An ITER-relevant magnetised hydrogen plasma jet

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An ITER-relevant Magnetised Hydrogen Plasma Jet

PROEFSCHRIFT

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Chapter 1

Introduction

The continued growth of the world population, and even more of the average standard of living, leads to a growing world energy demand. The annual world total primary consumption is more than 10000 Mtoe (megatons of oil equivalent) and expected to increase by a factor of 2 to 4 until 2100 ([1], chapter 2). Where today roughly 80 % of the energy demand is met by fossil fuel, it is clear that within a century the energy generation must essentially have been replaced by sustainable sources. Not only will fossil fuel run out (the world’s resources of fossil fuel are expected to last at maximum two centuries, with oil and gas having reserves for only a few tens of years ([1], chapters 3-5).), it is now clear that the emission of CO$_2$ associated with the use of fossil fuel bears unacceptable risks for climate and environment. Moreover, the global localization of the oil and gas reserves leads to global political tensions that are expected to become worse when the production is unable to follow the demand. The enormous, global problems associated with the growing energy demand make energy the defining political issue of the coming century. Resolving this issue relies to large extent on technological advances, especially in the development of sustainable energy sources that can be applied at large scale.

Nuclear fusion, a process in which light nuclei are merged under release of vast amounts of energy, is a safe, environmentally responsible and practically inexhaustible energy source that could give a substantial contribution to the future energy mix. As the development of fusion reactors started more than 50 years ago [2], clearly fusion is not a simple technology. The main problem is that the fuel, a mixture of the hydrogen isotopes deuterium and tritium, must be burned at a temperature of 150 million degrees. In present day fusion reactors of the so-called tokamak type temperatures like this and much higher can now routinely be sustained, by confining the hot fuel in a toroidal chamber by means of a magnetic field [3]. Of course, the fuel at these temperatures is fully ionized and forms a plasma.

In 2005 the agreement was reached on the construction of the next generation tokamak: ITER (Figure 1.1) [4, 5]. ITER will explore the scientific and technological feasibility of fusion energy, its target being the achievement of a
power multiplication factor of 10 (i.e. the power generated from the fusion reactions exceeds the power needed to sustain the high plasma temperature by a factor of ten) during pulses of 500 seconds or longer. ITER is a project of worldwide collaboration, carried by the participating parties Europe, Japan, the Russian Federation, the USA, China, South-Korea and India. ITER will be

![Figure 1.1: Schematical drawing of ITER (design). The system of magnetics windings surrounds the inner toroidal reactor chamber with divertor region in the bottom of the chamber.](image)

both a scientific experiment and a technology development and demonstration project. The two most important areas of scientific research to be tackled in ITER are the magnetic confinement of the hot plasma, and the exhaust of particles and energy from the plasma. For the predicted confinement in ITER there is a very well established theoretical and experimental basis, but ITER will take the plasma (in terms of dimensionless parameters) beyond the boundaries of parameter space that can be accessed by present day experiments. In particular,
in ITER the self-heating of the plasma by the alpha particles released in the
fusion reactions is essentially new and unexplored physics.

But possibly the largest uncertainty in ITER is that of the interaction of the
edge of plasma with the material wall. The heat generated in the plasma will be
conducted to the wall and leads to a very large power flux density on exposed
areas of the vessel. Moreover, there must be a continuous refueling of the
plasma, and the corresponding outflow of particles also constitutes a formidable
wall load. The fluxes foreseen for ITER are $\sim 10^{24} \text{ m}^{-2} \text{s}^{-1}$ plasma particles
at $1 - 10$ eV and energy fluxes of $\sim 10 \text{ MW m}^{-2}$ [4]. It is especially in the
power and particle fluences to the wall that ITER constitutes an extrapolation
of many orders of magnitude from present experiments. [4].

In present day experiments the exposed areas of the vessel are usually pro-
tected with cooled carbon tiles. For ITER, or fusion reactors beyond ITER, the
choice of wall material is very limited, with carbon and tungsten as the main
candidates, and beryllium as an option for areas that are not subjected to a
high heat load. The choice is so limited because under exposure to high heat
and particle fluxes the materials must not melt or erode (or if they erode, must
not introduce heavy particles into the plasma, since these are very detrimental
for the performance of the reactor), while they must also have an acceptable
life time under the neutron flux from the fusion reactions. It is not clear that
any material will satisfy these demands simultaneously. Therefore, the question
if the plasma-wall interaction processes can be controlled and manipulated in
such a way that erosion is minimised has come high on the research agenda for
ITER.

The erosion rates of carbon presently foreseen for ITER are a critical issue
for prolonged operation. The erosion itself is not the primary problem, but the
eroded carbon forms hydrocarbons that are redeposited. These deposits have a
high hydrogen content, i.e. there is significant codeposition of tritium which is
removed from the reaction cycle and cannot easily be recovered during operation
(this problem is usually referred to as tritium retention). Extrapolation of
present day tokamaks predicts that codeposits in ITER would reach the tritium
inventory limit of 350 g within 50 discharges [4, 6]. The complex processes of
erosion, migration and redeposition are not well understood. Progress in this
area requires detailed knowledge on the physics of sticking and re-erosion of car-
bon under ITER-relevant conditions. However, this information is not available
because these conditions are unprecedented.

It is not easy to obtain the physics understanding on plasma-surface in-
teractions from tokamak experiments. First of all, the primary site of carbon
erosion is the so-called divertor [7, 8] and its accessibility for \textit{in situ} diagnostics
is extremely limited. The alternative are \textit{ex situ} analyses, but these can be per-
formed only after venting the reactor and taking out parts of the interior wall.
This is usually only possible in between experimental campaigns that continue
for several months or even years. Consequently, only the integral effect of wide
ranges of experimental conditions on the reactor wall can be assessed.

Clearly, there is a need for magnetised plasma generators in which funda-
mental studies of plasma-surface interactions can be carried out in an accessible
laboratory experiment, under well-defined conditions matching those of ITER. This thesis addresses the development of such a plasma generator. In particular, it deals with the production of an intense hydrogen plasma jet and the transport of the plasma in a high magnetic field to a target.

1.1 Plasma-wall interaction in a fusion reactor

The magnetic field configuration of present tokamaks is such that the outer layer of the plasma flows into a special region in the reactor chamber, called "divertor" [7, 8] (Figure 1.2). Magnetic field lines form surfaces which are surfaces of constant temperature and pressure. The separatrix separates the hot core plasma (closed flux surfaces) from the low temperature divertor region in the bottom of the reactor: here the field lines cross the vessel wall, allowing plasma particles to leave the plasma.

Figure 1.2: Cross-section of the divertor region of a tokamak. The magnetic flux surfaces shown in projection are surfaces of constant temperature and pressure. The separatrix separates the hot core plasma (closed flux surfaces) from the low temperature divertor region in the bottom of the reactor: here the field lines cross the vessel wall, allowing plasma particles to leave the plasma.
is strongly inhibited by the field. With external coils a separatrix is induced, which separates the hot plasma core from the low-temperature outer layers near the wall. The field lines in this outer layer cross the vessel wall in the divertor. Helium (the "ash" of the fusion reaction) and impurities that originate also from plasma surface interaction are diverted into this zone following the magnetic field lines. Here the plasma particles collide with cold gas at higher pressure and subsequently radiate and cool down. It is possible to tune the conditions in the divertor such that the plasma is cooled to a temperature of 1-10 eV when it reaches the surface of the divertor. Such a regime of divertor operation is called the "detached regime". It minimizes the direct conduction of heat to the wall and the low impact energy of the particles prevents the potentially damaging physical sputtering. [4].

The particle and energy flux densities at the divertor plates foreseen for ITER represent an unexplored regime for the plasma facing materials. The physics understanding of processes at these conditions is lacking. The reference design of ITER has tungsten wall protection in the divertor, with carbon\(^1\) tiles in the places where the maximum heat load is expected (i.e. the strike zones of the diverted field lines). However, the data available on the interaction of such high fluxes of low-temperature hydrogen plasma with these materials is very limited. (In contrast, there are extensive databases on physical sputtering, obtained in early experiments e.g. [9]. This is relevant to the attached divertor operation, which is no longer the reference scenario for ITER). Moreover, the issue of tritium retention has been recognized as a major problem only in the last decade, and also on this issue the data is scarce. Hence, plasma surface interaction at these extreme conditions is of great interest and importance and needs to be explored with the highest priority.

1.2 Linear plasma generators for PSI

The FOM-Institute for Plasma Physics Rijnhuizen\(^2\) initiated the Magnum-PSI\(^3\) project to fill this niche. A linear plasma generator (called after the project, Magnum-PSI) will be built to provide a highly accessible laboratory experiment in which the interaction of a magnetised plasma (variable species mix) with surfaces of different materials can be studied in detail under the ITER-relevant conditions [10]. The design of the apparatus is presented in Figure 1.3.

Plasma parameters must be variable over a wide range, in particular covering the high-density, low-temperature conditions expected for the detached plasma regime in the divertor of ITER. A hydrogen plasma jet with a diameter at the target of around 10 cm and a particle flux density of $\sim 10^{24} \text{ m}^{-2}\text{s}^{-1}$ in a magnetic field of up to 3 T will be generated. With these parameters, the so-called strongly coupled regime of plasma surface interaction will be accessed.

---

\(^{1}\)specially developed carbon fibre composites - CFC
\(^{2}\)Partner in Triilateral Euregio Cluster (TEC)
\(^{3}\)MAgnetised plasma Generator and NUmerical Modelling for Plasma-Surface Interaction studies
Figure 1.3: Design of the Magnum-PSI set-up.
In this regime particles that come off the surface are trapped and remain part of the plasma surface interaction system, due to the combination of the ionization mean free path (short compared to the system dimension) and the small cyclotron radius in a high magnetic field. Moreover, the particle flux density is so high that the surface is modified to a depth of many atomic distances. Thus, the plasma in front of the surface and the top layers of the material form a strongly coupled system. This is the regime of plasma surface interaction that is typical of the divertor of a fusion reactor, and by virtue of the jet diameter of 10 cm and the high plasma density, this regime will be accessible in Magnum-PSI.

Several linear plasma generators are operational in the world (Table 1.1). The most well-known are PISCES (Plasma Interaction with Surface and Components Experimental Simulator) at the University of California, San-Diego [12], NAGDIS (Nagoya University Divertor Simulator) [20], PSI at the Max-Planck-Institute for Plasma Physics in Berlin [13, 14] and LENTA [15] at the Kurchatov Institute in Moscow. Studies at these devices have led to progress in the understanding of processes relevant for a tokamak divertor. Some phenomena were first discovered at linear plasma generators and later observed in tokamaks. For example, a detached regime in helium and hydrogen plasma as well as the appearance of plasma flow reversal were discovered and investigated at PISCES-A [16]. A series of plasma-surface interaction studies at moderate flux densities (see Table 1.1) was carried out at these linear apparatuses. To give an impression of the wide range of subjects investigated, we mention several of them. Measurements of erosion mechanisms from solid (carbon, tungsten) and liquid materials (gallium and lithium) were performed at the PISCES-B apparatus [17]. Experiments with hydrogen plasma at NAGDIS-II were devoted to the role of molecular activated recombination in the plasma detachment [19]. A series of experiments on the interaction of helium plasma with tungsten surfaces was conducted at the NAGDIS-I set-up [20, 21]. Studies at PSI-I focused on chemical sputtering of carbon based materials at high ion flux densities of deuterium plasma [22, 23]. Investigations on high-frequency and microwave radiation from the zone of interaction of hydrogen and helium plasma streams with neutral background gas targets were performed at LENTA linear plasma generator [24]. Also the development of a liquid lithium surface as a candidate for a reactor first wall [25] and imitation of deuterium plasma interaction with tungsten surfaces [26] and carbon materials [27] were carried out at this device.

All of the above mentioned plasma generators can produce hydrogen, deuterium and helium plasma with electron densities $10^{18} - 5 \cdot 10^{20}$ m$^{-3}$, flux densities in the range of $10^{21} - 10^{23}$ m$^{-2}$ s$^{-1}$, and operate in a magnetic field of 0.1 - 0.3 T. However, this does not cover the required range of parameters for ITER (indicated in the previous paragraph) at least by an order of magnitude in the flux density (Table 1.1). Moreover, in these devices, due to the type of plasma source, the higher fluxes are typically obtained at values of the electron temperature of tens of eV, rather than in the eV range.

Magnum-PSI will be unique in its ability to reach the high plasma density in front of the target in the temperature range typical of the detached divertor. The new experimental set-up that is being developed uses a high pressure plasma
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<td>10²</td>
<td>0.1</td>
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<td>6-15</td>
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<td>H₂, He</td>
<td>10⁻³</td>
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Table 1.1: Parameters of the different set-ups for PSI in the world
1.3. THIS THESIS: PILOT-PSI, A FORERUNNER OF MAGNUM-PSI

source (the cascaded arc [28, 29]) in combination with a strong magnetic field (up to B=3 T). In preparation of the design and construction of Magnum-PSI, a scaled-down experiment has been built at Rijnhuizen: Pilot-PSI. This thesis presents an analysis of the source performance and transport of the magnetised plasma jet in Pilot-PSI, and of the diagnostic equipment developed for these studies.

1.3 This thesis: Pilot-PSI, a forerunner of Magnum-PSI

As we have discussed, to study ITER-relevant PSI the Magnum-PSI set-up has to provide a well-defined high-flux density hydrogen plasma stream to a target in the presence of a high magnetic field. However, it is not easy to produce and transport a hydrogen plasma with such a high flux density at a temperature in the eV range, because at such low temperature the plasma vanishes quickly via molecular assisted recombination [18, 31, 32, 49].

Hence, to support the design of Magnum-PSI set-up a pilot set-up Pilot-PSI has been built within the project to investigate the following subjects:

- efficient production of high flux density hydrogen plasma jet
- plasma transport in high magnetic fields
- diagnostics to monitor the plasma parameters in Pilot-PSI and in the future machine Magnum-PSI

These are the central themes of the work presented in this thesis.

Pilot-PSI has been operational since 2001, and during this time a very successful development of the source and the transport and confinement of the plasma jet by means of a magnetic field has been carried out. Today, the plasma densities in the Pilot-PSI plasma jet reach world-record values for linear plasma generators, already achieving the target at a temperature of 2 eV. An overview of the Pilot-PSI results is presented in Sec. 1.4. Before turning to that, we discuss the applied plasma source.

For efficient production of a high-flux hydrogen plasma jet a DC cascaded arc plasma source (figure 1.4) has been chosen. The cascaded arc was introduced in the 1950s [28] and experiments and applications since have demonstrated that the arc can be operated in a wide range of pressures (10^4 to 10^8 Pa) and currents (5 to 2000 A) [29]. Operation of the arc on hydrogen is mainly explored by Schram and coworkers [31, 32, 49, 50, 82, 84, 89, 133]. The cascaded arc used at Pilot-PSI is based on the design of van de Sanden et al. [79, 92], which has been developed at the Eindhoven University of Technology for fast deposition and chemistry in a remote source configuration. In comparison with the plasma sources at other linear plasma generators, it operates at a higher pressure (0.1 –1 bar) and is able to produce a high-density (10^{20} – 10^{21} m^{-3}) hydrogen plasma with relatively low level of impurities present in the plasma,
CHAPTER 1. INTRODUCTION

Figure 1.4: The cascaded arc plasma source. The working gas enters the plasma source via the gas inlet (8). It is ionised in the arc channel (6) of the cascaded plates (3) due to the discharge between the cathodes (1) and the anode (7). The anode (7) is attached to the vessel wall (5). Plasma exits the source through the nozzle (4).

which is very important for plasma surface interaction studies. The high density in the source, much higher than in other sources, is necessary to reach the high flux density in the plasma jet. Moreover, due to the higher collision frequency of the particles in the source the plasma is in thermal equilibrium and is more homogeneous. For example a (low pressure) hollow cathode source, such as applied in e.g. the PSI-2 set-up leads to a hollow profile of electron temperature and density in the plasma jet [30].

The supersonic expansion of the neutral component from the source nozzle into the vacuum vessel, while the radial diffusion of the ionised plasma is limited by a strong magnetic field ($\sim 1$ T), can increase the ionisation degree significantly [93]. At the same time the energy dissipation from the jet as well as the plasma particle losses due to diffusion and recombination processes are reduced substantially. To separate the high-pressure source region from the low-pressure PSI region and to reduce recombination losses in the plasma jet due to collisions of plasma particles with cold background gas, differential pumping
1.4. OVERVIEW OF THE EXPERIMENTAL PROGRESS AT PILOT-PSI

can be applied. Finally, heating of the plasma jet can be achieved by drawing a longitudinal current in it, e.g. from the nozzle of the source to the target or a ring electrode [94, 95].

The important plasma parameters such as electron density, electron and ion temperature and plasma jet velocity were measured in the experiments by several diagnostic techniques: electric Langmuir probes, Thomson scattering and high-resolution spectroscopy. The diagnostics are described in detail in section 2.3.

In the following section an overview of the experiments conducted in Pilot-PSI and the main obtained achievements will be given.

1.4 Overview of the experimental progress at Pilot-PSI

Experiments in Pilot-PSI first focused on the optimisation of the cascaded arc plasma source for high flux operation. These experiments, aimed at maximizing the efficiency of the source, were carried out for values of the magnetic field varying from 0 to 1.6 T. Variation of the discharge channel diameter, from 1 mm to 4.5 mm, showed that the arc works better at bigger diameters. However, the current density and inflow of gas into the channel must be above certain minimal values for efficient and stable operation. Variation of the current in the arc channel showed increasing efficiency for increasing current density.

As hydrogen plasma was observed to recombine very quickly at the exit of the source [133], we experimented with an increased size of the nozzle diameter. This indeed led to a strong improvement of the efficiency of the source. Especially in combination with a magnetic field the effect of widening the nozzle was very pronounced. Melting of copper on the edges of the cascaded plates at plasma disruptions brought us to the idea of making the nozzle and source of tungsten (or an alloy of tungsten with copper that has higher thermal conductivity and is softer and easier to machine). The idea here is that a high surface temperature inside the source may be decreasing the wall passivation, which would therefore suppress the plasma recombination.

The application of a magnetic field of up to 1.6 T resulted in a spectacular change of the plasma jet. Whereas the hydrogen plasma, without magnetic field, extends only a few cm outside the nozzle a hazelnut-sized red glow with magnetic field a bright, finger-thick, confined plasma jet is obtained. The high brightness of this beam is evidence of a very high plasma density, as was confirmed by Thomson scattering measurements. The jet was studied with several diagnostic techniques, and turned out to exhibit a number of interesting phenomena. Among these are a very rapid axial rotation, a wobbling motion of the entire jet, and strongly asymmetric atomic line profiles. The analysis of these phenomena, presented in this thesis, led to a better understanding and description of the physics of the magnetically confined hydrogen plasma jet.

All together, these improvements made it possible to create the unprece-
dented hydrogen plasma jet with particle flux density of $\sim 10^{24} \text{ m}^{-2}\text{s}^{-1}$ at electron and ion temperature of 1-2 eV. With these results, the ITER-relevant regime of plasma-surface interaction can be accessed, and the foundation for the design and construction of Magnum-PSI has been laid.

1.5 Outline of this thesis

The experimental arrangement is presented in chapter 2. In this chapter we describe the Pilot-PSI set-up, the cascaded arc plasma source and its operational details as well as the plasma diagnostic techniques used for monitoring plasma parameters. The diagnostics used at Pilot-PSI are Langmuir probe, Thomson scattering on free electrons in the plasma and high-resolution emission spectroscopy on atomic lines. For each of them we give a short overview of methodology and the description of experimental set-up.

In chapter 3 we characterise the dependence of plasma parameters on discharge parameters such as discharge current, gas flow rate and on background pressure, in the absence of a magnetic field. The influence of the cascaded arc discharge channel diameter on the plasma production is also discussed in this chapter.

In chapter 4 the drastic change of the plasma jet due to confinement by a high magnetic field is described and discussed. The plasma diagnostics witness an increase of electron densities by at least two orders of magnitude and at the same time the electron and ion temperature remain around 1 eV or even grow up to 2 eV. Further, we report results on the effect of the nozzle geometry on the plasma jet parameters investigated by Thomson scattering. Measurements of the axial jet velocity derived from the Doppler shifts of atomic lines finalise this chapter.

In chapter 5 studies of the electron density, ion temperature and rotation velocity of the plasma jet by means of high-resolution optical emission spectroscopy are reported and discussed. Electron temperatures derived from the spectroscopic measurements are compared with results of Thomson scattering. A detailed mapping of the electron density and ion temperature by means of high-resolution optical emission spectroscopy in the magnetised plasma jet are also presented in this chapter.

The final chapter presents an overview and general discussion of all the results presented in this thesis.
Chapter 2

Experimental arrangement

The results described in this thesis were obtained with experiments performed on the Pilot-psi set-up. The experimental lay-out of Pilot-PSI, the cascaded arc plasma source and its operational details and the diagnostics used to measure plasma parameters are described in detail in this chapter.

2.1 Pilot-PSI instrumental layout

A schematic drawing of the Pilot-PSI set-up [70, 71] is shown in Figure 2.1.

Figure 2.1: Schematic drawing of the experimental setup Pilot-PSI.
A cascaded arc plasma source is mounted on a face plane of a cylindrical vacuum vessel of stainless steel. The vessel is 1 m long and 40 cm in diameter. A water-cooled copper target is installed at another side of it. Five coils are evenly distributed along the vessel to create an axial magnetic field inside.

The high pressure difference ($\sim 10^4$ Pa) between the gas inlet and the plasma source exit forces the produced plasma to expand supersonically through the nozzle into the vacuum vessel. The vessel is evacuated to $10^{-2}$ Pa (the base pressure) by a 3-stage pumping system. A forepumping is done by a Balzers pump DUO100 with a pumping speed 1 m$^3$/hour. A mechanical booster pump EH-500A (Edwards High Vacuum International) is installed in series as second stage pump. A EH-4200 mechanical booster pump with a pumping speed up to 4200 m$^3$/hour evacuates the vessel via a bent duct of 40 cm in diameter. The pump is protected by a small-meshed grid which limits the pumping speed to 3600 m$^3$/hour. The pressure in the vessel can be varied between 2 and 200 Pa by controlling the rotation velocity of the EH-4200 pump. The pressure depends also on the gas flow rate through the plasma source (Figure 2.2). Pressure is monitored by two gauges installed at the plasma source inlet (a membrane gauge PRAD D005.S70.C210 by Baumer sensopress, 1–1000 mbar) and at the end of the vacuum vessel (two Baratrons by MKS Instruments Inc: a 310BHS-1000 for 1-1000 mbar range and a 370HA-00001 for the range below 1 mbar).

Oil-cooled Bitter coils generate a magnetic field of up to 1.6 T inside the Vessel. The operational time is limited by the cooling of the coils, varying from 3 minutes at 0.4 T down to 3 s at 1.6 T. At 0.2 T continuous operation is possible. The current through the coils is 14.2 kA at 1.6 T. The calculated distribution of the magnetic field within the set-up for 14.2 kA current is plotted in Figure 2.3 [66]. Note that only a half of the vessel length is shown as the field is mirror-symmetric with respect to the central coil. It is seen that in the centre
of the vessel, where the plasma jet is located, the field varies monotonically along the axis of the vessel, there are no local minima. The field strength at the position of the source and the target (i.e. in the centre of the first and the last coils) is 20% lower than the maximum value. The magnetic field created by a single coil was measured directly (Figure 2.4) and the total magnetic field of all five coils was calculated on the basis of that (Figure 2.4). The measured axial magnetic field distribution is in a good agreement with the calculated one.

Five quartz windows are mounted at each side along the vessel for diagnostic access. The first and the last windows are rectangular (7x16 cm), and the others are round (10 cm in diameter). On the top and bottom of the vessel there are five circular ports with standard flanges assigned for installation of diagnostic equipment (probes etc.).

Controllers of vacuum pumps, gas flow-controllers, magnetic field system and controllers of the cascaded arc power-supply are connected to a Programmable Logical Controller (PLC) system (Siemens) which is in turn connected to a PC. This PC is operating on Windows NT 4.0 and uses the Wonderware FactorySuite 2000 package (Wonderware Corporation) to control the equipment via the PLC system, taking into account security schemes.
CHAPTER 2. EXPERIMENTAL ARRANGEMENT

(a) axial magnetic field of a single coil

(b) total axial magnetic field of five coils

Figure 2.4: The measured axial magnetic field of a single coil (a) and the total axial field of five coils (b) are presented.
2.2 Cascaded arc plasma source

The cascaded arc plasma was introduced by Maecker in 1956 [28] for spectroscopic studies in near equilibrium plasmas. At the University of Kiel, detailed studies were made among others on hydrogen line broadening by Helbig and coworkers [33]. In Eindhoven, the cascaded arc plasma was used to study deviations from thermal equilibrium [34]. Later, very high densities and temperatures were reached in pulsed operation of the cascaded arc with current pulses up to 2 kA and pressure pulses up to 14 bar [35]. The effect of these simultaneous pressure and current pulses was a rise in electron temperature from $\sim 1$ eV to $\sim 2$ eV and in electron density by more than a factor of 10 from $10^{23}$ to $10^{24}$ m$^{-3}$ at the source exit. This work was especially aiming at reaching a non-ideal plasma regime. In the work of the Haas et al. [36], the cascaded arc was for the first time used as a source of particles rather than photons, either to reach very high deposition rates [37, 38, 40, 41], to etch [42] or to modify the surface in another way [43, 44]. At the same time the dynamics and kinetics of the cascaded arc were studied in great detail: source development [45], expansion characteristics [46, 47, 48], and the influence of hydrogen [31, 49, 51] and nitrogen [52, 53, 54]. One particular hydrogen source related subject was the restoration and preservation of archeological artifacts [55].

The characteristics of the arc are the following: It can be operated in a wide range of pressures ($10^4 - 10^8$ Pa) and currents (5–2000 A) and can produce plasma with high electron densities ($10^{21} - 10^{24}$ m$^{-3}$ in argon plasma and $10^{19} - 10^{22}$ m$^{-3}$ in hydrogen plasma) at relatively low electron temperatures of $\sim 1$ eV [82, 92]. Another advantage of a cascaded arc in comparison with other arc discharges is that it is a wall-stabilised plasma source. The stability of the plasma parameters is usually within 15% or better [82].

2.2.1 Cascaded arc design at Pilot-PSI

In Pilot-PSI we use a design of the arc developed at the Eindhoven University of Technology [65] with some modifications for operation on hydrogen [82]. It consists of a cathode chamber with three cathodes, a set of copper plates 5 mm thick with a central discharge channel (nominally) 4 mm in diameter, and an anode plate with a nozzle (Fig. 2.5). All components are water-cooled. The 2 mm diameter cathodes are made of thoriated tungsten and are placed on a circle at 120° from each other and 45° with respect to the axis of the channel.

The plates are electrically insulated from each other, from the cathode part and from the anode plate, by boron-nitride plates of 1 mm thick. Rubber O-rings provide the vacuum sealing of the arc. The working gas (Ar, H$_2$) continuously flows into the arc at a pressure of $(1-2) \times 10^4$ Pa and is ionised there. Usual discharge parameters are: a gas flow rate of 1.5–3.5 slm (standard liters per minute, 1 slm corresponds to $4.5 \times 10^{20}$ particles per second), but can be up to 13.5 slm, and the arc current 30–100 A. The upper limit is determined by the power supply. Most experiments in Pilot-PSI were carried out with a power supply that could deliver up to 100 A, at up to 700 V, with a maximum power.
2.2.2 Plasma expansion from a cascaded arc

In the Pilot-PSI set-up a hot, partly ionised gas from subatmospheric pressure (10-20 kPa) expands from the nozzle of the plasma source into a low-pressure (1-10 Pa) vacuum vessel. In the absence of a magnetic field this high pressure difference leads to acceleration of the flow to supersonic velocities [117]. The expansion pattern in a similar device with the same plasma source working on argon has been studied by van de Sanden [79]. Peculiarities of the partly
ionised hydrogen expansion involving ro-vibrationally excited molecules were investigated by M. de Graaf [82], Zhou Qing [83], R. Meulenbroeks [84], S. Brussaard [85], P. Vankan [89] and others. It is important to investigate the expansion region as it influences the formation of the plasma jet. Moreover, the interaction of a magnetised (nonexpanding) plasma jet with the expanding neutral component will also have consequences for the temperature and density of the jet particles.

