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Spontaneous Formation of Ion Holes and Ion Beams in Expanding Plasmas

Evan M. Aguirre

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SPONTANEOUS FORMATION OF ION HOLES AND ION BEAMS IN EXPANDING PLASMAS

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Dissertation submitted to the Eberly College of Arts and Sciences
at West Virginia University

In partial fulfillment of the requirements for the degree of

Doctor of Philosophy
in
Physics

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Abstract

Spontaneous Formation of Ion Holes and Ion Beams in Expanding Plasmas

Evan M. Aguirre

Presented here are direct, spatially resolved measurements of the spontaneous formation of a steady-state, ion accelerating structure and an ion hole in the plume of an expanding, low-pressure plasma. Ion velocity distribution function (IVDF) measurements, obtained with laser induced fluorescence, are taken parallel and perpendicular to the background magnetic field in an argon helicon plasma. The parallel IVDFs show an ion beam with $v \approx 8,000$ m/s flowing downstream and confined to the center of the discharge. The ion beam is measurable for tens of centimeters along the expansion axis before the LIF signal fades, likely a result of metastable quenching of the beam ions. The parallel ion beam velocity slows in agreement with expectations for the measured parallel electric field. The perpendicular IVDFs show an ion population with a radially outward flow that increases with distance from the plasma axis. At the periphery of the ion beam, bipolar, magnetic field-aligned electric field structures, ion holes, form and accelerate ions across the ambient magnetic field. Multi-species plasmas provide further evidence that the ion acceleration mechanism is not a true double layer (DL). These measurements demonstrate that at least two-dimensional and perhaps fully three-dimensional models are needed to accurately describe the spontaneous acceleration of ion beams in expanding plasmas. A new paradigm for the spontaneous formation of ion accelerating structures in expanding plasmas is presented.

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Abbreviations

DL	Double Layer
CFDL	Current Free Double Layer
rf	radio frequency
LIF	Laser Induced Fluorescence
HPD	Helicon Plasma Device
IVDF	Ion Velocity Distribution Function
RFEA	Retarding Field Energy Analyzer
ANU	Australian National University
WVU	West Virginia University
MHD	MagnetoHydroDynamics
FAST	Fast Auroral Snapshot Explorer
HDLT	Helicon Double Layer Thruster
HELIX	Hot hELIcon eXperiment
LEIA	Large Experiment on Instabilities and Anisotropies
EEPF	Electron Energy Probability Function
EEDF	Electron Energy Distribution Function
FFT	Fast Fourier Transform
SNR	Signal to Noise Ratio
VDF	Velocity Distribution Function
RMS	Root Mean Square
PMT	PhotoMultiplier Tube
FWHM	Full Width Half Maximum
SCCM	Standard Cubic Centimeter per Minute
LASER	Light Amplification by Stimulated Emission of Radiation

Dedicated to my parents

Chapter 1

Introduction

What is a double layer (DL)? A classic plasma DL consists of two sheets of opposite charge, one negative and one positive, that appear in a plasma. As a result, quasi-neutrality is violated in the DL region. The thickness of a DL is on the order of a few Debye lengths. In laboratory plasmas, particularly in radio frequency (rf) driven, helicon source plasmas in an expanding magnetic field, DLs have been the subject of significant investigation since the discovery that current free DLs (CFDLs) form spontaneously in such systems when they are operated at very low neutral pressures.¹⁻³ The laboratory studies have focused on basic plasma science,⁴ the use of ion beams accelerated by the DL for plasma propulsion,⁵⁻⁷ and space-relevant plasma phenomena.⁸ Before these helicon source DL experiments appeared, creation of a DL in a laboratory plasma typically required creating an interface between two different plasma sources⁹ or by applying strong potentials to a boundary.¹⁰ How DLs are able to form in these single helicon source, expanding plasma experiments has been an open question.

In observations, theory, and computation, there is a long-standing correlation between the formation of DLs and phase space holes (ion or electron) in space plasmas.¹¹ Phase space holes are deficits of ions or electrons in a particular region of phase space that are sustained by

a localized minimum or maximum in the plasma potential. DLs and phase space holes typically arise in regions of space associated with accelerated particle populations, such as in the auroral zone¹² and where magnetic reconnection occurs.¹³

In a single species plasma, ions entering a sheath must have a minimum speed, the Bohm sound speed, to satisfy the conditions for a stable sheath at a boundary.¹⁴ The Bohm speed is defined as $\sqrt{k_B T_e / m_i}$, where k_B is the Boltzmann constant, T_e is the electron temperature and m_i is the ion mass. In plasmas with multiple ion species, sheath ion velocities have been investigated in multi-dipole devices¹⁵ but not in helicon devices that produce a DL. A recent theory,¹⁶ later confirmed in experiments,^{17,18} suggested that an ion-ion instability will cause ions in a two-species plasma to fall through a sheath at a common speed if the individual densities of the species are the same. If the ion densities are very different, the ions fall into the sheath at their own Bohm speed. This phenomena provides a unique test for understanding the sheath-like behavior of DLs observed in helicon plasmas. Extending the theory of *Baalrud et al.*¹⁶ to DLs suggests similar behavior should occur.

The remainder of this chapter presents a review of recent DL research. Chapter 2 presents a review of helicon plasma sources and the experimental facility used for these studies. Chapter 3 provides an overview of the diagnostics used for data collection. Chapter 4 details the two-dimensional structure of the plasma plume downstream of an expanding helicon source. Chapter 5 presents measurements of ion flows in multi-species plasmas and pressure effects on ion beam velocity.

The conclusions in Chapter 6 form the basis of a new explanation for ion beam formation in expanding helicon plasmas. The acceleration of the ions occurs over a distance much greater than the Debye length so the acceleration mechanism is not a true DL. This dissertation strongly suggests that the acceleration mechanism is specific to helicon plasma sources. The rf power

couples strongly to the electrons at the skin depth. These electrons then propagate from the plasma source to the expansion region as an annulus. The ions in the center of the upstream plasma are accelerated to supersonic speeds but the bulk plasma downstream flows towards the source. Therefore, the observed particle motion is inconsistent with expectations for a DL. A very similar picture was recently put forth by *Gulbrandsen* and *Fredriksen*¹⁹ although they continued to call the acceleration mechanism a DL.

1.1 DLs

The story of DLs begins in 1929 when Irving Langmuir first characterized them as double sheaths.²⁰ In its simplest description, a DL consists of two spatially separated charge layers, one positive and one negative. As a result, ions (electrons) will be accelerated as they travel from the high (low) potential side to the low (high) potential side. A conventional sheath forms at a plasma boundary or on the surface of an object inserted into the plasma, but a DL is not subject to this constraint. Prior to a sheath, a pre-sheath region accelerates ions to the Bohm speed. A simple pictorial of a DL is shown in Fig. 1.1, taken from a review of early DL studies by *Block*.²¹

Even in a quasi-neutral plasma, quasi-neutrality is violated within the DL. A sheath has a width of a Debye length given by

$$\lambda_D \equiv \left(\frac{\epsilon_0 k_B T_e}{n e^2} \right)^{1/2} \quad (1.1)$$

where ϵ_0 is the permittivity of free space, k_B is the Boltzmann constant, T_e is the electron temperature, n is the plasma density, and e is the elementary unit of charge.²² Beyond the

FIGURE 1.1: A cartoon of a simple DL. Fig. from Ref. [21].

Debye length, the potential of a charge in the plasma is shielded from the rest of the plasma. DLs can have widths of several²³ to thousands²⁴ of Debye lengths.

The next development involving DLs came in 1958 when Hannes Alfvén suggested that DLs were responsible for the aurora by accelerating electrons into the upper atmosphere.²⁵ Subsequent observations provided evidence for Alfvén's hypothesis.²⁶ DLs have been shown to be quite common in Earth's magnetosphere ranging from the auroral zone to the plasma sheet.^{27–30} Auroral phenomena observed on spacecraft has driven an interest in recent DL research.^{31,32} Laboratory studies have been used to understand space plasmas by mimicking the DL potential structure.³³

FIGURE 1.2: Color photograph of luminous plasma potential structure. The boundary of the light emitting region follows the diverging magnetic field lines and lies slightly to the inside of the potential jump. For this photograph the neutral argon pressure was $\sim 3 \times 10^{-3}$ torr. Fig. from Ref. [10].

FIGURE 1.3: Two-dimensional mapping (r, z) of equipotential contours. Note that the scale is 1:1. Numerical values of potential (in volts) are shown for various contours. The ion gyroradius varies from $\simeq 0.4$ cm at $z = 10$ cm to $\simeq 1.2$ cm at $z = 30$ cm. Fig. from Ref. [10].

In 1981 *Perkins*³⁴ predicted in an analytical study that CFDLs could form in diverging magnetic fields. A study by *Alport et al.*,¹⁰ although not current free, is especially illuminating because it showcases the curved, two-dimensional electric potential contours that appear in an expanding field geometry. Fig. 1.2 shows the plasma itself while Fig. 1.3 shows the two-dimensional plasma potential structure. The potential drop in z is greatest in the central region of the DL structure ($-3 < r < 3$ cm) implying a large acceleration in this direction. For radial positions $|r| > 3$, there is little acceleration in z . Early experiments that were not necessarily current-free were reviewed by *Hershkowitz*³⁵ in 1985. The research presented in this dissertation is only concerned with CFDLs.

1.2 DLs in Helicon Plasmas

Much of the recent work on DLs has been conducted in helicon plasmas. Although the individual sizes of experiments may vary, DLs in a helicon plasma device (HPD) share important features. A HPD consists of a Pyrex tube connected to a larger expansion chamber, thereby creating rapid plasma expansion. Additionally, the magnetic field magnitude decreases abruptly in the expansion region. Since helicon plasmas operate by coupling the power into the plasma by way of an antenna wrapped around the Pyrex tube, there is no net current flowing into the plasma. DLs formed by this process are therefore considered current free. A basic review of helicon plasma sources is presented later in Chapter 2.

1.2.1 DL Measurements at Australian National University (ANU)

The first report of a DL in a helicon plasma source was by *Charles* and *Boswell*¹ in 2003. A schematic of the experimental device known as Chi-Kung is shown in Fig. 1.4. The plasma originates in the source on the left in a uniform magnetic field before flowing into the weaker field

FIGURE 1.4: The Chi-Kung experimental device at ANU. Fig. from Ref. [1].

expansion chamber on the right. An axially movable retarding field energy analyzer (RFEA) and Langmuir probe provided measurements of the plasma potential, electron temperature, and plasma density. A key finding in the experiment was that the DL only formed for a low pressure of 0.2 mTorr. For a higher pressure of 3 mTorr, no DL was observed and the plasma obeyed the Boltzmann relation given by

$$n = n_0 \exp(-e\phi/k_B T_e). \quad (1.2)$$

For the lower pressure case of 0.2 mTorr shown in Fig. 1.5, a sharp discontinuity in the measured plasma potential at the source/expansion chamber junction extended over less than 50 Debye lengths. The pressure dependency was explained by calculating the energy gain by the ions in the span of a mean free path. The mean free path increased with decreasing pressure which means that ions acquired more energy before being scattered. *Charles and Boswell*¹ argued that the length of the helicon source chamber determined the pressure threshold. Longer sources would require a lower pressure because ions would be scattered before they could gain

FIGURE 1.5: A sharp discontinuity in the plasma potential in the Chi-Kung helicon plasma device. Fig. from Ref. [1].

enough energy and no DL or ion beam would form. They also proposed that this configuration could be the basis for a new type of plasma thruster with thrust values on the order of mN.

Of significant importance to this dissertation were a series of two-dimensional measurements of the DL structure by *Charles*.³⁶ Two-dimensional mapping of the plasma potential and the ion density just inside of the expansion chamber, shown in Fig 1.6, identified conic sections of high plasma density following the most diverging magnetic field line that lays wholly within the plasma chamber. It is important to note that both the density and plasma potential have been mirrored and radial symmetry assumed. By doing so, the DL takes on a U-shape, i.e. a parabolic potential contour. The potential upstream of the DL is 46 V and downstream it is 36 V. No measurements were taken in the DL and all the intermediate contours are interpolated. In the region immediately following the plasma source, large perpendicular electric fields

FIGURE 1.6: (a) 2D equipotential contours measured with the RFEA for 250 W and 0.053 Pa (0.4 mTorr). The DL extends between the 36 V and 46 V contours. The solid parabolic line represents a fit of the 36 V contour which is the low potential edge of the U-shaped CFDL. The solid diverging line shows the most diverging magnetic field line exiting the source at $(x, z) = (6.7, 30)$ cm. (b) 2D contours of the ion density measured with the RFEA. The density has been calculated from the total ion current from the RFEA and hence is only about ± 20 % accurate. Fig. from Ref. [36].

of 1000 V/m suggest the ions would be accelerated from the center outwards. Since RFEAs are completely incapable of measuring perpendicular particle distribution functions,³⁷ no such measurements were reported by Charles³⁶

Recent measurements in Chi-Kung also indicate the surprising presence of high energy electrons near the edge of the plasma source, which results in enhanced ionization at the edge and creates a radial potential barrier that appears to confine the ions.³⁸ Radial confinement thus appears to be important in establishing the necessary conditions for a DL.

FIGURE 1.7: The logarithm of the LIF amplitude vs. parallel velocity and axial position. Fig. from Ref. [39].

1.2.2 DL Measurements at WVU

Around the same time that *Charles* and *Boswell*¹ were conducting their experiments, the WVU group measured ion beams in a larger helicon device.³ Laser induced fluorescence (LIF) measurements were taken upstream, downstream and inside the DL. Fig. 1.7 shows the logarithm of the amplitude of the parallel ion velocity distribution function (IVDF) versus parallel velocity and axial position. The end of the plasma source is located at $z=150$ cm, which is approximately the same location as the DL. The ion acceleration takes place over a range of 20 cm with strong acceleration over a narrower region of ~ 5 cm. The 5 cm region, located at the maximum of the magnetic field strength gradient, can be regarded as the DL itself.

Both LIF and retarding field energy analyzers (RFEAs) have been used at WVU to characterize the ion beam that appears downstream of the DL.⁴⁰ It is important to note that while an RFEA has a lower detection threshold for ion beams than LIF, RFEAs are unable to reliably measure the details of the parallel ion velocity distribution function.³⁷ Inserting an RFEA in a plasma is perturbative while LIF is not. Also, because RFEAs do not directly measure the IVDF, ion temperatures cannot be measured. LIF has an unfortunate drawback as well. Previous studies of the ion beam amplitude decay in the plasma downstream of a DL

have attributed the decay to quenching of the metastable state probed by LIF due to collisions of the metastable ions with electrons. In other words, the decrease in LIF signal results from the particular requirements of the LIF measurement process and does not necessarily indicate actual decay of the ion beam. In fact, RFEA measurements in this plasma plume⁸ and in other experiments⁴¹ have found that the ion beam persists downstream with little reduction in beam density.

