoretical development of the RF-MOS, on the fundamental investigations of the electron heating mechanisms in the NLD and on the development of the planar-type Helicon discharge. The ideas resulting from the many discussions on the themes tackled in this work gave me valuable insights and unveiled for me a new view on the world of plasma discharges. His mentoring approach on solving problems or finding new ways of looking into old ones inspired and helped me to develop myself further both as a scientist and as a human being. Prof. Czarnetzki is and will forever remain for me an example to follow.

I would like to thank to my second supervisor Prof. Brinkmann for the insights given during the corrections of this thesis.

To Prof. Akihiro Kono I would like to thank for teaching me the first steps on building and operating a Triple Grating Spectrograph in my one month stay as a visiting student in his group at Nagoya University in Japan.

Many thanks to Dr. Dirk Luggenhöschler, his support during the time of building the Triple Grating Spectrograph at the Ruhr – University is much appreciated. The first measurements on the wave phenomena in the NLD were made together with Dr. Tatsuo Ishijima. His work on this subject is much valued.

Many thanks for the support during this project to Dr. Timo Gans and Dr. Deborah O’Connell.

I would also like to thank to Dr. Dinescu George, Dr. Bogdana Mitu and all the members of my formal group at I.N.F.L.P.R., Magurele, Romania for my first steps in plasma physics. Many thanks to my colleagues and friends who contributed along the years of my Ph.D. to the accomplishments of this work:
• Martin Brennscheidt worked together with me during his diploma on building the Triple Grating Spectrograph used in this work. Without him the new alignment procedure who allowed the electron drift detection would not have been possible.

• Tobias Kampschulte developed in his early diploma work the Hall measurement system. He is responsible for the smart way 😊 on how the magnetic field lines are computed from the measured magnetic flux density.

• Viktor Kadetov developed the analysis of the Phase Resolved Optical Emission Spectroscopy data. His simulations on the induced electric field had been of real help on the investigations of the wave phenomena in the NLD. His contribution to this work here is much valued.

• Erik Kieft from the Technical University in Eindhoven visited during the building time of the TGS. The talks with him are much appreciated.

• Julian Schultze and Brian Heil are the two most valuable colleagues I ever had. Thank you for all the good time that we had in our Ph.D. years.

Many thanks to Iris Nikas for her support that so often helped me to better integrate into the EP5 group during my stay at the Ruhr University Bochum.

I would like to thank also to Volker Koster, Oliver Post and Frank Hütter for their support at MT AG during the time of writing this thesis.

Last but not least I would like to thank Thomas Zierow, Bernd Becker and Frank Kremer for their technical support with the experimental setup. Their never-ending good will of trying to teach me to make the technical things right paid out in wonderful results.

To all my formal colleagues a BIG Thank You! The times spend together in the lab, conferences or outside activities will always be remembered.

This thesis is dedicated to my parents and to my first love. Without their support and understanding this work would have never been possible. Thank you!

This research was supported by the Deutsche Forschungsgemeinschaft in the frame of the SFB 591 and GRK 1051.
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