It is common to express the ratio between the flow velocity $v$ and the sound speed $c_s$ at given conditions in the Mach number $M$:

$$M = \frac{v}{c_s} = \frac{\sqrt{\gamma k_B T}}{m} \quad (2.1)$$

Here, $\gamma$ is the specific heat ratio $c_p/c_v$, $k_B$ is the Boltzmann constant, $T$ is the temperature of the gas with atomic mass $m$. At the same time the density and the temperature of electrons and ions decrease over the expansion. According to the description of expansion in [46] the density drops in the axial direction approximately as:

$$n(z) = n_0 \frac{1}{1 + z^2/z_{ref}^2} \quad (2.2)$$

where $z_{ref}$ is determined mainly by the source nozzle geometry and it is of the order of the radius of the source outlet. At some point downstream the pressure of the expanding gas becomes even lower than the background pressure in the vessel. This is the so-called valley of the expansion. The temperature of particles in the expanding flow drops as the thermal energy transforms into the kinetic energy of the directed collective movement of particles. The temperature decrease can be described to good approximation as adiabatic cooling.

Colliding with the background gas in the vessel, the supersonic flow forms a normal stationary shock front that is the front boundary of the barrel shock structure [118] (Figure 2.6). The position of the shock front $z_M$ depends on the pressure in the source, the nozzle diameter and the downstream pressure [119]. The expression can be rewritten in terms of the gas flow rate $\Phi$ (in standard m$^3$s$^{-1}$), the background pressure $p_{back}$ (in Pa), atomic mass number $A$ and the stagnation temperature $T_0$ (in K) [46]:

$$z_M = \sqrt{2.5 \Phi \frac{(1 + \gamma/2)}{\gamma^{1/2}} \frac{T_0^{1/2} A^{1/2}}{p_{back}}} \quad (2.3)$$

Here, $\gamma$ and $T_0$ are the specific heat ratio and the stagnation temperature respectively. The value of $\gamma$ that should be used in the case of plasma is still being discussed. Opinions vary from $\gamma=1.2$ [120] for a wide range of ionisation degree (0.03–0.3) up to $\gamma=5/3$ as for an isentropic flow of a monoatomic gas. Using Thomson scattering $\gamma_e=1.3$ has been measured for the electron gas [121], and Laser Induced Fluorescence (LIF) gave a value of $\gamma=1.3$ for atomic radicals [50]. In this thesis we will take $\gamma$ to be equal to 5/3, following [123].
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Figure 2.6: The scheme of gas flow supersonically expanding from a nozzle with the stationary shock front formation.

It is more convenient to use the expression 2.3 in terms of flow rate $\Phi$ in practical sccs units (standard cubic centimeters per second), source temperature at the axis $\hat{T}_s$ in eV, background pressure $p_{\text{back}}$ in Pa and atomic mass number $A$ as is done in section 3.2 of [89]:

$$z_M = 1.8 \times 10^{-2} \sqrt{\frac{\Phi}{p_{\text{back}}}} \sqrt{A\hat{T}_s}$$  \hspace{1cm} (2.4)

The width of the shock front is of the order of the mean free path of gas particles. Within the shock the flow velocity drops below the speed of sound and the temperature increases due to the collisions. The flow after the shock front is subsonic and does not diverge downstream, as the pressures in the flow and outside the flow are equalized. Thus the diameter of the jet depends on the background pressure, because this determines the position of the Mach disc and thus how far expansion continues.

There is always a recirculation of hydrogen in the expansion vessel ([82], section 3.4). Atomic hydrogen goes to the wall and associates there with atoms of hydrogen that have already accumulated there (the sticking probability can be up to 1 on a clean metal surface, and is between $10^{-2}$ and $10^{-1}$ for a passivated metal surface). As a result rotationally-vibrationally excited molecules may come off the surface [122] and a fraction of them diffuse back into the plasma jet. The presence of rotationally-vibrationally excited molecules in the plasma has consequences for the population of the $n = 4$ energy level and on the subsequent $H_\beta$-line radiation as will be discussed in section 5.2. This is an important issue, because the analysis of the $H_\beta$-line profile is used as a diagnostic tool.
to determine the ion temperature, the electron density and the jet rotational
velocity.

In a magnetic field, when plasma particles are confined, the expansion of the
plasma jet is limited while the neutral gas expands freely. The questions whether
plasma might be accelerated to supersonic velocities and how it interacts with
the shock front that is formed due to the expansion of neutrals, are still under
debate.

2.2.3 Operational details of the cascaded arc

In the experiments with the cascaded arc plasma source we worked at different
gas flow rates (1–3.6 slm), discharge currents (30–100 A) and we changed some
dimensions of the plasma source such as the bore of the channel, which was
varied from 1.5 to 4.5 mm (section 3.4).

The effective diameter of the plasma in the discharge channel is smaller
than the channel diameter because of the low temperature of the gas close to
the (water-cooled) walls. This leads to a peaked temperature, which in turn
leads to a peaked current density profile. Because the thermoconductivity of
hydrogen is higher than that of argon, the diameter of the plasma channel is
smaller with hydrogen operation ([82], p. 21 and [83], section 3.2.4).

An interesting peculiarity of the plasma source working on hydrogen is that
the discharge voltage decreases with increasing discharge current (Figure 2.7).
This fact can be explained by increasing cross-section of plasma in the discharge
channel while the conductivity $\sigma$ of the plasma remains constant and depends
on the electron temperature $T_e$ (in eV) [72] (Spitzer):

$$\sigma = \frac{2 \cdot 10^4 T_e^{3/2}}{\ln(\Lambda)},$$

with $\ln(\Lambda)$ is the Coulomb logarithm. $\Lambda$ is 9 times the number of charged
particles in the Debye sphere (2.6).

Without a magnetic field the arc voltage ranges from $\sim 40$ V for a pure argon
plasma to $\sim 140$ V for a pure hydrogen plasma. In a magnetic field the voltage
for a hydrogen plasma rises up to $\sim 220$ V. Typical potential distributions in
the source for the case of hydrogen are shown in the Fig. 2.8.

M. de Graaf [82] found that quick erosion of the tungsten cathodes can be
avoided by operating the arc a few of minutes in argon before switching to
hydrogen. This appears to passivate the cathode material.

As cathodes burn out with time, their tips gradually become further away
from the discharge channel. At some point they are so far from the discharge
channel that stable operation of the source is inhibited, and new cathodes should
be installed. Cathode tips must be carefully aligned and positioned as close
as possible to the axis of the discharge channel. Replacing the cathodes is a
minor operation, which takes about 30 minutes including 10 minutes to fill the
vessel with nitrogen to the atmospheric pressure and 10 minutes to pump it out
afterwards. Pilot-PSI is a very flexible and easy to handle device, that can be
made ready for experiment within 10–15 minutes after a complete shut-down.
2.3 Plasma Diagnostics

It is very important that the plasma that impinges on the target can be characterized accurately. The important parameters in the interaction region are the electron density \( n_e \) and temperature \( T_e \), and of course the composition. The electron temperature and density determine many processes, such as the balance between ionisation and recombination, the flux density to the target (through the Bohm criterion) and the sheath potential (and thereby, the impact energy of the ions), the collisionality, mean free paths for collisions and ionisation.

There are several methods to measure \( n_e \) and \( T_e \). We have applied Lang-
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Figure 2.8: A potential distribution over cascaded plates in hydrogen plasma. In a magnetic field the voltage over the arc is significantly higher. The extra voltage drops mainly between the anode and the first plate. The arc parameters are: gas flow rate 2.5 slm, $I_{arc} = 80$ A.

muir probes, Thomson scattering, and optical emission spectroscopy (the Stark broadening of atomic lines or radiation continuum analysis can be used) [60]. Each of them has certain characteristic ranges of application.

A Langmuir probe, in its basic form an electrode that is inserted in the plasma, is technically easy to implement and does not require expensive equipment. However, the theory of electrical probes requires several assumptions about plasma parameters and thus the interpretation of probe data is not always straightforward. This is especially so if the measurements are done in a magnetised plasma, where the flow of plasma particles to the electrode may be influenced by the field. Moreover, a probe inserted into the plasma disturbs it. Also the reverse can be true: a probe does not survive long in a high energy density plasma. In our experiments, the application of the probe in the magnetized plasma beam was therefore strongly limited.

Thomson scattering requires expensive equipment: a high-power pulsed laser and sensitive light-detectors. But it gives direct information about the free electron distribution function in the plasma, and hence $T_e$. The analysis does not require any assumptions on the plasma. Only a relative calibration is required to deduce $T_e$, while an absolute calibration will also give $n_e$. The observed photons come from a volume defined by the intersection of the laser beam and the viewing optics, which allows the measurements to be strongly localized. This is an advantage in comparison with the line-of-sight integrated passive emission spectroscopy. Thomson scattering normally does not disturb the plasma. Its main drawback, apart from the investment cost of the equipment, is the low cross-section of Thomson scattering. In our application, this meant that the resulting detection threshold makes Thomson scattering measurements only
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well possible if the plasma beam is confined by a magnetic field. Furthermore, Thomson scattering calls for a well collimated laser beam and a view line perpendicular to it. This implies a requirement on diagnostic access which can be limiting.

High-Resolution Optical Emission Spectroscopy (HiRES) is a good method to determine the electron density, ion temperature and velocity components of a plasma jet. A line profile can be analyzed with respect to the Stark broadening, the Doppler broadening and the Doppler shifts of the spectral lines at once. If the Stark width of the line is of the same order as the Doppler width, then the electron density and the ion temperature can be obtained simultaneously. Like Thomson scattering, HiRES does not disturb plasma. In our experiments, the intensity of the atomic hydrogen lines is usually orders of magnitude higher than the Thomson scattering signal (section 2.3.3), which makes the detection of signals much easier. Moreover, as only a single view line is required for this measurement, diagnostic access is much simpler than for Thomson scattering.

In summary, the three methods applied are complementary. The absolute values of $T_e$ and $n_e$ measured by Thomson scattering provide a very important calibration of the Langmuir probe and HiRES measurements. These, in turn, can be applied at lower density and (HiRES) without the spatial limitations of Thomson scattering. Thomson scattering, again, provides the best spatial resolution. Together these diagnostics cover the full range of measurements on the plasma that were needed for the studies described in this thesis. Apart from these, there are of course the simple measurements of the voltage applied to the source, the discharge current, etc.

2.3.1 Basics of Langmuir Probe

One of the simplest techniques is an electrical Langmuir probe. The original theory was developed already in 1926 by Langmuir [73]. In the most general case, a Langmuir probe is just an electrode inserted into a plasma, which can be electrically biased with respect to the plasma potential and due to that collects charged particles from the plasma. (See, for example, [58, 59]). Information on the electron velocity distribution function and density is revealed in the relation between the measured flux (current) to the probe and the bias potential. From this relation, $n_e$ and $T_e$ can be derived (see below). It is obvious that we can speak about the electron temperature only if the velocity distribution function is Maxwellian, e.g. the electrons are in local thermal equilibrium. Further, the electron density in the plasma must be equal to the ion density (the plasma is quasineutral). The accuracy of the Langmuir probe measurements, performed at expanding plasmas similar to those in Pilot-PSI, as well as the correct interpretation of the data has been established by Brussaard [85] using the Theory of Sperical and Cylindrical Langmuir Probes in a Collisionless, Maxwellian Plasma at Rest by Laframboise [74] and results of Peterson and Talbot [75].

Electric probes require only a minimum of simple equipment and provide a measurements that is local. The spatial resolution depends only on the probe dimensions. These, in turn, are constrained by the Debye radius, i.e. the
characteristic length of the local charge separation in the plasma ([59], pp. 25-30:

\[ r_D = \sqrt{\frac{\epsilon_0 k_B T_e}{n_e e^2}} \approx 7.4 \cdot 10^3 \sqrt{\frac{T_e}{n_e}} \] (2.6)

Here, \( \epsilon_0 \) is the dielectric permittivity, \( k_B \) is the Boltzmann constant, \( e \) is the charge of electron, and \( \hat{T}_e = k_B T_e/e \) is the electron temperature expressed in eV.

Another important length scale for the Langmuir probe theory is the mean free path of charged particles, i.e. the average distance a particle can move without a collision. The mean free path for different sorts of collisions (electron-ion and ion-ion) can be estimated using [76]:

\[ \lambda_e \approx \frac{\lambda_{ee}}{\sqrt{2}} = 1.4 \cdot 10^{17} \frac{T_e^2}{n_e \ln \Lambda} \] (2.7)

\[ \lambda_i = 2.04 \cdot 10^{17} \frac{T_i^2}{n_e \ln \Lambda} \] (2.8)

where \( \ln(\Lambda) \) is the Coulomb logarithm. In our expanding and recombining downstream hydrogen plasma with electron and ion temperatures in the range of 0.1 to 0.2 eV and electron densities from \( 10^{16} \) to \( 10^{17} \) m\(^{-3} \), the Debye radius is around \( 10^{-5} \) m while the mean free path at these conditions ranges from \( 10^{-3} \) to \( 10^{-1} \) m and is mainly much larger than the Debye radius. In the case of the magnetised plasma jets produced in Pilot-PSI, the temperatures are, at 1-2 eV, an order of magnitude higher, while the densities increase by up to four orders of magnitude, up to \( 7 \cdot 10^{20} \) m\(^{-3} \). In these conditions the Debye length and collision mean free paths are even shorter: \( r_D \approx 3 \cdot 10^{-7} \) m and \( \lambda_{e,i} \approx 5 \cdot 10^{-5} \) m.

The main disadvantage of the probe technique is the fact that a probe disturbs the plasma because of its physical presence in the plasma. Therefore the probe should be made as small as possible but at the same time must satisfy the condition \( R_{\text{probe}} \gg r_D \). For the probe used at Pilot-PSI (see below for details) this condition is well met but for the plasma edges, where the density falls below \( 10^{15} \) m\(^{-3} \). Here the basic assumption for the application of the probe is not valid, which may lead to incorrect values of \( T_e \) and \( n_e \) (see section 3.2).

A single Langmuir probe can be applied to determine the plasma floating potential as well as \( T_e \) and \( n_e \). However, in an electrodeless downstream plasma there is no proper reference potential, which makes the measurement hard to interpret. For such cases a double Langmuir probe has been proposed [61]. In a double probe a sweeping voltage is applied between two identical electrodes and the current in the circuit is measured. The total current to the system can never be greater than the ion saturation current, since any electron current to the total system must always be balanced by an equal current of ions. Thus a measured voltage-current characteristic (or Volt-Ampere Characteristic - VAC) of a double probe is usually symmetric except for the cases when the electrodes
are not identical. A typical voltage-current characteristic of a double Langmuir probe measured at Pilot-PSI is shown in Figure 2.9.

Figure 2.9: A typical voltage-current characteristic of the double Langmuir probe.

The distance between the two electrodes should be as small as possible for the measurements to be local but to prevent an overlapping and an interference, it should be at least two times larger than the sheath around the electrodes, i.e. a few Debye lengths. The slope of the VAC at zero voltage gives a measure of the electron temperature in electronvolts $k_B T_e/e$ ([58], pp. 178-183):

$$\frac{k_B T_e}{e} = \left(\frac{dI}{dV}\right)^{-1} \frac{I_{\text{ion.sat}}^+ - I_{\text{ion.sat}}^-}{I_{\text{ion.sat}}^+ + I_{\text{ion.sat}}^-}$$

(2.9)

Here $k_B$ is the Boltzmann constant, $T_e$ is the electron temperature in K, $e$ is the charge of electron, and $I_{\text{ion.sat}}^+$ and $I_{\text{ion.sat}}^-$ are the ion saturation currents in the VAC. The ion saturation current can be also used to calculate the ion density $n_e$ from the expression

$$I_{\text{ion.sat}} = jA_s = kn_e eA_s \sqrt{\frac{k_B \cdot T_e}{M_i}}$$

(2.10)

Here, $j$ is the current density to the probe, $A_s$ is the area of the probe in $m^{-2}$, $k$ is the factor that depends on the temperature ratio $T_i/T_e$ [59] and for $T_e = T_i$
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The electron temperature is obtained from the VAC slope in the vicinity of zero voltage. The density then is:

\[ n_e = \frac{2 \cdot I_{\text{ion sat}}}{eA_1 \sqrt{\frac{k_B T_e}{M_e}}} \]  

(2.11)

2.3.2 Probe measurements at Pilot-PSI

A double Langmuir probe was installed at about 30 cm from the nozzle (Figure 2.10), in the second diagnostic port. It could be moved over 130 mm perpendicularly to the plasma axis. Thus, radial profiles of \( T_e \) and \( n_e \) in the plasma jet could be measured. It was not possible to place the probe closer to the source, in the first diagnostic port, because there it would be in the region of the shock front (see chapter 2.2.2) and would strongly disturb the jet pattern.

![Figure 2.10: A double Langmuir probe was installed at about 30 cm from the nozzle in the second diagnostic port and could be moved perpendicularly to the plasma axis within 130 mm to obtain radial profiles of the electron density and temperature in the plasma jet.](image)

The probe consists of two cylindrical tungsten wires of 4.5 mm length and 0.2 mm in diameter. A digital SourceMeter-2400 (by Keithley) was used to apply
CHAPTER 2. EXPERIMENTAL ARRANGEMENT

a sweeping voltage between the wires and at the same to record the measured current to the probe. The data were transferred to a PC via GPIB interface for further treatment of the obtained VAC. The sweeping of the voltage as well as the data transfer via GPIB usually took 10–20 seconds depending on the voltage range (usually -7..+7 V) and the number of steps in the range (usually ∼ 200).

The obtained VACs were stored in a database created in Microsoft Access 2000 and were treated by a simple macros written in Visual Basic for Applications. The ion saturation current regions of the graph were fitted by straight lines and the intersection points of the fitted lines with the Y-axis gave the values of the ion saturation currents (negative and positive). In the vicinity of zero voltage we also fitted the graph by a straight line with the same method and thus the slope of the characteristic \( \frac{dI}{dV} \) was obtained. The error bars for the calculated \( T_e \) and \( n_e \) were derived as a result of the fitting procedure of the saturation currents and the slope of VAC. Usually the error was less than 20 % in the center of the plasma jet, increasing towards the edges of the plasma jet where the electron density is much lower.

2.3.3 Thomson scattering to determine electron temperature and density

This chapter is devoted to the electron temperature and density measurements obtained with Thomson scattering. This diagnostic [92, 129, 88] is based on detection of laser radiation scattered by free electrons in the plasma. The intensity of the scattered light is proportional to the electron density. Usually the light is collected in the direction perpendicular to the initial laser beam. Due to the particle velocities the scattered light is Doppler shifted and represents the electron velocity distribution function. If the distribution function is Gaussian the electron temperature \( T_e \) (or \( \hat{T}_e = k_B T_e / e \) in eV) can be calculated from the width \( w_g \) of the Gaussian:

\[
    w_g = \frac{\lambda}{c} \sqrt{\frac{8 k_B T_e \sin^2(\theta/2)}{m_e}} \approx 1.48 \times \sqrt{T_e} \quad (2.12)
\]

Here, \( \lambda \) is the wavelength of the laser, \( k_B \) is the Boltzmann constant, \( e \) is the electron charge, \( m_e \) is the electron mass, \( \theta \) is the angle at which the scattered light is collected with respect to the initial laser beam (in our experiment \( \theta = 90^\circ \)), and \( c \) is the speed of light (all values are in SI-units).

To suppress stray-light, the laser beam in a Thomson scattering system must be very well aligned throughout the optical path containing lenses, mirrors and a special system of apertures. The last optical surface lens or window must be placed far from the plasma. In the Pilot-PSI setup, the light enters the vacuum chamber through a tube of 2 m length. Clearly, this entire optical system cannot easily be moved, so the measurement point is fixed. To make a scan of \( T_e \) and \( n_e \) profiles along the plasma plume, it would be necessary to make the source movable by placing it on a bellows section. This was not implemented on Pilot-PSI, but is seriously considered for future experiments.
2.3. PLASMA DIAGNOSTICS

2.3.4 Details of the Thomson scattering system

The schematic of the Thomson scattering system is shown in Figure 2.11. For

![Diagram of the Thomson scattering system](image)

the measurements the frequency-doubled green 532 nm line of a Nd:YAG-laser (LAB-170, Spectra-Physics) was used, which produced 450 mJ pulses with a pulse-width of 7 ns at a repetition frequency of 10 Hz. A singlet laser lens (focal length of 2500 mm) focuses the beam at the centre of the vessel to a beam waist of 0.33 mm by virtue of the small divergence of the beam (0.13 mrad). This gives a beam diameter of less than 0.7 mm at the ends of the full observational chord of 50 mm. The scattered light was collected by a lens and imaged on a linear array of 40 optical fibres (the same as we used for the high-resolution emission spectroscopy measurements, see section 2.3.6). Thus the light of the full observational chord was detected in 40 channels and analysed, providing profiles of $n_e$ and $T_e$ across the plasma jet with a spatial resolution of 1.3 mm.

An in-house constructed spectrometer in a Littrow arrangement [106], with the focal length of 1 m, dispersed and detected the light from all fibers simultaneously. A gated Instaspec V ICCD-camera (Andor) was used for light detection, consisting of a generation II image intensifier and a CCD chip ($12.67 \times 8.47 \text{ mm}^2$) of 512 spatial by 385 spectral pixels (22 mkm size). By gating the power supply of the image intensifier for only 50 ns at around each laser pulse, effective plasma light reduction is achieved. This is essential, because otherwise the plasma light
would saturate the CCD detector. To eliminate the effect of laser jitter on the gating, the laser pulse itself is used to trigger the gate. The minimal detectable electron temperature is determined by the width of the fibre, i.e. 0.4 mm. That corresponds to 12 CCD pixels or 0.31 nm at 532 nm and yields $T_{e,\text{min}} = 0.18$ eV. The spectral range is determined by the width of the CCD and is approximately equal to 10 nm. As a result, the maximum electron temperature that can be detected is about 10 eV.

To obtain the absolute electron density the intensity of scattered light was calibrated [130] at a known pressure of pure nitrogen, by measuring the Rayleigh scattering.

A typical CCD image of a Thomson scattering signal is presented in Figure 2.12. The intensity distribution in every fibre was fitted by a gaussian (Figure 2.13) and the $T_e$ was derived using the expression 2.12.

![Figure 2.12: A typical CCD image of Thomson scattering signal. The light from separate fibres across the plasma jet is seen.](image)
2.3. Plasma diagnostics

In this subsection we describe in detail the relationships between plasma parameters such as the electron density, the ion temperature and velocity components with characteristics of atomic line shapes: the Stark width, the Doppler shift and width of the line. Phenomena such as the Zeeman effect and hydrodynamical instabilities that can disturb the line shapes are discussed as well.

Methodology: Features of an atomic line profile and the related plasma parameters

Doppler shift and the velocity components

Light that is emitted by a moving particle is shifted in wavelength due to the Doppler effect. If there is a velocity component $v$ of the particle directed parallel to the line of observation then the shifted wavelength $\lambda$ is:

$$\lambda = \lambda_0 \left( 1 \pm \frac{v}{c} \right),$$  \hspace{1cm} (2.13)

Here, $\lambda_0$ is the original wavelength, and $c$ is the speed of light. Plus in this expression corresponds to the case when the particle moves away from the observer and minus corresponds to the opposite case, when the particle approaches the observer.
If there is a directed collective motion of radiating particles, then the whole line is shifted in the spectrum and we can speak about drift velocity components of the population. In the case of random movement of particles, the line is broadened. The natural line width is usually much smaller than the Doppler broadening, and we will neglect it in our measurements.

To estimate possible Doppler shifts of atomic lines in our experiments let us take the sound speed (2.1) for hydrogen at the temperature $T \approx 2\text{ eV}$ as the higher limit. Even in the supersonic expansion region the stream velocities are not expected to be much higher than the sound speed (Mach number $M = 1–2$, see chapter 2.2.2). The sound speed of up to $2 \cdot 10^4 \text{ m/s}$ corresponds to a Doppler shift $\Delta \lambda/\lambda$ of up to $7 \cdot 10^{-5}$, leading to a shift of 0.03 nm of the $H_\beta$ line and 0.04 nm of the $H_\alpha$ line. It introduces a necessary lower resolution limit of the spectrometer.

**Doppler broadening and the ion temperature**

When the radiating particles have a thermal velocity distribution, then due to the Doppler effect an atomic line shape represents this distribution function (if there are no other line-broadening effects). The distribution function of particle velocities (or, rather, projections of the velocities to a certain coordinate axis) which are in thermal equilibrium is a Gaussian. In this case we can derive the temperature of the radiating particles from the line shape. For a purely Doppler-broadened line the intensity distribution of a Gaussian shape is:

$$I(\lambda) = \frac{I_{tot}}{\sqrt{\pi w_D}} \exp \left( -\frac{(\lambda - \lambda_0)^2}{w_D^2} \right)$$  \hspace{1cm} (2.14)

Here, $I_{tot}$ is the total line intensity, $w_D$ is the width of the Gaussian (at $1/e$ of the maximum). The relation between the Doppler width of a line $w_D$ at a wavelength $\lambda$ and the temperature $T$ (in K) of the particles with mass $M$ is:

$$\frac{w_D}{\lambda} = \sqrt{\frac{2k_B T}{M c^2}}$$  \hspace{1cm} (2.15)

Here, $k_B$ is the Boltzmann constant and $c$ is the speed of light (all values are in SI-units). The ion temperature in electronvolts $\hat{T}_i = k_B T_i / e$ is:

$$\hat{T}_i = \frac{M c^2}{2e} \left( \frac{w_D}{\lambda} \right)^2$$  \hspace{1cm} (2.16)

where $e$ is the charge of electron.

The ion temperatures in our plasma are expected to be up to 2-3 eV. The corresponding Doppler widths are 0.04 nm for the $H_\beta$ line and 0.05 nm for the $H_\alpha$ line.

**Stark broadening and the electron density**

Another mechanism that broadens the line profile is the Stark effect. In atomic hydrogen the energy levels that represent different orbital quantum numbers $l$
are degenerated due to the Coulomb field of the nuclei. In an external electric field those energy levels are perturbed and the degeneracy is partly cancelled. The energy levels with different \( l \)-numbers have different shifts. Due to the high density of charged particles in the plasma, the radiating particles are always influenced by a superposition of the local microfields. This leads to the Stark broadening of spectral lines. For hydrogen atoms and hydrogen-like ions the Stark effect is linear with respect to an external electric field, while for other atoms it is quadratic (i.e. it is much smaller). Stark broadening in certain approximations can be described by different models. See, for example ([100], [101]). A Stark-broadened line shape is approximated by a Lorentz profile \( L(\lambda) \):

\[
L(\lambda) = \frac{w_L/\pi}{w_L^2 + (\Delta \lambda - d)^2}
\]

Here, \( w_L \) is the width of the Lorentz line profile (at half of the maximum), \( \Delta \lambda = \lambda - \lambda_0 \), \( \lambda_0 \) is the centre of the line and \( d \) is the Stark shift as the Stark mechanism leads not only to a broadening of lines but also to a shift of lines [102]. Griem states [100] that widths, shifts and profiles of suitable spectral lines are very insensitive to both electron and ion temperatures and gives a simple relation between the Stark width and electron density ([101], p. 305):

\[
n_e = C(n_e, T_e) \cdot \frac{w_{FWHM}^{3/2}}{2}
\]

where \( w_{FWHM} \) is the Full Width at Half Maximum (FWHM) of the Lorentz profile in angstrom, and \( C(n_e, T_e) \) is a coefficient that is only a weak function of the electron density. In addition, it has a slight temperature dependence. For the hydrogen H\( \beta \) line at electron densities of \( 10^{20} \) \( m^{-3} \) this coefficient varies from \( 3.84 \cdot 10^{14} \) \( \text{Å}^{3/2} \text{cm}^{-3} \) for \( T_e = 0.5 \) eV to \( 3.72 \cdot 10^{14} \) \( \text{Å}^{3/2} \text{cm}^{-3} \) for \( T_e = 2 \) eV (in that book wavelengths are expressed in angstrom and the densities are expressed in \( \text{cm}^{-3} \)). In the \( 10^{15} \text{cm}^{-3} \) range of densities this coefficient changes from \( 3.68 \cdot 10^{14} \) to \( 3.55 \cdot 10^{14} \) for the same range of temperatures. We used these widely accepted coefficients of Griem in all our calculations.