Considerable effort at WVU was directed towards understanding the temporal evolution of the DL. Time-resolved measurements in pulsed plasmas demonstrated that large electrostatic instabilities arise at the same time as strong DLs.⁴² For the first 40 ms of the pulse, the ion beam velocity increases and then maintains a constant velocity throughout the rest of the pulse. A strong correlation exists with a previously identified ion acoustic wave at 17.5 MHz. The appearance of the wave does not completely destroy the DL, rather it limits the current that passes through the DL before the layer becomes unstable. Thus, any spontaneously formed DL will be subject to this constraint. A detailed description of time-resolved experiments at WVU is found in Ref. [43].

1.3 Simulations of DLs

A one-dimensional, unmagnetized, hybrid simulation (particle ions and fluid electrons) of an expanding plasma in a diverging magnetic field showed that a DL will form at the location of rapid plasma expansion.⁴⁴ Downstream of the DL, the plasma expansion was simulated with a simple, spatially dependent loss mechanism. In addition to the accelerated ion beam, a low energy ion population created by ionization and charge exchange collisions was also observed. Shown in Fig. 1.8 is the ion velocity distribution in phase space for a neutral pressure of 1 mTorr. The ion acceleration takes place over many centimeters in the simulation domain; in

FIGURE 1.8: (a) Ion velocity distribution in phase space, showing a supersonic ion beam. Increased brightness indicates increased density. (b) Ion velocity distribution function f_v in arbitrary units in different locations. Fig. from Ref. [44].

the presheath and the sheath. One might then ask, where is the DL? The region of strongest acceleration is deemed to be the DL. The results of the simulation are consistent with experiments summarized earlier (compare Fig. 1.7 and Fig. 1.8).^{3,4} Some recent computational models have included two-dimensional magnetic fields⁴⁵ or tracked particle motion in three-dimensional phase space.⁴⁶ None of those models, however, have been fully self-consistent.

1.4 Models of DL Formation

Initial theory and computation of DL formation focused on one-dimensional models. Those models yield predictions of the neutral pressure threshold for DL formation,^{1,2,41,47} of the potential drop across the DL, of the DL thickness in multiples of Debye lengths,^{5,48,49} and of the relative densities of the ion and electron populations needed to sustain the CFDL.⁵⁰

One-dimensional models are clearly unable to reproduce “U” shaped potential structures, electric fields perpendicular to the magnetic field near the throat of the expansion region and the large variations in plasma density transverse to the magnetic field observed downstream of some DLs. Soon after the initial discovery of DLs in HPDs, an MHD treatment was proposed in 2006.⁴⁸ The physical mechanism proposed focused on the evolution of the sheath (DL) until it runs out of energy. Unfortunately, this theory was one-dimensional and did not take into account the kinetic effects of the electrons near the throat of the HPD.

At the same time, a one dimensional diffusion controlled theory coupled the dynamics of the particles in the non-neutral DL to the diffusive flows of the quasi-neutral plasma in the source and expansion chambers.⁵¹ In addition to the classic four DL particle groups described by *Andrews* and *Allen*,⁵⁰ this new theory added a group of counter streaming electrons formed by the reflection of almost all of the accelerated electrons from the sheath at the insulated end wall of the source chamber to make the DL current free. This would cause additional ionization upstream of the DL, which was corroborated by previous experiments at WVU.^{52,53} Studies that varied the neutral pressure and measured the pressure threshold for DL formation were also consistent with the model.^{1,2,47} The diffusion-model correctly predicts the potential drop of the DL. However, this model did not account for the wide spatial range of acceleration (the pre-sheath region) or the multi-dimensional structure of the DL itself. An in-depth review of the theory with its connection to previous WVU research can be found in Ref. [52].

The discovery of a U-shaped DL in the laboratory suggests an opportunity to compare ground based experiments with auroral data obtained by the FAST satellite.⁵⁴ However, *Singh*⁵⁵ pointed out some errors arising in making such comparisons. In a helicon plasma, the scale length of the diverging magnetic field is smaller than the ion gyroradius radius, while in the aurora the opposite is true. A review article by *Singh*⁵⁵ provided a different explanation for the 2D structure of the DL in the HPD. *Singh*⁵⁵ uses the magnetization and transit times of electrons and ions as the basis for his argument. A particle in a plasma gyrates around magnetic field lines with a cyclotron period of:

$$\tau_\sigma = \frac{2\pi}{\omega_\sigma} \quad (1.3)$$

where $\omega_\sigma = \frac{|q_\sigma|B}{m_\sigma}$, m is the mass, q is the charge of the species, B is the magnetic field strength, and σ is the plasma species. The particle's radius, the Larmor radius, is:

$$\rho_\sigma = \frac{mv_\perp}{q_\sigma B} \quad (1.4)$$

where $v_\perp$ is the perpendicular velocity; assumed here to be the thermal velocity. Comparing the time it takes to traverse the DL to the cyclotron period, the magnetic field's influence on the particles is understood. Because the ions are supersonic and have a larger Larmor radius than the electrons, the ions are unmagnetized while the electrons are highly magnetized. Therefore, much like ambi-polar diffusion,²² the electrons leave the source and follow the rapidly diverging field lines while the ions do not. The subsequent charge separation sets up an ambi-polar electric field that drags the ions outwards, which can explain the perpendicular potential drop measured by *Charles*.³⁶ *Singh* then proposed that the perpendicular field is the source of the parallel potential drop of the DL. The perpendicular electric fields would then be shorted out by conducting boundaries in the plasma.

1.5 Two Dimensional Plasma Potential Measurements

A two-dimensional study of the downstream potential structure by *Saha et al.*⁵⁶ in a helicon source revealed large perpendicular electric fields, as large as 20 V/cm, at the junction of the source tube and expansion chamber (see Fig. 1.9).⁵⁶ The most diverging continuous magnetic field line that lays within the plasma chamber, indicated in Fig. 1.9 by an arrow, sets a boundary for the large perpendicular potential drop. Just outside of this magnetic field line *Saha et al.*⁵⁶ describe secondary lobes in the plasma potential contours. In this region, $(r, z) = (5, 5)$ there is a minimum in the plasma potential. In other words, these structures resemble ion holes aligned to the indicated magnetic field line. The parallel ion acceleration takes place over approximately 5 cm, providing further evidence that this region may not be worthy of the “DL” designation because $5 \text{ cm} \gg \lambda_D$.

Additionally, two dimensional RFEA measurements were performed over the same spatial range. As the RFEA is moved through the DL into the downstream plasma, the background plasma returns and the ion beam current density decays. The beam current density decays radially as well. The ion beam radial extent was approximately 6 cm for a source tube of radius 7.5 cm. The ion beam was found to expand with increasing axial distance suggesting that it is not initially detached from the magnetic field.

The thrust generated by the ion beam has been a subject of detailed study because of its potential for commercial application and tests of a prototype Helicon DL Thruster (HDLT) were recently reported.⁵⁷ In previous studies using RFEAs, detachment of the ion beam from the magnetic field lines has been observed,⁵⁸ with beam divergences of less than 6 degrees.⁵⁹ RFEAs are used for detachment measurements by moving the RFEA further downstream and measuring the radial extent of the beam. When the radial extent of the beam does not increase in radius but the magnetic flux tubes do, the beam is detached. For LIF, metastable quenching

FIGURE 1.9: The 2-D contour plots of the plasma equipotential at 0.1 mTorr. The labels of some contours are shown. The two magnetic field lines marked by arrows just pass through the junction between the quartz tube and the flange of the expansion chamber. Fig. from Ref. [56].

makes measuring detachment impossible. Further downstream the LIF signal decreases but the ion beam is still present.⁴³ Therefore, different magnetic geometries must be compared at the same downstream locations in order to measure detachment of the ions.

1.6 Sheath Experiments in Multi-Species Plasmas

For single ion species plasmas, the Bohm criterion requires that ions entering a sheath must have a velocity $v = c_s = \sqrt{k_B T_e / m_i}$. A presheath forms before the sheath to accelerate the ions to the Bohm speed. The Bohm criterion for multi-ion species plasmas was presented by

Baalrud and *Hegna*.⁶⁰ Two solutions are possible. In the first solution, the ions will reach their individual Bohm speeds at the sheath edge. In the second solution, the ions reach a common sound speed given by

$$c_s = \sum_i \left[c_{s,i}^2 \left(\frac{n_i}{n_e} \right) \right]^{1/2} \quad (1.5)$$

where $\sum_i$ is the summation over the gas species i . Previous experiments in weakly collisional, weakly ionized plasmas containing two ion species with comparable ion densities agreed well with the latter solution.^{17,61} No mechanism to achieve the second solution was established until *Baalrud et al.*¹⁶ developed a theory that included collisional friction from ion-ion instabilities. Recent experiments in multi-dipole devices measuring the sheath infall speed agreed with this theory.^{15,18,62}

DL experiments in multi-ion plasmas at WVU were conducted to test theoretical predictions regarding the infall speed for different ion species.⁶³ It is important to note that the study at WVU did not measure the individual ion densities by launching ion acoustic waves but relied on the neutral gas composition fractions. The total gas flow was maintained at 10 SCCM. Ar-Xe and Ar-He gas mixtures were studied but only argon and xenon are accessible to current LIF schemes. The argon ion beam velocity increased linearly with the addition of xenon from ~ 6.7 km/s in pure argon to ~ 8 km/s for a 4% xenon fraction in the plasma source. Measurements of the argon IVDF for various argon-helium compositions are shown in Fig. 1.10. As helium was added, for fractions up to 30%, LIF signal decreased and the argon ion beam increased in velocity. A higher helium concentration caused the argon beam velocity to be defined more by the lighter ion mass (helium). These results were consistent with the predictions by *Baalrud* and *Hegna*.⁶⁰

FIGURE 1.10: Effect of increasing helium fraction on the argon parallel IVDF in HELIX. Measurements were obtained at $z = 146$ cm. Fig. from Ref. [63].

Chapter 2

Experimental Setup

All the experiments presented in this work were conducted in the Hot hELIcon eXperiment and Large Experiment on Instabilities and Anisotropies (HELIX and LEIA). For these experiments, an in-situ, scanning, mechanical probe was manufactured to allow high resolution LIF measurements of the plasma immediately after the HELIX-LEIA junction. A detailed description of the probe and LIF is reserved for Chapter 3.

2.1 Helicon Plasmas

Helicon plasmas are magnetized, inductively coupled rf plasmas characterized by high efficiency and high density. The first helicon plasma source was developed by Rod Boswell⁶⁴ with densities on the order of 10^{13} cm^{-3} and the signature “blue core” for argon.⁶⁵ The ability to run continuously makes helicon plasmas an attractive source for a variety of plasma research including: space plasma research,^{66,67} plasma propulsion,^{5,57–59} plasma processing,^{68,69} and basic plasma physics.^{70,71} It is still unclear exactly what mechanism is responsible for the efficient coupling of rf power into the plasma. Possible explanations include: collisional processes,^{72,73} helicon

wave penetration,⁷⁴ mode conversion near the lower hybrid frequency,⁷⁵ Landau damping,^{76,77} antenna localized acceleration and^{78,79} non-linear trapping of fast electrons.^{80,81} Some reviews of helicon research can be found in *Boswell and Chen*,⁸² *Chen and Boswell*,⁸³ and *Scime et al.*⁸⁴

2.1.1 Helicon Waves

Helicon waves are low frequency electromagnetic waves confined by cylindrical geometry. Whistler waves are the unbounded brethren to the helicon wave; both are right hand circularly polarized waves. Whistler waves derive their name from their descending tones, first reported during World War 1 by an engineer in the German army.⁸⁵ It was later determined that lightning creates the waves in the atmosphere.⁸⁶ *Aigrain*⁸⁷ was the first to use the term “helicon” to describe the propagation of bounded right hand circularly polarized waves in a solid rod of sodium. Additional exploration of helicon waves in the 1960s took place in gaseous plasmas⁸⁸ and solid state plasmas.⁸⁹ Some of the first published studies on the basic theory of helicon waves include: Woods,⁹⁰ Klozenbreg *et al.*,⁹¹ and Davies *et al.*⁹²

The dispersion relation for a helicon wave propagating at an angle θ relative to the magnetic field in an infinite magnetized plasma is

$$n_{ref}^2 = \frac{\omega_{pe}^2}{\omega(\omega_{ce} \cos \theta - \omega)} \quad (2.1)$$

where n_{ref}^2 is the index of refraction defined as $n_{ref} = k_{||}c/\omega$, $k_{||}$ is the wave number parallel to the magnetic field, c is the speed of light, ω is the wave frequency, ω_{pe} is the electron plasma frequency and ω_{ce} is the electron cyclotron frequency. For $\theta=0$ and $\omega \ll \omega_{ce}$, the dispersion relation reduces to

$$n_{ref}^2 = \frac{\omega_{pe}^2}{\omega\omega_{ce}}. \quad (2.2)$$

The frequency range of propagation is sufficiently high that the ions cannot respond to the wave electric field and sufficiently low that electron inertia is small, i.e., $\omega_{ci} \ll \omega \ll \omega_{ce}$. A calculation of the group velocity reveals that $v_g \sim \sqrt{\omega}$.⁹³ Therefore, high frequency waves travel faster than low frequency waves⁸⁶ giving rise to the “whistling” effect described earlier. A very good introduction of helicon plasmas including wave propagation, mode transitions and antenna design can be found in Chapter 8 of *Chabert and Braithwaite*.⁹⁴

2.2 HELIX-LEIA Experimental Facility

HELIX-LEIA consists of two separate regions: a small diameter plasma source (HELIX) and a large aluminum expansion chamber (LEIA). Originally developed for studying magnetospheric relevant plasmas in the laboratory, HELIX-LEIA is capable of producing high beta ($\beta = \frac{P}{B^2/2\mu_0}$) plasmas. The abrupt transition from HELIX to LEIA is ideal for studying the spontaneous formation of CFDLs at low neutral pressures. A photograph of HELIX-LEIA is shown in Fig. 2.1.

2.2.1 Plasma Chamber

HELIX consists of a 61 cm long, 10 cm diameter PyrexTM tube connected to a 91 cm long, 15 cm diameter stainless steel chamber. The other end of the Pyrex tube is connected to a four-way glass cross. The other three legs of the glass cross are terminated with a turbo-molecular drag pump, an ion gauge, and a 12 inch stainless steel flange fitted with a 4 inch viewport. The stainless steel chamber has one set of four 6 inch ConflatTM crossing ports in the center of the chamber and four sets of four 2.75 inch ConflatTM crossing ports. The 2.75 inch ConflatTM crossing ports are spaced evenly on either side of the set of the 6 inch ConflatTM crossing ports.

FIGURE 2.1: Photograph of HELIX-LEIA.

FIGURE 2.2: A schematic view of the HELIX chamber. The positions marked A, B, D and E are the 2.75 inch crossing ports while the larger 6 inch ports are at location C. Figure adapted from Ref. [39].

The stainless steel chamber opens up into a 1.8 m inner diameter, 4.5 m long expansion chamber (LEIA).