**Voigt profile**

The profile of a line which undergoes Doppler broadening and Stark broadening simultaneously is the convolution of a Gaussian \( L_G(\lambda) \) and a Lorentzian \( L(\lambda) \):

\[
L_{\text{conv}}(\lambda) = \int_{-\infty}^{+\infty} L_G(\Delta \lambda') L(\Delta \lambda - \Delta \lambda') d(\Delta \lambda'),
\]

where \( \Delta \lambda = \lambda - \lambda_0 \) and \( \lambda_0 \) is the centre of the line.

Such a profile is called a Voigt profile. Of course, this requires that the processes of broadening must not influence each other. When both widths \( w_D \) and \( w_L \) are of the same order and the measurement is sufficiently accurate, it
is possible to obtain both the ion temperature and the electron density from a single Voigt profile.

We consider our plasma to be optically thin, i.e. self-absorption is negligible, no perturbations of a line profile due to this effect are taken into account.

**Zeeman effect**

Finally we consider the Zeeman effect [103] as a phenomenon that can influence the shape of spectral lines. The physical background of the effect is as follows. Degeneracy of atomic energy levels with respect to the orbital quantum number \( l \) is cancelled in a magnetic field and the spectral line splits into components\(^1\) with different polarisation. The number of the components and their position in the spectrum with respect to one another depends on the quantum states between which the transitions occur, on the magnetic field strength and on the angle between the magnetic field and the line of sight [104]. In a strong magnetic field, when the perturbation of energy levels from the magnetic field is higher than the one from spin-orbit interaction, in the direction perpendicular to the magnetic field, three components (the Lorentz triplet) are seen. The central, non-shifted \( \pi \)-component, and two \( \sigma \)-components equally shifted on either side of the central one (\( \lambda_0 \)):

\[
\lambda = \lambda_0 \pm \Delta \lambda,
\]

where

\[
\Delta \lambda = \frac{2\pi ce}{eB}
\]  

This shift in the wavelength numerically corresponds to the inverse frequency of electron gyration in a magnetic field.

The electric field vector in the \( \pi \)-component is parallel to the magnetic field and for the \( \sigma \)-components it is perpendicular to the magnetic field. The intensity of the \( \pi \)-component is two times higher than that of each \( \sigma \)-component. Experimentally, the central component can simply be isolated by placing a polarizer in the optical path. This greatly simplifies the analysis. The influence of Zeeman broadening on temperature measurements in fusion plasmas was the subject of the Ph.D thesis of Anders Blom [105].

**Rotation and "wobbling" of the plasma jet in a magnetic field**

As we study plasma in a high magnetic field some other possible phenomena such as rotation of the plasma column or magnetohydrodynamic instabilities (i.e. "wobbling" of the plasma jet) should be taken into account.

Rotation of a plasma jet in a magnetic field is a very common and natural phenomenon, that has observed in many set-ups, for example [107, 108, 109, 110, 111]. In these papers and in [112] and [116] (Chapter 8) the authors also present physical models describing this phenomenon. The radial electric field \( \vec{E} \) that is perpendicular to the axial magnetic field \( \vec{B} \) is proposed as the main

\(^{1}\)In the case of the \( H_\beta \) line, (transitions from \( n = 4 \) to \( n = 2 \)) there are 54 components.
reason for the rotation of the plasma. This is due to the $E \times B$-drift, with the drift velocity $\vec{v}_{\text{drift}}$

$$\vec{v}_{\text{drift}} = \frac{\vec{E} \times \vec{B}}{B^2} \quad (2.21)$$

In section 4.7.2 we will describe the rotation of the plasma jet in Pilot-PSI and present a physical model for it.

The origin of the jet wobbling in a magnetic field is related to the current in the plasma jet and the Lorentz force is considered to be the main driving mechanism [113]. To estimate the frequency of the "wobbling" one can use the expression from [116] (Chapter 8, p. 357):

$$\omega = \sqrt{\frac{\pi I B}{n_e m_i \Lambda L^2}} \quad (2.22)$$

where $I$ is the current in the jet, $B$ is the magnetic field strength on the axis, $\Lambda$ is the $1/e$ radius of the plasma density, and $L$ is the length of the plasma column. Measurements on the wobbling and its influence on the HiRES and Thomson scattering data interpretation are the subject of section 4.6.

### 2.3.6 Details of the spectroscopy experimental technique

#### High-resolution spectrometer

The optical scheme of the HIRES measurements is depicted in 2.14. A single-pass imaging spectrometer in Littrow arrangement was designed and constructed at IPP Juelich.

The grating dimensions are $11 \times 11 \text{ cm}^2$ and it has a groove density of 1200 per mm. The light is focused by a lens of 15 cm in diameter with a focal length of 2.25 m. The blazed diffraction grating is optimised for the second diffraction order (blaze angle $17.45^\circ$). The main equation for a diffraction grating with the Littrow condition is [106]:

$$km\lambda 10^{-6} = 2 \sin \alpha \quad (2.23)$$

The angular dispersion for our spectrometer can be calculated as:

$$\frac{d\lambda}{d\alpha} = \frac{10^6 \cos \alpha}{km} \left( \frac{nm}{\text{rad}} \right) \quad (2.24)$$

and the linear dispersion is:

$$\frac{d\lambda}{dx} = \frac{10^6 \cos \alpha}{kmL} \left( \frac{nm}{\text{mm}} \right) \quad (2.25)$$

where $\alpha$ is the angle between the normal to the grating plane and the output ray of light with wavelength $\lambda$ (in nm), $k$ is the groove density (1200 per mm of the grating), $m$ is the diffraction order ($m = 2$ in our case), and $L$ is the focal length of the spectrometer ($L = 2250 \text{ mm}$).
CHAPTER 2. EXPERIMENTAL ARRANGEMENT

Figure 2.14: Schematical drawing of the optical emission spectroscopy system at Pilot-PSI. The light from the plasma is coupled to the high-resolution Littrow-spectrometer via a fibre-bundle preserving also the spacial information.

Rotation of the grating is provided by a precision mechanism with accuracy better than 0.0025 degrees. For the \( H_\beta \)-line (\( \lambda = 486.133 \) nm) this corresponds to \( \Delta \lambda = 3 \cdot 10^{-5} \). We used a hydrogen spectral lamp for wavelength calibration. Hydrogen lines from the lamp also served to find the optimal slit width. For this purpose the width of the \( H_\beta \)-line was measured versus the slit width (Figure 2.15). The instrument function due to the finite slit width was considered to be a single Gaussian. To determine the instrument function of the spectrometer, the 632.8 nm red line of a He-Ne laser was used. The width of this line is much smaller than the resolution of the spectrometer and the width of the instrumental function was found to be equal to or less than 0.003 nm, which corresponds to a CCD pixel size at 486.13 nm (\( H_\beta \)-line). In the case of pure Doppler broadening of the line shape and Gaussian shape of the instrumental function the resulting line width is:

\[
w_D = \sqrt{w_{\text{measur.}}^2 - w_{\text{instr.}}^2}.
\] (2.26)

Most of the measurements were performed on the \( H_\beta \) hydrogen line (486.133 nm). This line is recommended for the electron density measurements because atoms at higher energy levels are stronger coupled to the ions via collisions. Furthermore, the Doppler width of this line is smaller than the Doppler width of the \( H_\alpha \) line at the same particle Temperature, making the \( H_\beta \) line more
2.3. PLASMA DIAGNOSTICS

Figure 2.15: The measured width of the $H_\beta$-line versus the slit width. The real width is close to 0.01 nm (2.9 pixels). The optimal slit width is around .. mkm (5 points on the micrometer scale).

sensitive to Stark broadening. The calculated linear dispersion (2.25) in this wavelength range is 0.00341 nm/pixel and the linear dispersion measured from the resolved $H_\beta$ and the deuterium $D_\beta$ (486.003 nm) lines is around 0.0032 nm/pixel. The resolving power $\frac{\lambda}{\Delta\lambda}$ is around 20000. Ions in our hydrogen plasma in a magnetic field of up to 1.6 T have temperatures between 0.5–3 eV and the density is in the range of $n_e = 10^{19} - 10^{21} \text{ m}^{-3}$. The corresponding Doppler widths of the $H_\beta$ line are 0.016 to 0.040 nm (from 0.022 to 0.052 nm for $H_\alpha$ line) and the Lorentz widths are up to 0.022 nm.

The optical system

Light from the plasma is collected by a lens (6 cm in diameter with a focal length of 20 cm) in the direction perpendicular to the jet axis. A polariser was installed in front of the collecting lens to eliminate $\sigma$-components of Zeeman-splitted spectral lines. The collected light is coupled into the spectrometer via an array of 40 individual quartz fibers (0.4 mm in diameter). In this way, spatial information is preserved and a certain spectral range (about 1 nm in the $H_\beta$ line vicinity) can be investigated over the entire plasma jet profile (2–3 cm) at once. A CCD-camera with 298x1152 pixels was attached to the spectrometer to detect the signal.
HiRES to determine the axial velocity component

The method to determine velocity components of a plasma stream from Doppler shifts (2.13) of atomic lines we have described in section 2.3.5. To measure the axial velocity component the optical scheme was modified: a prism was placed in the vacuum vessel as shown in Figure 2.16. With this setup we measured the velocity component at approximately 15° from the jet axis. Due to sensitivity of the spectrometer to even small vibrations a spectral line can shift in a CCD image by several pixels. To ensure that the shift of line occurs due to the Doppler effect and not because of some other effects, a control measurement of the same spectral line (Hβ) from a spectral lamp was done just before and just after the measurement of radiation from the plasma jet. At the same time it gave an estimation of the error bar for the measurements. The error bar was found to be around 5%.

Plasma light is collected from a certain volume along the line of sight. It is important to note here that it contains information that is averaged over this volume and thus can indicate somewhat lower velocity values. The Zeeman effect caused another inconvenience for the measurements, as the light was detected in the direction parallel to the magnetic field thus only two σ-components of the line are seen in the spectrum (Figure 2.17). However, as the components are symmetrically shifted from the line centre, we were interested in the Doppler shift of this line centre only.

Figure 2.16: The optical scheme of the experiment with a prism to measure the axial velocity of the plasma jet derived from the Doppler shift of Hβ line. The measurement is made at the source exit in a magnetic field of 0.4 –1.6 T. Discharge current is 80 A, gas flow rate is 2.5 slm of H2.
Figure 2.17: In the direction parallel to the magnetic field only two shifted \( \sigma \)-components and no central \( \pi \)-component of the \( H_\beta \)-line are seen.
Chapter 3

Efficient hydrogen plasma production by a cascaded arc in the absence of a magnetic field

Abstract

We investigated the influence of discharge current and discharge channel diameter on the efficiency of plasma production in the cascaded arc source, for operation in both hydrogen and argon. Moreover, the influence of the pressure in the vacuum tank on the plasma plume was investigated. Radial profiles of the electron temperature ($T_e$) and density ($n_e$) were measured with a double Langmuir probe at an axial position of 30 cm downstream from the nozzle. This was done for discharge currents ranging from 50 to 80 A. The plasma expanded into a low pressure vessel, of which the pressure was varied from 6 to 60 Pa. At the position of the probe, measured electron densities were in the range of $3 - 5 \times 10^{18} \text{ m}^{-3}$ for argon and $2 - 6 \times 10^{16} \text{ m}^{-3}$ for hydrogen plasma. Electron temperatures were 0.20–0.25 eV for argon and 0.15–0.20 eV for hydrogen plasma. The background pressure was found to determine the position of the shock, and thereby the width of the plasma jet, in good agreement with theoretical expectations. The I-V characteristic of the source was measured for scans of the plasma current (30-100 A) and plasma channel diameter (2.5-4.5 mm) showed that the surface averaged resistivity of the plasma is independent of the channel diameter. The averaged resistivity scales with the averaged current density as $\bar{\eta} \propto \bar{j}^{-0.6}$ for argon and as $\bar{\eta} \propto \bar{j}^{-1.3}$ for hydrogen. A simple physical model of the plasma in the arc is presented, which assumes based on earlier observations and numerical modeling that the central $T_e$ in the channel is in good approximation independent of the discharge current and channel width, so that the only parameter determining the resistance of the plasma column is
its effective width. On the basis of this model, the stability of the source can be understood as a stable equilibrium between Ohmic dissipation and power loss. The most striking difference between operation in hydrogen and argon is that the argon discharge fills the channel almost to the wall, whereas in hydrogen operation a hot core forms that extends to only half the radius of the channel. From generic scaling arguments it is concluded that the power loss from the plasma channel is determined mainly by the cold layer surrounding the hot plasma core. The much larger heat conductivity of hydrogen then explains why in hydrogen the power equilibrium requires a narrow plasma column.

### 3.1 Introduction

Plasma production by a cascaded arc without a magnetic field has been studied extensively in the literature. Its operation in argon, with a focus on the supersonic expansion of the plasma into a low-pressure vessel, was studied in depth by van de Sanden [79]. Elaborating from this work, others have added hydrogen to the argon working gas, starting at a level of a few percent and gradually moving to pure hydrogen operation [82, 83, 84, 85, 87, 89]. These studies covered a wide range of subjects, including the supersonic expansion of hydrogen plasma [50], anomalous fast recombination in hydrogen plasmas involving rovibrationally excited molecular hydrogen [49], the influence of surface chemistry on the transport of hydrogen atoms in a supersonic hydrogen plasma jet [69], the kinetics of hydrogen atoms and rotationally-vibrationally excited molecules in the supersonic expansion and the production of atomic hydrogen radicals [89]. The work presented in this chapter continues in the line of these investigations, but for the first time places the focus on the efficient production of hydrogen ions.

The production of hydrogen plasma with a cascaded arc has not been fully characterized from the viewpoint of maximal plasma output. Qing et al. [83] did investigate the source in terms of hydrogen plasma output, but did not push the arc to its operational limits. For example, the arc current did not exceed 50 A and the gas flow rate was kept at around 1 slm. With respect to the arc geometry, the influence of the total length of the discharge channel was explored. They found that for higher plasma yield it is to be kept as short as possible to minimize power losses to the walls. However, a certain length is required for pressure build-up in the cathode region.

We intend to use the cascaded arc to produce extremely high hydrogen plasma fluxes. The first step in our approach is to expand the operational parameters. We push the arc performance by going to the higher current and gas flow range of $60 - 100$ A and $2 - 3.5$ slm, respectively. In particular, the effect of the current variation on the plasma production was investigated in detail and is described in section 3.2.

The expansion of the plasma was probed by varying the pressure in the vessel as is discussed in section 3.3.

With respect to the geometry of the arc, we focused on the influence of the
3.2 Influence of arc current on the plasma density

The discharge current is probably the most direct knob the experimentalist has with which the source operation can be changed. Changing the discharge current may affect the plasma inside the channel in many ways, but the most obvious influence is to change the total power that is dissipated in the source. However, since the I-V characteristic of the source is not very simple - e.g. its slope is positive for operation in argon, but negative for operation in hydrogen (see Figure 3.10) at the current densities used - we will present results both as function of discharge current and of input power. In a series of experiments, the current through the cascaded arc was varied from 50 A to 80 A in steps of 10 A. The gas flow was kept at 2 slm throughout these measurements. In these conditions the pressure in the vessel was typically around 5 Pa for both argon and hydrogen operation. The plasma was characterised in terms of electron density and temperature with double Langmuir probe measurements (see sections 2.3.1 and 2.3.2 for details on the probe measurements). The probe was positioned at an axial distance of 30 cm from the exit of the plasma source and could be moved radially from 130 mm above the plasma center to 30 mm below the center. The results are summarized in Figure 3.1.

Figure 3.1(a) shows that $n_e$ increases significantly with increasing arc current in argon operation. This is corroborated by visual inspection of the plasma: the plasma light emission does also increase significantly for increasing arc current. However, there is little influence of the discharge current on the shape of the $n_e$ profile.

The $T_e$ profiles in Figure 3.1(a) are much broader than the $n_e$ profiles. The measurements suggest a slight increase of $T_e$ with increasing current. However, we cannot exclude that this small increase is an artifact of the fit procedure, and due to the much stronger increase of $n_e$.

To evaluate the effect of the current on the arc output, the peak density of the profiles in Figure 3.1(a) are plotted versus the discharge current in Figure 3.2(a). This shows that, for operation in argon, $n_e$ grows approximately linearly from 3 to $5 \times 10^{18} \text{ m}^{-3}$ as the arc current is increased from 50 to 80 A. The corresponding arc voltage increases from 45 to 64 Volt, so that the resulting total power input to the source increases from 0.9 to 6.3 kW.

From these experiments it is clear that increasing the discharge current is an
(a) Argon plasma

(b) Hydrogen plasma

Figure 3.1: Electron temperature and density profiles determined from double Langmuir probe measurements for a scan of the arc current. The upper panel (a) shows the results for argon and the lower panel (b) for hydrogen plasma. In both cases, the gas flow rate was 2 slm, which results in a background pressure of around 5 Pa.
3.2. INFLUENCE OF ARC CURRENT ON THE PLASMA DENSITY

Effective means to increase the total ion flux emitted by the source, at constant input gas flow. Hence, it increases the efficiency of ion production expressed in terms of fraction of the gas input. With operation in argon, it does not, however, lead to increased energy efficiency, i.e. the energy needed to produce an ion at the exit of the source. We do not, in this stage, compute the absolute values of the efficiencies, as in the absence of a confining magnetic field the loss of plasma between nozzle and probe is large. This will be taken up in Chapter 4.

We return to Figure 3.1(b) and will now discuss the influence of the discharge current on the hydrogen plasma production. It is seen that increasing the current from 50 to 60 A has no effect at all on the measured $n_e$ profile. A further current increase does improve the $n_e$. It should be noted here that the IV-characteristic of the cascaded arc is negative for hydrogen operation. In the range 50 – 60 A, a current increase is accompanied by a voltage decrease such that the total power deposited into the arc remains almost constant - within the error bar (figure 3.3). Increasing the current from 60 A to 80 A does increase also the total power from approximately 7.5 kW to 9.8 kW. In this range $n_e$ grows from 3 to $6 \times 10^{16}$ m$^{-3}$, i.e. much faster than proportional.

Comparison of the hydrogen and argon results shows that the absolute $n_e$ levels differ by two orders of magnitude. Again, this is corroborated by visual inspection of the plasma. In hydrogen operation, there is no plasma light visible at the position of the probe measurement, whereas for argon light emission is still strong. This difference can partly be attributed to the higher sound speed in hydrogen than in argon, due to the lower particle mass. As a consequence, for the same ion flux density, the electron density should be roughly a factor 6

Figure 3.2: The peak electron density versus the arc current for argon (a) and hydrogen (b). The data are extracted from the profiles of 3.1(a) and 3.1(b).
smaller (assuming the same $T_e$ for both cases). But according to literature there is also a much faster loss process in the case of hydrogen, i.e. the efficient Molecular Activated Recombination (MAR) of hydrogen plasma [18, 19, 49], which is a two-particle process. In the case of argon plasma, volume recombination occurs via a three-particle process, which is much slower at the densities at hand [49, 92]. The issue of volume recombination will be discussed in more detail in section 3.3, that deals with experiments in which the background pressure was systematically varied. The rates for different ways of recombination of argon and hydrogen plasma are given in a table 3.1.

The electron temperature of the hydrogen plasma is observed to be in the range of 0.1 - 0.2 eV. The $T_e$ profile is very broad, for all cases except for the 60 A scan we measure a uniform $T_e = 0.17$ eV across the plasma.

Figure 3.3, in which the same measurements of the electron density are plotted versus the total input power to the source, clearly shows that for hydrogen operation both the ionization efficiency and the energy needed per ion improve strongly with increasing input power.

In summary, increasing the power to the source by increasing the discharge current is a very effective way to produce a greater ion flux. However, the effect is much more pronounced in hydrogen than in argon. We will return to this point after the presentation of the results on variation of the discharge channel width.
Table 3.1: Rates of recombination processes in hydrogen [136] plasma. (* the rate for this reaction is in m^6s^{-1})

<table>
<thead>
<tr>
<th>Reaction</th>
<th>rate (m^6s^{-1})</th>
</tr>
</thead>
<tbody>
<tr>
<td>( H^+ + e^- \rightarrow H^* )</td>
<td>( \sim 10^{-20} )</td>
</tr>
<tr>
<td>( H^+ + e^- + e^- \rightarrow H + e^- )</td>
<td>( 6.3 \cdot 10^{-38} )</td>
</tr>
<tr>
<td>( H^+ + H_2 \rightarrow H + H_2^* )</td>
<td>( 3 \cdot 10^{-15} )</td>
</tr>
<tr>
<td>( H_2^+ + e^- \rightarrow H^* + H )</td>
<td>( 8 \cdot 10^{-14} )</td>
</tr>
<tr>
<td>( H^+ + H^- \rightarrow H + H )</td>
<td>( 5 \cdot 10^{-14} )</td>
</tr>
<tr>
<td>( H_2^+ + H^- \rightarrow H^* + H_2(r,v) )</td>
<td>( 3.1 \cdot 10^{-13} )</td>
</tr>
<tr>
<td>( H_3^+ + e^- \rightarrow H + H + H )</td>
<td>( 1.4 \cdot 10^{-14} )</td>
</tr>
<tr>
<td>( H_3^+ + e^- \rightarrow H_2(r,v) + H^* )</td>
<td>( 1.1 \cdot 10^{-14} )</td>
</tr>
</tbody>
</table>

3.3 Influence of pressure in the vessel on the plasma flux

In the previous section we quantified the arc operation on the basis of probe measurements at \( \sim 30 \) cm downstream from the nozzle. By doing so, we must realize that the plasma expands and shocks between the nozzle and the probe. To get more insight in the extent to which this might influence our results, we consider the expansion in more detail in this section. In our approach, we used the background pressure in the vessel as the experimental parameter to probe the expansion.

The vessel pressure was varied from 5 to 60 Pa by tuning the pumping speed of the roots pump. The gas flow was kept constant at 2 slm and the discharge at 60 A. The effect on the expansion is dramatic as can be seen in the photographs shown in Figure 3.4. These clearly show that the plasma is expanding over the entire region where it emits light if the vessel is at low pressure. However, when the background pressure is higher, the expansion is suppressed and the plasma

Figure 3.4: Photo of argon plasma at 5 Pa (a) and 100 Pa (b) background pressure.
is confined. This was observed before for argon [68]. It is explained by the fact that the shock front is formed closer to the plasma source exit.

To consider this effect in more detail, we start from the expression (2.4) for the position of the shock front \( z_M \) in terms of flow rate \( \Phi \) (in \( \text{sccs} \) units - standard cubic centimeters per second), source temperature at the axis \( \hat{T}_s \) in eV, background pressure \( p_{\text{back}} \) in Pa and atomic mass number \( A \) [89]:

\[
 z_M = 1.8 \cdot 10^{-2} \sqrt[3]{\frac{\Phi}{p_{\text{back}}}} \sqrt{A\hat{T}_s} 
\]  

This shows that the shock position is closer to the source for increasing pressure. Because the plasma expands until it undergoes a shock and subsequently does not increase in diameter anymore, the diameter of the jet is determined by the position of the shock front. In other words, the expansion is reduced by increasing the pressure, as is schematically shown in Figure 3.5.

![Diagram showing the plasma expansion from low to high pressure](image)

**Figure 3.5:** Schematic of the plasma expansion from the plasma source nozzle to explain the influence of the background pressure on the jet diameter. The position of the stationary shock front depends on the pressure. At low pressure (a), the Mach disk is formed further away from the exit which leads to a larger diameter plasma jet compared to the case of high pressure (b).

The effect of the vessel pressure on the plasma expansion was studied by measuring the \( n_e \) and \( T_e \) profiles with the double Langmuir probe, for a scan of vessel pressure settings. The results are shown in Figure 3.6, for both argon and hydrogen operation.

The argon profiles (upper panel) show the expected trend of narrower \( n_e \) profiles for higher background pressures. The \( T_e \) profiles do not show a clear dependency on the background pressure.

In the case of hydrogen, the \( n_e \) profiles show the same trend: the diameter of the beam becomes smaller for increasing pressure. However, the \( T_e \) measurements show a decrease for better confined jets.
3.3. INFLUENCE OF PRESSURE IN THE VESSEL

Figure 3.6: The radial profiles of electron density and temperature in (a) argon and (b) hydrogen plasma. The profiles were measured with the double Langmuir probe for several settings of the background pressure in the vessel. The pressure was set via the rotation speed of the roots pump. The source was operated at 60 A and 2 slm.
The electron density profiles in Figure 3.6(a) were fitted with Gaussians. The width of the plasma jet is then characterized by the $1/e$ (half-)width of the profile. The results are plotted versus the background pressure $p$ (actually, versus $p^{-1/2}$) in Figure 3.7. This plot shows that the jet (half-)width $w$ varies with the pressure as $w \propto p_{\text{back}}^{-1/2}$. This is in perfect agreement with the shock front being positioned according to the vessel pressure following expression (3.1).

The total flux of argon plasma particles as a function of pressure was calculated by integration over the density profile as done in [68] (the flow velocity $v$ is believed to be approximately constant after the shock front i.e.,

$$
\Phi = v \int_0^\infty n_e(r) 2\pi r dr
$$

(3.2)

data not shown). This demonstrated that the total argon ion flux is not influenced by the background pressure. This result was expected, since there is no efficient volume recombination mechanism for argon plasma in the conditions of our experiments, as was shown by van de Sanden et al. [92, 68]. According to their analyses, the candidate pathways for recombination are two particle recombination

$$
Ar^+ + e^- \rightarrow Ar^* \rightarrow Ar + h\nu
$$
and three particle recombination

\[ \text{Ar}^+ + e^- + e^- \rightarrow \text{Ar}^* + e^-, \]

which are both slow processes (see table 3.1).

However, this is not the case for hydrogen. Again, we integrated the \( n_e \) profiles to estimate the plasma flux as a function of the vessel pressure. This integrated density is plotted versus pressure in Figure 3.8.

\[ \text{Figure 3.8: } n_e \text{ integrated over the plasma jet cross-section in hydrogen plasma versus the background pressure. Note the double logarithmic scale.} \]

The plot shows that the integrated density initially decreases with pressure. This is in agreement with observations of de Graaf et al. [49]. They attributed this anomalous fast process of hydrogen plasma extinguishing to Molecular Activated Recombination (MAR) [18, 19, 49]. In line with their analyses, the recombination rate for a hydrogen plasma grows significantly with the collision frequency due to the resonant charge exchange between a hydrogen ion and a hydrogen molecule:

\[ H^+ + H_2 \rightarrow H + H_2^+ \]

and the subsequent dissociative recombination of the molecular ion:

\[ H_2^+ + e^- \rightarrow H + H \]

(see also table 3.1.) The second step of MAR has much higher rate than the first one [49]. That means that the two-step recombination is limited by the rate of the first reaction. The rate of the charge exchange is proportional to the
ion density and to the density of molecular hydrogen:

\[
\frac{d n_i}{dt} = -k_{H^+H_2} \cdot n_i \cdot n_{H_2}
\]  

(3.3)

The Figure 3.8 shows the strong decrease of the integrated density for increasing background pressure. The data is in fair agreement with the exponential behaviour predicted by equation 3.3. We note that the data point at 6.5 Pa is not reliable, because at this low pressure the beam is so wide that it extends significantly beyond the range of the probe. A further uncertainty in the interpretation is the flow velocity of the hydrogen plasma, of which no measurements are available. For a proper assessment of the loss mechanism, the profile of the ion flux density, rather than that of \(n_e\) should have been integrated. However, given the limitations of the measurements, we can draw the following conclusions from the scan of the background pressure:

- increasing the background pressure leads to a narrowing of the \(n_e\) profiles, both in argon and hydrogen operation

- this narrowing is ascribed to the shift of the position of the Mach disk, and the measurements are in good agreement with the theoretical prediction based on this effect

- for argon operation, increasing the background pressure has almost no effect on the integrated \(n_e\). From this, it is concluded that the loss channel of Ar ions is slow and/or does not involve collisions with the background plasma. This is in agreement with the theory, in which the background gas only enters in a three-particle process, of which the rate is negligible in the prevailing densities. Similarly, the constancy of \(T_e\) under variation of the background pressure indicates that the dominant power loss processes do not involve the background gas. Also this is in agreement with the expectation based on previous work.

- for operation in hydrogen, there is a very strong decay of the integrated \(n_e\), hence ion density, with increasing background pressure. This is in agreement with results published in the literature, where the loss is attributed to the MAR process. Our measurements are in agreement with the exponential decrease of the ion flux as function of the background pressure that would follow from this loss process.

### 3.4 Variation of the arc channel diameter

The diameter of the discharge channel is a very important parameter of the cascaded arc source, in particular with an eye on the envisaged upscaling of the present design to a large, high-power source. The physics questions that play a role are the stability of the arc on the one hand, and the power efficiency of the source, i.e. the balance of input power and power loss, on the other. To study
the effect of the discharge channel diameter on the operation of the cascaded arc source, we performed a series of experiments in which the diameter of the channel in the cascaded plates was varied from 1.5 to 4.5 mm in steps of 0.5 mm. In these experiments [90], we measured the arc voltage as a function of arc current and gas flow. The upper limit of the range of channel diameters was given by experimental constraints: the available power supplies and the input gas flow. Operation at too low current density resulted in an unstable discharge, which also led to damage to the wall of the channel. The lower limit was determined experimentally: the arc only operates for a channel width of at least 1.5 mm in argon, and 3.5 mm in hydrogen. There are also limits for the gas flow rate. With the present equipment, not more than 1.9 slm of argon could be driven through the smallest diameter channel of 1.5 mm. The bigger bores work in the whole range from 0.8 up to 3.6 slm that could be covered with the equipment available in these experiments. The arc current for the diameter of 1.5 mm cannot be higher than 60 A. For larger channel diameters, the arc becomes unstable for arc currents below 20-30 A. The scans of flow rate at fixed bore were mostly conducted at a current of 60 A, with a variation of the gas flow rate from 1 slm to 3.6 slm (with 0.2 slm steps). The scans of arc current were conducted at a constant flow rate of 2 slm, with currents varied from 20-100 A (the limits depending of the stability of the arc), in steps of 5 A. The complete range of experimental parameters we investigated is summarized in table 3.4. Appendix 3.9 gives a full table of the measurements.

<table>
<thead>
<tr>
<th>Gas</th>
<th>Diameter (mm)</th>
<th>Gas flow rate (slm)</th>
<th>Arc current (A)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Argon</td>
<td>1.5</td>
<td>1.0</td>
<td>10-60</td>
</tr>
<tr>
<td></td>
<td>1.9</td>
<td>10-60</td>
<td></td>
</tr>
<tr>
<td></td>
<td>2.0</td>
<td>2.0</td>
<td>60</td>
</tr>
<tr>
<td></td>
<td>2.5</td>
<td>2.0</td>
<td>10-95</td>
</tr>
<tr>
<td></td>
<td>1.0-3.6</td>
<td>60</td>
<td></td>
</tr>
<tr>
<td></td>
<td>3.0</td>
<td>1.1</td>
<td>30-40</td>
</tr>
<tr>
<td></td>
<td>2.0</td>
<td>40-95</td>
<td></td>
</tr>
<tr>
<td></td>
<td>1.0-3.6</td>
<td>60</td>
<td></td>
</tr>
<tr>
<td></td>
<td>3.5</td>
<td>2.0</td>
<td>30-95</td>
</tr>
<tr>
<td></td>
<td>1.0-3.6</td>
<td>60</td>
<td></td>
</tr>
<tr>
<td></td>
<td>4.0</td>
<td>2.0</td>
<td>20-100</td>
</tr>
<tr>
<td></td>
<td>0.8-3.6</td>
<td>60</td>
<td></td>
</tr>
<tr>
<td></td>
<td>4.5</td>
<td>2.0</td>
<td>20-100</td>
</tr>
<tr>
<td></td>
<td>0.8-3.6</td>
<td>60</td>
<td></td>
</tr>
<tr>
<td>Hydrogen</td>
<td>3.5</td>
<td>2.0</td>
<td>40-100</td>
</tr>
<tr>
<td></td>
<td>4.0</td>
<td>2.0</td>
<td>30-100</td>
</tr>
<tr>
<td></td>
<td>1.0-3.6</td>
<td>60</td>
<td></td>
</tr>
<tr>
<td></td>
<td>4.5</td>
<td>0.8-3.6</td>
<td>30-100</td>
</tr>
</tbody>
</table>
The voltages required to maintain a stabilised current through the arc were measured. The electric field in the channel is assumed to be uniform and equal to the voltage $V$ divided by the channel length ($l = 40\ mm$): $E = V/l$. To arrive at a unified representation of the data, we introduce the averaged current density $\bar{j}$, i.e. the current density averaged over the channel cross-section of radius $a$: $\bar{j} = I/\pi a^2$. The average resistivity $\bar{\eta}$ is computed from the measured voltage and discharge current according to

$$\bar{\eta} = E/\bar{j} = \frac{V \pi a^2}{II} \quad (3.4)$$

Figure 3.9 summarises the results of the different parameter scans, both for argon and hydrogen. In this plot, the average resistivity is plotted versus the

![Figure 3.9](image)

Figure 3.9: Average resistivity of argon (Ar) and hydrogen (H) plasma in the cascaded arc of different diameters (2.5 – 4.5 mm) versus the averaged current density, for arcs of different bore. The different markers correspond to different channel widths. Points with different flow rate are not separately marked (the flow rate has little influence on the resistivity (see Figure 3.10)) but can be recognized as small vertical clusters. The measurements show that, at constant current density, the average resistivity of argon plasma does not depend on the arc channel diameter. For hydrogen the picture is less clear: between the channel diameters 4.0 and 4.5 mm no change of resistivity is observed, but at 3.5 mm, the smallest diameter at which operation could be sustained, the resistivity is higher.
3.4. VARIATION OF THE ARC CHANNEL DIAMETER

Averaged current density. This representation of the data was chosen because it was found to organise the data points in well distinguishable clusters. The figure is essentially a normalised plot of the I-V characteristic of the source. Note that in this double-log plot a slope of -1 corresponds to a flat I-V characteristic.

The behaviour of the arc source, represented in condensed form in 3.9, can be described as follows:

- for argon, at fixed averaged current density \( \bar{j} \) the averaged resistivity \( \bar{\eta} \) does not depend on the bore radius \( a \).

- for argon, the averaged resistivity \( \bar{\eta} \) depends on the averaged current density \( \bar{j} \) as \( \bar{\eta} \propto \bar{j}^{-0.8} \) for \( \bar{j} \) at the low end of the operational window, weakening to \( \bar{j}^{-0.6} \) at the highest values of \( \bar{j} \) applied (corresponding to a positive slope of the I-V characteristic).

- for hydrogen, the specific resistivity is higher than for argon. The difference is a factor 4 at small \( \bar{j} \), decreasing to a factor 2 at the highest values of \( \bar{j} \).

- for hydrogen, \( \bar{\eta} \) depends on \( \bar{j} \) as \( \bar{\eta} \propto \bar{j}^{-1.3} \) (exponent less than -1 corresponds to a negative slope of the I-V characteristic).

- for hydrogen the curves for the bore radius \( a = 4 - 4.5 \) mm do line up, but at \( a = 3.5 \) mm the resistivity is higher by a factor of \( \approx 1.2 \). It is noted that this is the smallest radius for which operation in hydrogen was possible, \( a = 2.5 - 3 \) mm was tried but there was no sustained arc operation with the power source available.

The above is not more than a description of the experimental data. Even without any interpretation or model forming, these results indicate that under variation of the channel diameter the source behaviour is described by the universal curves in Figure 3.9. This will serve as an empirical guide for further upscaling of the source in the frame of the development of Magnum-PSI. The relative simplicity of Figure 3.9, which nonetheless shows salient differences between operation in hydrogen and argon, suggests that the behaviour of the source may be dominated by simple, global effects. To investigate whether such a simple picture can be constructed, we will propose a single-parameter model of the source in the next section.

In Figure 3.9 the points with variation of the gas flow rate were not separately marked. The effect of the variation of the gas flow, which is rather small and difficult to distinguish in Figure 3.9, is brought out in Figure 3.10. This shows that the voltage over the hydrogen arc increases slightly with increasing gas flow rate, however it decreases with increasing current for all used channel diameters (Figure 3.10).
CHAPTER 3. HYDROGEN PLASMA IN $B=0$

Figure 3.10: The voltage over arc on hydrogen slightly increases with increasing gas flow rate at constant current $I = 60 \text{ A}$ (left plot). The discharge voltage decreases with increasing discharge current (at constant gas flow rate of 2 slm) when operating on hydrogen (right plot).

3.5 A one-parameter physical model of the arc

The fact that by plotting the averaged resistivity versus the averaged current density led to essentially single curves describing all data of current, gas flow and channel diameter scans (Figure 3.9), suggests that the behaviour of the source may be dominated by simple, global effects. In this section we will investigate if this behaviour can be understood on the basis of a simple physical model.

The basic assumption for this model will be that the electron temperature $\hat{T}_e$ (in eV) in the source is practically fixed, as follows from expression (3.5) [46]:

$$\hat{T}_e \approx \frac{\hat{E}_{\text{ion}}}{\ln(10p_\text{s}\sqrt{A}) - \ln(\hat{T}_h)}$$

(3.5)

where $\hat{E}_{\text{ion}}$ is the ionisation energy in eV of the working gas (13.6 eV for hydrogen) with atomic mass $A$ (1 a.m.u. for hydrogen), $p$ is the pressure in the discharge channel ($p \approx 10^4 \text{ Pa}$) of the length $l_s$ (approximately 40 mm), and
3.5. A ONE-PARAMETER PHYSICAL MODEL OF THE ARC

$\hat{T}_h$ is the temperature of heavy particles.

At the basis of this is the fact that the source is operated at high pressure, in relatively low ionization. The plasma therefore lives at the low end of the SAHA equilibrium, where a temperature increase is associated with a steep increase of the ionization degree, and therefore requires a lot of energy. In [99] this has been verified experimentally, while also numerical simulations [98] of the cascaded arc support the hypothesis that $T_e$ has very narrow margins. Typically, for the arc under consideration, $T_e$ is expected to lie in the range 1.1-1.3 eV.

Based on these considerations, we will take as central assumption that the electron temperature $T_e$ is invariant.

The arc conductivity according to Spitzer (3.6) depends only on the electron temperature $T_e$ (in eV):

$$\sigma = 2 \cdot 10^4 \frac{T_e^{3/2}}{\ln(\Lambda)}$$  (3.6)

Here, $\ln(\Lambda)$ is the Coulomb logarithm that is a very slow varying function of the electron temperature and density (see, for example, [116], p. 317). Thus the resistivity in the current carrying channel (where the plasma exists and the electron temperature is high enough) is constant as well.

For the radial profile of $T_e$, and thereby of the resistivity, we will assume a generic profile shape which is characterized by a single parameter, which generically would be a peaking factor or effective width of the profile. The strategy will then be to determine on the basis of the experiments how this parameter depends on the channel radius and current density.

For the model it is in fact not very important what we choose for the generic profile shape. Based on considerations in the literature, especially [68] in which a similar reasoning is followed, we use here the flat-top profile, in which $T$ is uniform out to a radius $r < a$. For convenience we introduce the 'filling factor' $\alpha = r/a$. Thus, we now have a single parameter $\alpha$ which describes the effective conductivity of the plasma channel. As in the model the temperature in the hot part of the channel is uniform, we can define the resistivity $\eta_0$ in the current carrying channel and relate it to the averaged resistivity $\bar{\eta}$, and similarly we write the local current density $j_0$ and its trivial relation to the averaged current density $\bar{j}$:

$$\eta_0 = \bar{\eta} \alpha^2$$  (3.7)

$$j_0 = \bar{j} / \alpha^2$$  (3.8)

In terms of the current density and resistivity in the current carrying channel $j_0$ and $\eta_0$, the total current and resistance of the arc are $I_{arc} = j_0 \pi a^2 \alpha^2$ and $R = \eta_0 l / (\pi a^2 \alpha^2)$ with $l$ denoting the length of the arc. The power dissipated in the arc is then:

$$P_{in} = I_{arc}^2 \cdot R = j_0^2 \pi a^2 \alpha^2 \cdot \eta_0 l$$  (3.9)

However, $j_0$ depends on $\alpha$ and is not known. Hence, we rearrange the expression in terms of the known averaged current density $\bar{j}$ to reveal a clear
dependence on the parameter $\alpha$, at other values being constant:

$$P_{in} = \frac{\gamma^2 \pi a^2 \eta_0 l \alpha^{-2}}{2}$$  \hspace{1cm} (3.10)

The first characteristic of this model immediately becomes apparent: at given total current $I_a$ (this is how the source is operated experimentally), the voltage $V$ and hence the power dissipated by the plasma scales as $\alpha^{-2}$.

Figure 3.11 depicts this behaviour. We note that this strong dependency provides for stable operation of this model arc. Whatever the precise processes are that are at play, in the end the width of the channel is determined by the power balance. In Figure 3.11 also a curve is sketched that represents the power loss as function of $\alpha$. We do not know this curve, but for all reasonable loss mechanisms this curve will be an increasing function of $\alpha$. Hence, from a power balance point of view, equilibrium occurs where the power dissipation and power loss curves cross. We note that this is a stable equilibrium: if at any time the dissipated power is larger than what is needed to sustain the discharge as it is, the current channel will widen (it must in this model, since the temperature is fixed), which immediately reduces the input power. Conversely, if the power falls short, the edge of the plasma cools, i.e. the channel contracts and the dissipated power increases. The relation between the filling factor and the dissipated power is so strong that in this way it will stabilize very effectively fluctuations in the arc operation.

Figure 3.11: Schematic plot of the power dissipation and power losses in the arc versus the filling factor $\alpha$ according to the single-parameter model described in the text. $P_{in} \propto 1/\alpha^2$, $P_{loss} \propto const \cdot \alpha/(1 - \alpha)$ with different constants for hydrogen and argon.
3.6 Interpretation of the measurements in terms of the single-parameter model

We shall now analyse the observed behaviour of the arc in terms of the single-parameter model, starting with operation in argon.

First, the fact that at given averaged current density $\bar{j}$ the averaged resistivity $\bar{\eta}$ does not depend on the channel bore implies that $\alpha$ does not vary with the channel bore. ($\bar{\eta} = \eta_0 \cdot \alpha^{-2}$ is independent of the bore radius $a$, and since $\eta_0$ is a constant in this model, the filling factor $\alpha$ must be also independent of $a$. This is a very important result. It implies a scale invariance of the source operation at a given averaged current density. Below we will analyse what this means for the power loss channel.

Second, the empirical dependence of $\bar{\eta}$ on $\bar{j}$: $\bar{\eta} \propto \bar{j}^{-0.6 \ldots -0.8}$. Translates into a relatively weak dependence of $\alpha$ on $\bar{j}$: $\alpha \propto \bar{j}^{0.3 \ldots 0.4}$. Thus, for a given channel diameter, increasing the discharge current leads to an increase of the filling factor, but the effect is weak. Nonetheless, the filling factor cannot exceed unity, and by raising the current density it must approach this value. It is therefore important to estimate the absolute value of the filling factor.

Using the expression for the conductivity (3.6) (Spitzer’s formula), taking $T = 1.2$ eV, we can estimate the absolute value of $\alpha$. For the series of the measurements with the arc channel diameter of 4.5 mm and currents between 30 and 80 A the filling factor $\alpha$ for argon is between 0.64 and 0.96, increasing with the current. This is an important observation. Clearly, the applicability of the model breaks down when $\alpha$ approaches unity. There, the current channel cannot expand anymore and the only mechanism to absorb more power is by increasing $T_e$, and thereby the degree of ionisation (and in the case of hydrogen, dissociation). We shall return to this consideration in the discussion (section 3.8).

The same analysis applied to the hydrogen data leads to the following conclusions:

- The resistivity is up to a factor 4 higher than in argon. Applying Spitzer’s conductivity (3.6), this leads to absolute values of $\alpha$ in the range 0.43 to 0.97 for arc current from 30 to 60 A. Hence, the hydrogen arc fills only a narrow tube inside the available channel. The probable explanation for this higher heat loss is considered in the discussion (section 3.8). Whatever the detailed physics behind the small filling factor is, it does imply that there is much scope for increasing the current in the hydrogen arc. This will increase $\alpha$, and thereby improve the plasma output and possibly the efficiency of the source.

- The dependence of the average resistivity on the current density clearly is much stronger than in the case of argon. The empirical relation $\eta \propto \bar{j}^{-1.3}$ translates into $\alpha \propto \bar{j}^{0.65}$.

- as with argon, $\alpha$ appears to be independent of the channel diameter. However, the evidence is less clear in the hydrogen case, as effectively the
channel bore was only varied from 4.0 to 4.5 mm. The 3.5 mm bore did have 20% higher resistivity, but this was at the limit of the operational window and therefore not a reliable point for a scaling study.

3.7 Interpretation of the empirical dependence of the filling factor on the current density and the channel diameter; a study of power loss using generic scaling arguments

In the single-parameter model developed above, the plasma temperature in the arc is assumed constant leaving only the filling factor \( \alpha \) as the sole parameter determining the behaviour and performance of the source. As indicated before, \( \alpha \) is the result of a power balance: it adjusts itself such that the Ohmic dissipation in the source equals the power losses. The Ohmic dissipation can be expressed, keeping all relevant parameters \( (\bar{j}, a \text{ and } \alpha) \), as:

\[
P_{in} \propto \bar{j}^2 a^2 \alpha^{-2}
\]

For the power loss we will write the generic scaling with variables in the exponents:

\[
P_{loss} \propto \alpha^{2x} a^{2y}
\]

To place this scaling in context, we consider a few typical loss mechanisms:

a. Volume losses from the hot plasma channel, e.g. losses on ionisation and radiative losses from an optically thin medium, are described by \( x = y = 1 \):

\[
P_{loss} \propto \alpha^2 a^2
\]

b. Conductive losses from the hot cylinder, at constant temperature, scale with \( x = y = 0 \) as:

\[
P_{loss} \propto \alpha^0 a^0
\]

c. If the loss is determined by the conduction of heat through the layer outside the current carrying channel, the generic form becomes:

\[
P_{loss} \propto \alpha/(1 - \alpha)
\]

For small values of \( \alpha \) this constitutes a relatively weak dependence of \( P_{loss} \) on \( \alpha \), but for \( \alpha \) close to unity, the dependence is very steep (see Figure 3.11). The equilibrium condition \( P_{in} = P_{loss} \) yields a parametric dependence for \( \alpha \):

\[
\alpha \propto \bar{j}^{1/(1+x)} a^{(1-y)/(1+x)}
\]

For argon, the experimental result \( \alpha \propto \bar{j}^{0.3-0.4} a^0 \) leads to the following values for the exponents \( x \) and \( y \): \( x = 1.5 - 2 \), \( y = 1 \).

For hydrogen, the empirical relation \( \alpha \propto \bar{j}^{0.65} a^0 \) results in \( x = 0.5 \), \( y = 1 \).
3.8 DISCUSSION OF THE SINGLE PARAMETER MODEL

From these global scaling exercises, it becomes clear that power losses that are associated with the hot plasma core are not easily reconciled with the experimental results. By assuming that the power loss is limited by the conduction through the layer that surrounds the hot plasma, the different behaviour of hydrogen and argon can be understood. The fact that experimentally \( \alpha \) is found to be independent of the channel radius is not readily understood from the power balance analysis.

3.8 Discussion of the single parameter model

A systematical study of the arc I-V characteristics at different diameters of discharge channel revealed certain empirical regularities:

- the arc resistivity \( \bar{\eta} \) (averaged over the channel cross-section) occurred to be independent of the channel diameter \( a \) for both argon and hydrogen
- \( \bar{\eta} \) scales with the averaged current density \( \bar{j} \) as \( \bar{\eta} \propto \bar{j}^{-0.6-0.8} \) for argon plasma (that corresponds to a positive I-V characteristic) and as \( \bar{\eta} \propto \bar{j}^{-1.3} \) for hydrogen plasma (a negative I-V characteristic).
- \( \bar{\eta} \) for hydrogen is 2–4 times larger than that for argon (the difference is less at higher current density).
- a filling factor \( \alpha = r_{\text{plasma}}/a_{\text{bore}} \) is large for argon (approaches 1) and is significantly smaller for hydrogen (0.4–0.8).

A model on the arc power balance with a single parameter \( \alpha \) clearly explains stability of the arc operation. The model assumes a constant electron temperature in the plasma channel with only the filling factor varied for self-stabilisation. Power input scales as \( P_{\text{in}} \propto \alpha^{-2} \). Power losses are considered to be of two different types: the volume losses (ionisation and plasma heating, radiation) that is proportional to \((\alpha \cdot a)^2\); losses to a cooled wall due to heat conduction through a layer of cold gas at the wall which scales as \( \alpha/(1 - \alpha) \). The volume losses increase with \( a \) while conductive losses are independent of \( a \).

The significant difference in \( \alpha \) between argon and hydrogen arc is due to a difference in heat conduction of these two gases (a factor of 11 higher for hydrogen). It is also consistent with a considerable difference in the input power at the same arc current. If we take for example \( \alpha = 0.9 \) for argon and consider a power balance as:

\[
P_{\text{in}} = C_1 \cdot \alpha^{-2} = P_{\text{loss}} = C_{\text{Ar}} \cdot \alpha/(1 - \alpha)
\]

then \( C_1/C_{\text{Ar}} \approx 7.3 \). For hydrogen this ratio is less by a factor of 11 due to thermal conductivity as mentioned above. Thus, for hydrogen:

\[
\frac{\alpha^3}{1 - \alpha} \approx 0.66
\]
and we find $\alpha = 0.63$ for hydrogen that is in a good agreement with our results. The model is valid till $\alpha$ approaches to 1. The layer of the cold gas along the wall must have a finite non-zero thickness. At the limit the electron temperature starts to increase rather than $\alpha$ increases. Increase of electron temperature leads to increase of ionisation degree. In spite of the limitations, the model provides a good way to consider power balance and stability of the arc and explains the arc behaviour and differences in operation in argon and hydrogen.

The model suggests a more efficient plasma production at higher current densities. For hydrogen it is also because the dissociation degree increases and the gas can behave more like a monoatomic gas. Preliminary experiments with arc current up to 300 A at 4 mm bore showed that even hydrogen I-V characteristic becomes positive at higher current densities. That means that the filling factor of the arc approaches to its limit - unity.

Based on the model for power losses, we find that operation in the flat range of the I-V characteristic is the most efficient in energy per ion. Thus, to produce more ions it is better to increase the arc diameter rather than to push the current harder. There is a very wide range of the arc current with approximately the same efficiency that provides operational flexibility. Increase of the arc diameter to produce more ions is also essential for upscaling the plasma source for the future Magnum-PSI set-up.

The last point that is not clear at the moment is the question why the filling factor $\alpha$ is independent of the bore $a$. Maybe, if we push up $\alpha$ high enough the electron temperature increases, and that would give a dramatic improvement of the ion output.

3.9 Appendix:

Arc voltages measured for a range of gas flow rate (1.0 - 3.6 slm), arc current (20-100 A) at different arc channel diameter (2.5 - 4.5 mm for argon and 3.5 - 4.5 for hydrogen).
Figure 3.12: Arc voltages measured for a range of argon flow rate (1.0 - 3.6 slm) and arc current (20-100 A) at different arc channel diameter (2.5 - 4.5 mm).

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Figure 3.14: Arc voltages measured for a range of hydrogen flow rate (1.0 - 3.6 slm) and arc current (20-100 A) at different arc channel diameter (3.5 - 4.5 mm).

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Chapter 4

Hydrogen plasma in a high magnetic field

Abstract

We consider the effect of a high (0.4–1.6 T) axial magnetic field on the production of hydrogen plasma in a cascaded arc and on the transport of the plasma jet. Moreover, we varied the geometry of the nozzle of the arc. The electron temperature ($T_e$) and density ($n_e$) profiles were measured by Thomson scattering at 35 mm from the nozzle. The jet velocity components were measured by means of high-resolution emission spectroscopy. The combination of a wide nozzle and the magnetic field was found to lead to effective confinement of the plasma, producing a hydrogen plasma jet with typically ~1 cm diameter that extends deep into the vessel, up to the target at 1 m from the nozzle. $n_e$ was found to increase linearly, as B, at B=1.6 T peak values of $n_e = 7.5 \times 10^{21}$ m$^{-3}$ are observed. $T_e$ was around 1.9 eV, almost independent of B. The forward velocity measured at the position of Thomson scattering is 3 km/s, also independent of B. With this result, the ITER-relevant flux density ($10^{24}$ m$^{-2}$s$^{-1}$) has been achieved in Pilot-PSI. The variation of the nozzle diameter had a strong influence on both $n_e$ and $T_e$. Going from a 5 mm nozzle to 8 mm, $T_e$ increases from 0.8 to 1.9 eV. The peak $n_e$ remains the same, but the (half-)width of the $n_e$ profile increases from 6 to 10 mm. The effect on the forward velocity is weak. The higher $T_e$ is explained by the longer, cross-field return path of the arc current. This is corroborated by the higher arc voltage for larger nozzle diameters. The corresponding cross-field electric field is consistent with the measured rotation velocity of the plasma jet. Finally, a 0.7 MHz wobble of the plasma jet is observed. The frequency is in agreement with the theoretical prediction of this wobble. The amplitude was observed to decrease with increasing B.
4.1 Introduction

The aim of this thesis is the development of a high-flux magnetized linear plasma generator for plasma surface interaction (PSI) experiments. The magnetic field has a dual purpose in those experiments. First, it is an essential element of the physics of plasma-surface interaction. The field traps particles that come off the surface as soon as they are ionized, so that the plasma in front of the surface is determined by the PSI rather than the incoming plasma. This defines the strongly coupled limit of PSI. Thus, to mimic PSI conditions as they occur in the divertor of a fusion reactor, a strong magnetic field is a condition. Second, the field is necessary to confine the plasma jet and transport it from source to target. It was already seen that in the absence of a magnetic field, the plasma density decays rapidly as function of distance to the source. This was understood as the result of supersonic expansion of plasma (see section 2.2.2) and of Molecular Activated recombination (MAR) [18, 19, 49]. Application of an axial magnetic field is expected to confine the plasma in a jet and thus improve energy and particle confinement. This chapter reports on experiments in Pilot-PSI in which axial magnetic fields of up to 1.6 T were applied. The questions that could be addressed with these experiments are:

• is there an effect of the field on the operation of the source itself
• what is the effect of the field on the temperature and density of the plasma measured at a distance from the source, and
• can we understand this effect and deduce an extrapolation to higher field.