2.2.2 Vacuum System

The vacuum pressure in HELIX-LEIA is maintained by three turbo-molecular drag pumps, each backed by a diaphragm roughing pump. Each pump is separated from the vacuum chamber by

a gate valve. The gate valves are part of an interlock system that closes the valves automatically if the pressure rises beyond a threshold value. One pump is located under the glass cross at the front of HELIX while the other two are located at the back end of the LEIA chamber. The HELIX turbo pump is set to a constant rotational frequency of 549 Hz while the two LEIA turbo pumps can be each be set to 400 Hz or 600 Hz. Therefore, by changing only the LEIA turbo pump frequency and or closing a gate valve, there are five pumping frequencies available (400 Hz, 600 Hz, 800 Hz, 1000 Hz and 1200 Hz). A higher pumping frequency reduces the ion-electron and ion-neutral collisions that cause metastable quenching.^{2,4} With the LEIA turbo pumps each set to 400 Hz, the three pumps maintain a base pressure of 10^{-7} Torr. For the experiments conducted in this work, the two LEIA turbo pumps are set to 400 Hz and the HELIX turbo pump is set to 549 Hz with all gate valves open.

The pressure in HELIX-LEIA is measured with a series of gauges placed at various locations. A Balzers PKR250 pressure gauge is located on the top of the glass cross and another is located at the back end of LEIA. The Balzers gauges achieve a full range of pressure readings by combining a cold cathode gauge for pressures below 10^{-2} Torr and a Pirani gauge for pressures above 10^{-2} Torr. Unfortunately, the Balzers gauges are not gas species independent so a Baratron[®] capacitance manometer is used for experiments involving multiple gas species. The Baratron[®] is attached to a 2.75 inch port on the HELIX chamber at position D in Fig. 2.2.

The gas flow into HELIX is regulated by two MKS1179 mass flow valves controlled by a single PR-4000 power supply. HELIX has two locations used for gas injection. The first, known as the “end feed”, is adjacent to the Balzers pressure gauge at the glass cross. The second, known as a “mid feed” is located on the underside of the stainless steel chamber of HELIX at position D in Fig. 2.2. Operating pressures in argon range from 0.1 mTorr to 10 mTorr. All the experiments in this work utilized the end feed location because it allows easier access to low neutral pressure plasmas.

2.2.3 Magnetic Field

Ten water cooled electromagnets produce a steady-state, nearly uniform axial magnetic field of 0-1200 G in HELIX. Each magnet has 46 internal copper windings with a resistance of $17\text{ m}\Omega$ and an inductance of 1.2 mH. A maximum current of 400 Amperes is provided by two Xantrex XFR power supplies connected in parallel. The magnets rest on a rail system so each magnet can be moved to access various flanges and diagnostics. The magnets are chilled by a Neslab System III Heat Exchanger which is cooled using the building chilled water.

Seven custom built 9' diameter water cooled electromagnets produce a steady-state, uniform axial magnetic field of 0-150 G in LEIA. Each magnet has five sets of aluminum tubing wound into five two-coil pancakes of four layers each, for a total of 40 turns per magnet. The aluminum tubing is hollow with a square cross section of 0.5" x 0.5" and is wrapped in insulating paper. A maximum current of 200 Amperes is provided by a DC EMHP power supply resulting in a magnetic field of 150 Gauss. The LEIA magnets are chilled by building water.

Typical operating conditions result in an axial magnetic field gradient of 10 Gauss/cm over a distance of 70 cm. Fig. 2.3(a) shows a diagram of the HELIX-LEIA system drawn to scale. Fig. 2.3(b) shows the on axis magnetic field strength and magnetic field gradient versus axial position. The magnetic field is calculated with a two-dimensional modeling code and confirmed repeatedly with measurements from a Gaussmeter. Fig. 2.3(c) shows contours of constant magnetic flux for the rapidly diverging magnetic field. The HELIX field is kept constant at 700 G but two different LEIA field strengths, 70 G (dashed red line) and 14 G (solid blue line), are plotted to show the effect of only changing the downstream magnetic field.

FIGURE 2.3: (a) Diagram of HELIX-LEIA showing the axial position of important features. (b) The magnetic field strength and magnetic field gradient in HELIX-LEIA versus axial position. (c) Contours of constant magnetic flux for LEIA magnetic fields of 70 G (dashed red line) and 14 G (solid blue line). The field in HELIX is 700 G. Figure adapted from Ref. [43].

FIGURE 2.4: Diagram of the $m = +1$ helical antenna and the actual design of the antenna.⁹⁵ Figure adapted from Ref. [39].

FIGURE 2.5: Matching circuit for the helicon antenna for HELIX. C_L is the load capacitor and C_T are the three tuning capacitors. The copper rods are surrounded by hollow pieces of Macor[®] with copper sheets on the outside which are grounded. Figure adapted from Ref. [39].

2.2.4 RF Antenna and Matching Network

Up to 2.0 kW of rf power is coupled through a 19 cm $m = +1$ helical antenna over a frequency range of 6-18 MHz. The antenna is wrapped around the pyrex tube as shown in Fig. 2.4. The rf system consists of a 50 MHz wavetek model-80 function generator that provides the rf signal to an ENI 1000, 30 dB amplifier. The rf power is then transmitted from the amplifier to a π matching network with a high frequency co-axial cable. The matching network matches the 50 Ω output impedance of the amplifier to the antenna/matching network system. The matching

network consists of one load capacitor and three tuning capacitors. The load capacitor has a range of 20-2000 pF. Two of the tuning capacitors have a range of 4-250 pF and the third has a range of 5-500 pF. All the capacitors are Jennings high voltage vacuum capacitors. The three tuning capacitors are connected in parallel with each other and in series with the antenna. The load capacitor is connected in parallel with the tuning capacitors and the antenna as shown in Fig. 2.5. The capacitors are connected to the antenna with rods of silver plated copper. In order to minimize power reflected to the amplifier and maximize the forward power coupled to the antenna, the real impedance of the matching network must equal the 50Ω output impedance of the amplifier while the imaginary part of the combined matching network-antenna circuit must be zero. *Chen*⁹⁶ calculated the required load (C_L) and tuning (C_T) capacitances for an inductive load to be

$$C_L = \frac{1}{2\omega R} \left[1 - \left(1 - \frac{2R}{R_0} \right) \right]^{\frac{1}{2}} \quad (2.3)$$

and

$$C_T = \left[\omega X - \frac{1 - \frac{R}{R_0}}{C_L} \right]^{-1} \quad (2.4)$$

where R is the real resistance of the antenna, $R_0 = \sqrt{L/C}$ is the normalized impedance, and $X = \omega L$ is the reactive impedance of the antenna.

After the discharge is initiated, the effect of the inductive load of the plasma on the antenna must be considered. For a typical helicon plasma source in the “inductive” or “helicon” mode, Eq. 2.4 becomes

$$C_T^{-1} = \omega^2 L - \frac{\left(1 - \frac{R}{R_0} \right)}{C_L} \quad (2.5)$$

where L is the total inductance in the antenna portion of the circuit.

2.2.5 HELIX-LEIA Plasma Parameters

Table 2.1 shows typical values for various plasma parameters in HELIX and LEIA.

Parameters (units)	HELIX	LEIA
B (Gauss)	500-1200	5-140
$p_{neutral}$ (mTorr)	0.1-10	0.01-1.0
n ($\times 10^{12}$ cm $^{-3}$)	0.1-10	0.001-0.1
T_e (eV)	4-12	2-20
T_i (eV)	<1	<1
λ_{debye} (cm)	5×10^{-3} - 8×10^{-4}	3×10^{-2} - 1.6×10^{-1}
r_e (cm)	4×10^{-2} - 2×10^{-2}	2.5×10^{-2} - 2.5
r_i (cm)	3×10^{-1} - 6×10^{-1}	4-60
ω_{pe} (10^{10} rad/s)	11.1-111	1.1-11.1
ω_{ce} (10^9 rad/s)	8.7-21	0.087-2.4
ω_{ci} (10^6 rad/s)	0.12-0.3	1.2×10^{-3} - 3.2×10^{-2}

Chapter 3

Diagnostic Methods

Plasma diagnostic techniques fall into three broad categories: *ex situ*, *in situ* intrusive and *in situ* non-intrusive. *Ex situ* measurement techniques involve removing a portion of material from a system and transferring it elsewhere for examination. All *in situ* techniques are intrusive to some extent. However, in cases where the perturbation to the plasma is negligible, the technique is considered *in situ* non-intrusive. For *in situ* intrusive measurement techniques, the perturbation is considerable and must be properly accounted for in the analysis.

The three diagnostics used in this work are: Langmuir probes, electrostatic probes and laser induced fluorescence (LIF). All three are considered *in situ* intrusive because they are located inside the vacuum chamber. Langmuir probes and electrostatic probes are simply conductors inserted into the plasma. LIF is accomplished with an *in situ*, scanning, mechanical probe (a.k.a. the superprobe). All of these techniques are well established and have been tailored to meet the specific requirements of the experiments in this work.

3.1 Langmuir Probe

The use of a current collecting probe in a plasma was first discussed by Irving Langmuir in 1929,²⁰ hence it is widely known as the Langmuir probe. In its simplest form, a Langmuir probe consists of a single metal wire (usually tungsten) that is inserted into the plasma. The Langmuir probe is used for determining the electron temperature, density and occasionally the plasma potential. Manufacturing a Langmuir probe and collecting data are quite simple. The analysis of Langmuir probe data can be extremely complex especially in cases of drifting, non-Maxwellian or collisional plasmas which require comparison with complex theoretical models. The perturbative effects of the probe itself can be significant. Thankfully, the available literature is quite extensive and some excellent reviews can be found in *Chen*,⁹⁷ *Merlino*,⁹⁸ *Hutchinson*,⁹⁹ *Hershkowitz*,¹⁰⁰ *Demidov et al.*,¹⁰¹ *Sheridan et al.*¹⁰² and others. The use of Langmuir probes in this dissertation is limited so only a brief overview of the theory, design and data analysis relevant to the conducted experiments will be presented here.

3.1.1 Langmuir Probe Theory

A typical Langmuir probe measurement involves applying a voltage bias to the conductor inserted into the plasma and measuring the current drawn. The collected current is plotted as a function of the applied voltage and the result is known as an I-V trace. An ideal I-V trace is shown in Fig. 3.1. A Langmuir probe inserted into typical plasmas will charge up negatively because the electrons have a higher flux than the ions due to their lighter mass and higher mean velocities.

A high impedance voltage source is used to control the probe's bias voltage. The floating potential V_f , labeled in green in Fig. 3.1, is the voltage to which the probe charges such that the current drawn by the probe vanishes. For applied voltages increasingly less than the floating

FIGURE 3.1: An example of an ideal Langmuir probe I-V trace with the ion saturation current, electron saturation current, floating potential and plasma potential indicated. Fig. from Ref. [103].

potential, the probe repels electrons and attracts ions resulting in the ion saturation current. Eventually all the electrons are repelled and the ion current reaches a maximum. Biasing the probe more positive than the floating potential has the opposite effect, leading to the electron saturation current. It is important to note that V_f is not the potential of the plasma. The inflection point (2nd knee) where the slope of the I-V trace begins to decrease near the electron saturation current is the plasma potential. There is no potential difference between the probe and the plasma when the applied bias voltage is at the plasma potential. The magnitude of the electron saturation current is much greater than the ion saturation current because electrons are more mobile than ions.

Certain assumptions are required to determine the plasma parameters from a Langmuir

probe. Assuming a collisionless plasma with no magnetic field and Maxwellian particle distributions, the current near the floating potential can be written as

$$I(V_0 - V_p) = n_e e A_p \left(\frac{k_B T_e}{m_i} \right)^{1/2} \left[\left(\frac{1}{2} \right) \left(\frac{2m_i}{\pi m_e} \right) e^{\frac{e(V_0 - V_p)}{k_B T_e}} - \left(\frac{A_s}{A_p} \right) e^{-\frac{1}{2}} \right] \quad (3.1)$$

where V_0 is the externally applied voltage, V_p is the plasma potential, n_e is the plasma density, e is the elementary charge, T_e is the electron temperature, m_i is the ion mass, m_e is the electron mass, A_s is the surface area of the sheath around the probe, and A_p is the surface area of the probe.⁹⁹ When the probe size is much larger than the thickness of the surrounding sheath, $\frac{A_s}{A_p} \approx 1$ is a good approximation. The sheath surrounding the probe is on the order of the Debye length. As V_0 becomes more negative, the first term in the brackets of Eqn. 3.1 becomes negligible. Keeping only the second term inside the brackets yields

$$I_{si} = -0.61 e n_e A_p \sqrt{\frac{T_e}{m_i}} \quad (3.2)$$

where I_{si} is the ion saturation current in an unmagnetized plasma.

The electron temperature and plasma density are the two unknowns in Eqn. 3.2. The electron temperature is obtained by taking the derivative of Eqn. 3.1 with respect to the voltage $V = V_0 - V_p$ which yields

$$\frac{dI(V)}{dV} \cong \frac{e}{T_e} (I - I_{si}) + \frac{dI_{si}}{dV}. \quad (3.3)$$

The second term can be neglected because $dI_{si}/dV \gg dI(V)/dV$ in the saturation regime.

Solving Eqn. 3.3 for T_e yields

$$T_e = \frac{e[I(V) - I_{si}]}{\frac{dI(V)}{dV}}. \quad (3.4)$$

The electron temperature is obtained by performing a linear fit to the semi-logarithmic plot of $\ln(I - I_{si})$ versus V and taking the inverse of the slope of the fit. The plasma density in Eqn.

3.2 can be obtained by using the measured ion saturation current and the calculated electron temperature.

Unfortunately, high density plasmas such as a helicon source do not produce an ideal I-V trace. In a helicon plasma a cylindrical Langmuir probe cannot achieve true electron saturation because the sheath surrounding the probe continues to expand and collects more electrons as the voltage is increased.¹⁰⁴ Without reaching electron saturation the plasma potential cannot be directly measured and must be approximated.

For plasmas with $T_i < T_e$ a useful relationship between the plasma potential and floating potential exists. The ion current at V_f is

$$j_i = \frac{1}{4} n e \sqrt{\frac{8 k_B T_e}{\pi m_i}} \quad (3.5)$$

and the electron current at V_f is

$$j_e = \frac{1}{4} n e \sqrt{\frac{8 k_B T_e}{\pi m_i}} e^{\left(\frac{e(V_f - V_p)}{k_B T_e} \right)} \quad (3.6)$$

The total current to the probe at V_f is zero, in other words $j_i = j_e$. After some algebra of the above terms, the plasma potential can be expressed as¹⁰⁵

$$V_p = V_f + \frac{k_B T_e}{2e} \ln \left(\frac{T_e m_i}{T_e m_e} \right) = V_f + \frac{k_B T_e}{2e} \ln \left(\frac{m_i}{m_e} \right) \quad (3.7)$$

For argon ions, $m_i = 40m_p$, where m_p is the proton mass and Eqn. 3.7 reduces to

$$V_p = V_f + 5.6 T_e. \quad (3.8)$$

The difference in the floating potential and the plasma potential for argon is thus 5.6 times the

electron temperature. Therefore, the slope of dI/dV can be well approximated by $dI/dV_{applied}$ for a reasonable estimate of the electron temperature.