In these experiments the essential measurements are: current and voltage of the cascaded arc (effect on the source operation) and $n_e, T_e$ profiles and forward velocity measured at 40 mm from the source.

In a separate but linked series of experiments the shape of the nozzle of the source was varied. The motivation for this experiment is that it has been observed in the literature that in hydrogen operation, the shape of the nozzle can have a strong influence on the plasma production [133]. This is attributed to the high conductivity of hydrogen, which leads to a high heat loss to the nozzle in a region where it is not balanced by a high input power density. The combination of a strong axial magnetic field, which reduces the perpendicular electrical and heat conductivity, and a variation of the nozzle diameter, could be expected to introduce new phenomena and has not been researched before.

In addition to the plasma confinement a magnetic field introduces several other effects in the plasma jet in comparison with the case without a magnetic field. The radial electric field perpendicular to the axial magnetic field causes rotation of the plasma column due to the $E \times B$-drift. Rotation of a plasma column in a magnetic field is a well-known phenomenon, that has been studied in magnetic fields up to 0.2 T by e.g. [107, 108, 109, 110, 111, 112, 116]. In our case we extend the range to fields of 0.4–1.6 T, so that much higher rotation velocities are anticipated. The measurements of the rotation velocity
4.2. INFLUENCE OF MAGNETIC FIELDS ON THE PLASMA JET

were obtained from the Doppler shifts of hydrogen atomic lines. Section 4.4.1 presents the results of these measurements. The high-resolution spectroscopic measurements that are at the basis of these rotation velocity measurements are quite complex and detailed description and discussion of these measurements is given later, in chapter 5.

Also magnetohydrodynamic instabilities are often associated with the presence of a magnetic field and currents in the plasma [113]. Results on detection of such a high-frequency movements of the plasma jet are presented in section 4.6.

In our studies we build on the results of Zhou Qing [83] and a theory and experiments on transport of plasmas in a magnetic field described by Schram in chapter 8 of [116]. Our investigations continue those studies and enter an unexplored range of parameters for a cascaded arc. In particular, we worked with arc currents of 60 to 100 A in a continuous magnetic field of 0.4–1.6 T., whereas the quoted experiments were done at arc currents up to 50 A in a magnetic field up to 0.2 T.

4.2 Influence of magnetic fields on the plasma jet

A drastic change in the expansion of the plasma jet happens already in a magnetic field of 0.4 T. Instead of hazy, faintly radiating, low-density plasma \( n_e \sim 10^{16} \text{ m}^{-3} \) as shown in chapter 3, an extremely bright narrow plasma jet appears (Figure 4.1).

![Figure 4.1: In comparison with a faint expanding plasma without a magnetic field (a) plasma in a magnetic field of 0.4 T (b) is a confined narrow bright jet.](image)

In a magnetic field charged particles move helically, following the magnetic field lines. The so-called Larmor radius \( r_L \) of the helix depends on the magnetic field strength \( B \) and the perpendicular to the field velocity component \( v_\perp \) of a
particle with mass $m$ and charge $e$:

$$ r_L = \frac{mv_{\perp}}{eB} \tag{4.1} $$

For example, the Larmor radius for electrons at the temperature of 1 eV in a magnetic field of 0.4 T is around $6 \cdot 10^{-6} \text{ m}$ and for hydrogen ions of the same temperature it is around $2.6 \cdot 10^{-4} \text{ m}$. Hence, charged particles are bound to a magnetic field and can diffuse across it only due to collisions with other particles. Hence, the ratio between the gyro-frequency

$$ \omega_c = \frac{eB}{m} \tag{4.2} $$

and the electron-ion and ion-ion collision frequency $\nu_{e,i}$

$$ \nu_e = \frac{1}{\tau_e} = 2.9 \cdot 10^{-12} n_e \ln(\Lambda) \frac{\hat{T}_e^{3/2}}{\hat{T}} \tag{4.3} $$

$$ \nu_i = \frac{1}{\tau_i} = 4.8 \cdot 10^{-14} n_i \ln(\Lambda) \frac{\hat{T}_i^{3/2}}{\hat{T}} \tag{4.4} $$

(here $\ln(\Lambda)$ is the Coulomb logarithm see also section 2.3.1), the Hall parameter $H$

$$ H_{e,i} = \frac{\omega_c}{\nu_{e,i}} = \omega_c \tau_{e,i} \tag{4.5} $$

is important. If $H > 1$ the plasma particles are magnetised (their mobility across the magnetic field is reduced). For our hydrogen plasma with electron and ion temperatures around 1 eV and density of the order of $10^{20} \text{ m}^{-3}$ in a magnetic field of 0.4–1.6 T the Hall-parameter for electrons is in the range of 10 to 50 (they are magnetised). The Hall-parameter for ions (if we consider ion-ion collisions) is lower by a factor of $\sqrt{M_i/2m_e}$. However, ion-ion collisions do not lead to diffusion in the first order approximation because there is no momentum transfer in a collision of equal particles. Only collisions with neutral gas contribute to diffusion. More detailed consideration of this is given in the discussion of this chapter (section 4.7.1), on the basis of the presented results of our measurements. After leaving the plasma source the ions expand due to collisions until the density has decreased to the value at which the ion gyro-frequency is of the same order as the collision frequency. However, as positive ions and electrons have to diffuse together the magnetisation of electrons leads to plasma confinement.

It is manifest from the photograph that a magnetic field of 0.4 T or higher reduces the expansion of the plasma jet.

We may therefore expect that the plasma density does not decay so rapidly as a function of the distance from the source, and, as there is no expansion cooling, that $T_e$ and $T_i$ can stay high for a longer distance, too.

The $T_e$ and $n_e$ profiles in the hydrogen plasma jet were measured by the Thomson scattering system (see sections 2.3.3 and 2.3.4 for description).
4.3. Effect of the nozzle geometry on the plasma jet parameters

The experiments with variation of the nozzle diameter showed two clear effects. First, increasing the nozzle diameter leads to an increased voltage (for the same arc current) and hence an increased power dissipation in the arc. Second, there is a clear effect on the \( n_e \) and \( T_e \) profiles measured at 35 mm from the nozzle.

The effect on the voltage is illustrated in Figure 4.3.

Figures 4.4 and 4.5 present the \( n_e \) and \( T_e \) profiles, for \( B = 0.4, 0.8, 1.2 \) and 1.6 T and for nozzle opening diameters of 5, 6, 7 and 8 mm.
CHAPTER 4. HYDROGEN PLASMA IN A HIGH MAGNETIC FIELD

Figure 4.3: Arc voltage in a magnetic field of 0.4–1.6 T for a nozzle diameter of 5, 6, 7 and 8 mm. The voltage increases together with both the field and diameter.

It is seen that the peak density does not depend on the nozzle diameter but scales only with the magnetic field. The full width of the \( n_e \) profile remains almost constant for all nozzle opening diameters in a magnetic field of 0.4 T but in higher magnetic fields increases from approximately 7 mm for 5 mm nozzle opening to approximately 10 mm for 8 mm nozzle opening. Another interesting observation is that flattening of the \( n_e \) profile is seen at increasing nozzle diameter. This flattening is not observed for the \( T_e \) profiles.

\( T_e \) increases from about 0.8 to 1.8 eV going from the 5 to the 8 mm nozzle, but interestingly, most of this increase is concentrated at the step from 6 to 7 mm. For different values of \( B \), the variation of \( T_e \) stays within the error bar of roughly 10%.

The combined effect of the wide nozzle and the application of a strong magnetic field has allowed the record values of \( n_e \), \( T_e \) and ion flux density, that bring Pilot-PSI in the ITER relevant range of plasma surface interaction.

### 4.4 Measurements of jet velocity components

Using the high-resolution optical emission spectroscopy we can derive the plasma jet velocity components measuring Doppler shifts of atomic lines. We used the hydrogen \( H_\beta \) line in our measurements (see sections 2.3.5 and 2.3.6). The radial or azimuthal velocity components are measured if the signal is detected in the direction perpendicular to the jet axis. The axial jet velocity can be evaluated by determining the Doppler shift in a direction close to the jet axis.
4.4. MEASUREMENTS OF JET VELOCITY COMPONENTS

Figure 4.4: The electron density profiles in Hydrogen plasma in a magnetic field of 0.4, 0.8, 1.2 and 1.6 T for nozzle opening diameters of 5 (a), 6 (b), 7 (c) and 8 mm (d). The peak density does not depend on the nozzle diameter but scales only with the magnetic field. The broadening and flattening of the plasma density profile is seen at increasing nozzle diameter.

4.4.1 Rotation velocity of the plasma jet

Figure 4.7 presents the measured rotation velocity profiles for B = 0.4-1.6 T. For a detailed description of the model that we use to derive the rotation velocity we refer to chapter 5. As the obvious candidate cause for the observed rotation is the $E \times B$-drift, it is useful to investigate the dependencies of the observed rotation velocity on B and the radial E field. For the latter, there is no direct measurement, but it as the radial E-field will be correlated to the voltage drop between the last plate and the nozzle, we shall take that as the observable.

Figure 4.7 shows that the rotation velocity increases approximately linearly with the B-field. The series of measurements in different magnetic fields conducted for different nozzle opening diameter show that the velocity increases approximately linearly with the nozzle diameter. Figure 4.7(a) shows the rela-
Figure 4.5: The $T_e$ profiles in hydrogen plasma for $B = 0.4$, 0.8, 1.2 and 1.6 T and for nozzle opening diameters of 5 (a), 6 (b), 7 (c) and 8 mm (d). $T_e$ grows significantly at the transition from approximately 6 to 7 mm. The difference in $T_e$ for different magnetic fields is insignificant.

4.4.2 The axial velocity of the plasma jet

The axial velocity component of the plasma jet in combination with the ion density determines the particle flux density in the jet, which is a very important quantity for PSI studies. We deduce the axial jet velocity from the Doppler shifts of the hydrogen $H\beta$ line using expression (2.13). The description of the measurements is presented in section 2.3.6.

It is important to note that the plasma light is collected at a small angle (of around 15°) from the axis of the plasma jet. First of all it means that it contains information on the Doppler shifts over a plasma length of the order of a few centimeters. Thus the derived velocities are averaged over this long
distance and over the cross-section. In addition, a projection of the rotation velocity component \( v_{\text{rot}} \) to the line of sight is added to the axial velocity \( v_{\text{axial}} \) (taking into account the sign of the projection depending on the direction of the rotation in the top and in the bottom of the jet as shown in a schematic figure 4.8). In fact, the measured value \( v_{\text{meas}} \) is:

\[
v_{\text{meas}} = v_{\text{axial}} \cos 15^\circ \pm v_{\text{rot}} \cos 75^\circ
\]  

(4.6)

These complications all together allow only an estimate of the axial velocity component of the plasma jet.

The axial velocity component was found to be in the range 2 to 5 km/s in a magnetic field of 0.4 T. Its radial profiles are shown in Figure 4.9.

The profiles clearly show where the projection of the rotation velocity is added or subtracted in the top and in the bottom of the plasma jet.

The axial variation of the velocity measured in the middle of the jet within approximately 10 cm downstream from the plasma source is shown in Figure 4.10. It decreases from approximately 5 km/s to 2 km/s downstream.
4.5 Integrated fluxes of ions and energy and efficiency of the source.

Now that we have measured the $n_e$, $T_e$ and flow velocity profiles, we can compute the integrated flux of electrons and ions for different values of B and nozzle diameter. These results are summarized in Figure 4.11. Dividing by the flux of hydrogen atoms that is fed into the source (2.5 slm corresponds to $2.25 \times 10^{21}$ hydrogen atoms/s), the total ionisation efficiency evaluated at 35 mm outside the nozzle is obtained (Figure 4.12). Other useful figures of merit are the energy efficiencies: energy per ion, and the total power in plasma beam over the total power fed into the source. To evaluate the latter, we take ionization potential of hydrogen (13.6 eV), a half of the dissociation energy of $H_2$ (4.4/2 eV), and thermal energy $3n_e T_e$ (equal for electrons and ions). This adds up to approximately 20.3 eV per ion, at 1.5 eV. The energy efficiencies are plotted in figures 4.13(a) and 4.13(b).

4.6 (In-)stability of the plasma jet in a magnetic field

The wobbling frequency was measured in the Pilot-PSI set-up with a photodiode (Thorlabs PDA55 - the Amplified Si detector, DC 10 MHz, 400-1100 nm). The light from a cross-section of the plasma jet at approximately 35 mm from the
4.6. (IN-)STABILITY OF THE PLASMA JET IN A MAGNETIC FIELD

Figure 4.8: Schematic drawing of the jet velocity components (top view). The measured velocity component is a sum of projections of the axial velocity and of the rotational velocity of the plasma jet. Note that the direction of the rotational velocity is opposite in the top and in the bottom of the jet.

Table 4.1: The frequencies of the wobbling of argon and hydrogen plasma jet in a magnetic field of 0.4 to 1.6 T.

<table>
<thead>
<tr>
<th>Gas</th>
<th>Magnetic field (T)</th>
<th>Arc current (A)</th>
<th>Wobble freq. (s⁻¹)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Argon</td>
<td>0.4</td>
<td>80</td>
<td>9.8 \cdot 10⁴</td>
</tr>
<tr>
<td>Hydrogen</td>
<td>0.4</td>
<td>80</td>
<td>7.7 \cdot 10⁵</td>
</tr>
<tr>
<td>Hydrogen</td>
<td>1.6</td>
<td>80</td>
<td>7.5 \cdot 10⁵</td>
</tr>
<tr>
<td>Hydrogen</td>
<td>1.6</td>
<td>100</td>
<td>8.5 \cdot 10⁵</td>
</tr>
</tbody>
</table>

nozzle was collected with a lens into the linear array of 40 fibres (the same that was used for HROES measurements, section 2.3.6) and the light from the edge of the jet was focused on the photodiode. The time-resolved signal was detected by an oscilloscope PC-card (National Instruments PCI-5112, 0-100 MHz) running with National Instruments software (NI Scope-SFP 1.5.1.). The detected signals for argon and hydrogen in a magnetic field of 0.4 T at the arc current of 80 A are shown in Figure 4.14.

The signals look quite periodical. All frequencies were found to be in the range $10^5 - 10^6$ s⁻¹ (Table 4.1). That is in a good agreement with estimations of the frequencies using expression (2.22). It is necessary to note that when the magnetic field is higher, the electron density is also higher and other parameters that determine the frequency also change. Finally, the frequency changes not significantly. The big difference in frequency between the argon and the hydrogen plasma jet is the mass of ions. In fact the difference in frequencies is approximately equal to the square root of the mass ratio: $\sqrt{m_{Ar}/m_H} \approx 6.3$.

To find out whether the plasma jet really moves at that high frequency a hundred shots of the plasma jet cross-section were recorded by a short-time gated CCD camera attached to a spectrometer. The spectrometer was tuned
CHAPTER 4. HYDROGEN PLASMA IN A HIGH MAGNETIC FIELD

Figure 4.9: The axial velocity profiles of the plasma jet derived from the Doppler shift of $H_\beta$ line. The measurement is made at several points along the jet (at 2, 25, 45, 70 and 90 mm downstream from the source exit) in a magnetic field of 0.4 T. The discharge current is 80 A, the gas flow rate is 2.5 slm of $H_2$.

to the wavelength range around the $H_\alpha$ line that is the strongest line in the spectrum of our plasma. The exposure time of a single shot was $1 \times 10^{-7}$ s (the gating time of the camera) that is approximately 10 times shorter than a wobbling period. The profiles of the jet in the highest and in the lowest detected positions are presented in Figure 4.15.

The top of the profiles is flat and their sides are steeper than that of the fitted gaussians. That means that the jet slightly moves within the time of shots.

To find an averaged shift of the plasma jet due to the wobbling these profiles were fitted by Gaussians and positions of the centres of the Gaussians were gathered as a histogram (Figure 4.16). The histogram was fitted by a Gaussian with the half-width at 1/$e^2$-level of around 0.8 mm. This half width we consider as the maximal wobbling radius (95% of the shots show the shift of the jet within this range). The wobbling radius is bigger than the error bar in determining the width of every jet profile that we get from the Gaussian fit ($\Delta w_p = 0.2$ mm) and bigger than the error bar in determining the position of the centre of the jet profile fitted by the Gaussian ($\Delta x_c = 0.1$ mm).

We consider this as a sufficient proof of the wobbling of the plasma jet in a high magnetic field. However the radius of the wobbling can be bigger as the jet moves within the time of shots and consequently the Gaussian fits of the profiles give understate values of the shifts.
4.7 Discussion of the results on the hydrogen plasma in a high magnetic field

The model that explains all the above mentioned results assumes that in a strong magnetic field the arc current can not close the electrical circuit from the plasma channel to the anode within the plasma source due to the reduced electrical conductivity across the field [123]. Hence the current exits the source, flows axially with the jet and expands slowly radially all along the plasma jet and only then returns to the anode as it is schematically shown in Figure 4.19. Thus the effect becomes stronger with increasing magnetic field and with increasing nozzle opening radius (Figure 4.7). This effect leads to the significant increase of the arc voltage at constant current which has been observed at higher magnetic fields and larger nozzle opening diameters (Figure 4.17). Higher power input into plasma ends up in production of more electrons and higher electron temperature as it has been observed by the Thomson scattering measurements (Figures 4.4 and 4.5). The arising radial electric field leads to plasma rotation due to the $E \times B$-drift. Eventually, the energy is transferred to ions via viscosity. The increase of the electron temperature at larger nozzle opening diameters is also explained by this.

The radial electron density profile in a magnetic field of 0.8 T is higher...
and narrower than the one at 0.4 T. That means simply a better confinement. However, in a higher magnetic field the profiles become higher but not narrower. The total number of electrons is higher in a higher magnetic field (Figure 4.11).

Comparing the density profiles at different nozzle opening diameters but at constant magnetic field we see that the profiles at higher diameters become broader but never higher. It means that this is the limit of the density that can be confined by this magnetic field and due to collisions the plasma jet expands until the density drops to the value at which the collision frequency of the ions is approximately equal to the ion gyro-frequency or in other words till ions are no longer magnetised.

A flattening of the plasma density profile is seen at increasing nozzle diameter. It is not clear yet whether it is a result of integrating the scattered light over many periods of wobbling of the plasma jet (subsections 2.3.5 and 4.6).

Relying on the expression (3.5) we find that a higher electron temperature is possible in the jet where the density of electrons and neutrals is much lower than in the high-pressure plasma source.

### 4.7.1 Diffusion of charged particles across the magnetic field

Diffusion of charged particles across a magnetic field is determined by the Hall parameter (4.5) that is the ratio between the gyro-frequency of particles (4.2) and the collision frequency of the particles (expressions (4.3) and (4.4)). For electrons in the first order approximation we have to take the elastic electron-ion collision frequency as dominating the electron momentum transfer. Collisions
4.7. DISCUSSION OF THE RESULTS IN A MAGNETIC FIELD

Figure 4.12: The ionisation efficiency - ratio of the total flux of produced ions and the flux of hydrogen atoms fed into the plasma source - in a magnetic field of 0.4–1.6 T for a nozzle diameter of 5 and 8 mm. The measurement is made at 35 mm from the source exit, a discharge current is 100 A, gas flow rate is 2.5 slm of $H_2$.

with other electrons do not contribute to the momentum transfer and elastic collisions of electrons with neutrals are much less frequent. Thus we estimate the electron-ion collision frequency at an average electron temperature $\hat{T}_e$ of 1 eV and electron density $n_e$ of around $3 \cdot 10^{21} m^{-3}$ and (taking the Coulomb logarithm $\ln(\Lambda) \approx 7$) from the expression (4.3):

$$\nu_{ei} = \frac{1}{\tau_{ei}} = 2.9 \cdot 10^{-12} \frac{n_e \ln(\Lambda)}{\hat{T}_e^{3/2}} \approx 6 \cdot 10^9$$  \hspace{1cm} (4.7)

The Hall parameter for electrons in a magnetic field from 0.4 to 1.6 T is in the range of approximately 10 to 50. Hence electrons are always magnetised.

The gyro-frequency of ions is $M_i/m_e$ times less than that of electrons. At equal temperatures ($T_i = T_e$) the ion-ion collision time is $\sqrt{M_i/m_e}$ times larger. Thus the ion-ion Hall parameter is smaller by this factor. Estimation of the ion-ion Hall parameter from expressions 4.4 and 4.5 gives values in a range 0.5–1.5 for our conditions (the ion densities of $2–7 \times 10^{20} m^{-3}$, the ion temperature of 1–2 eV in a magnetic field of 0.4–1.6 T). For ion diffusion we need also to consider ion-neutral collisions as ion-ion collisions do not lead to diffusion in the first order approximation. This is because ion-ion collisions do not contribute to momentum transfer as they have equal masses. Electron-ion collisions do change the momentum of the ions only very slightly due to the negligibly small mass of electron. Collisions of atomic ions $H^+$ with molecular ions $H_2^+$ are rare due to the significantly lower density of the latter. The ion-neutral collision frequency was estimated for a neutral density of around $3 \cdot 10^{20} m^{-3}$ and a
collision constant $C_{ii}^0$ for hydrogen of the order of $10^{14}$ [116, 125, 126]:

$$\nu_{i0} = \frac{1}{\tau_{i0}} = \frac{n_0}{C_{ii}^0} \approx 3 \cdot 10^6$$  \hspace{1cm} (4.8)

From this we find the Hall parameter for ions in a magnetic field of 0.4 to 1.6 T also in the range 10–50. That is occasionally of the same order as the one for electrons. Thus, ions are also magnetised.

Diffusion of electrons and ions across a magnetic field is reduced depending on their Hall parameter:

$$D_\perp = D_\parallel \cdot \frac{1}{1 + \omega_c^2 \tau^2}$$  \hspace{1cm} (4.9)

Here, $D_\parallel$ is diffusion coefficient in the absence of a magnetic field or parallel to it. For electrons the diffusion coefficient is:

$$D_{e\parallel} = \frac{kT_e}{m_e \tau_{e\parallel}} \approx 30$$  \hspace{1cm} (4.10)

and in a magnetic field of 0.4 to 1.6 T it is in the range 0.3 down to 0.01. For ions it can be estimated in a similar way:

$$D_{i\parallel} = \frac{kT_i}{M_i \tau_{i\parallel}} \approx 30$$  \hspace{1cm} (4.11)

As we see it is also accidentally very close value as for electrons. As the Hall parameter occurred to be the same for electrons and ions the diffusion coefficient for ions at our preset magnetic fields of 0.4 to 1.6 T is also in the range 0.3 down
4.7. DISCUSSION OF THE RESULTS IN A MAGNETIC FIELD

(a) argon

(b) hydrogen

Figure 4.14: The detected signal of a photodiode from the edge of the wobbling argon (a) and hydrogen (b) plasma jet. Both measurements are made at 35 mm from the source exit in a magnetic field of 0.4 T. Discharge current is 80 A, gas flow rate is 2.5 slm of argon or hydrogen.
Figure 4.15: The CCD shots ($H_{\alpha}$-line radiation) of the wobbling plasma jet with the exposure time of $10^{-7}$ s. The profiles fitted by gaussians show clear shifts with respect to each other. The measurements are made at 35 mm from the source exit in a magnetic field of 0.4 T. The discharge current is 80 A, the gas flow rate is 2.5 slm hydrogen.
4.7. **DISCUSSION OF THE RESULTS IN A MAGNETIC FIELD**

Figure 4.16: The histogram: positions of the centre of the plasma jet profile on the CCD shots. The histogram was fitted by a gaussian with the half-width at $1/e^2$-level of around 0.8 mm.

It is interesting to note that in such a case there is no significant charge separation!

As the diffusion of both electrons and ions is strongly reduced across a magnetic field it becomes insignificant within the life-time of the particles in the jet as at an average axial jet velocity of about 5000 m/s particles are lost at the end of the 1 m long vessel in less than a millisecond.

### 4.7.2 Rotation of the plasma jet in a magnetic field

Rotation of the plasma jet is caused by the drift of charged particles in crossed electric and magnetic fields: the radial electric field and the axial magnetic field (Figure 4.18):

$$\vec{v}_{drift} = \frac{\vec{E} \times \vec{B}}{B^2}$$  \hspace{0.5cm} (4.12)

A high $E \times B$-drift velocity means a high radial electric field in the plasma column. In our plasma two reasons for the arising radial electric field can be
CHAPTER 4. HYDROGEN PLASMA IN A HIGH MAGNETIC FIELD

Figure 4.17: The discharge voltage grows with increasing nozzle opening diameter and increasing magnetic field at constant current. Discharge current is 100 A, gas flow rate is 2.5 slm of $H_2$.

Figure 4.18: Scheme of rotation of the plasma jet due to the ExB-drift.

considered. They are the ambipolar field that has internal causes and the external electric field of the power supply that is related to the arc current due to the possible charge separation because of the difference in mobilities of positive ions and negative electrons.

The ambipolar electric field $\vec{E}$ (4.13) [116]

$$\vec{E} = \frac{D_i - D_e}{\mu_i + \mu_e} \nabla n_e$$

(4.13)

is related to the charge separation in the plasma due to the difference in mobility of electrons $\mu_e$ and ions $\mu_i$. Here, $D_i$ and $D_e$ are the diffusion coefficients of ions and electrons respectively, and $n_e$ is the electron density. The diffusion coefficients are related to the mobility of the particles via Einstein’s relation:

$$\frac{D_{i,e}}{\mu_{i,e}} = \frac{k_B T_{i,e}}{e} = \hat{T}_{i,e}$$

(4.14)
4.7. DISCUSSION OF THE RESULTS IN A MAGNETIC FIELD

Thus at \( \tilde{T}_e \approx \tilde{T}_i \approx 1 \text{ eV} \) in such a magnetic field the ambipolar radial electric field is of the order of 200 V/m while the radial electric field must be approximately of 16000 V/m to cause the "\( \mathbf{E} \times \mathbf{B} \)" drift velocity of up to 10000 m/s, as it has been detected (Figure 4.7).

The electric field can be related also to macroscopic radial currents in the plasma. According to our model, the arc current can not close the electrical circuit from the plasma channel to the anode within the plasma source due to the reduced electrical conductivity across the field [123]. Hence the current exits the source, flows radially all along the plasma jet and only then returns to the anode as it is schematically shown in Figure 4.19. This radial current is related to a radial electric field in the plasma jet (the Ohm’s law): \( \mathbf{j} = \sigma \mathbf{E} \). The arising electric field induces the \( \mathbf{E} \times \mathbf{B} \)-drift of charged particles (independent of the charge sign) leading to the plasma rotation. In a magnetic field of 1.6 T an extra voltage of approximately 90 V drops over the plasma in the vessel (see Figure 2.8), mainly over the radius of the jet (\( \approx 5 \text{ mm} \)) because the conductivity of plasma perpendicular to a magnetic field \( \sigma_\perp \) is smaller than that along the magnetic field \( \sigma_\parallel \) due to reduced mobility \( \mu_\perp \) of charged particles across a magnetic field ([128], pp. 242-247):

\[
\mu_\perp = \mu_0 \frac{1}{1 + \omega_c^2 \tau^2}
\]  

(4.15)

Here, \( \omega_c = eB/m \) is the cyclotron frequency, and \( \tau \) is the characteristic time between collisions.

A larger diameter of the nozzle opening decreases the number of electrons that reach the anode within the plasma source. It leads to increase of the current exiting the plasma source and thus causes a higher potential difference over jet radius and consequently strengthening of the plasma rotation. When

![Figure 4.19: The electrons cannot close the circuit within the source due to their limited mobility across the magnetic field but they exit the source, diffuse radially and return back to the anode along the magnetic field lines.](image-url)
Table 4.2: Characteristic time (CT) and mean free path (MFP) of collisional, 
gyronal and radional processes in the hydrogen plasma with the neutral 
density of $10^{21} m^{-3}$, the electron density of $5 \times 10^{20} m^{-3}$ the electron and ion 
temperatures of 1 eV in a magnetic field of 1.6 T

<table>
<thead>
<tr>
<th>Process</th>
<th>CT</th>
<th>scale (s)</th>
<th>MFP</th>
<th>scale (m)</th>
</tr>
</thead>
<tbody>
<tr>
<td>ion-neutral collisions</td>
<td>$\tau_{i0}$</td>
<td>$3 \cdot 10^{-6}$</td>
<td>$\lambda_{i0}$</td>
<td>$3 \cdot 10^{-2}$</td>
</tr>
<tr>
<td>$n = 4 \rightarrow n = 2$ transition</td>
<td>$1/A_{ik}$</td>
<td>$1.2 \cdot 10^{-7}$</td>
<td>$\lambda_{4\rightarrow 2}$</td>
<td>$1 \cdot 10^{-3}$</td>
</tr>
<tr>
<td>Electron-ion energy transfer</td>
<td>$\tau_{ei}^\epsilon$</td>
<td>$5 \cdot 10^{-8}$</td>
<td>$\lambda_{ei}^\epsilon$</td>
<td>-</td>
</tr>
<tr>
<td>Gyration of ions</td>
<td>$2\pi/\Omega_i$</td>
<td>$4.1 \cdot 10^{-8}$</td>
<td>$\rho_i$</td>
<td>$1.2 \cdot 10^{-4}$</td>
</tr>
<tr>
<td>Ion-ion collisions</td>
<td>$\tau_{ii}$</td>
<td>$6 \cdot 10^{-9}$</td>
<td>$\lambda_{ii}$</td>
<td>$6 \cdot 10^{-5}$</td>
</tr>
<tr>
<td>Gyration of electrons</td>
<td>$2\pi/\Omega_e$</td>
<td>$2.2 \cdot 10^{-11}$</td>
<td>$\rho_e$</td>
<td>$3 \cdot 10^{-6}$</td>
</tr>
<tr>
<td>Electron-ion collisions</td>
<td>$\tau_{ei}^m$</td>
<td>$6 \cdot 10^{-11}$</td>
<td>$\lambda_{ei}^m$</td>
<td>$3 \cdot 10^{-4}$</td>
</tr>
<tr>
<td>Debye radius</td>
<td>-</td>
<td>-</td>
<td>$r_D$</td>
<td>$3 \cdot 10^{-7}$</td>
</tr>
</tbody>
</table>

the rotational velocity approaches the local thermal velocity the ion viscosity 
leads to ion heating and (for positive rotation) to higher densities.

Comparison of collision scale for electron-ion energy transfer and for expan-
sion and rotation (Table 4.2) shows that electrons and ions are fully thermally 
coupled on the length scale of plasma jet radius ($\sim 0.5$ cm)

Another interesting detail about $E \times B$-drift of particles is the fact that ions 
have an upper limit of the rotational velocity due to the viscosity and inertia 
while electrons do not have this limit. The momentum balance equations for 
ions and for electrons differ significantly ([116], section 8.4, equations (8.85) and 
(8.86)):

$$n_i m_i (\vec{v}_i \cdot \nabla) \vec{v}_i + \nabla p_i + \nabla \cdot \Pi_i^{jk} = \rho_i (\vec{E} + \vec{v}_i \times \vec{B}) - \vec{R}_{ie} - \vec{R}_{i0} - \vec{M}_s$$  \hspace{1cm} (4.16)

$$\nabla p_e = - \rho_n (\vec{E} + \vec{v}_e \times \vec{B}) + \vec{R}_{ie}$$  \hspace{1cm} (4.17)

Here, $p_i,e$ denotes the pressure, $\Pi_i^{jk}$ is the viscosity tensor, $\vec{R}$ stands for the 
momentum transfer due to collisions, $\vec{M}_s$ is the momentum transfer due to the 
mass source term. In the momentum equation for electrons the inertia term 
$n_e m_e (\vec{v}_e \cdot \nabla) \vec{v}_e$ is neglected due to $m_i \gg m_e$, as well as the 
viscosity term. The 
term due to collisions with neutrals $\vec{R}_{i0}$ is neglected either, as compared to the 
$\vec{R}_{ei}$ (as $\tau_{ei} \ll \tau_{e0}$) whereas $\vec{R}_{i0}$ can not be neglected compared to $\vec{R}_{ie}$. The 
reason is that in the ion momentum equation $\vec{R}_{i0}$ has no contribution: one must have $\vec{R}_{ie}$ which is equal to (apart from the "minus" sign) $-\vec{R}_{ei}$. Therefore $\tau_{e0}$ should not only be longer than $\tau_{ei}$ (it has no relevance for diffusion in the first 
order, but only for ion heat conduction) but also be larger than:

$$\tau_{ie}^m = \frac{M_i}{m_e} \tau_{ei}^m$$  \hspace{1cm} (4.18)

These differences in the momentum equations for electrons and ions mean 
that that electrons are thermal but ions may seriously be affected by inertia, 
viscosity, collisions with neutrals.
4.8. RESISTANCE OF THE PLASMA COLUMN

The difference in the rotational velocity of ions and electrons can cause the azimuthal current in the jet. Note, that the $\mathbf{E} \times \mathbf{B}$-drift depends neither on mass of the particles nor on the charge sign.

In fact, the measured profile of the rotational velocity represents also the distribution of the radial electric field in the jet. As the magnetic field is known and varies not significantly within the plasma jet volume (Figure 2.3) the electric field can be estimated from the formula (4.12).

4.8 Resistance of the plasma column

To complete our consideration of the model of a magnetised plasma jet let us now estimate resistance of the jet with respect to the current that exists outside the plasma source in a magnetic field. We choose the most extreme case with the diameter of the nozzle opening of 8 mm in a magnetic field of 1.6 T. The peak electron temperature in this case is almost 2 eV and the arising extra potential difference between the last cascaded plate and the anode is around 90 V at the arc current of 80 A.

For estimation of resistance of the plasma column we use the same Spitzer’s expression for plasma conductivity as we have used in chapter 3 [72]:

$$\sigma = \frac{2 \cdot 10^4 T_e^{3/2}}{\ln(\Lambda)}$$

(4.19)

The arc conductivity according to Spitzer depends only on the electron temperature $T_e$ (here, $\ln(\Lambda)$ is the Coulomb logarithm that is a very slow varying function of the electron temperature and density). We consider the plasma jet as consisting of two coaxial cylinders: one in the centre of the jet and another one (hollow) around it. The first cylinder has the same diameter as the arc channel ($d_1 = 2r_1 = 4$ mm) and the electron temperature is high in it ($T_1 = 2$ eV). The inner diameter of the second cylinder coincides with the outer diameter of the first one (4 mm) and the outer diameter approximately corresponds to the nozzle opening (for the current to close the loop along magnetic field lines) ($d_2 = 2r_2 = 8$ mm). The electron temperature in this outer layer is lower ($T_2 = 1$ eV). With the plasma column length $l$ of about 0.5 m we obtain:

$$R_1 = \frac{l}{\sigma_1 \pi r_1^2} \approx 4.9 \text{ Ohm}$$

(4.20)

and

$$R_2 = \frac{l}{\sigma_2 \pi (r_2^2 - r_1^2)} \approx 4.6 \text{ Ohm}$$

(4.21)

Resistances are close in value because the electron temperature in the first cylinder is higher but its cross-section is smaller.

The value of the current outside the source is not known and it is very difficult to measure directly. However, it is possible to estimate it now in frames of our simple consideration from the known potential difference of 90 V and the estimated resistance. Thus the current is around 10 A.
CHAPTER 4. HYDROGEN PLASMA IN A HIGH MAGNETIC FIELD
Chapter 5

Plasma Jet Rotation diagnosed with High-Resolution Optical Emission Spectroscopy

Abstract

The rotation of the hydrogen plasma jet in Pilot-PSI is clearly revealed in the line shape of hydrogen Balmer-β light. Quantification of this rotation is obtained by exploring the observed asymmetric line shape. A physical model is proposed that explains the asymmetry. It assumes that the line is composed of a contribution from $H(n = 4)$ atoms coupled to the plasma ions and $H(n = 4)$ atoms that have collided with neutral particles. This model is used to determine the rotation velocity profiles. The result was a peak rotation velocity that increases proportionally to the magnetic field strength up to 10 km/s at $B=1.6$ T. It corresponds to the rotation frequency at the axis of $\sim 5 \cdot 10^6$ rad/s. These high values are explained by the development of an additional potential drop between the last source plate and anode due to the arc current extending from the plasma source into the magnetised jet. This potential drop is proportional to $B$ and it is shown to be in agreement with a radial electric field that causes the rotation via $E \times B$ drift. The axial variation of the rotation showed that the rotation persists over 0.5 m. This is interpreted as a map of the discharge current that runs outside of the source to cross the magnetic field and returns to attach to the nozzle. The ion temperature derived from the Doppler broadening was found to be systematically higher than the electron temperature. Viscous heating of the ions of the ions due to their rotation is shown to be significant and may cause this temperature difference.
5.1 Introduction

A characteristic property of a magnetised plasma column is that it will rotate [116]. Any radial electric field perpendicular to the magnetic field will induce a drift motion in the gyration of the charged particles and because of cylinder symmetry, this will cause rotation. A radial electric field always exists in a plasma. Even if the jet is current free, charge separation will occur due to the difference in mobility for ions and electrons across the magnetic field.

In the previous chapter we explained additional ionization from a modified nozzle geometry by discharge current crossing the magnetic field outside of the plasma source. If this current is indeed running beyond the nozzle, this must induce significant electric fields and thus rotation of the plasma jet. In this chapter we will probe the rotation and determine to what extent the discharge current leaves the nozzle.

Emission spectroscopy is the obvious diagnostic to probe at collective movements of particles via the Doppler shift in the light emission. However, the protons in a hydrogen plasma do not emit light and only the neutral atoms (and molecules) can be assessed via spectroscopy. Thus, the first question is whether the radiating atoms are coupled to ions and reflect their characteristics (e.g. rotation). To answer this question let us consider the history of radiating particles: processes that populate them.

5.2 Origin of H\(_{\beta}\) light

Atomic H\(_{\beta}\) light emitted by the plasma jet originates from neutral atoms. In order to use this light as a probe for the properties of the ions, it is obviously required that in some way the emitting atoms and ions are coupled. This is actually the reason that we have chosen to use the H\(_{\beta}\) line that originates from the \(n = 4\) level. As will be discussed later, the cross section of an excited atom increases roughly with the square of the main quantum number \(n\). Atoms in higher excited states have therefore a better probability to be coupled via collisions with ions. We didn't consider higher lines because these appeared to be at too low intensities to be practical. We investigate whether the coupling may be expected and start from the mechanisms that populate the \(n = 4\) level of hydrogen.

The simplest production pathway of H\((n = 4)\) is direct excitation by electron impact of ground state atoms H\((1s)\):

\[
H(1s) + e^- \rightarrow H(n = 4) \rightarrow H(1s) + h\nu \quad (5.1)
\]

Also important can be dissociative excitation by electron impact:

\[
H_2 + e^- \rightarrow H_2^+ + e^- \rightarrow H(n) + H + e^- \quad (5.2)
\]

Before explaining the creation of the atom, we note that the energy difference between H\((1s)\) and the first excited level H\((n = 2)\) is already 10.2 eV. This means
that the electrons in the 1 eV Pilot-PSI plasma do not carry sufficient energy to excite significantly H(1s). Recombination of ions may be a more effective way to produce excited atoms (H\(^*\)) [83]. For example, radiative recombination of a hydrogen ion with an electron is the possibility:

\[ H^+ + e^- \rightarrow H^* \rightarrow H(1s) + h\nu \] (5.3)

This process has a very low rate (\(\sim 10^{-20} \text{ m}^3/\text{s}\)) because the total angular momentum of the system must be conserved while the reaction ends with only one particle. In three-particle recombination, the conservation of angular momentum is taken care of by a third particle:

\[ H^+ + e^- + e^- \rightarrow H^* + e^- \] (5.4)

However, the rate of three particle recombination is again low as it requires three particles to collide at once (see Table 5.1).

In Molecular Activated Recombination (MAR) [18, 19, 49] an excited atom is created in two consecutive steps. First, a molecular ion is produced via charge exchange between an ion and a low-temperature background gas molecule:

\[ H^+ + H_2 \rightarrow H(1s) + H^+_2 \] (5.5)

This reaction is endothermic (-2.1 eV), which at the observed temperatures of above 1 eV and probable excitation rotational and vibrational excitation will not have a large effect on the rate. Subsequently, fast dissociative recombination of the molecular ion results in two hydrogen atoms:

\[ H^+_2(r, v) + e^- \rightarrow H^* + H(1s) \] (5.6)

of which one is carrying an excess of internal energy and ends up at least in \(n = 2\) or most probably in \(n = 3\). If the intermediate molecular ion was rotationally-vibrationally excited \((v \geq 4)\), this reaction would produce directly H\((n = 4)\). Alternatively, H\((n = 4)\) is produced by direct electron excitation from the \(n = 2\), \(n = 3\) levels.

At the prevalent conditions with \(n_e \approx n_i\) of the same order as \(n_{H_2}\) the first step (5.5) in the MAR sequence is the slowest one (\(<\sigma v> = 5 \cdot 10^{-15} \text{ m}^3/\text{s}\)) and determines the rate of MAR (see Table 5.1).

For completeness we should also consider recombination of negative ions as a population mechanism for H\(^*\)(\(n\)). If negative ions are produced by dissociative attachment to H\(_2\)(\(r, v\)):

\[ H_2(r, v) + e^- \rightarrow H^- + H \] (5.7)

This may also lead to excited atoms by mutual recombination of H\(^-\) and H\(^+\) resulting in excited H atoms. In this case H\((n = 4)\) can be reached easily. However, at the plasma densities encountered in Pilot-PSI the production channel of negative ions is probably too low to be significant so we neglect it.
Table 5.1: Reaction rate constants [136, 137] of the processes that populate \( n = 4 \) excited hydrogen atoms (at \( T_{\text{e}} = 1 \text{ eV} \)). (* the rate for this reaction is in \( m^6s^{-1} \))

<table>
<thead>
<tr>
<th>Reaction</th>
<th>rate constant (( m^6s^{-1} ))</th>
</tr>
</thead>
<tbody>
<tr>
<td>( H^+ + H_2 \rightarrow H + H_2^* )</td>
<td>( 5 \times 10^{-15} )</td>
</tr>
<tr>
<td>( H_2^*(v, v) + e^- \rightarrow H(n \geq 2) + H )</td>
<td>( 8 \times 10^{-14} )</td>
</tr>
<tr>
<td>( H^+ + e^- \rightarrow H(n \geq 2) )</td>
<td>( \sim 10^{-20} )</td>
</tr>
<tr>
<td>( H^+ + H(n = 4) \rightarrow H(n = 4) + H^+ )</td>
<td>( 3 \times 10^{-13} )</td>
</tr>
</tbody>
</table>

An important loss channel of \( H(n = 4) \) at low electron densities is radiation (Paschen-\( \alpha \), Balmer-\( \beta \) and Lyman-\( \alpha \)):

\[
H(n = 4) \rightarrow H(n = 3) + P_{\alpha} \tag{5.8}
\]

\[
H(n = 4) \rightarrow H(n = 2) + H_{\beta} \tag{5.9}
\]

\[
H(n = 4) \rightarrow H(n = 1) + L_{\gamma} \tag{5.10}
\]

with transition probabilities \( A_{h_i} \) of \( 0.9 \cdot 10^7 \text{ s}^{-1} \), \( 0.84 \cdot 10^7 \text{ s}^{-1} \) and \( 1.3 \cdot 10^7 \text{ s}^{-1} \) respectively. A similar loss is encountered for the \( n = 3 \) \( (H_{\alpha}) \), transition probability of \( 4.4 \cdot 10^7 \text{ s}^{-1} + L_{\beta}, 5.6 \cdot 10^7 \text{ s}^{-1} \) and the \( n = 2 \) levels \( (L_{\alpha}, 4.7 \cdot 10^8 \text{ s}^{-1} \). Thus the radiative loss time of \( n = 4 \) is almost \( 3 \cdot 10^7 \text{ s}^{-1} \) if \( L_{\gamma} \) is optically thin and \( 1.7 \cdot 10^7 \text{ s}^{-1} \) if it is optically thick.

De-excitation rates for atomic hydrogen can be found in van der Mullen [138]. These are in the range of \( 10^{-12} \text{ m}^3/\text{s} \) for \( n = 3 \) at \( T_{\text{e}} \approx 1 \text{ eV} \). So the balance of population and de-population processes can be written as \( n_{H^+} \equiv n_i \approx n_e \):

\[
n_{H^+} n_{H_2} k_{cx} = n_{H_2^*} n_e k_{dr} \tag{5.11}
\]

\[
n_{H_2^*} n_e k_{dr} = n_{H(n=3)} (n_e K_3 + \sum_{i<3} A_{3i}) \tag{5.12}
\]

with \( cx \) and \( dr \) denoting "charge exchange" and "dissociative recombination" respectively, and \( K_3 \equiv \sum_{i \neq 3} k_{3i} \).

\[
n_{H(n=4)} = \frac{n_{H(n=3)}}{n_e K_4 + A_4} \tag{5.13}
\]

with \( K_4 \equiv \sum_{i \neq 4} k_{4i} \) and \( A_4 \equiv \sum_{i<4} A_{4i} \).

If \( n_e > (\sum_{i} A_{3i})/K_3 \) (and thus also \( n_e > (\sum_{i} A_{4i})/K_4 \)) and \( n_e/n_{H_2} \sim n_i/n_{H_2} > 10^{-2} \) then we find for \( H(n = 4) \) a balance between underlying processes:

\[
n_{H^+} n_{H_2} k_{cx} \simeq n_{H(n=4)} n_e K_4 \tag{5.14}
\]

The important consequence of 5.14 is that the \( H_{\beta} \) line emission is only proportional to \( n_{H_2} \) and not to \( n_e \).
5.2. ORIGIN OF H\(\beta\) LIGHT

At sufficiently high electron densities \((n_e > 10^{19} \text{ m}^{-3})\) de-excitation of \(H(n = 4)\) level overweightes radiation and the line emission becomes independent of \(n_e\) and depends only on \(n_{H_2}\). Thus, the highest light emission can occur not at the highest electron density in the centre of the jet. On the contrary, 5.14 predicts a hollow emission profile because the neutral molecular density \(n_{H_2}\) will have a hollow profile. The electron temperature peaks at axis and thus equal partial pressure means lower density. On top of this, burn out of neutrals in the centre can occur, as will be discussed in more detail below.

Let us now consider which temperature and velocity is to be expected on the basis of the formation process. The first reaction (5.5) gives an \(H_2^+\) ion with initially averaged energy (and velocity) of intermediate those from \(H_2\) molecules and ions. But within \(\mu\)s the \(H_2^+\) ions will be thermalised in ion-ion collision and approach the ion temperature and \(E\times B\)-drift velocity. In the subsequent steps the \(H_2^+\) ion will dissociatively recombine to \(H(n = 3)\) (and \(H(n = 1)\)), and still have ion-like behaviors. The excitation to \(n = 4\), will not change this. Thus the radiating \(H(n = 4)\), that we observe, will have a velocity and temperature resembling that of the ion population. Even if temperatures and velocity are lower due to the partial neutral history, they are coupled very effectively to the atomic ions by resonant charge exchange:

\[
H^+ + H^* \rightarrow H^* + H^+ \quad (5.15)
\]

In one collision, the excited atoms will thermalise with the ions and have the ion temperature and velocity, thus also the rotational velocity component corresponding to the \(E\times B\)-drift. The charge exchange process has a high reaction rate. Not only it is resonant but also the cross-section of an excited atom increases roughly with the fourth power of the quantum number as the Bohr radius \(r_b\) of the atom grows as \(n^2\):

\[
r_b = \frac{e_0 h^2}{\pi c^2 m_e n^2} \quad (5.16)
\]

Thus the cross-section is \(\sim 2 \cdot 10^{-18} \text{ m}^2\) for \(n = 4\) in comparison with \(\sim 1 \cdot 10^{-20} \text{ m}^2\) for the ground state.

This process will couple even more the ion temperature and velocity to the \(H(n = 4)\) atoms; this will remain true for the rotation velocity if the mean free path of charge exchange of \(H(n = 4)\) and \(H^+\) ion \(\lambda_{mfp} (n = 4, i)\) is smaller than the radius.

As the radiative lifetime of \(H(n = 4)\) is sufficiently short \((0.3\text{--}0.6 \cdot 10^{-7} \text{ s})\), depending on whether the \(L_\gamma\) is optically thin), the radiation will give a direct vision on the \(n = 4\) distribution which is (initially) coupled to the ion distribution. However, in particular further outside, there are encounters with neutrals: \(H\) atoms and \(H_2\) molecules. Collisions with \(H\) atoms are also resonant and have also large rates \((\sim 3 \cdot 10^{-14} \text{ m}^3/\text{s})\). Thus part of the \(H(n = 4)\) atoms will adopt a neutral, colder distribution. As neutrals in the ground state themselves have much lower collision rates they cannot follow ion rotation: they will have zero or small velocity and the low temperature of \(\sim 2000\text{--}4000 \text{ K}\) of the background...
CHAPTER 5. PLASMA JET ROTATION

gas. These latter temperatures are a result of a balance between energy input from the source and heat conduction to the walls of the vessels. In conclusion, the $H_\beta$ light will carry information both on the hot rotating ions and on the cold background gas.

The dissociation degree is high in the plasma source. Thus the molecules are formed mostly at the wall of the vacuum vessel and then return to the plasma jet. They are assumed to be rotationally-vibrationally excited [89] and to have temperatures of about 0.2 eV [50]. The penetration depth of the molecules to the plasma jet can be estimated as the mean free path of neutral-ion charge exchange collisions (reaction 5.5). With the ion density at the edge of the jet of larger than $2 \cdot 10^{20} \text{ m}^{-3}$ (as measured with Thomson scattering, see section 4.2) and an average velocity of $3 \cdot 10^3 \text{ m/s}$ the mean free path is smaller than 3 mm. Hence, the plasma jet is not transparent for molecules. As radiation originates mainly via MAR, we expect a hollow emissivity profile because the $H_2$ density profile is hollow. This means that emissivity is maximal, where $H_2$ is higher and the electron density has decayed already to a substantially lower value, but not so low that still electrons can dissociatively recombine and excite to the $H_2(n = 4)$ level. This requires at least that $n_e k_{34} \geq \sum_i A_{3i}$, which with appropriate values for $A_{3i} (\sim 10^8 /\text{s})$ and $k_{34} (\sim 10^{-12} \text{ m}^3 /\text{s})$ gives a value around $10^{20} \text{ m}^{-3}$. Hence, we expect to see this value rather independent of the central value of $n_e$.

5.3 The $H_\beta$ line measured at Pilot-PSI

The High-Resolution Emission Spectroscopy (HiRES) measurements are described in detail in section 2.3.6. Here we repeat in short that light is collected perpendicularly to the plasma jet and relayed to the spectrometer via an array of 40 individual fibers. In this way the entire jet cross-section is covered simultaneously. Figure 5.1 shows part of a raw CCD-image of the spatially (vertical axis) as well as spectrally (horizontal axis) resolved $H_\beta$ line. The measurement was done at $B = 1.6 \text{ T}$, 80 A discharge current, and 2.5 slm hydrogen flow. The horizontal axis in the figure represents wavelength. The complete width of the CCD-image corresponds approximately to 1 nm (only the central part of $\sim 0.35 \text{ nm}$ is shown in the figure) at the wavelength of $H_\beta$ line (486.13 nm). The vertical axis is the spatial coordinate (radius of the jet). Taking into account a magnification factor of two due to the collecting lens it covers approximately 40 mm (of which only $\sim 20 \text{ mm}$ is shown in the figure). The different bands correspond to the different fibres in the bundle. The finer structures within these bands are due to imperfections in the spectrometer slit.

The intensity profile shows a broad peak that corresponds to the visually observed plasma jet profile. It is clearly seen that the line is Doppler shifted to the red in the bottom of the plasma jet, to the blue in the top, which reveals rotation of the jet. However, it turns to be unshifted at the jet edges that means no rotation there. It follows from the sign of the Doppler shifts in the upper and in the lower parts of the plasma jet that the rotation is clockwise looking
5.4 Radial emissivity profile of the $H_\beta$ line

The relative emission intensities of the $H_\beta$ line emission across the plasma jet is obtained by integrating the CCD image over the wavelength axis. This was done for a magnetic field from 0.4 to 1.6 T and the results are shown in Figure 5.2 (the profiles are plotted on the same scale). It is evident that the profiles become narrower in a higher magnetic field due to the better confinement of the plasma jet. The peak intensity remains almost constant while the electron density grows significantly (Fig. 4.2(a)). This confirms that the emission does
CHAPTER 5. PLASMA JET ROTATION

Figure 5.2: The measured profiles of relative emission intensity of the $H_\beta$-line in a magnetic field of 0.4 to 1.6 T for the 8 mm nozzle diameter. The profiles become narrower that is in agreement with better confinement. The peak intensity does not increase. Note also the top hat shape.

not depend on the density anymore but depends on the $H_2$ density. Excitation of $n = 4$ atoms to higher levels due to collisions with electrons become faster than spontaneous line emission as discussed in the previous paragraph (equation (5.14)).

More important is that the emission profiles have a flat top, which is in agreement with a hollow emission profile. This follows from the fact that optics collect light over the entire line of sight through the plasma column. Thus a lateral intensity profile $I(y)$ is measured instead of a radial emissivity $\epsilon(r)$ distribution. These are related as:

$$I(y) = \sqrt{R^2 - y^2} \int_{-\sqrt{R^2 - y^2}}^{R^2} \epsilon(r) dr = 2 \int_y^R \frac{\epsilon(r) r dr}{\sqrt{r^2 - y^2}}$$  

(5.17)

Here, $R$ is the radius of the light emitting column. The observation that the measured distribution is flat means that the emissivity profile of the plasma column is hollow. The radial distribution of the plasma emissivity $\epsilon(r)$ can be obtained by Abel inversion [114] of the measured lateral intensity distribution $I(y)$ according to:

$$\epsilon(r) = -\frac{1}{\pi} \frac{1}{r} \int_r^R \frac{\partial I(y)}{\partial y} \frac{dy}{\sqrt{y^2 - r^2}}$$  

(5.18)
5.5. **ASYMMETRY OF THE Hβ LINE**

Figure 5.3 shows the Abel inversion of the intensity profile measured in a magnetic field of 0.4 T. It is observed that the maximum emissivity is at the edge of the jet as predicted in the previous section: a hollow emissivity profile. Note that this has consequences for the determination of plasma parameters from the line emission. A hollow emission profile means that predominantly the edges of the jet are sampled. To indicate how important this may be, we return to the Thomson scattering data of the previous chapter (see Figures 4.5 and 4.4). These show that at \( r \approx 6 \) mm (maximum emissivity), the temperature is dropped to \( \sim 25\% \) and the density to \( \sim 30\% \) of their central values. Again this agrees with the prediction that the emission is maximal at \( n_e \approx 1 \times 10^{20} \text{ m}^{-3} \).

The existence of a hollow emission profile indicates that the jet is not transparent for the background gas molecules. After all, these molecules are the rate determining step in the production of the atomic light in the Pilot-PSI jet.