Notably absent in the discussion of Langmuir probe theory up until this point are magnetic fields. Ions and electrons gyrate around magnetic field lines, limiting transport across field lines and thus restricting the particle flux to the probe. The effects of the magnetic field are determined by the ratio of the gyroradius to the dimensions of the probe. If the probe is much larger than the gyroradius for a given species, that species will be impeded from reaching the Langmuir probe. Eqns. 3.1 - 3.4 must be modified to account for cross-field transport and collisions.⁹⁹ Including magnetic field effects on the ions, *Hutchinson*⁹⁹ showed that Eq. 3.2 must be adjusted to

$$I_{si} = -0.49en_e A_p \sqrt{\frac{T_e}{m_i}}. \quad (3.9)$$

The effect of strong rf fields have also been ignored up to this point. RF fields cause both acceleration and deceleration (sloshing) of the electrons towards the probe. The result is an error in the floating potential and a broadening of the electron distribution function. To minimize these effects, the Langmuir probe must be rf compensated. *Sudit* and *Chen*¹⁰⁶ developed a method to compensate for the rf fields in helicon plasmas involving a floating electrode. The floating electrode is exposed to the plasma potential fluctuations and connected to the probe tip through a small capacitor, forcing the probe tip to follow the potential oscillations, thereby reducing the sheath impedance. The Langmuir probe used in these experiments has a floating electrode, but it is not directly exposed to the plasma. A set of rf chokes are connected in between the probe tip and the voltage source, increasing the impedance of the current measurement circuit at the rf frequency and its harmonics.

3.1.2 Langmuir Probe Design

Two different designs of Langmuir probes were used in the experiments presented in this dissertation. A cylindrical Langmuir probe was used for measurements in HELIX while a planar Langmuir probe was used for measurements in LEIA.

3.1.2.1 HELIX Cylindrical Probe

Fig. 3.2 shows a photograph of the cylindrical Langmuir probe used in HELIX. The Langmuir probe electrode is a $500\ \mu\text{m}$ carbon tip with an exposed length of 2 mm protected by an alumina (Al_2O_3) shield. A set screw fixes the alumina inside a brass connector. The carbon electrode is wrapped in a strand of silver-coated copper wire which is then soldered to the inside of the brass connector. A 10 nF RF capacitor is soldered to the brass connector. A MacorTM cap surrounds the brass connector assembly. A set of rf chokes (inductors), designed to block a particular range of frequencies, is attached to the probe wire. The rf chokes, in order from the brass connector are, in MHz: 23, 47, 23, 7.3 and 6.8. The end of the rf choke chain is soldered to a shielded coaxial wire which is attached to a vacuum feedthrough. A RG-58 BNC cable carries the probe signal to a Keithley 2400 SourceMeter. The source meter is controlled by custom software written in LabWindowsTM.

3.1.2.2 LEIA Planar Probe

Fig. 3.3 is a schematic highlighting the important features of the Langmuir probe used in LEIA. It is important to note that the figure is not to scale. The probe tip is a circular, tungsten sheet with diameter 0.635 cm. The tungsten disk is spot-welded to a 0.5 mm in diameter tungsten rod. The tungsten tip faces the flowing plasma from HELIX with the backside of the disk coated with aluminum oxide. The tungsten rod is inserted into a hollow alumina shield and attached by

FIGURE 3.2: Photograph of the Langmuir probe used in HELIX. A protective MacorTMcap that surrounds the brass connector has been removed to reveal probe tip internal components. Fig. from Ref. [107].

FIGURE 3.3: A schematic of the Langmuir probe used in LEIA including a) tungsten probe tip, b) alumina shaft, c) MacorTMcap, d) set screw, e) copper base, f) capacitor, g) stainless steel probe shaft and h) chain of rf chokes. Figure not to scale.

a set screw to a copper base. A 10 nF capacitor is attached to the copper base. This assembly is surrounded by a 3.81 cm long, 0.953 diameter MacorTMcap. A 0.635 cm diameter stainless steel tube carries the probe signal, via a coaxial wire, to a vacuum feedthrough. A set of rf chokes is attached to the probe wire just before the feedthrough. The rf chokes, in order from the probe tip to the feedthrough are, in MHz: 26,53,26,13.2,6.8. The rest of the LEIA Langmuir probe setup is the same as the HELIX Langmuir probe described in the previous section.

3.1.3 Langmuir Probe Analysis

For electron distributions that are not Maxwellian, the basic Langmuir probe theory of Section 3.1.1 is insufficient. Cylindrical Langmuir probes cannot distinguish the approach direction of the particles they collect, but they can resolve the energies of those particles. Therefore, it is possible to construct the electron energy distribution function (EEDF) f_e or the electron energy probability functions (EEPF) f_p . The electron temperature and density are then calculated by taking moments of the distribution function. The electron density is

$$n_e = \int_0^\infty f_e(E) dE, \quad (3.10)$$

where f_e has units of $\text{eV}^{-1}\text{m}^{-3}$ and the electron temperature is

$$T_e = \frac{2}{3} \langle E \rangle = \frac{2}{3} \left[\frac{1}{n_e} \int_0^\infty E f_e(E) dE \right]. \quad (3.11)$$

Note that T_e (in eV) is defined relative to the mean energy of the distribution as opposed to the distribution width for a Maxwellian distribution. It is common in the literature for authors to interchangeably use the EEDF or EEPF. The two are related by

$$f_p(E) = \frac{f_e(E)}{\sqrt{E}} \quad (3.12)$$

with $f_p(E)$ having units of $\text{eV}^{-3/2}\text{m}^{-3}$.

The first to use the analysis to be described here was Druyvesteyn, hence it is known as Druyvesteyn analysis.¹⁰⁸ The electron current from the Langmuir probe is given by

$$I_e = \frac{eA_p}{\sqrt{8m_e}} \int_{eU}^\infty (E - eU) f_p(E) dE \quad (3.13)$$

FIGURE 3.4: An example of Druyvesteyn analysis showing the EEPF. The black curve describes a Maxwellian distribution over several decades, whereas the red curve is not Maxwellian. Fig. from Ref. [107].

where A_p is the probe area, $f_p(E)$ is the EEPF and $U = |V - V_p|$ such that $E = eU$. The plasma potential V_p is the voltage of maximum $\frac{dI_e}{dV}$ or the zero crossing of $\frac{d^2I_e}{dV^2}$. The EEDF and EEPF are then defined as

$$f_e(E) = \frac{\sqrt{8m_e}}{e^3 A_p} (eU)^{1/2} \frac{d^2 I_e}{dU^2} \quad (3.14)$$

and

$$f_p(E) = \frac{\sqrt{8m_e}}{e^3 A_p} \frac{d^2 I_e}{dU^2}, \quad (3.15)$$

respectively. In an EEPF, a Maxwellian distribution appears as a straight line. Fig. 3.4 shows an example of Druyvesteyn analysis for a Maxwellian and Non-Maxwellian EEPF. The Maxwellian distribution is nicely represented by a linear fit while the Non-Maxwellian distribution has a smaller group of low energy electrons.

The Druyvesteyn analysis described above is for low pressure, unmagnetized plasmas,

where collisions inside the sheath are negligible and diffusion is isotropic. For magnetized plasmas such as HELIX with anisotropic diffusion across and along field lines, an extended Druyvesteyn analysis must be used. The generalized probe current is¹⁰⁹

$$I_e = \frac{eA_p}{\sqrt{8m_e}} \int_{eU}^{\infty} \frac{(E - eU)f_p(E)}{1 + [(E - eU/E)]\Psi(E)} dE \quad (3.16)$$

where $\Psi(E)$ is the diffusion parameter defined as^{110,111}

$$\Psi(E) = \frac{r_p}{\gamma\rho_e(E)} \ln \left(\frac{\pi l_p}{4r_p} \right). \quad (3.17)$$

Eqn. 3.17 depends on the length of the probe tip l_p and a unitless geometric factor, $\gamma = 4/3 - 0.62 \exp(-\lambda_e/2r_p)$, based on the electron momentum loss scale length λ_e and the probe radius r_p where ρ_e is the energy dependent electron gyroradius. There are three regimes depending on the value of Ψ . If $\Psi \ll 1$, then Eqn. 3.17 reverts to Eqn. 3.13 and the conventional Druyvesteyn method is acceptable. For $\Psi \sim 1$, a more complex Druyvesteyn analysis is required. For the experiments in this work, $\Psi \gg 1$ and $f_p \propto dI_e/dU$. In this regime, the diffusion parameter is very large and the probe collects electrons tied to the field lines that intersect the probe surface. The EEPF is found from the first derivative of the probe current given by

$$\frac{dI_e}{dU} = \frac{eA_p}{\sqrt{8m_e}} \left[\frac{eU}{\Psi} f_p(E) + \int_{eU}^{\infty} \frac{E f_p(E)}{(1 + \Psi)[(1 + \Psi)E - \Psi eU]} dE \right]. \quad (3.18)$$

The EEPF from Eqn. 3.15 then becomes^{109,112}

$$f_p(E) = \frac{3\gamma\sqrt{2m_e}}{2e^3A_p} \frac{\Psi}{U} \frac{dI_e}{dU}. \quad (3.19)$$

3.2 Electrostatic Probe

The electrostatic probe used in these experiments is an uncompensated multi-tip Langmuir probe. The particular structure of the electrostatic probe consists of three tips of tungsten arranged in the shape of a triangle, hence the more colloquial name, triple probe. The tips are separated spatially to measure the electric field and to study waves propagating in the plasma.

3.2.1 Electrostatic Probe Theory

Without an applied voltage, the conducting tips of the electrostatic probe charge up to the floating potential at which point no current flows to the probe. Electrostatic waves propagating in the plasma cause changes in the particle flux to the probe, which result in fluctuations in the probes' measured floating potential. The local DC electric field and the wavelength of the travelling electrostatic wave are determined from the triple probe measurements. Assuming that the variation in electron temperature is ignorable over the scale of the probe tip separation, the DC electric field is calculated from $E = \frac{\Delta\phi}{d}$ where $\Delta\phi = \phi_1(x_1) - \phi_2(x_2)$ is the difference in floating potential between two probe tips and $d = x_1 - x_2$ is the spatial separation of the two probe tips. The wave number of a travelling wave is given by $k = \frac{\Delta\theta}{d}$, where $\Delta\theta$ is the measured phase difference and d is the distance between two probe tips. A power spectrum of the probe signal yields a measure of the strength of the fluctuations at the frequencies of interest. Knowing the frequencies (ω) and the wavelengths (k), the dispersion relation (ω versus k) for ambient waves is determined.

3.2.2 Electrostatic Probe Design

The electrostatic probe used to collect data in LEIA consists of three tungsten 500 μm diameter tips. Each tip has an exposed length of 3 mm and is placed in a separate alumina tube. The

FIGURE 3.5: Dimensions and orientation of the triple probe to the coordinate axes. Each tip has a reference number in blue.

alumina tubes are placed in a MacorTM cap which is locked in a 3/16 inch stainless steel tube with a set screw. Fig. 3.5 shows the dimensions of the triple probe and the orientation of the tips in the lab frame.

Underneath the MacorTM cap, the tungsten filaments are spot-welded to silver-plated copper wire. Each spot-welded connection is wrapped in electrical shrink wrap to keep the tips from touching each other and the stainless steel tube. The stainless steel tube is connected to the superprobe which will be described later in this chapter. The signals from each tip are sent to a digital oscilloscope for processing and storage.

3.2.3 Electrostatic Probe Analysis

Waves in a plasma are identified by measuring the wave frequency, propagation direction, and wavelength. For waves having sinusoidal behavior $\sin(\mathbf{k} \cdot \mathbf{r} - \omega t)$ and traveling from one tip to another, a phase difference $\theta = \mathbf{k} \cdot \mathbf{r}$ occurs because of the finite transit time of the wave.

The phase difference is directly measured from the time series measurements of the two signals. However, typical timeseries include noise and multiple wave signals. The phase difference for a specific wave frequency is determined by using the cross power spectrum of two measured time series. The cross power spectrum of two time series is defined as the product of the Fast Fourier transform (FFT) of one time series and the complex conjugate of the FFT of the other time series. Let $f_1(x_1, \omega)$ and $f_2(x_2, \omega)$ represent the two individual time series. The FFT of each time series is defined as

$$\Phi_1(x_1, \omega) = \int_{-\infty}^{\infty} f_1(x_1, t) \cos(\omega t) dt - i \int_{-\infty}^{\infty} f_1(x_1, t) \sin(\omega t) dt \quad (3.20)$$

and

$$\Phi_2(x_2, \omega) = \int_{-\infty}^{\infty} f_2(x_2, t) \cos(\omega t) dt - i \int_{-\infty}^{\infty} f_2(x_2, t) \sin(\omega t) dt, \quad (3.21)$$

where x_1 and x_2 are the respective probe tip locations. The cross power spectrum $P_{12}(\Delta x, \omega)$ is the product of Eqn. 3.20 and the complex conjugate of Eqn. 3.21;

$$P_{12}(\Delta x, \omega) = \Phi_1(x_1, \omega) \Phi_2^*(x_2, \omega) \quad (3.22)$$

where $\Delta x = x_1 - x_2$ is the spatial separation of the probe tips. Expanding Eqn. 3.22, the cross power spectrum is written as a function of the real and imaginary parts of each FFT;

$$\begin{aligned} P_{12}(\Delta x, \omega) = & (Re\{\phi_1\}Re\{\phi_2\} + Im\{\phi_1\}Im\{\phi_2\}) \\ & + i(Im\{\phi_1\}Re\{\phi_2\} - Re\{\phi_1\}Im\{\phi_2\}). \end{aligned} \quad (3.23)$$

In the complex plane, the phase difference is simply the angle between the real and imaginary vectors of the cross power spectrum which is:

$$\theta(\omega) = \tan^{-1} \left(\frac{(Im\{\phi_1\}Re\{\phi_2\} - Re\{\phi_1\}Im\{\phi_2\})}{(Re\{\phi_1\}Re\{\phi_2\} + Im\{\phi_1\}Im\{\phi_2\})} \right). \quad (3.24)$$

A large signal to noise ratio (SNR) has been assumed in these calculations. When the SNR is small, ensemble averaging of many cross-power spectra significantly improves the accuracy of the phase measurements because the random errors decrease as $1/\sqrt{M}$, where M is the number of data samples recorded.¹¹³ Large values of M require significant computing capability as well as vast amounts of storage space. To accomplish these tasks, a LeCroy WaverunnerTM604Zi oscilloscope was used to perform real-time averages of the FFTs and cross power spectra. The averages were digitally stored for later analysis.

A challenge with FFT measurements is aliasing, i.e., under sampling, in both the frequency and spatial domains. Aliasing in the frequency domain occurs when a periodic signals' frequency is larger than the Nyquist frequency ($f_{Nyquist} = 0.5 \times \text{sample rate}$) of the data acquisition system. In an FFT of an aliased time series measurement, the frequency of a high frequency signal will appear downshifted. Harmonics of the artificially downshifted frequency could also appear in the FFT. For a sampling rate of 50 MHz, the Nyquist frequency will be 25 MHz, which is double the frequency of the rf antenna for these experiments. The oscilloscope was bandwidth limited to 20 MHz to minimize temporal aliasing concerns.