5.5 Asymmetry of off-centre measured Hβ line

The spectral line shape is considered in more detail by integrating over the spatial axis for the individual fibres. Figure 5.4 shows the result for the central and the second fibre up and down from the jet axis. It is seen that the spectral line is clearly asymmetric outside the jet axis. This asymmetry is flipped comparing the profiles recorded at the top and bottom of the plasma jet.

Before trying to quantify the rotation that is hidden in the line profiles we should understand what is determining the shape of the line. On the basis of the considerations in section 2.3.5, we expect Doppler and Stark broadening
Figure 5.4: The measured $H\beta$-line profiles at the upper, central and lower parts of the plasma jet cross-section. This reveals the Doppler shift due to the rotation of the plasma jet. The line shape becomes asymmetric outside the axis of the jet and this asymmetry is flipped for the top and bottom.
which together lead to a Voigt profile. Rotation would come in as a shift of the entire line. First we see how this works for a symmetric line measured in the central fibre. Figure 5.5 shows the line profile that was measured for a magnetic field of $B=1.2$ T, at $z=35$ mm from the plasma source. The profile is fitted with one Voigt profile, which yielded an ion temperature of 0.5 eV and an electron density of $5 \cdot 10^{20} \text{ m}^{-3}$. However, it is seen that the data are not well described by a Voigt profile. The wings and thus the Lorentz components are overestimated whereas the Gauss width is underestimated by the fit; the center is underestimated as well.

The underlying problem becomes more clear when we consider a line profile observed at the top and bottom of the jet. Here the line is asymmetric as clear from Figure 5.4. First we also fit the asymmetric line shapes with a single Voigt. This is shown in Figure 5.6, where the profile measured two fibers aside from the central one is treated. One single Voigt can not describe the asymmetry of the

![Figure 5.5: A measured π-component of $H_β$ line profile fitted by one Voigt profile. The light was collected at 35 mm from the exit of the plasma source nozzle at a magnetic field of 1.2 T and the arc operating at 80 A arc current and 2.5 sml hydrogen flow. In this way, the Gaussian and Lorentzian contribution are separated and the ion temperature and electron density are determined from the respective widths to be $5 \cdot 10^{20} \text{ m}^{-3}$ and 0.5 eV.](image-url)
profile. The systematic residues have become a factor of two higher. Moreover, the top of the fit is shifted with respect to the top of the measured profile. This means that the interpretation of the Doppler shift of a single Voigt as the jet rotation becomes questionable.

We have looked at several possible explanations for the asymmetry. In particular, simulations of the effect of line of sight averaging over the hollow emissivity profile and jet wobble (see section 4.6) were carried out. Line of sight averaging effects do not introduce significant asymmetries. Fast rotational motion of the plasma column, the wobble [113], was measured but was found not to have a sufficient amplitude to be of significance. Only at low fields below 0.1 T wobble becomes important. None of these could explain the observed asymmetries and are therefore not further discussed here.
5.6. THE TWO-COMPONENT MODEL

Inspired by the population mechanisms of H($n = 4$), that predict the excited atoms to carry information of both the plasma ions and the colder background gas, we assume that the line profile accordingly consists of two components.

5.6 The two-component model for the asymmetric line profile description

The measured off-center profile was fitted with the sum of two Voigt profiles. We used the two-component model in four slightly different approaches to determine the ion temperature and the rotational velocity from the asymmetric line shape.

In the first type fitting procedure (denoted as “fit (a)” throughout), 8 parameters were left free: these are the baseline (offset), the amplitudes (or areas) of the two Voigt components, their Doppler $w_D$ widths and one Lorentz $w_L$ width (the same for both), the exact position of the “cold” component, the shift of the “hot” component with respect to the “cold” one. The Lorentz width is assumed to be the same for both, because the light was expected to come from the same volume with a certain local $n_e$ (see the last fitting approach with both Lorentz widths independent for comparison). This procedure led to a remarkable improvement of the fit in comparison with one Voigt (Figure 5.6). The residue was much smaller and is more stochastic. A two-component fit of an asymmetric line profile with the two components and with a residue (a difference between the measured and the fitted profiles) is shown in Figure 5.7. The profile was measured at 45 mm from the source exit in a magnetic field of 1.6 T for 80 A discharge current and 2.5 slm hydrogen flow. The Voigt that describes the “cold” component reveals a temperature of $T_{\text{low}} = 0.54$ eV and the “hot” one gives a temperature of $T_{\text{high}} = 4.3$ eV. The electron density derived from the Lorentz width is $9.5 \cdot 10^{19}$ m$^{-3}$. The Doppler shift of 0.016 nm reveals an (azimuthal) rotational velocity of 9.8 km/s for the “hot” component and 0.9 km/s for the cold component ($\sim 0.0015$ nm shift). The shift of the cold one with respect to the calibrating signal from a hydrogen spectral lamp is of the order of 1/10 of the shift of the “hot” component. We observe also that $T_i$ is higher than $T_{\text{TS}}$ (measured by Thomson scattering - TS) and that $n_{\text{HiRES}}^e$ is lower than $n_{\text{TS}}^e$. The fact that $n_{\text{HiRES}}^e$ is lower is in retrospect not unexpected in view of the hollow emission profile. In fact, the obtained values for $n_e$ are not far from the values we expect for the location of the maximum of emissivity ($r \approx 6$ mm). The ion temperature is significantly higher than $T_{\text{TS}}$. This was unexpected at first. It may point, however, to ion heating by viscous effects at high rotation velocities; this could be a mechanism to cause $T_i$ to be larger than $T_e$. Because initially $T_i$ was expected to be close to $T_e$, also fits were made with prefixed values of $T_i = T_e$. These will be used to investigate the sensitivity of the rotation velocity to the particulars of the two-component model.

To reduce the number of free parameters in the fitting procedure we considered a possibility to fix some of them on the basis of results from other measuring techniques (i.e. Thomson scattering). The component that is cou-
Figure 5.7: A measured asymmetric $H\beta$-line profile fitted with two Voigt profiles according to fit $a$. This gives an unshifted Voigt with a low temperature $T_{\text{low}} = 0.54$ eV and a shifted one with a high temperature $T_{\text{high}} = 4.3$ eV. The electron density derived from the Lorentz width is imposed to be equal for the two Voigts and is $7.6 \cdot 10^{20}$ m$^{-3}$. The Doppler shift of 0.016 nm reveals an (azimuthal) rotational velocity of 9.8 km/s. The residues are less than 1 %. The profile was measured at 45 mm from the source exit in a magnetic field of 1.6 T for 80 A discharge current and 2.5 slm hydrogen flow.

pled to the ions was first expected to have a temperature of the order of the electron temperature due to a low electron-ion energy transfer time ($\sim 10^{-7}$ s in the middle of the jet). A temperature of the other component that is coupled to the colder background gas, was expected to be $\sim 0.2$ eV [50]. Thus, in the second approach (fit $(b)$), we fixed the ion temperature at the electron temperature measured with Thomson scattering ($T_{\text{e}}^{\text{TS}} = 1.9$ eV) and the gas temperature at 0.2 eV. The baseline was fixed such that the residues become zero at the sides of the spectral window. In addition, also the ratio between the Lorentz widths of the two components was fixed to 2 : 1 between the hot and the cold component, respectively. The reasoning behind this ratio originates from the measured hollow emission profiles. The electron density should refer to the locations where the hot and cold components are maximal. The latter
neutral based one will probably peak more outside than the ion-related component. The ratio was chosen 3 : 1 for the densities of hot and cold components (maybe an exaggerated ratio) which This corresponds approximately to a 2 : 1 ratio in Lorentz widths. The fitted line profile is shown in Figure 5.8.

The results differ mainly at the higher magnetic fields by giving a rotational velocity to approximately 8.5 km/s — a value that is close to the one obtained in the fit with 8 free parameters (the first approach). The derived value for \( n_{\text{hot}} \) is of 1.4 \( \cdot \) 10^{20} m\(^{-3}\). The plot of the residues reveals now a systematic discrepancies of up to \( \sim 10 \) %. This is a remarkable result. Apparently, fixing the temperature at the electron temperature causes the Gauss broadening to be too small and the Lorentz broadening too high as is evident from the character of the residue. This means that the temperature of ions must be higher than the electron temperature.

In the third approach (denoted as (c)), the temperatures of the "cold" and "hot" components were fixed again at 0.2 eV and 1.9 eV, respectively, but the electron density was considered to be the same for both components, assuming the radiation to originate from the same volume. The fit with the two components and the residue is shown in Figure 5.9. The derived \( n_e \) is 1.3 \( \cdot \) 10^{20} m\(^{-3}\). The rotational velocity is of 14 km/s that is higher than the one obtained with the first two fitting approaches (a and b). The residue is larger than in the case a but smaller than in the case b showing however the same trend of overestimation of the Lorentz component.

To investigate whether two different independent values for \( n_{\text{hot}} \) and \( n_{\text{cold}} \) would improve the fit we analyzed the data again with all the physical parameters free (fit d), and thus with an additional free parameter in comparison with fit a. The result shown in Figure 5.10 again demonstrates a smaller and more stochastic residue. The plasma parameters proved to be very similar to those of the fit a also with free temperatures. Moreover, the densities \( n_{\text{hot}} \) and \( n_{\text{cold}} \) are very close to each other. Hence we conclude that the most satisfactory fits are a and d and that the fits b and c show systematic derivations pointing to too large Lorentz and too small Gauss. The hot component is more broadened than the value corresponding to the electron temperature. Also we expect the values for the rotation velocity for the fit methods a and d to be the most trustworthy.

In figure 5.11 the results are given for the maximum rotation velocity for the four fit approaches a, b, c, d. Figure 5.11 summarises the results of the four approaches. It shows that the rotation velocity is not influenced significantly by the assumptions in the fit approach for the lower fields. At the larger fields, approach c gives still increasing velocities whereas the other approaches suggest an upper limit for the rotation velocity at around 10 km/s.

A radial profile of the rotation velocity was calculated from the results of approach a. In Figure 5.12 a plot of the rotation profile for the "hot" ion-related and the "cold" background neutral gas-related component is given. The position of the "neutral" component was not fixed and exhibits a distinguishable rotation. It was found to be impossible to fix the position of the low-temperature component at the wavelength calibrated with a hydrogen spectral lamp with an acceptable fit. The rotational velocities of the low-temperature component
Figure 5.8: A measured asymmetric $H_\beta$-line profile fitted a sum of two Voigt components. According to fit $b$ the Doppler widths are fixed to be corresponding to a temperature of 1.9 eV (equal to the electron temperature as measured with Thomson scattering) and 0.2 eV (the estimated gas temperature), respectively. The electron densities of the two components are at ratio 1 : 3 (the Lorentz widths 1 : 2). The derived $n_{e}^{hot}$ is $1.4 \cdot 10^{20}$ m$^{-3}$. The rotational velocity is of 8.5 km/s. The profile was measured at 45 mm from the source exit in a magnetic field of 1.6 T for 80 A discharge current and 2.5 slm hydrogen flow for the 8 mm arc.
Figure 5.9: A measured asymmetric $H_\beta$-line profile fitted with a sum of two Voigt components. According to fit c, the Doppler widths are fixed to be corresponding to temperatures of 1.9 eV and 0.2 eV, respectively. The Lorentz width is the same for both and is free in the fitting procedure. The derived $n_e$ is $1.3 \cdot 10^{20} \text{ m}^{-3}$. The rotational velocity is of 14 km/s. The profile was measured at 45 mm from the source exit in a magnetic field of 1.6 T for 80 A discharge current and 2.5 slm hydrogen flow for the 8 mm arc.
Figure 5.10: A measured asymmetric $H_β$-line profile fitted with a sum of two Voigt components. According to fit $d$, the Doppler widths, the Lorentz widths, the amplitudes of the two components as well as the wavelength of the "cold" component, the shift between them and the baseline were free in the fit. The derived $T^{\text{hot}} = 4.2$ eV, $T^{\text{cold}} = 0.6$ eV, $n_e^{\text{hot}} = 8.8 \cdot 10^{19}$ m$^{-3}$, $n_e^{\text{cold}} = 8.7 \cdot 10^{19}$ m$^{-3}$. The rotational velocity of 10 km/s. The profile was measured at 45 mm from the source exit in a magnetic field of 1.6 T for 80 A discharge current and 2.5 slm hydrogen flow for the 8 mm arc.
Figure 5.11: Four different approaches to determine the jet rotation speed from a measured asymmetric $H_\beta$-line profile. 

- **a**: Doppler widths of the two components, their areas and offset, Lorentz width (the same for both), wavelength and the Doppler shift between them were free in the fit (8 parameters).
- **b**: the Doppler widths are fixed at a temperature of 1.9 eV (equal to the electron temperature as measured with Thomson scattering) and 0.2 eV (the estimated gas temperature), respectively. The electron densities for the non-shifted and the shifted component are fixed at a ratio 1 : 2 (5 parameters).
- **c**: The temperatures are fixed as in **b**, the electron density is assumed to be the same for both components (5 parameters).
- **d**: all 9 physical parameters (including two independent Lorentz widths) are free in the fit. The measurement was done for a scan of $B$ at 45 mm from the source exit for an 8 mm nozzle, 80 A discharge current, and 2.5 slm hydrogen flow.
CHAPTER 5. PLASMA JET ROTATION

Figure 5.12: The ion and neutral rotation velocity profiles in the plasma jet determined from the two-component model. The rotation velocity of the neutral component typically 10% of the ion rotation and is explained as the result of a drag force from the rotating ions. The measurement was done at 45 mm from the source exit in a magnetic field of 1.6 T. Discharge current is 80 A, gas flow rate is 2.5 slm of $H_2$.

Profiles of the $H_\beta$ line were measured across the plasma jet at each setting of the magnetic field (i.e., 0.4, 0.8, 1.2 and 1.6 T) for a nozzle diameter of 8 mm. Profiles of the rotation velocity, ion temperature of each component and electron density were determined by fitting these line profiles according to the two-population model (approach a). The results for a measurement at 45 mm downstream from the nozzle are presented in Figure 5.13. It is observed that for all four magnetic fields $T_i$ is higher than $T_e$ however the ratio $T_i/T_e$ is the highest for the high magnetic field (it varies from 1.4 at 0.4 T to 2.2 at 1.6 T) at the centre of the plasma. Part of it may be due to the line of sight integration but still the factor between $T_i$ and $T_e$ is higher than expected.

5.7 Discussion of the two-component model

The electron density derived from the analysis of the $H_\beta$ line shape was typically five times less than the peak density as determined with Thomson scattering. We explain this by the fact that the signal is integrated over a line of sight crossing the plasma jet giving a lateral profile. As the radial profile is hollow (fig. 5.3), also the central cord measurements give information of the plasma at the periphery. This is in accordance with the expectation from the balance for
5.7. DISCUSSION OF THE TWO-COMPONENT MODEL

Figure 5.13: Radial profiles of rotation velocity, temperature of the "hot" component, electron density and temperature of the "cold" component determined according to the two-population model (approach a) from the $H_\beta$ line shape. The measurement was done for the 8 mm nozzle at 45 mm from the exit of the plasma source at an arc current of 80 A and a hydrogen flow rate of 2.5 slm, in a magnetic field of 0.4, 0.8, 1.2 and 1.6 T.
H\textsubscript{β} emission, which peaks where \( n_{\text{H}_2} \) is appreciable and still \( n_e \) is high enough to populate the level. Hence, the emission is mainly produced at the edges of the plasma jet where the electron density drops below 2 \( \times 10^{20} \) m\(^{-3} \) (more than a factor of 3 lower than in the center of the jet, as can be seen from the electron density profiles measured with Thomson scattering, e.g., fig. 4.4 in chapter 4).

Fixing the ion temperature in the fitting procedure equal to the electron temperature seemed justified by a high electron-ion collision frequency and thus a good coupling between them. This is still very likely at the centre of the plasma where the electron density is high and e-i energy equilibration times:

\[
\tau_{ei}^\text{m} \simeq \frac{\tau_{ei}}{2 m_e} \simeq 3 \cdot 10^{14} \frac{T_e^{3/2}}{n_e \ln \Lambda} \tag{5.19}
\]

are in the sub \( \mu \)s. However at the periphery this coupling is substantially weaker due to lower \( n_e \) and there the ion temperature and electron temperature could differ.

The rotational velocity obtained from the line shape with the two-component model were found to be up to 14 km/s. This value is close to the thermal velocity of hydrogen atoms at 2 eV. We note, that if \( v_{\text{rot}}^{\text{ion}} \) approaches the ion thermal velocity then friction, viscosity and inertia effects start to limit the ion rotation significantly (therewith by the way producing a radial current). The results underline this by showing that indeed for all conditions \( v_{\text{rot}}^{\text{ion}} \) remains smaller than the thermal velocity. They follow roughly the scaling indicated in [116], section 8.4. Hence we observe the \( \mathbf{E} \times \mathbf{B} \)-drift rotation, which at larger values is limited by the effects of friction and viscosity. This can be observed in figure 5.17 in which the ratio is given (velocity profile divided by local value for \( v_{\text{th}}^{\text{i}} \) here defined as \( \sqrt{k_B T_i/M_i} \simeq 10^4 \sqrt[4]{T_i(eV)} \)). It is clear that in all cases the value for the ratio approaches 0.4. Apparently there is a mechanism which tends to accelerate the ions to values close to the thermal velocity. The occurrence of viscous heating of ions, causing \( T_i \) to be higher than \( T_e \) is a finite possibility.

Next the accelerating mechanism should be discussed. It is the axial build up of \( E_z \) because the current continues into the vessel, due to the presence of the magnetic field. Because of the resistance of this current channel a potential is built up in the axial direction (see also section 4.7). This potential gives at the same time a radial electric field, which of course is the built up potential divided by the radius. This field is several tens of volts over the radius of a few mm and is much larger than the ambipolar fields (section 4.7.2). The presence of strong currents in the jet and the associated electric fields can be such a mechanism.

This was already discussed in section 4.7. The voltage between the anode and the last cascaded plate in a magnetic field corrected for the value at zero field (see Figure 5.14) increases up to 80 V over the maximum nozzle radius of 4 mm. This corresponds to an electric field of 20 kV/m and would give, in the absence of limiting effects, a \( \mathbf{E} \times \mathbf{B} \)-drift velocity of 13 km/s in a magnetic field of 1.6 T. Apparently, this is the explanation of the observations: the current continues into the plasma vessel and returns to the anode nozzle, therewith building up an appreciable potential. Hence a strong radial E-field arises, which causes the
Figure 5.14: Potential difference between the anode and the cascaded plate that is the closest to it versus a magnetic field for the nozzle diameter of 5, 6, 7 and 8 mm. The voltage increases with both the magnetic field and the nozzle diameter. The measurement was done at an arc current of 80 A and a hydrogen flow rate of 2.5 slm.
plasma to rotate strongly. As a consequence the velocity is limited by among others viscosity effects and the ions may undergo viscous heating. In the next section we will evaluate the data on rotation and confront them with estimates of the radial electric field.

5.8 Effect of the magnetic field and the nozzle diameter on the jet rotation

In order to investigate the relation between the rotation and radial electric field built up, caused by current continuation into the vessel, we will analyze further the data on rotation velocity and radial electric field. First, we will derive the local rotation frequency profile, $\Omega_i(r) = v_{\text{ion}, \text{rot}} / r$ from the rotation profiles shown in Figure 5.13(a) for the four fields. First we construct symmetrised rotation velocity profiles from the measured profiles with result presented in Figure 5.15.

In Figure 5.16 the maximum values of the rotation velocity are plotted as function of $B$ for the four nozzle radii. In Figure 5.17(a) the profiles for the rotation frequency in the jet center are then given for all four fields. It is evident that the central value of the rotation frequency does increase with $B$ and that the radial extent decreases somewhat masked by the spatial integration caused by a finite fibre dimension. In Figure 5.17(b) we show this increase of rotation

![Graph](image-url)
5.8. EFFECT OF THE MAGNETIC FIELD AND THE NOZZLE

frequency as function of $B$, together with the value for the maximum velocity divided by the radius $r_{max}$ where $v_{max}^{rot}$ takes place.

![Figure 5.16: The rotation velocity determined according to the two-population model from the $H_\beta$ line shape. The measurement was done at 45 mm from the exit of the plasma source nozzle at an arc current of 80 A and a hydrogen flow rate of 2.5 slm, in a magnetic field of 0.4, 0.8, 1.2 and 1.6 T for the nozzle diameter of 5, 6, 7 and 8 mm.](image)

The potential increases with magnetic field, which is an indication of the longer length of the current carrying column. For a larger nozzle diameter it is more difficult for the arc current to cross the magnetic field and therefore the potential drop over the jet radius increases with the nozzle diameter as well (figure 5.14). That means a high radial electric field and consequently high rotational velocities. In that way, an increase of both the magnetic field and the nozzle opening diameter increases the rotational velocity.

For the estimate of the radial electric field we need the radius of the potential distribution, for which we take the radius deduced from the rotation frequency profile. We should remark here that the measured potential refers to the position zero at the exit of the arc source, whereas the measurement of rotation is performed at $z = 40$ mm. This causes an overestimation of $E_r$ by a factor of 1.5 for the lowest field. It will probably not have a significant impact at the highest fields as there the length of the current continuation is much longer.

To compare the measured rotation frequency $\Omega$ with $E \times B$ predicted values we must take $E/BR_{rot}$ which is estimated as $V/BR_{rot}^2$ where $V$ is the potential difference between the anode and the closest to it cascaded plate. This comparison is given in Figure 5.17(b). We conclude that the agreement is satisfactory in absolute values. A more detailed picture including the calculation of
Figure 5.17: The rotation frequency profiles versus a magnetic field (a) and the rotation frequency together with the value for the maximum velocity divided by the radius $r_{\text{max}}$ where it occurs versus a magnetic field (b). The measurement was done at 45 mm from the exit of the plasma source nozzle at an arc current of 80 A and a hydrogen flow rate of 2.5 slm, in a magnetic field of 0.4, 0.8, 1.2 and 1.6 T for the nozzle diameter of 8 mm.

the potential by Ohmic resistance, requires more precise measurements of the radial extent and the axial dependence of the plasma parameters. For example a potential drop over the distance between the nozzle and the position of the spectroscopy measurements 45 mm downstream can presently not be accounted for.

5.9 Axial development of the rotation to probe field crossing currents

We used the HiRES measurements and the two-component model also to determine the axial variation of the rotational velocity of the plasma jet and to evaluate the $z$-dependence of the radial electric field in the jet. Spectra were measured along the plasma jet over 50 mm in steps of 4 mm from the plasma source. The results of the measurements - the peak rotational velocity along the jet for 8 mm nozzle opening in a magnetic field of 1.6 T are presented in Figure 5.18. It is seen that the velocity slightly decreases with the distance from the source but then a sudden rise is observed at approximately 40 mm from the source. We associate this jump with the shock front of the neutral gas that expands from the source. The position of the shock front calculated with expression (2.3) (with the following parameters: Mach number at the exit of the plasma source is around 1, hydrogen flow rate is of 2.5 slm, the specific
5.10. **IS THERE ION VISCOUS HEATING?**

![Graph of Rotational Velocity vs Jet Axis](image)

Figure 5.18: The axial variation of the rotational velocity determined with the two-population model from the $H_β$ line shape. The spectra were measured within 50 mm from the exit of the plasma source nozzle. Source parameters: 80 A arc current, 2.5 slm hydrogen flow rate, 8 mm nozzle diameter, 1.6 T magnetic field.

The heat ratio is 5/3, and the background pressure of about 7 Pa) is around 60 mm from the nozzle. If we take into account the fact that a pre-expansion occurs already in the nozzle opening (diameter of 8 mm within first 10 mm of the nozzle and then 10 mm of the nozzle with diameter of 14 mm) then the density jump corresponds to the shock front of the neutral gas expansion pattern.

To investigate the length scales over which the arc current extends outside the source into the plasma jet, we conducted the HiRES measurements also in the second and in the third window (counting from the plasma source). For example, for a magnetic field of 0.4 T the rotational velocity in the first window (at around 45 mm from the source) is about 4000 m/s; in the second window (at around 300 mm from the source) it drops to around 2000 m/s and in the third window (at around 550 mm from the source) it occurred to be undetectable within the accuracy of our measurements. It should be noted that the plasma jet is almost extinguished here as well. We conclude that for low fields the arc current in the magnetized plasma jet reaches up to 40-50 cm, roughly corresponding to the end of the visible plasma jet.

**5.10 Is there ion viscous heating?**

In the preceding sections we have analysed the data on electron density and temperature from Thomson scattering and the ion temperature and rotation
velocity from HiRES. It appeared from the data that the electrons are magnetised for all investigated conditions, i.e. for four values of the applied magnetic field and for four nozzle diameters. The values for $\omega_{\text{ce}} \tau_{\text{ee}}$ are all around 20–40 and that values for $\omega_{\text{ci}} \tau_{\text{ii}}$ are all around 1. So in the discussion of the results we can assume the electrons to be confined to the field lines, which is the reason for the current to continue in the vessel, building up in this way the potential along the column and with that leads to a strong radial electric field. We have seen that this radial electric field leads to a rotation, which at the edge approaches the ion thermal velocity. Hence viscous effects become probable and we will estimate the magnitude from the experimental values obtained.

The viscous heating term is equal to [124]:

$$Q_{\text{visc}} = \Pi_{\alpha \beta} \frac{\partial v_\alpha}{\partial x_\beta} \tag{5.20}$$

In our case $v_\theta$ is the largest and for $\omega_{\text{ci}} \tau_{\text{ii}} \sim 1$, the largest contribution is of the order:

$$Q_{\text{visc}} \sim C \cdot \eta_1 \left( \frac{\partial v_\theta(R)}{R} \right)^2 = C \cdot f_1 n_i k_B T_i \tau_{\text{ii}} \left( \frac{\partial v_\theta(R)}{R} \right)^2 \tag{5.21}$$

with

$$f_1 = \frac{2.33 + 4.8 \xi^2}{16 \xi^4 + 16.02 \xi^2 + 2.23}, \quad \xi \equiv \omega_{\text{ci}} \tau_{\text{ii}} \tag{5.22}$$

in which $\xi$ is the ion Hall parameter. Note that the estimates given in Braginskii’s review ([124], pp. 219–220) are relevant for magnetised plasmas with $\omega_{\text{ci}} \tau_{\text{ii}} \gg 1$. The quantity $f_1 \sim 1$ for $\omega_{\text{ci}} \tau_{\text{ii}} < 1$, $f_1 \sim 0.3$ for $\omega_{\text{ci}} \tau_{\text{ii}} = 1$ and $f_1 \sim 0.3/(\omega_{\text{ci}} \tau_{\text{ii}})^2$ for $\omega_{\text{ci}} \tau_{\text{ii}} \gg 1$. Here we take $C f_1 \sim 1$. Note that $Q_{\text{visc}}$ is determined by the rotation frequency $\Omega_\text{rot}$ and its radial dependence rather than by the velocity itself.

Hence this ion heating can be roughly estimated as:

$$Q_{\text{visc}} \sim n_i k_B T_i \tau_{\text{ii}} \Omega_i^2, \tag{5.23}$$

in which $\Omega_i \equiv v_{\theta i}/r$ is the radially dependent ion rotation frequency. Now it appears that the value of $f_1$ is 1 for very small fields and is still about 0.3 for for $\omega_{\text{ci}} \tau_{\text{ii}} = 1$. For larger values of the ion Hall parameter it decreases with $1/(\omega_{\text{ci}} \tau_{\text{ii}})^2$ and becomes very small. One could speculate whether the fact that the ion Hall parameter is consistently close to 1 is in fact connected with the optimum for viscous heating.

In order to investigate whether the ion heating by viscosity is significant we can compare it with ion heating by energy transfer from the electrons to ions and by comparing it to ion heat loss by ion heat conduction.