Spatial aliasing occurs when the wavelength of the wave is smaller than the separation distance between the two tips. Fig. 3.5 depicts a situation where the measured phase differences are the same even for two different waves. The probe separation distance places a lower bound on the wavelength λ of waves that can be resolved. The minimum wavelength is simply $\lambda_{min} = 2d$. The probe tip separation in r is 0.291 cm so wavelengths of the order ~ 0.6 cm are measurable.

FIGURE 3.6: An illustration of spatial aliasing for a pair of fixed probes. The black dots and dashed lines represent the spatial location of the probe tips while the color sinusoids represent the wave amplitudes at an instant in time. Fig. from Ref. [114].

Measured wavelengths in these experiments are approximately 20-30 cm, eliminating concerns about spatial aliasing.

3.3 Laser Induced Fluorescence

Laser induced fluorescence (LIF) is the primary diagnostic method used in this dissertation. The LASER (Light Amplification by Stimulated Emission of Radiation), invented in the early 1960s, has spawned numerous laser based spectroscopic measurement techniques. LIF employs a laser tuned to a natural absorption line of an atom or ion to induce emission from the upper pumped state to either the same initial state (resonant LIF) or a different third state (non-resonant LIF). In 1966, *Yardley* and *Moore* were the first to observe LIF from molecules other than the laser medium.¹¹⁴ Using a single frequency argon laser, *Stern* and *Johnson* provided the first experimental evidence of LIF in a plasma.¹¹⁵ *Hill et al.*¹¹⁶ were the first to establish the capability of tuning the laser wavelength to obtain velocity selective LIF measurements. *Hill et al.*¹¹⁶ used a resonant LIF scheme and measured the absorption line shape with a velocity resolution determined by the natural linewidth of the laser transition rather than the emission line shape, which is limited by the resolution of the spectrometer used for detection. LIF is routinely accomplished with either diode lasers or dye lasers. Non-resonant LIF employing a

tunable ring dye laser was used for the experiments in this work. A detailed review of LIF with tunable dye lasers can be found in Ref. [117]

The width of an absorption line depends on an assortment of broadening mechanisms including: the natural line width, Stark broadening, power broadening, Doppler broadening and Zeeman broadening. Only Doppler broadening and Zeeman broadening are significant for the argon ion lines used in HELIX-LEIA. The velocity distribution function (VDF) is measured with LIF because of Doppler broadening of the absorption line.

As the laser's frequency is scanned, the intensity of the fluorescent emission is measured as a function of the laser wavelength. Typically, only metastable excited states or ground states are populated enough to be used for the initial state. In a plasma, moving ions or neutrals absorb the Doppler shifted laser light, pumping an electron to a higher energy, excited state which then spontaneously decays. LIF measurements are capable of being highly localized by engineering the optical paths for laser injection and collection to overlap only in one location. Spatial resolutions on the order of 1 mm^{-3} are easily accomplished. Analysis of the VDF measurement yields three quantities; temperature, velocity and density of the absorbing state (if the LIF system is fully calibrated). Assuming a Maxwellian distribution, the Doppler broadening of an ion absorption line is

$$I(v) = I_0 \exp\left(\frac{-m(\nu - \nu_0^*)^2 c^2}{2k_B T \nu_0^2}\right) \quad (3.25)$$

where $I(v)$ is the absorbed photon flux as a function of frequency v , I_0 is the maximum photon flux absorption and $\nu_0^* = \nu_0 + \nu\nu_0/c$ is the proper frequency of the transition when viewed from the laboratory frame.

Most of the measurements presented in this work are conducted in LEIA with a low magnetic field ($<150 \text{ G}$), rendering all forms of broadening negligible except Doppler broadening.¹¹⁸

However, a few measurements were conducted in HELIX or close to HELIX with a strong magnetic field, which requires a correction due to the Zeeman effect. The Zeeman effect is the splitting of the spectral lines of an atom in the presence of a strong magnetic field. Zeeman splitting yields linearly polarized π lines ($\Delta m = 0$) and circularly polarized lines ($\Delta m = \pm 1$) for absorption between the initial and the upper states. The details of typical Zeeman splitting for the primary 611.6616 nm absorption line of ionized argon is shown in Fig. 3.7. The π lines are symmetrically distributed around the zero magnetic field transition. The σ lines consist of two clusters denoted by σ^+ and σ^- . The amplitude envelope of the σ^+ or σ^- clusters is asymmetric, but each cluster is distributed symmetrically around the central line which denotes the zero magnetic field transition. The frequency shift of each σ cluster from the central line depends linearly on the magnetic field strength. The strength of the magnetic field at the measurement location is determined from the measured shift of the σ clusters.

Each Zeeman component is also Doppler broadened by thermal effects. To obtain an accurate measurement of the temperature from the measured LIF signal, the signal has to be de-convolved into the individual Zeeman components.¹¹⁸ For typical LIF measurements in HELIX-LEIA, the Zeeman splitting of each of the two sets of σ lines and the set of π lines is much less than the Doppler broadening for each of the individual σ and π lines.³⁹ Therefore, each cluster of σ lines and the cluster of π lines can be treated as a single Doppler broadened line shifted from the proper frame wavelength by the statistically weighted average Zeeman shift of the individual lines in the cluster.

A polarizer is typically used to excite only one set of lines. For perpendicular injection, the laser is polarized parallel to the magnetic field, allowing only the π lines to be excited. For parallel injection, the laser is circularly polarized with a linear polarizer followed by a quarter wave plate so that only one cluster of σ lines is excited.

FIGURE 3.7: Schematic of the σ , and π transitions for the 611.6616 nm argon ion absorption line. The height of each line corresponds to the statistical weighting of each transition as a function of wavelength. Fig. from Ref. [39].

3.3.1 LIF of Argon Ions

A Matisse-DR tunable ring dye laser, with a line width of less than 20 MHz RMS (root mean square), was used for the LIF measurements presented in this work. A Millennia Pro 10s diode laser pumps the Matisse-DR ring dye laser with Rhodamine 6G dye that fluoresces from 550 nm to 630 nm. The Matisse-DR is tuned to 611.6616 nm (vacuum wavelength) to pump the Ar II $3d^2G_{7/2}$ metastable state to the $4p^2F_{7/2}$ state, which then decays to the $4s^2D_{5/2}$ state by emitting 461.086 nm photons. A diagram of the LIF scheme is shown in Fig. 3.8.

The Matisse-DR dye laser beam is split and 99% is modulated with a mechanical chopper at 5 kHz. After the chopper, the main laser beam is coupled into a multimode, non-polarization-preserving, optical fiber for transport from the laser laboratory to HELIX-LEIA in the adjacent room. As the laser's wavelength is scanned, the intensity of the fluorescence is

FIGURE 3.8: LIF scheme for Ar II using the ring dye laser. Figure from Ref. [43].

collected and transported by an optical fiber to a filtered (1 nm bandwidth around the fluorescence wavelength) narrowband, high-gain, Hamamatsu photomultiplier tube (PMT). The PMT signal consists of electronic noise, background spectral radiation and electron-impact-induced fluorescence radiation. A Stanford Research SR830 lock-in amplifier is used to isolate the modulated signal. The lock-in amplifier is necessary because electron-impact-induced emission is orders of magnitude larger than the fluorescence signal.

The other 1% of the laser beam is split into two beam paths for diagnostic purposes. One beam is sent to a Bristol Instruments 621-VIS wavelength meter for laser wavelength measurements. The second beam is sent through an iodine cell for a consistent zero-velocity calibration. The iodine spectral emission is detected with a photodiode and sent via a BNC cable to the data acquisition system. The *Salami*¹¹⁹ reference iodine spectrum was compared to experimentally obtained iodine spectra in the ranges of interest of each LIF scheme in order to identify appropriate lines to be used as a zero velocity reference. Additional information on the iodine calibration is found in Ref. [120]. The iodine line closest to the zero velocity reference with a sufficiently strong intensity is the 16348.94 cm^{-1} line, which corresponds to a frequency

FIGURE 3.9: Iodine spectra for approximately ± 15 GHz relative to 611.6616 nm. This spectrum was obtained from two partially overlapping 20 GHz wide laser scans (solid line). The iodine spectrum in the spectral ranges of interest for Ar II, as given in Ref. [119], is shown as the dotted line. Figure adapted from Ref. [121].

shift of 1.08 GHz from the Ar II absorption line at 16348.91 cm^{-1} (661.6616 nm). Figure 3.9 shows the iodine spectra surrounding the Ar II line.

Figure 3.10 shows an IVDF taken at position E in Fig. 2.2, 126 cm downstream of the antenna. The bulk velocity V of the Ar II IVDF is determined from

$$V = \lambda_0 \Delta \nu_{total} \quad (3.26)$$

where λ_0 is the rest frame wavelength in nm and $\Delta \nu_{total}$ is the total frequency shift in GHz. The total frequency shift is equal to the difference between the peak of the LIF signal and the iodine reference line (1.46 GHz for this example) plus the shift in the peak of the iodine reference line

FIGURE 3.10: Typical LIF measurement of an argon IVDF, measured in LEIA, 126 cm downstream of the antenna (position E in Fig. 2.2). The black line is the raw LIF signal, while the red line is a single Gaussian fit to the data and the pink line is the iodine reference spectrum. Figure from Ref. [121].

from the non-shifted absorption line at 611.6616 nm (1.08 GHz) and the Zeeman shift due to the magnetic field at the measurement location (1.03 GHz for this example). Therefore, the total frequency shift is 1.51 GHz, yielding a bulk velocity of ~ 925 m/s. The raw data is fit with a single Gaussian function (red line) and the ion temperature of 0.16 eV is calculated from the full width at half maximum (FWHM).¹¹⁸

3.4 Superprobe

A newly manufactured, in-situ, scanning probe (hereafter referred to as the superprobe) allows LIF measurements parallel and perpendicular to the background magnetic field in a 2-D plane inside LEIA. The superprobe is essentially a system of gears and shafts that rotate to move an aluminum structure, housing LIF optics, to a particular point in space.

3.4.1 Exterior Apparatus

Shown in Fig. 3.11 is a photograph of the stand used to support the exterior portion of the superprobe. A 1-meter long VELMEXTMstepping motor assembly provides the radial movement of the entire superprobe apparatus inside of LEIA. The VELMEXTMassembly has a measuring tape, 80 cm long, attached to its side to determine the radial location r of the LIF measurement. Two black lines are located at 59.5 cm and 65.5 cm on the tape which correspond to a LIF measurement location of $r=0$ cm and the width of the vertical T-mount shown in Fig. 3.12. A 1.5 m shaft is coupled to the VELMEXTMassembly on the vertical T-mount and passes through a double O-ring seal into the LEIA chamber. The vertical mount also contains a second VELMEXTMstepping motor, a rotational motor that rotates the 1.5 m shaft, providing the axial motion in the chamber. To keep track of the axial position of the probe assembly, a mechanical counter is attached to the end of the shaft and zeroed (00000) at $z=164$ cm. Fifty turns on the counter equal 1 cm of motion with counterclockwise rotation moving the measurement location towards HELIX. The rotating motor is limited to 2 cm of axial motion at a time because it has a tendency to overheat and get stuck. The axial measurement region extends from $z=163$ cm (00050) to $z=200$ cm (98200). Because the LEIA chamber is curved near the connection to HELIX, the radial scan range depends on axial location. The combined axial and radial limits are incorporated in the scanning program in the data acquisition system.

FIGURE 3.11: Photograph of the stand used for the exterior portion of the superprobe.

A 0.305 m custom made aluminum flange is attached to the LEIA chamber and houses the double O-ring seal and a Y-adaptor, a three-flange vacuum feedthrough. The entire assembly is shown in Fig. 3.12. The top flange contains two SMA optical fiber connections for laser injection. Two marks, one for parallel injection ($\parallel$) and one for perpendicular injection ($\perp$) are by the respective connections. The middle flange contains an extension for the collection fiber which leads to the PMT. The PMT is surrounded by magnetic shielding because the PMT is located inside the LEIA magnets. PMTs must be shielded from strong ambient magnetic fields because magnetic fields suppress the electron avalanche process inside the PMT. The bottom flange is a five feedthrough BNC flange for the triple probe and Langmuir probe. The central BNC connects to the Langmuir probe while the three connectors for the triple probe are marked with the numerals 1, 2, and 3 (see Fig. 3.5).

FIGURE 3.12: The counter used for verifying the axial location of the superprobe. A counter value of 00000 corresponds to $z=164$ cm.

3.4.2 Interior Apparatus

A second 1 meter linear motion stage inside LEIA is aligned parallel with the chamber axis. The slide assembly provides the axial movement and is driven by the 1.5 m shaft through a series of gears. The entire internal assembly is supported by two horizontal bars of extruded aluminum that extend the width of LEIA and two vertical bars that rest on the bottom of the chamber. The two horizontal bars terminate in the boxport of the LEIA chamber and are locked into place by two vertical braces of 309 high temperature stainless steel. Fig. 3.14 shows the interior superprobe assembly with the two 309 stainless steel bars just outside the field of view to the left of the photo. Fig. 3.15 is a schematic of the superprobe showing the interior and exterior slide assemblies.

A 0.61 m long, 0.0254 m square cross section aluminum arm connects a lightweight,

FIGURE 3.13: Photograph of the aluminum flange which houses the Y-adaptor. The PMT is wrapped in magnetic shielding material.

aluminum, CNC machined optics head to the linear motion stage. The optics head, shown in Fig. 3.16, contains all the necessary hardware for LIF; collimators, focusing lenses, mirrors and fibers. The three optical fibers are routed from the Y-adaptor to the optics head through flexible stainless steel conduit. A planar Langmuir probe and a triple probe are attached to the optics head but offset in radius to eliminate any interference of the probes on the LIF measurements. The triple probe is offset 2.5 cm in r from the LIF location and the Langmuir probe is offset an additional 2 cm (4.5 cm in total). The Langmuir probe and triple probe are offset 1 cm towards HELIX relative to the LIF measurement location (see Fig. 3.17).

It is important to note that the superprobe measurement location is subject to hysteresis or backlash. The 1.5 m shaft couples to the interior linear motion assembly by way of a flexible shaft coupler. The purpose of the coupler is to absorb misalignment of the shaft. As

FIGURE 3.14: Photograph of the interior superprobe assembly.

the superprobe is scanned in the radial direction, the coupler compresses or expands, thereby leading to backlash. To avoid uncertainty in the measurement location, every radial scan should start at the same location and move in the same direction. Confirmation of the radial position is accomplished by peering down the axis of HELIX-LEIA, and identifying when the triple probe crosses the machine axis (accounting for the offsets described earlier).

FIGURE 3.15: Schematic of the superprobe showing the slide assemblies.

FIGURE 3.16: Photograph of the optics head which contains the LIF hardware. The Langmuir probe and triple probe are offset from the LIF measurement location.