Electron-ion heat transfer is given by:

$$Q_{\text{ei}} = \frac{3}{2} n_e k_B (T_e - T_i) / \tau_{\text{ei}} \tag{5.24}$$
which can be compared with the above given estimate of the viscous heating. If one takes the values in the centre of the jet at high fields: 
\[ \ln \Lambda \approx 7, \ n_i \approx 7 \times 10^{20} \ \text{m}^{-3}, \ T_e \sim 2 \ \text{eV}, \ T_i \sim 4 \ \text{eV}, \ B = 1.6 \ \text{T}, \ \Omega \sim 5 \cdot 10^6 \ \text{rad/s} \] and filling in the relevant quantities it shows that the ion heating contributions are comparable at a level of several \( 10^8 \ \text{W/m}^3 \) and that thus ion heating above the electron temperature is possible.

Another way to estimate the significance of ion heating by viscosity is by comparing it Ohmic dissipation (which is the usual electron heating mechanism), which can be given as:

\[
Q_{\text{ohmic}} = j^2 \sigma || = \frac{I^2}{\sigma || (\pi R_j^2)^2} \quad (5.25)
\]

With the values of current, radial extent, and conductivity (electron temperature) we find again values of several times \( 10^8 \ \text{W/m}^3 \). Hence, we conclude that ion viscous heating is important, the occurrence of which agrees with the observed high ion temperatures. More work, with local measurements of ion temperature and velocity by LIF (Laser Induced Fluorescence) is needed to further unravel this interesting mechanism.

## 5.11 Summary and discussion on the plasma rotation

High-resolution optical emission spectroscopy measurements revealed the rotation of the plasma jet confined by the axial magnetic field of Pilot-PSI. From these measurements we concluded:

- The asymmetry of the measured line shapes can be composed with two Voigt distributions.
- The component with the largest shift corresponds to a temperature of at least the electron temperature. The axial variation of the shift agrees with jet rotation caused by \( \mathbf{E} \times \mathbf{B} \)-drift of charged particles in a radial electric field that is related to the discharge current continuing from the source into the vacuum vessel.
- The second component is also shifted, however only slightly, and has a typical background gas temperature of \( \sim 2 \cdot 10^3 \ \text{K} \). The shift is typically one tenth of the full rotation.
- The atomic emission is observed to be independent of the electron density. Furthermore, the emissivity profile is hollow, which is explained by a decreased molecular hydrogen density in the center of the jet. We expect that maximum emissivity is typically reached where the electron density is \( \sim 1 \cdot 10^{20} \ \text{m}^{-3} \). The consequence is that especially the jet edges are probed by the emission spectroscopy.
The combination of integrating light over the entire line of sight, mixing of plasma and background gas properties, and the hollow emissivity profile makes it difficult to use emission spectroscopy as a diagnostic for plasma properties in the center of the jet. Despite these complications, we have demonstrated that the technique is very useful to assess the rotation of the plasma column. We have compared different approaches in the analysis of the line profile, which showed good agreement at the lower magnetic fields and only started to disagree at the maximum field of 1.6 T. From these rotation measurements we were able to conclude that a strong magnetic field causes the discharge current to continue outside the source into the free jet. This was supported by the agreement in absolute values between the estimated values of the radial electric field (by the voltage drop over the nozzle radius) and those estimated from the measured rotation velocity and known magnetic field. Furthermore, while the rotation angular frequency on axis was observed to increase with for example magnetic field and nozzle diameter, the maximum velocity is always below the thermal velocity of the ions. We interpret this as a consequence of friction, viscosity and inertia effects that prevent approaching the thermal velocity.

Very remarkable are the high ion temperatures that were determined ($T_i$ up to 4 eV) in comparison with the electron temperatures ($T_e$ is up to 2 eV). Of course, the complexity of the line shape analysis behind this temperature determination might raise doubts on the significance of this discrepancy. However, given the two population model, it was not possible to reproduce the line shapes without these high temperatures. Moreover, the high temperatures are supported by the additional heating that is introduced by the fast jet rotation. Remember here also the counterintuitive finding that it is not necessarily the rotational velocity that approaches the thermal speed but already the high rotational frequency on axis that drives this heating. Finally, such a decoupling of the electron and ion temperatures is possible given the electron-ion energy transfer time being in the sub $\mu$s range. In conclusion, we presently do believe that the ion temperature in the Pilot-PSI hydrogen plasma jet is slightly higher than the electron temperature due to the jet rotation, especially at the higher (1.6 T) magnetic fields.
Chapter 6

General discussion

The different chapters in this thesis are more or less separate investigations with conclusions mostly specific for that chapter. Returning to the introduction, there was a shared aim that drove this work: the efficient production of an intense hydrogen plasma jet for the linear plasma generator Magnum-PSI, in which this will be used to mimic the plasma conditions of the ITER divertor. The following issues were formulated:

- efficient production of a high flux density hydrogen plasma jet
- plasma transport in high magnetic fields
- diagnostics to monitor the plasma parameters in Pilot-PSI and in the future machine Magnum-PSI

In this chapter we summarize the results of the different chapters in view of these themes and draw general conclusions about the cascaded arc operated on hydrogen in strong magnetic fields.

6.1 Hydrogen plasma production with the cascaded arc

The high-pressure wall-stabilised cascaded arc was chosen as a plasma source for production of high-speed dense hydrogen plasma jets. The arc was operated at a discharge current in the range of 20 to 100 A in magnetic fields up to 1.6 T. In the most recent experiments (not described in this thesis), currents up to 300 A were successfully applied. The general conclusion on the basis of the results that are repeated below is that the cascaded arc performs well under these conditions and produces the desired hydrogen plasma jets.

If there is no magnetic field applied, the plasma expands supersonically into the vacuum vessel due to the pressure difference (three orders of magnitude) between the source inlet and the vessel. The electron density $n_e$ in a hydrogen
plasma was measured at 30 cm from the exit of source with a double Langmuir probe. It increased with the arc current from $2 \times 10^{16}$ m$^{-3}$ at 50 A up to $6 \times 10^{16}$ m$^{-3}$ at 80 A.

Higher pressures reduce the expansion, which leads to a free jet with an increased $n_e$. However, this is not the way of plasma confinement that is suitable for our purposes. Hydrogen plasma recombines anomalously quickly via a two-step molecular activated recombination [49] and the recombination rate is proportional to the pressure. Consequently, no plasma reaches the remote target.

The current-voltage characteristic of the arc is anomalous for hydrogen operation: the discharge voltage decreases with increasing discharge current. This inverse behavior is only weak so that the total power consumption still increases with increasing current. It is explained by an effective plasma channel diameter that is smaller than the arc channel diameter and increases with higher powers. The plasma resistivity is determined by the electron temperature $T_e$, which is roughly constant and in the range 1.1–1.3 eV for a wide range of operational parameters (as it follows from experiments and modelling). An increasing effective plasma diameter of constant resistivity means a decreasing channel resistance.

Variation of the arc channel diameter demonstrated that for argon the resistivity of the arc scales with the discharge current density as $\eta \propto \bar{j}^{-0.6}$ and is independent of the diameter. For hydrogen, this relation is $\eta \propto \bar{j}^{-1.3}$ and it does depend weakly on the channel diameter. However, the range of channel diameters that could be investigated in hydrogen was too small to allow more detailed conclusions here.

A single-parameter model was developed that describes the stability of the discharge and the difference between operation on argon and hydrogen. It uses the filling factor $\alpha$ (the ratio between the plasma and the arc channel cross section) as main parameter, which is determined by the power balance. The input power must scale with $\alpha$ as $P_{in} \propto \alpha^{-2}$. The power losses are assumed to be set by a narrow layer of cold gas between plasma and the wall and to scale as $P_{loss} \propto \alpha/(1-\alpha)$. It follows that $\alpha$ must slightly increase with the current density in the arc in order to balance input and losses: $\alpha \propto \bar{j}^{0.3-0.5}$ for argon and $\alpha \propto \bar{j}^{0.65}$ for hydrogen. This shows that the effect is stronger for hydrogen. When $\alpha$ approaches 1, the model becomes unapplicable: plasma can not exist very close to the cooled wall. In that case, $T_e$ must increase to provide higher conductivity.

The model predicts a more efficient plasma production (gas efficiency) at higher current densities. Pushing alpha to its maximum would require the electron temperature to increase. This would give a dramatic improvement of the ion output and thus the source efficiency.

Preliminary experiments with arc currents up to 300 A in a 4 mm bore (not presented in this thesis) demonstrated that the I-V characteristic of the hydrogen discharge becomes also positive at higher current densities. We conclude that than the filling factor approaches its maximum (unity). Narrowing of the cold layer at the wall and thus growth of the dissociation degree decreases the importance of molecular processes and gives the discharge an atomic (argon-like)
character.

So far, it is not entirely clear why the filling factor $\alpha$ would be independent of the bore $a$.

Based on the model for power losses, we find that operation in the flat range of the I-V characteristic is the most efficient in consumed power per ion. Thus, to produce more ions it is better to increase the arc diameter rather than to push the current density. A wide range in arc current leads to approximately the same efficiency, which gives operational flexibility. Thus, the primary line of approach for up-scaling the plasma source for the future Magnum-PSI experiment will have to focus on wider arc channels. An enlargement of the discharge channel diameter (to $\sim 10$ mm) in combination with an increased discharge current (to $1 \sim 2$ kA) should give a plasma output that is specified for Magnum-PSI.

6.2 Plasma transport in strong magnetic fields

Pilot-PSI offers considerably higher magnetic fields in comparison to other linear plasma generators. It can be varied in steps of 0.4 up to 1.6 T. The plasma expansion changes drastically in such fields: the plasma is confined into a narrow ($\sim 1$ cm), high-density and high-temperature jet.

According to the present understanding, the effect of the field is not limited to the expansion characteristics but also effects the discharge current. A significant part of this current continues outside the plasma source into the free jet. A strong electric field (up to several teens kV/m) is related to this current and is directed inward. This causes together with the axial magnetic field an azimuthal $\vec{E} \times \vec{B}$-drift of the charged particles which makes the jet rotate. Rotation velocities up to 10 km/s (at a jet radius of $\sim 2$ mm) and frequencies up to 5 MHz have been determined from emission spectroscopy (discussed below). We presently believe that this rotation is a heating mechanism for the ions. This has as consequence that $T_i$ can be even higher than $T_e$. At the edges of the jet the electron-ion energy transfer time is of the order of $10^{-6}$ s while the ion-ion collision time and charge exchange with neutrals has a characteristic time of $3 \cdot 10^{-8}$ s.

The outer part of the arc current causes a "wobbling" of the plasma jet in a magnetic field. The Lorentz force is believed to be a driving mechanism of the wobbling. We detected frequencies of the wobbling up to 1 MHz depending on $n_e$, magnetic field and the length of the arc. The detected frequency is in a good agreement with the predicted values for hydrogen and argon. It appears that for magnetic fields of 0.4 T and higher, the wobble amplitude is negligibly small.

The arc voltage increases with the magnetic field. We interpret this as an additional voltage drop that exists in the nozzle region, just outside the plasma source. This causes extra power input in a region where there is less heat loss to the walls of the arc. Indeed, the electron temperature is observed to be increased under these conditions ($T_e$ up to 2 eV has been detected).

The plasma density was found to increase with field, current, and nozzle di-
ameter. The maximum $n_e$ as detected with Thomson scattering was $7 \cdot 10^{20} \text{ m}^{-3}$. The particle flux density in the magnetised plasma jet of Pilot-PSI is calculated by multiplying this $n_e$ with the axial velocity $v_{ax}= 3 \text{ km/s}$:

$$\Gamma_{H^+} = n_e v_{ax} = 2 \cdot 10^{24} \text{ m}^{-2}\text{s}^{-1}$$

It is a record value for linear plasma generators in the world and it is sufficient for ITER-relevant PSI studies. This answers the first theme of the thesis: efficient production of a high flux density hydrogen plasma jet.

Preliminary experiments on additional heating of the plasma by feeding the jet with extra current demonstrated that this is a promising way to control the temperature and density of the particles in the plasma jet. More studies are foreseen on this subject in order to determine an optimal regime for this Ohmic heating.

### 6.3 Development of diagnostics

In order to investigate the plasma transport in the high magnetic fields, we supported the Thomson scattering measurements by High-Resolution Emission Spectroscopy (HiRES). It provided us with information on the jet velocities (rotational and axial) and heavy ion temperatures. The rotation of the plasma jet was determined from the Doppler shift of atomic lines (in this work we mainly used the $H_\beta$ line), which was opposite in direction for the top and bottom of the plasma jet. The line shapes occurred to be asymmetric. The explanation of this asymmetry starts from molecular activated recombination as the main excitation mechanism of atomic hydrogen. Charge exchange occurs between a proton and a background molecule and subsequent recombination of the molecular ions produces an excited atom. Because the ion-ion collision time ($3 \cdot 10^{-8} \text{ s}$) is short in comparison with the $H_\beta$ transition time, this atom will quickly equilibrate with the plasma ions before it will radiate. In addition, also charge exchange with other protons (at a rate of 5–10 times faster than transition probability of $H_\beta$ line) will take care of this equilibration. However, some of these “hot” excited atoms may undergo elastic collisions with low-temperature background gas before they radiate. The rate of these collisions is between the charge exchange time and the transition probability. This produces a colder component in the emitted light, but still with some rotation and somewhat higher temperature than one would expect for the colder background gas. According to these views, a model was developed that was used for the analysis of the asymmetric line profiles.

The rotation velocity was determined to be between 6 and 10 km/s, depending on the exact operational conditions. It increased with magnetic field (0.4–1.6 T) and nozzle diameter (5–8 mm). These velocities approached the thermal velocity but were never found to exceed them.

Although we analyzed the line shapes as consisting of two separate components, we realize that a continuous non-thermal asymmetric distribution function of the excited atoms is more probable. We base this on the comparable
rates of the charge exchange, the elastic collisions and radiation decay of excited states.

It occurred to be difficult to separate the Doppler broadening from Lorentz broadening and to derive $T_i$ and $n_e$ with high accuracy. Applying no fit constraints yielded $T_i$ up to 5 eV, which seems too high for $T_e$ up to 2 eV. Similarly, the Lorentz width gave values for $n_e$ that were typically 5 times lower than the densities obtained from Thomson scattering ($1.4 \cdot 10^{20} \text{ m}^{-3}$ versus $7 \cdot 10^{20} \text{ m}^{-3}$). The latter is explained with the hollow emissivity profile: light originates mainly from the edges of the plasma jet, where $n_e$ is lower. However, we do not exclude that the Lorentz width is underestimated in the fit due to an overestimation of the Doppler width (which is vice versa a possible reason for the too high $T_i$).

Although, the Doppler shift between the two components allows to determine rotation velocity of ions with a moderate accuracy, it occurs not to be easy to use this technique as a reliable diagnostic tool on an every-day-basis. It requires extreme care in the interpretation and needs to be supported by data from other independent techniques, such as $T_e$ and $n_e$ from Thomson scattering. Still, the HiRES data gave important input for the understanding of the processes that take place in a magnetised hydrogen plasma jet as is encountered in Pilot-PSI. A comprehensive and consistent model of the plasma jet has been developed by virtue of these data.

The variation of the axial velocity of the plasma jet was also estimated from the Doppler shifts of atomic lines. The axial velocity was measured to decrease from about 5 km/s close to the source to 2 km/s further downstream. However, these measurements have a limited accuracy due to a long line of sight integration. Accurate axial velocity measurements require high spatial resolution, as is offered by Two-Photon Laser Induced Fluorescence (TALIF). Such data are inevitable to yield in combination with Thomson scattering precise information on the particle flux density in the jet.

6.4 Concluding remarks

In summary, our investigations were highly successful in reaching the targets that were set. A stable magnetised hydrogen plasma jet was produced with $n_e$ up to $7 \cdot 10^{20} \text{ m}^{-3}$, $T_e$ around 1–2 eV, and $T_i$ up to 4 eV. A detailed view on the transport of the plasma jet towards a target was produced. Rotation of the plasma jet with a frequency of $\sim 5$ MHz in a magnetic field of 1.6 T was detected. Axial jet velocities of around 3 km/s together with the achieved high densities provide the required for ITER-relevant PSI studies particle flux density of more than $10^{24} \text{ m}^{-2} \text{s}^{-1}$, which is unique for linear plasma generators. The electron density and temperature can be controlled by the current in the plasma source and the nozzle geometry in combination with the external confining magnetic field. Extra current through the plasma jet also showed promising results on post-heating of the plasma for temperature and density control. Moreover, we have developed a comprehensive understanding of the processes in the plasma jet. All together this work gives us confidence that a cascaded arc with a
discharge channel of \( \sim 10 \text{ mm} \) diameter operated at a discharge current of \( 1 - 2 \) kA will produce a hydrogen plasma jet as specified for Magnum-PSI.
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Summary

The interaction between the hydrogen plasma and the divertor wall is yet an unresolved issue in the design of ITER. Especially the erosion rates and retention of tritium (the fuel of a fusion reactor) presently foreseen for the ITER divertor are critical issues for its prolonged operation. The large linear plasma generator Magnum-PSI, is presently being build at our institute to study the underlying processes. This requires intense hydrogen plasma jets to mimic the plasma conditions of the ITER divertor.

The research described in this thesis was carried out at Pilot-PSI, the forerunner of Magnum-PSI, and focused on the efficient production of such intense hydrogen plasma jets with a wall-stabilized cascaded arc operated in a strong magnetic field.

The cascaded arc that was used here operates at a relatively high pressure (∼0.1 bar). Coupled to a vacuum vessel, the hydrogen plasma is observed to expand in a way similar to gas expansion. The investigations start with Langmuir probe measurements to determine the plasma density and temperature under these conditions. The results confirmed the importance of molecular assisted recombination.

The performance of the arc operating on hydrogen was characterized and compared with argon operation by power measurements. An important aspect observed in these measurements is a decreasing discharge voltage for increasing discharge currents for hydrogen operation. This is a consequence of an increasing effective plasma channel with increasing power input resulting in a decreased average resistivity. Such an IV-characteristic is not observed for argon. Results on argon for different diameters of the discharge channel demonstrate that the resistivity of the arc (i.e. its resistance divided by the area of the channel cross section and the channel length) scales predominantly with the discharge current density as $\eta \propto \bar{j}^{-0.6}$ and is independent on the diameter. For hydrogen, this relation is $\eta \propto \bar{j}^{-1.3}$. The independence on the channel diameter is not entirely clear for hydrogen and at most true for channels wider than 4 mm. A model is developed from a power balance for the discharge channel. It explains the experimental data on the basis of the higher heat conduction of hydrogen, which leads to a smaller hot plasma channel compared to argon. The model predicts higher efficiencies if higher discharge currents are combined with larger channel diameters.

Pilot-PSI offers magnetic fields up to 1.6 T, which is unique for a linear
plasma generator. This confines the otherwise expanding plasma into an intense jet and reduces the recombination losses that are typical for hydrogen plasma. Thomson scattering was applied to determine the record plasma parameters, unprecedented in a linear plasma generator: an electron density $n_e = 7 \cdot 10^{20} \text{ m}^{-3}$ and temperature $T_e = 2 \text{ eV}$ (at $B=1.6T$).

The average forward velocity of the plasma was determined from the Doppler shift in the light emission and amounts typically $\sim 3 \text{ km/s}$ at the position of the Thomson scattering experiment. Together with the measured electron density, this yields a proton flux density that is expected in the ITER divertor:

$$\Gamma_{H^+} = 2 \cdot 10^{24} \text{ m}^{-2}\text{s}^{-1}$$

The magnetic field also effects the operation of the source. This is concluded from the significant increase of the potential difference at the exit of the source induced by the magnetic field. The additional potential is dependent on the inner diameter of the nozzle (with respect to the discharge channel) and amounts up to 90 V for an 8 mm nozzle at $B=1.6 \text{ T}$. The potential increase from a wider nozzle causes a significantly improved source output, up to a factor of 2, as was quantified by Thomson scattering. The effect is explained in a physical picture where a significant part of the discharge current continues outside the source before it attaches to the nozzle.

The consequence of current continuing into the free jet is that a large potential is built up, which gives rise to an appreciable radial electric field. The radial electric field is perpendicular to the axial magnetic field and causes strong rotation of the jet via $E \times B$ drift of the plasma particles. High-resolution optical emission spectroscopy (HiRES) was performed to investigate this rotation from the Doppler shift in the atomic light emitted perpendicular to the plasma jet.

The measured line shapes were asymmetric, which was explained by the existence of two populations in the radiating atoms. One is coupled to rotating ions and has the ion temperature and velocity. The other is coupled to colder background gas and rotates at most slightly. This picture was implemented in a fitting procedure that yields the ion temperature and rotation velocity, the background gas temperature and rotation velocity, and the electron density. In this way, peak rotation velocities up to $10^4 \text{ m/s}$ were determined, probably limited by ion-neutral friction to below the thermal velocity. These rotation velocities correspond to electric fields larger than $10^4 \text{ V/m}$. The rotation frequency of the central part of the plasma jet was observed to scale with the potential difference between the last plate and the nozzle, which confirms the physical picture on currents continuing beyond the source. The ion temperature $T_i$ that followed from the fitting procedure was found to be systematically larger than $T_e$. Although the accuracy in the temperature determination is expected to be too limited to quantify the ratio $T_i/T_e$, we do conclude that it is larger than unity. This is in line with additional viscous heating of the ions due to the rotation of the jet.

On the basis of the results presented in this thesis we conclude that the cascaded arc can serve as the basis for a future Magnum-PSI source. Scaling of the arc will be based on an enlargement of the discharge channel diameter (to $\sim 10 \text{ mm}$) in combination with an increased discharge current (to $1 - 2 \text{ kA}$).
Samenvatting

De wisselwerking tussen het waterstofplasma en de divertorplaten is een vooral nog onopgelost vraagstuk in het kader van het design van ITER. Met name de erosie van de wanden en de opname van tritium, de brandstof van de kernfusiereactor, zijn nog niet onder controle. De lineaire plasmagenerator Magnum-PSI wordt momenteel gebouwd om de processen die hieraan ten grondslag liggen te onderzoeken. Hiervoor is een bron voor intense bundels waterstofplasma nodig waarmee de plasmamomstandigheden zoals die in de ITER divertor verwacht worden kunnen worden nagebootst.

Dit proefschrift beschrijft het onderzoek naar de mogelijkheden om dergelijke bundels te maken door een boogontlading in een zogenaamde cascadeboog te combineren met een sterk magnetveld.

De cascadeboog zoals deze gebruikt is in het hier gepresenteerde onderzoek produceert plasma onder relatief hoge druk (0.1 bar). Door deze te koppelen aan een vacuümvat strekt het plasma als in een gasexpansie het vat binnen. Het onderzoek dat beschreven is in dit proefschrift start met de bepaling van plasmadichtheid en -temperatuur van waterstofplasma dat op deze wijze in het vacuümvat van Pilot-PSI is geïntroduceerd. Deze metingen bevestigen dat transport van waterstofplasma gehinderd wordt door verliezen ten gevolge van ladingsruil met waterstofmoleculen gevolgd door recombinatie.

Onder dezelfde omstandigheden is de werking van de cascadeboog op waterstofgas gekarakteriseerd aan de hand van vermogensmetingen en vergeleken met zijn gedrag voor argongas. Een belangrijke bevinding hierbij is het negatieve inverse verband tussen de spanning door de bron en de hiervoor benodigde spanning. Dit is het gevolg van de verwijding van het effectieve plasmakanaal bij hogere vermogenstoevoer welke leidt tot een verlaging van de weerstand van het kanaal. Bij argongas is dit niet het geval. Metingen bij verschillende diameters van het gaskanaal tonen voor argon aan dat de stroomdichtheid in het kanaal de soortelijke weerstand van het plasmakanaal bepaalt, volgens \( \eta \propto j^{-0.6} \) onafhankelijk van de kanaaldiameter. Voor waterstof is deze relatie \( \eta \propto j^{-1.3} \). De onafhankelijkheid van de kanaaldiameter is niet geheel duidelijk en hooguit geldig voor kanalen met een diameter groter dan 4 mm. Op grond van deze metingen is een vermogensbalans voor de cascadeboog opgesteld die de resultaten voor zowel waterstof als argon beschrijft. Dit model voorspelt een hogere efficiëntie voor de cascadeboog indien het verhogen van de boogstroom samen gaat met het vergroten van de kanaaldiameter.
Het magneetveld van Pilot-PSI is maximaal 1.6 T, uniek voor een lineaire plasmagenerator. Met behulp van dit magneetveld wordt het plasma dat zonder magneetveld expandeerde opgesloten op de as van het vat en wordt een plasmabundel gevormd. Op deze manier worden de recombinatieverliezen die type­rend zijn voor waterstofplasma belangrijk verlaagd. Met behulp van Thomson verstrooiing is voor deze omstandigheden plasmadichtheid van $7 \cdot 10^{20} \text{m}^{-3}$ bij een elektronentemperatuur van 2 eV bepaald: een record voor waterstofplasma in een lineaire plasmagenerator.

Op grond van de Dopplerverschuiving in het geëmis­teerde licht is een voor­waartse snelheid van 3 km/s van het plasma in de bundel bepaald. Samen met de plasmadichtheid geeft dit de plasmasnelheid die ook in de divertor van ITER verwacht wordt:

$$\Gamma_{H^+} = 2 \cdot 10^{24} \text{m}^{-2} \text{s}^{-1}$$

Het magneetveld beinvloedt ook de werking van de bron. Dit volgt allereerst uit de significante verhoging van de potentiaalval aan de bronuitgang die evenredig toeneemt met het magneetveld. Deze potentiaalval is verder ook afhanke­lijk van de toename in inwendige diameter tussen het bronkanaal en de uitstroomopening (welke tevens de anode van de boog is). Bijvoorbeeld bij een 8 mm uitstroomopening en een 4 mm kanaaldiameter ontstaat 90 V extra potentiaalverschil door een langere plasmabundel bij het het aanleggen van 1.6 T magneetveld. De productie van waterstofplasma verbetert hierdoor met een factor 2. Het effect is uitgelegd als het gevolg van bronstroom die doorloopt in de bundel en pas in het vat terugbuigt naar de uitstroomopening.

Het gevolg van de stroom die pas in het vacuüm­vat het magneetveld kruist is een aanmerkelijk radiaal elektrisch veld. Dit veroorzaakt via een $\mathbf{E} \times \mathbf{B}$ drift­beweging rotatie van de plasmabundel om haar as. Hoge resolutie emissiespec­troscopie is toegepast om deze rotatie te meten. De gemeten lijnvormen waren asymmetrisch. Dit is geïnterpreteerd als het gevolg van twee verdelingen stralende atomen. De eerste heeft de temperatuur en rotatiebeweging van de protonen, de tweede is afgekoeld door interactie met het achtergrondgas en roteert nauwelijks. Op deze manier werd een maximale rotatiesnelheid van 10 km/s bepaald, waarschijnlijk gelimiteerd tot onder de geluidssnelheid. Dit komt overeen met elektrische velden groter dan $10^4$ V/m. Verder bleek dat de omwentelingsfrequentie van het centrale deel van de plasmakolom schaalt met de potentiaalval aan de bronuitgang. Dit bevestigt het beeld dat de boogstroom zich uitstrekt buiten de bron. Een andere opmerkelijk resultaat van deze analyse was een ionen­temperatuur die systematisch hoger was dan de temperatuur van de elektronen. De nauwkeurigheid van de fitprocedure schoot tekort om dit verschil te kwantificeren, maar was voldoende om te concluderen dat de ionen aan de rand van de plasmakolom heter zijn dan de elektronen. Dit is verklaard als het gevolg van viskeuze verhitting van de ionen ten gevolge van de rotatie van de bundel.

Uit de resultaten in dit proefschrift blijkt dat de cascadeboog kan dienen als plasmabron voor Magnum-PSI. Verdere opschaling zal met name gebaseerd zijn op vergroting van de diameter van het ontladingskanaal.
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11.04.2006
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## Curriculum Vitae

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