FIGURE 3.17: A cartoon showing the dimensions between the various diagnostics on the superprobe. The three positions marked L, T and LIF represent the locations of the Langmuir probe, triple probe and the LIF location respectively. The colored lines are the two injection laser paths (red) and the collection path (blue).

Chapter 4

Spatial Structure of Ion Beams

Classic DL theory requires four populations to establish a steady-state DL: trapped ions downstream of the DL, accelerated ions flowing downstream from upstream, accelerated electrons flowing upstream, and trapped electrons upstream of the DL.⁵⁰ Shown in Fig. 4.1 is a schematic of a typical DL with the various particle populations and the measurement region for these experiments. Using the scanning probe, the accelerated and background ion populations are measured throughout the downstream plasma plume. Two cases of the magnetic field mirror ratio B_{HELIX}/B_{LEIA} were studied, a high LEIA field case with a mirror ratio of 8 and a low LEIA field case with a mirror ratio of 28.

4.1 High Field Case, Mirror Ratio=8

The measurements in this section were obtained with an upstream magnetic field in HELIX of 860 G and an expansion magnetic field in LEIA of 108 G.

FIGURE 4.1: A schematic showing the four populations in the classical picture of a DL. The crosshatched region is the region in which the IVDF and probe measurements were obtained.

4.1.1 Parallel IVDF Measurements

Fig. 4.2 shows a typical parallel IVDF obtained in the center of the plasma at $z = 164$ cm (the HELIX-LEIA junction occurs at $z = 159$ cm). Evident in the IVDF is a large amplitude ion beam population at 8.0 km/s and a lower density background ion population centered around 0 m/s. A negative velocity means that the ions are traveling downstream from the source into the expansion chamber (towards the source of laser light). Further downstream, at an axial location of $z = 175$ cm, the background ion population with an upstream directed bulk velocity of approximately +500 m/s, appears and is larger in magnitude than the ion beam population (see Fig. 4.3). At this location, the temperature of the beam and bulk are virtually identical, 0.23 ± 0.01 eV and 0.24 ± 0.01 eV, respectively. The increase in background ion density relative to the ion beam and the average upstream flow of the background ions is consistent with the assumption that the measurement region is downstream of a DL.

FIGURE 4.2: The IVDF at $r = 0$ cm and $z = 164$ cm showing an ion beam at 8.0 km/s with little background plasma. The temperature of the ion beam population is 0.23 eV.

Measurements were obtained every 2 cm from $r = -12$ cm to $r = 8$ cm at axial locations $z = 164$ cm, $z = 170$ cm, $z = 175$ cm and $z = 180$ cm, as shown in Figures 4.4-4.7. At radial locations greater than $r = 8$ cm the body of the probe blocks the plasma from flowing into the LEIA chamber, therefore the measurement region is asymmetric around the plasma axis. The IVDFs are stacked in an array and plotted as a contour map with arbitrary units. Each individual IVDF measurement is scaled to account for the different lock-in gain settings used for each measurement location. In the LIF signal at $z = 164$ cm, the dominant ion beam population seen in Fig. 4.2 is noticeably absent. The ion beam population and the background ions (those at zero velocity) do not appear in the plots in the center of the plasma in Figures 4.4-4.7 because all the IVDF measurements have been plotted with a common color bar and there is a dramatic decrease in overall ion density towards the center of the discharge. As will be

FIGURE 4.3: The parallel IVDF at $r = 0$ cm and $z = 175$ cm showing an ion beam of 7.6 km/s and a population of background plasma moving upstream towards the source with $v = +500$ m/s. A third population with $v = -1.7$ km/s is flowing away from the DL.

shown by the Langmuir probe measurements later, the plasma density drops at least a factor of five from the edge to the center of the discharge. Therefore, although the ion beam population dominates the IVDF at $r = 0$ cm and $z = 164$ cm, the overall plasma density is much lower than the background ion density at the plasma edge and therefore the beam does not appear in the center of the plasma in Figures 4.4-4.7. Other helicon source groups have also reported a hollow plasma density profile in the plume of an expanding plasma.³⁶ The black dashed lines in the panels of Figures 4.4-4.7 lie along a common tube of constant magnetic flux that expands with downstream distance. The IVDF data in Figures 4.4-4.7 show that the hollow portion of the background plasma density profile expands with the expanding magnetic field.

FIGURE 4.4: The parallel IVDF as a function of radial location at $z = 164$ cm scaled to account for detector sensitivity.

FIGURE 4.5: The parallel IVDF as a function of radial location at $z = 170$ cm scaled to account for detector sensitivity.

FIGURE 4.6: The parallel IVDF as a function of radial location at $z = 175$ cm scaled to account for detector sensitivity.

FIGURE 4.7: The parallel IVDF as a function of radial location at $z = 180$ cm scaled to account for detector sensitivity.

To better visualize the spatial structure of the ion beam, each parallel IVDF is normalized to its peak value and plotted in an array of normalized IVDFs in Figures 4.8-4.11. In the normalized plots there is a well-defined region (from $r = \pm 5$ cm) that is dominated by the ion beam. Outside of that central region, the IVDF is dominated by the background ion population. For downstream distances beyond $z = 164$ cm, the background ion population appears in the central core region of the plasma. Also shown in Figures 4.8-4.11 are the same two dashed lines from Figures 4.4-4.7 which mark the edges of a cylinder of constant magnetic flux that maps to a flux tube of radius 2 cm in the helicon source. Both the amplitude and the speed, $\approx 8,000$ m/s, of the metastable ion beam population decrease slightly with increasing downstream distance. There is little to no change in the parallel ion temperatures of the beam and bulk populations with downstream distance or radial location.

FIGURE 4.8: The same parallel IVDF as function of radial location at $z = 164$ cm as shown in Fig. 4.4, but with each IVDF normalized to the peak value in the IVDF.

FIGURE 4.9: The normalized parallel IVDF as a function of radial location at $z = 170$ cm.

FIGURE 4.10: The normalized parallel IVDF as a function of radial location at $z = 175$ cm.

FIGURE 4.11: The normalized parallel IVDF as a function of radial location at $z = 180$ cm.

A faint, low velocity, downstream-flowing, third ion population also appears in the center of the plasma with increasing distance from the DL. Complex IVDFs similar to this have been observed in LEIA before.⁸ Shown in Fig. 4.3 are Maxwellian fits to the ion beam and bulk populations, along with a third Maxwellian population fitted to the residue of the IVDF after the beam and bulk populations are subtracted. The third population in Fig. 4.3 has a net flow of 1.7 km/s directed downstream from the DL. The effective temperature of this third population is 0.42 eV, significantly hotter than the beam and bulk ion temperatures. These ions are most likely beam ions that have slowed down through collisions with background ions or energetic neutrals created through charge-exchange between the beam ions and background neutrals.

FIGURE 4.12: The same corrected, but unnormalized, parallel IVDFs shown in Fig. 4.4 as a function of radial location at $z = 164$ cm, but only for velocities above 4,000 m/s.

FIGURE 4.13: The same corrected, but unnormalized, parallel IVDFs shown in Fig. 4.5 as a function of radial location at $z = 170$ cm, but only for velocities above 4,000 m/s.

FIGURE 4.14: The same corrected, but unnormalized, parallel IVDFs shown in Fig. 4.6 as a function of radial location at $z = 175$ cm, but only for velocities above 4,000 m/s.

FIGURE 4.15: The same corrected, but unnormalized, parallel IVDFs shown in Fig. 4.7 as a function of radial location at $z = 180$ cm, but only for velocities above 4,000 m/s.

The spatial structure of the ion beam portion of the parallel IVDF develops over the region of study as the ion trajectories respond to magnetic forces and electric fields. To gain a better understanding of the evolving spatial structure of ion beam, just the portion of the unnormalized parallel IVDFs above 4,000 m/s, well above the bulk thermal velocity, are shown in Figures 4.12-4.15. Because the IVDFs in Figures 4.12-4.15 are not normalized to the background plasma density, they highlight variations in the total amplitude of the beam population, i.e., the variations in the plots include the effects of the strongly radially varying plasma density. The radial profile at $z = 164$ cm shows a hollow ion beam amplitude profile with peaks on either side of the central axis. The radial profiles from further downstream show the same hollow structure expanding radially outward. In the time it takes for the ion beams to traverse 16 cm along the axis of the plasma, the peaks in the radial profiles shift radially outward by approximately 3 cm. Such radial expansion of the hollow ion beam profile exceeds the radial expansion of the magnetic field. Therefore, these parallel IVDF measurements suggest the action of significant additional radial forces pushing the beam ions radially outward as they travel downstream. The axial evolution of the radial ion beam profiles in the IVDFs shown in Figures 4.12-4.15 provides additional evidence of the beam ion slowing mechanism noted in the previous discussion of the third ion population. With increasing downstream distance, the ion beam distributions elongate and flatten in velocity space, stretching out from a beam peak at 8 km/s to include slower beam ions extending to 4 km/s.

To investigate the effects of collisions on the ion beam velocity and the total parallel IVDF, high spatial resolution measurements were performed at $r = 0$ cm in 1 cm steps from $z = 170$ cm to $z = 191$ cm (see Fig. 4.16). With increasing downstream distance, the relative ion beam to background intensity decreases and there is a slight decrease in the ion beam velocity. Note that Fig. 4.16 begins at $z = 170$ cm, well after the background ion population appears in the parallel IVDF. Previous studies of ion beam amplitude decay in expanding plasmas have

FIGURE 4.16: The normalized parallel IVDF at $r = 0$ cm as a function of downstream distance.

attributed the decay to quenching of the metastable state probed by LIF due to collisions of the metastable ions with electrons.⁴ In other words, the decrease in LIF signal results from the particular requirements of the LIF measurement process and does not necessarily indicate actual decay of the ion beam amplitude. In fact, RFEA measurements in this plasma plume⁸ and in other experiments⁴¹ have found that the ion beam persists downstream with little reduction in beam density. An exponential fit to the decaying LIF amplitude yields a $1/e$ folding distance of 11.4 cm (see Fig. 4.17). Assuming the $1/e$ folding distance is the effective mean-free-path for the metastable ions, these measurements yield a quenching cross-section of $2.7 \times 10^{-18} \text{ m}^2$, consistent with results from previous measurements of the effective metastable quenching cross section.^{2,4,41,122}

Over an 18 cm distance, there is a slight decrease in the ion beam velocity from 8,320 to 8,040 m/s shown in Fig. 4.16. Typical error in these measurements is 50 m/s, smaller than

FIGURE 4.17: LIF signal amplitude for the peak of the ion beam in the IVDF and exponential fit versus axial position from data of Fig. 4.16.

the symbol marker. Although the measured slowing is slight, it is statistically significant. A collisional process would be expected to produce an exponential velocity decrease given that ion momentum loss due to collisions with background neutrals is described by

$$\frac{d\mathbf{v}}{dt} = -\nu\mathbf{v}, \quad (4.1)$$

where ν is the relevant collision frequency. Therefore, some non-collisional process must contribute to the slowing of the ion beam. In other words, the nature of the slowing of the beam ions suggests the existence of an upstream directed electric field, completely inconsistent with what would be expected for a DL.

Although the parallel IVDF measurements presented so far were obtained for a relatively strong downstream magnetic field of 108 G, most DL studies in LEIA and elsewhere have

FIGURE 4.18: Ion beam velocity as a function of axial distance for the data of Fig. 4.16.

FIGURE 4.19: The parallel IVDF as a function of the expansion magnetic field at $r = 0$ cm and $z = 171$ cm.

employed weak downstream magnetic fields or none at all.^{6,43,53} By simple magnetic moment conservation ($\mu = kT_{i\perp}/2B$ for thermal ions gyrating around a magnetic field), a weaker downstream magnetic field should yield an increase in the ion beam velocity in addition to any DL acceleration effects. Faster ion beams are of particular importance for spontaneous ion beam generation applications such as plasma thrusters.

Shown in Fig. 4.19 is the effect on the on-axis ion beam velocity due to changing the downstream magnetic field strength from 8 to 108 G. The ion beam velocity is measured well downstream of the DL at $z = 171$ cm. All other plasma source parameters were held fixed. The ion beam velocity drops from 12,100 m/s for a downstream magnetic field of 8 G to roughly 8,800 m/s for a field of 54 G. As the downstream field increases from 50 G to 110 G, there is a modest, linear decrease in the beam velocity.

Assuming that the ions flow downstream slow enough that μ is an adiabatic invariant in these experiments, energy conservation and measurements of the upstream and downstream perpendicular ion temperatures are enough information to calculate the maximum possible increase in parallel ion flow speed due to μ conservation. For these experiments, the upstream perpendicular ion temperature was 0.55 eV. Conversion of all the perpendicular thermal energy into parallel flow kinetic energy would only accelerate stationary argon ions up to a parallel flow speed of 1,150 m/s. However, given that for a downstream magnetic field of $B = 108$ G the perpendicular ion temperature at $z = 171$ cm is measured to be 0.45 eV (nearly unchanged from the upstream value), it is clear that energy conservation and magnetic moment conservation (if magnetic moment is even conserved in this system) are insufficient to explain the observed ion acceleration, i.e., the existence of additional ion acceleration from upstream to downstream is implied by the parallel IVDF measurements. For a downstream magnetic field of $B = 31$ G, the downstream perpendicular ion temperature was 0.51 eV. Given the small change in perpendicular ion temperature, μ conservation yields at most a 56 m/s increase parallel ion

velocity. Therefore, the sharp decrease in ion beam velocity as the downstream magnetic field increases from 20 to 50 G must result from a substantial change in the potential difference across whatever electric field structure is responsible for the ion acceleration. For downstream magnetic fields of 50 to 110 G, it appears that the potential difference across the ion accelerating structure remains relatively constant. Note that at speeds of 8,000-10,000 m/s, beam ions only complete 1/4 of a gyro-orbit while traveling 16 cm in the axial direction. Such ion motion is simply too fast for adiabatic constraints such as μ conservation to hold.

It is important to note that the parallel IVDFs shown in Fig. 4.19 are not self-normalized. The data shown are raw IVDF measurements and therefore Fig. 4.19 indicates that there is a critical downstream magnetic field for which the background and beam ion densities are largest, approximately 31 G, and above which the ion beam velocity and ion beam density start to decrease. For downstream magnetic fields less than 20 G, the ion beam velocity also decreases — additional confirmation that the observed parallel ion beam velocities can not be a result of μ conservation. The implications of these measurements are profound with regard to potential use of these systems as plasma thrusters. Clearly, some downstream field enhances both the specific impulse and the thrust of such a thruster. For the weakest downstream magnetic fields, the collisionality of the plasma also appears to play less of a role. There is a clear reduction in LIF signal for ion velocities between the beam velocity and the background at the weakest downstream fields (visible as a purplish region around 5,000 m/s and 10 G in Fig. 4.19). As there should be no effect of the changing magnetic field on neutrals, it appears that poor ion confinement (and therefore fewer ion-ion collisions) at the smallest downstream magnetic field strengths may reduce the number of ions that are scattering into the slowest velocity ranges in the plasma.

4.1.2 Perpendicular IVDF Measurements

The perpendicular IVDFs were measured at the same locations as the parallel IVDFs and under the same conditions. While simultaneous parallel and perpendicular IVDF measurements using multiplexing with the scanning LIF probe are possible,¹²³ the parallel and perpendicular LIF measurements were obtained sequentially in these experiments. Switching from parallel to perpendicular measurements is accomplished by moving the injection fiber to a different fitting on the external interface of the probe.

Figures 4.20-4.23 shows the self-normalized, perpendicular IVDFs at axial locations of $z = 164$ cm, $z = 170$ cm, $z = 175$ cm and $z = 180$ cm, respectively. There are a number of significant features in these perpendicular IVDF measurements. In the center of the plasma, in the same central region where the parallel IVDFs show clear evidence of an ion beam, there is an ion population with a finite radial flow that switches sign across the plasma axis. It is important to note that in the expansion region the local magnetic field is not purely axial. At the most upstream locations measured, the magnetic field has a significant radial component. These perpendicular IVDF measurements are in the laboratory frame. Thus, any ion beam flowing along the local magnetic field will have velocity components in the radial and axial directions. This projection effect is evident in Fig. 4.20. Outside of the central core of the plasma, the perpendicular IVDFs show quite complex behavior. Outside of $r = \pm 5$ cm, the perpendicular IVDF cannot be described with a single Maxwellian velocity distribution.

Overplotted on the panels in Figures 4.20-4.23 is a dashed black line marking what the projected perpendicular velocity of the ion beam should be given the measured parallel velocity at that location and the expected magnetic field angle. Since the magnetic field angle relative to the axial direction is only a few degrees (depending on radial location), the radial velocity should be no larger than a few hundred meters per second. The perpendicular IVDF in the core

FIGURE 4.20: The normalized perpendicular IVDF at $z = 164$ cm as a function of radial location.

FIGURE 4.21: The normalized perpendicular IVDF at $z = 170$ cm as a function of radial location.

FIGURE 4.22: The normalized perpendicular IVDF at $z = 175$ cm as a function of radial location.

FIGURE 4.23: The normalized perpendicular IVDF at $z = 180$ cm as a function of radial location.

of the plasma at $z = 180$ cm (Fig. 4.23) is generally consistent with the predicted structure. There is a slight positive flow offset that is likely instrumental in nature.

However, moving upstream the complexity of the perpendicular IVDF increases dramatically. At $z = 170$ cm there is a clear increase in the difference between the measured radial flow and what is predicted based on the projection of the measured parallel flow. Towards the plasma edge, the perpendicular IVDF becomes broader (hotter) and shows evidence of multiple ion populations. Shown in Fig. 4.24 is a typical perpendicular IVDF obtained at $z = 170$ cm and $r = -8$ cm. The perpendicular IVDF is well fit with two Maxwellian distributions with temperatures of 0.35 eV and 0.30 eV and relative normalized densities of 1.0 and 0.30 for the bulk and flowing populations, respectively.

By $z = 164$ cm, the discrepancy between the radial flow in the core plasma and the projected parallel flow has further increased. In fact, the perpendicular IVDFs yield a radial speed of 2,000-3,000 m/s, suggesting a significant additional ion accelerating mechanism in the perpendicular direction - reminiscent of the curved DL structures mentioned earlier and also consistent with the radial expansion of the beam ions identified in Figures 4.12-4.15. At the edge of the plasma at $z = 164$ cm, the velocity spread has increased so much that a single (and naive) Maxwellian fit to the perpendicular IVDF yields an ion temperature of 1-2 eV. Using a single temperature to describe the perpendicular IVDF is clearly inappropriate, but the substantial spread in velocities suggests the presence of a significant source of particle energization in the perpendicular direction to the local magnetic field. It is important to remember that these IVDF measurements are obtained over long time intervals and therefore the particle energization mechanism is a steady-state phenomenon that is also clearly multi-dimensional.

FIGURE 4.24: The perpendicular IVDF at $r = -8$ cm and $z = 170$ cm.

4.1.3 Probe Measurements

To measure the electric field structure throughout the plasma plume for the same plasma conditions of the IVDF measurements (downstream magnetic field of 108 G), the triple probe was scanned through the same locations while measuring the two-dimensional (radial and axial), steady-state, electric field. The vector electric field in the plasma plume is shown in Fig. 4.25. Also shown in Fig. 4.25 as solid black lines is the expansion of magnetic flux tubes in the downstream region. The two outermost flux tubes map back to a radial location approximately 1 cm from the inner surface of the Pyrex vacuum chamber under the rf antenna in the source.

Past studies of expanding helicon plasmas have employed emissive probes⁵⁶ or interpretations of RFEA measurements to determine the local plasma potential throughout the DL region.¹ Most studies have only measured the plasma potential along the system axis. Those

FIGURE 4.25: The average DC electric field for a downstream magnetic field of 108 G. Fifty measurements were taken at each location. Because the triple probe is offset in the negative radial direction, the measurement range is asymmetric about the plasma axis.

that have measured the radial and axial structure of the plasma potential have reported curved equipotentials.^{10,36,56} Very recent measurements have reported regions of electric fields pointed upstream¹²⁴ and significant gradients in plasma potential towards the plasma edge.³⁶ Here the steady-state electric field is measured directly. From $E = -\nabla\Phi$, where Φ is the plasma potential, the plasma potential structure can be obtained from these measurements.

In the center of the plasma the magnitude of the axial electric field is small, between 1 and 10 V/m. For the average axial electric field at $r = 0$ cm, which is 4.78 V/m and points upstream, the ion beam will slow from 8,320 m/s to 8,070 m/s over 0.18 m. According to the measurements shown in Fig. 4.16 and Fig. 4.18, the velocity drops to 8,040 m/s, over this range. Therefore, the electric field is sufficient to explain all (to within measurement error) of the observed ion beam slowing. Ion-neutral collisions¹²⁵ appear to have little effect on the ion

beam velocity.

Further downstream, the axial electric field on axis decreases, consistent with other experiments.⁵⁸ Moving outwards, the axial electric field and the radial electric field increase. Both field components then abruptly switch sign across the outermost flux tube shown in Fig. 4.25. The switch in sign of the field components maps along the field line over the entire downstream region sampled. This electric field structure is clearly field aligned and is a region of ion density depletion, i.e., an ion hole as $\nabla \cdot E < 0$.

The large scale radial flows and broad perpendicular IVDFs, particularly at $z = 164$ cm, are entirely consistent with these measured electric fields. The average radial electric fields from $r = 0$ cm to just before the ion hole are 91.2 V/m, 25.5 V/m, 14.0 V/m and 29.3 V/m at $z = 164$ cm, $z = 170$ cm, $z = 175$ cm, and $z = 180$ cm, respectively. These field strengths are more than sufficient to accelerate the ions to the perpendicular velocities observed in Figures 4.20-4.23.

The last few centimeters of the HELIX chamber are unable to be accessed to perform a measurement of the electric field through the ion acceleration region. However, plasma potential measurements in the helicon source upstream are possible. Shown in Fig. 4.26 are measurements of the electron energy probability function (EPPF) and the plasma potential at $z = 112$ cm. While the plasma potential measurements shown in Fig. 4.26 extend beyond 5 cm, the actual plasma source tube is only 5 cm in radius. The measurements were performed in the larger diameter stainless steel chamber downstream of the plasma source tube. Along the plasma axis, the upstream plasma potential is 35 V. The plasma potential measurements indicate that within the plasma source there is a radially outward electric field due to a potential drop of 10-15 V. Measurements of the downstream plasma potential on axis yield an upstream-downstream total plasma potential difference of 35 V. This number falls within the range of plasma potential

FIGURE 4.26: (a) Electron energy probability function and (b) The plasma potential as a function of radial location at $z = 112$ cm, inside the plasma source.

differences reported in other experiments.^{36,124} Strong acceleration of ions outwards into the walls of the glass chamber is consistent with our observations of significant sputtering of the glass walls when the plasma source is operated at the low pressures required to create ion beams. The sputtering is severe enough to create small holes in the glass tube (2-3 mm in diameter) or to completely etch through the glass walls. Similarly strong radial electric fields have been reported in other expanding plasmas at the junction of the plasma source and the expansion chamber.⁵⁶

The upstream EEPF measurements transition from a single Maxwellian energy distribution in the plasma core to a plasma with a significant energetic, “fast”, electron component by $r = 3$ cm. A calculation of the electron skin depth $\delta = c/\omega_{pe}$ for these plasma conditions yields a value of ~ 2 cm, consistent with location of the potential dip at 3 cm in Fig. 4.26b. In other words, the expected radial location for peak rf absorption matches the upstream region with energetic electrons and a strong electric field that then maps along the expanding magnetic field downstream. The same explanation for an observed annulus of fast electrons in an expanding helicon plasma was independently proposed by *Takahashi et al.*¹²⁶

Significant electrostatic wave activity is also observed in the time-resolved electric field measurements. Shown in Fig. 4.27 is the low frequency power spectrum as a function of radial location at $z = 180$ cm. The same measurements are shown in Fig. 4.28 for $z = 164$ cm. Frequencies below 300 Hz and above 50 kHz are omitted because they show little spatial variation. There is clearly stronger wave activity in the center of the plasma. At the plasma edge, i.e., $r = -12$ cm, there is effectively no wave activity. The radial region with significant wave activity expands in radius further downstream of the DL, suggesting that waves are excited upstream and damp with distance downstream. Previous measurements in this system have identified a number of instabilities associated with ion beam formation.^{42,127} The 17 kHz wave, propagating primarily axially, was previously determined to be an ion acoustic wave.⁵² The width of the

peaks in the power spectra are consistent with an active damping mechanism in these plasmas. Detailed analysis of the wave activity is left for future work.

Shown in Fig. 4.29 is the electron density in the plasma plume. The field-aligned, high-density structure occurs in the same region of the plasma plume as the strong electric fields. Fig. 4.29 is consistent with Figures 4.4-4.7 showing a hollow density profile. A comparison with Fig. 1.6b shows a similar feature. There is little electrostatic wave activity outside this region of high density, suggesting that drift waves are not present. Drift waves typically arise near regions with a density gradient.

FIGURE 4.27: The low frequency power spectrum for a single tip of the triple probe at $z = 180$ cm as a function of frequency and radial position.

FIGURE 4.28: The low frequency power spectrum for a single tip of the triple probe at $z = 164$ cm as a function of frequency and radial position.

FIGURE 4.29: The electron density (in units of 10^{12} cm^{-3}) throughout the plasma plume as measured with an rf compensated planar Langmuir probe on the scanning probe.

4.2 Low Field Case, Mirror Ratio=28

To study the issue of ion beam detachment from the magnetic field, the spatial structure of the ion beam was studied for a lower expansion magnetic field (higher mirror ratio). Fig. 4.19 established that an expansion magnetic field of 31 G produces a sharp increase in the ion beam velocity and the LIF signal. Plasma parameters in this section are identical to those described in Section 4.2 except that the downstream magnetic field was reduced to 31 G.

4.2.1 Parallel IVDF Measurements

Shown in Figures 4.30-4.33 are the IVDFs scaled for detector sensitivity. Compared with Figures 4.4-4.7, the maximum background plasma density occurs further from the plasma axis ($r = 0$). At $z = 180$ cm the background plasma profile has expanded to $r = -10$ cm. Because of the asymmetric scanning range, the hollow plasma profile cannot be measured completely as it was for the higher magnetic field. It is likely that that hollow profile is symmetric and simply diverges more rapidly due to the lower magnetic field.

FIGURE 4.30: The parallel IVDF as a function of radial location at $z = 164$ cm scaled to account for detector sensitivity.

FIGURE 4.31: The parallel IVDF as a function of radial location at $z = 170$ cm scaled to account for detector sensitivity.

FIGURE 4.32: The parallel IVDF as a function of radial location at $z = 175$ cm scaled to account for detector sensitivity.

FIGURE 4.33: The parallel IVDFs as a function of radial location at $z = 180$ cm scaled to account for detector sensitivity.

The normalized IVDFs shown in Figures 4.34-4.37 show a well defined ion beam with $v \sim 11$ km/s in the center of the plasma. At $z = 164$ cm only, there is a clear decrease in the ion beam

speed in the central region of the plasma ($r = \pm 2$ cm). In this region the ion beam speed is ~ 9 km/s, close to the value in the high field case. The ion beam dominates over the background plasma for $z = 164$ cm and $z = 170$ cm whereas, in the high field case, this only occurred for $z = 164$ cm. At $z = 175$ cm, the third ion population reappears in the plots because the ion beam signal starts to decrease relative to the other populations. By $z = 180$ cm the ion beam signal has decreased so much that it is hardly distinguishable from the noise. The third ion population has also diminished, suggesting that rather than being beam ions that have slowed down, this population is more likely background ions that have gained energy.

FIGURE 4.34: The normalized parallel IVDF as a function of radial location at $z = 164$ cm.

FIGURE 4.35: The normalized parallel IVDF as a function of radial location at $z = 170$ cm. FIGURE 4.36: The normalized parallel IVDF as a function of radial location at $z = 175$ cm.

FIGURE 4.37: The normalized parallel IVDF as a function of radial location at $z = 180$ cm.

Shown in Figures 4.38-4.41 is the radial profile of the ion beam only (above 4000 m/s) at axial locations of $z = 164$ cm, $z = 170$ cm, $z = 175$ cm, and $z = 180$ cm respectively. Similar to Fig. 4.12 there exists a hollow ion beam amplitude profile with peaks on either side of the central axis. Figures 4.12-4.15 are mostly symmetric but Figures 4.39-4.41 are not. By comparing Fig. 4.13 to 4.39 it is clear that the beam ions move outwards in radius rather quickly compared to the high field case. It is important to note that the signal to noise ratio for any measurements beyond $z = 170$ is quite poor.

FIGURE 4.38: The normalized parallel IVDF as a function of radial location at $z = 164$ cm.

FIGURE 4.39: The normalized parallel IVDF as a function of radial location at $z = 170$ cm.

FIGURE 4.40: The normalized parallel IVDF as a function of radial location at $z = 175$ cm.

FIGURE 4.41: The normalized parallel IVDF as a function of radial location at $z = 180$ cm.

4.2.2 Perpendicular IVDF Measurements

Figures 4.42-4.45 show the self-normalized, perpendicular IVDFs at axial locations of $z = 164$ cm, $z = 170$ cm, $z = 175$ cm, and $z = 180$ cm, respectively. The perpendicular IVDFs show mostly the same trends for both mirror ratio cases. In the center of the plasma there is an ion population with a finite radial flow that switches sign across the plasma axis. The radial velocity of the ions approaches 5 km/s around $r = -10$ cm, slightly higher than the low mirror ratio case. Outside of $r = \pm 6$ cm, the perpendicular IVDF cannot be described with a single Maxwellian velocity distribution. The projected velocity, based on the field aligned parallel flow, shown by the overplotted black line, is much less than the measured velocities from LIF at axial locations $z = 164$ cm and $z = 170$ cm. There is no projected velocity line at $z = 180$ cm in Fig. 4.45 because there is no discernible ion beam in the parallel IVDFs. At the most upstream locations and towards the plasma edge, the perpendicular IVDFs are broadened which

suggests either localized heating or two Maxwellian populations. These results further reinforce the conclusion that the accelerating surface is curved. The perpendicular and parallel IVDFs suggest the strength of the accelerating surface is greater in the lower mirror ratio.

FIGURE 4.42: The normalized perpendicular IVDF at $z = 164$ cm as a function of radial location.

FIGURE 4.43: The normalized perpendicular IVDF at $z = 170$ cm as a function of radial location.

FIGURE 4.44: The normalized perpendicular IVDF at $z = 175$ cm as a function of radial location.

FIGURE 4.45: The normalized perpendicular IVDF at $z = 180$ cm as a function of radial location.

4.2.3 Probe Measurements

The vector electric field in the plasma plume is shown in Fig. 4.46. Also shown in Fig. 4.46 as dashed lines is the expansion of magnetic flux tubes in the downstream region. The ion hole perfectly follows the outermost flux tube that maps back to a radial location approximately 1 cm from the inner surface of the Pyrex vacuum chamber under the rf antenna in the source. In the center of the plasma the magnitude of the axial electric field is 23.92 V/m and points upstream. This value is larger than the axial electric field for the lower mirror ratio case. The average radial electric fields from $r = 0$ cm to just before the ion hole are 76.9 V/m, 80.38 V/m, 112.5 V/m and 59.0 V/m at $z = 164$ cm, $z = 170$ cm, $z = 175$ cm, and $z = 180$ cm, respectively. The perpendicular IVDFs show a bulk flow that is faster than the low mirror ratio case and the larger radial electric fields are consistent with these observations.

FIGURE 4.46: The average DC electric field for a downstream magnetic field of 31 G. Fifty measurements were taken at each location. Because the triple probe is offset in the negative radial direction, the measurement range is asymmetric about the plasma axis.

Shown in Fig. 4.47 is the low frequency power spectrum as a function of radial location at $z = 164$ cm. The same measurements are shown in Fig. 4.48 for $z = 175$ cm. Frequencies below 300 Hz and above 50 kHz are omitted because they show little spatial variation. There is clearly stronger wave activity in the center of the plasma just as there was for the low mirror ratio. The 17 kHz wave and its harmonic are still present regardless of the downstream magnetic field. After $z = 175$ cm the spectral measurements begin to get noisy, similar to what was seen in IVDF measurements in Sections 4.2.1-4.2.2.

FIGURE 4.47: The low frequency power spectrum for a single tip of the triple probe at $z = 164$ cm as a function of frequency and radial position.

FIGURE 4.48: The low frequency power spectrum for a single tip of the triple probe at $z = 175$ cm as a function of frequency and radial position.

Chapter 5

Multi-Species Plasmas

The theoretical description by *Baalrud* and *Hegna*⁶⁰ for ion infall speeds into sheaths for multi-ion plasmas is tested in the DL structures in HELIX-LEIA. The ion beam speed downstream of the DL is measured in the same location, $(r, z) = (0, 164)$ cm, for three different gas mixtures; argon and xenon, argon and helium, xenon and helium. A helium ion LIF scheme does not exist so only argon and xenon IVDFs are measured. The *Baalrud* and *Hegna*⁶⁰ theory predicts that adding a lighter ion mass will result in an increase in the ion beam speed when the ion densities are equal. Conversely, adding a heavier ion mass should slow the ion beam.

5.1 Single-Ion Species Baseline

Pressure has a well-known and significant effect on DL formation in single-ion species plasmas.^{51,128} The ion beam speed versus pressure for an argon plasma was measured to establish a baseline for the effects of pressure (see Fig. 5.1). The argon ion beam speed decreases as pressure is increased until the DL disappears (around 2.5 mTorr). Two explanations are possible. The first is that pressure increase changes the accelerating structure, in other words, the pressure

introduces a fundamental change in the nature of the ion beam accelerating structure. The second is that the mean free path for ion-neutral collisions (charge exchange and elastic collisions) decreases with increasing pressure, thereby scattering the ions before they gain significant energy. The result in Fig. 5.1 is likely a combination of both processes.¹²⁸

Shown in Fig. 5.2 is a similar baseline for xenon. The xenon ion beam disappears at a lower pressure threshold compared to argon. IVDFs with a distinct xenon ion beam and background ion population also disappear at a few mTorr. Assuming that the accelerating structure is the same, the difference in ion beam speed for xenon and argon at the same pressure will be related by

$$\frac{v_{B,1}}{v_{B,2}} = \sqrt{\frac{m_2}{m_1}} = \sqrt{\frac{131}{40}} \sim 1.81 \quad (5.1)$$

where species 1 is argon with a mass of 40 amu and species 2 is xenon with a mass of 131 amu. For a pressure of 0.1 mTorr, the measured argon and xenon ion beam speeds are, respectively, 4.5 km/s and 8.3 km/s yielding a measured ratio of $v_{Ar}/v_{Xe} = 1.84$. Therefore, the electric field structure appears to be remarkably similar for different ion species at the same pressure.

The ion beam velocities decrease exponentially with increasing pressure for argon and xenon as shown in Fig. 5.3. Collisional processes are expected to produce an exponential velocity decrease given that ion momentum loss due to collisions with background neutrals is described by

$$\frac{d\mathbf{v}}{dt} = -\nu\mathbf{v}, \quad (5.2)$$

where ν is the relevant collision frequency. Therefore, these baseline cases demonstrate that the ion beam is created upstream from the measurement location and slows down as a result of collisions during transit of the ions to the measurement location.

FIGURE 5.1: The normalized parallel argon IVDF at $(r, z) = (0, 164)$ cm as a function of pressure.

FIGURE 5.2: The normalized parallel xenon IVDF at $(r, z) = (0, 164)$ cm as a function of pressure.

FIGURE 5.3: The bulk velocity of argon (blue) and xenon (red) as a function of pressure from Figures 5.1-5.2 with exponential fits.

5.2 Argon LIF in Ar-He and Ar-Xe Plasmas

Fig. 5.4 shows the argon IVDF for a case in which the argon partial pressure and helium partial pressure were changed, but the total pressure was held constant like the conditions in the *Biloiu* and *Scime* study. As the argon partial pressure is lowered, the ion beam speed increases, consistent with Sec. 5.1. At the same time, the helium partial pressure increases. The *Baalrud* and *Hegna*⁶⁰ model predicts that if sheath physics creates the ion beam, the argon ion speed should increase up to the average ion sound speed given by Eq. 1.5 at whatever neutral pressures result in equal argon and helium plasma densities (which has to occur at some ratio of argon and helium pressures between essentially 100% argon and essentially 100% helium). The mechanisms responsible for the change in argon ion beam speed in Fig. 5.4 are unclear because both the argon and helium partial pressures are changing simultaneously.

As shown in Fig. 5.5, increasing the neutral argon pressure for a constant value of helium

partial pressure also results in slowing of the argon ion beam. The helium partial pressure in Fig. 5.5 is 0.71 mTorr while the argon partial pressures are significantly lower than the pressures in Fig. 5.1. The maximum argon pressure in Fig. 5.5 is 0.1 mTorr, nearly the same pressure as the lowest pressure in Fig. 5.1. Such low argon partial pressures are used for two reasons.

First, a high partial pressure of helium, along with high rf power, is required to produce helium ions. Raising the total pressure further by adding an equal pressure of argon nearly destroys the argon ion beam. The second reason is best illustrated with a third case in which the total argon pressure is held constant. The argon IVDF was measured for a constant argon pressure of 0.17 mTorr while adding various amounts of helium to the plasma. Fig. 5.6 shows that the argon ion beam slows linearly with increasing amounts of helium. This suggests that helium acts as a drag force on the argon ions thereby slowing them, a very different result than the multi-dipole experiments discussed earlier. The measured argon ion beam speed for a pressure of 0.1 mTorr (0.8 mTorr) in an argon only plasma was 8.3 km/s (5.3 km/s). With an argon pressure of 0.1 mTorr and a helium pressure of 0.71 mTorr the argon ion beam speed is 6.7 km/s (see top end of Fig. 5.5). Therefore, the argon ion beam speed is faster, with the addition of helium, than it would have been at the same total pressure under an argon only plasma. Assuming hard sphere collisions, the total neutral cross section is simply $\sigma = \pi(2r_1)^2$ for like atoms and $\sigma = \pi(r_1 + r_2)^2$ for different atoms where r_1 and r_2 are the atomic radii of species 1 and 2 respectively. The cross section for an argon only plasma is $4.44 \times 10^{-19} \text{ m}^2$. With the addition of helium the average cross section is $3.38 \times 10^{-19} \text{ m}^2$ (31%) smaller than the argon only plasma. The ion beam speed is 26.4% higher with the addition of helium, which is comparable to the reduction in the cross section. In other words, the increase in ion beam speed is explainable by just the reduced collisionality of the argon-helium mixture. There is no need to invoke the effect described by *Baalrud* and *Hegna*.⁶⁰

The argon ion beam speed was also studied as a function of xenon partial pressure. Xenon

has an ionization potential of 12.13 eV, lower than the ionization potential of argon (15.76 eV). Therefore, adding xenon to an argon plasma quickly changes the argon plasma into a xenon plasma. Fig. 5.7 shows the argon IVDF as a function of xenon partial pressure. The argon partial pressure was kept constant at 0.17 mTorr yielding an ion beam with $v = 8.2$ km/s. For a xenon partial pressure of 0.04 mTorr, the argon ion beam has slowed to 6.7 km/s. Increasing the xenon partial pressure further does not result in further slowing of the argon ion beam. In fact, the argon ion beam speed increases slightly with the addition of xenon from 6.9 km/s at a xenon pressure of 0.04 mTorr to 7.1 km/s at a xenon pressure of 0.09 mTorr. This result is completely at odds with the prediction from *Baalrud* and *Hegna*.⁶⁰ Eventually, with enough xenon, the argon LIF signal is lost and the plasma turns from a predominantly argon plasma (purple) to a xenon plasma (light blue). The multi-dipole sheath experiments discussed earlier measured the individual ion densities by launching ion acoustic waves and using the dispersion relation to solve for n_i . The present experiment does not use this method. Instead, the individual ion densities are inferred by the color change of the plasma. In Fig. 5.7, the xenon and argon ion densities must be equal somewhere between the two extremes of pure argon and pure xenon. The calculated average neutral cross sections for an argon only plasma and a Xe-Ar plasma, at a gas ratio of 35-65, are $4.44 \times 10^{-19} \text{ m}^2$ and $4.93 \times 10^{-19} \text{ m}^2$, respectively. The increase in the average neutral cross section of 11.0% is comparable to the decrease in ion beam speed of 12.7% shown in Fig. 5.7.

FIGURE 5.4: The normalized parallel argon IVDF at $(r, z) = (0, 164)$ cm as a function of both helium and argon pressure. The total pressure, 0.47 mTorr, was kept constant.

FIGURE 5.5: The normalized parallel argon IVDF at $(r, z) = (0, 164)$ cm as a function of argon pressure with a constant helium pressure of 0.71 mTorr.

FIGURE 5.6: The normalized parallel argon IVDF at $(r, z) = (0, 164)$ cm as a function of helium pressure with a constant argon pressure of 0.17 mTorr.

FIGURE 5.7: The normalized parallel argon IVDF at $(r, z) = (0, 164)$ cm as a function of xenon pressure with a constant argon pressure of 0.17 mTorr.

5.3 Xenon LIF in Xe-Ar and Xe-He Plasmas

Shown in Fig. 5.8 is the xenon IVDF as a function of argon partial pressure with a constant xenon pressure of 0.09 mTorr. Argon is the lighter mass in this case and the xenon speed decreases rather than increases. It is difficult to describe the xenon IVDF as having a clear ion beam because the total pressure is high enough that the ion beam is no longer present. Instead, the measured xenon flow is more similar to ambipolar diffusion obeying the Boltzmann relation.¹²⁸ The neutral pressure of argon (higher ionization potential) must be large in order to produce any argon ions in a xenon plasma.

The large discrepancy between the xenon and helium masses, 131 amu and 4 amu respectively, presents an interesting case. Fig. 5.9 shows the xenon IVDF as a function of helium partial pressure with the xenon partial pressure kept constant at 0.09 mTorr. The xenon ion beam is clearly slowing with the addition of helium. It is unlikely, given the experimental setup, that any significant amount of helium ions would be produced in a xenon-helium plasma unless the partial pressure of helium was significantly high. The ionization energy of helium is 24.59 eV, double that of xenon. By the time any helium ions are produced, the pressure threshold of ion beam formation has been surpassed. The measured xenon ion beam speed for a pressure of 0.09 mtorr in a xenon only plasma was 4.5 km/s. With a xenon pressure of 0.09 mTorr and a helium pressure of 1.0 mTorr the xenon ion beam speed is 3.3 km/s (see Fig. 5.9). For the data shown in Fig. 5.2, the xenon ion beam disappears above pressures of 0.50 mTorr. The xenon ion beam persists for higher total pressures with helium than it would have in a xenon only plasma.

FIGURE 5.8: The normalized parallel xenon IVDF at $(r, z) = (0, 164)$ cm as a function of argon pressure with a constant xenon pressure of 0.09 mTorr.

FIGURE 5.9: The normalized parallel xenon IVDF at $(r, z) = (0, 164)$ cm as a function of helium pressure with a constant xenon pressure of 0.17 mTorr.

Chapter 6

Conclusions

These measurements suggest a new model of the ion beam formation process in expanding, low-density, helicon source plasmas.¹²⁹ At low neutral pressures, the rf power couples strongly to electrons within one skin depth of the chamber wall. The electrons are strongly heated, forming a high energy, low collisionality energetic tail. These energetic, magnetized¹³⁰ electrons exit the source by streaming out along the expanding magnetic field. As they pass through the neutral gas, these energetic electrons create an annulus of increased plasma density through enhanced ionization upstream and downstream of the source-expansion chamber junction. The annulus of increased plasma density appears as a hollow plasma density profile. Hollow plasma density profiles downstream of the expansion region were recently reported by other researchers^{36,131} and similar hollow profiles are observed in these experiments.

The ring of high-energy, magnetized, electrons streaming out along the expanding magnetic field naturally creates a significant upstream-downstream charge imbalance. The resulting ambipolar electric field²² pushes out a centrally confined cylinder of energetic ions; an ion beam. As noted previously, these ions, with bulk speeds of 8-10 km/s, complete only 1/4 of a gyro-orbit while traveling 16 cm in the axial direction. The ions are essentially unmagnetized and

are effectively detached from the magnetic field,¹³⁰ i.e., their motion is too fast for adiabatic constraints such as μ conservation to hold. It also appears that the beam ions respond to the radial component of the ambipolar electric field established by energized electrons flowing along magnetic field lines near the plasma radial periphery. Since the beam ions get a boost in perpendicular velocity, there is a radial displacement of the hollow beam radial profile as beam ions are tugged along with the expanding electron rich annulus at the plasma periphery. The perpendicular energization of beam ions is strongest at the furthest upstream locations measured because the annulus of energetic electrons is closer to the central axis upstream and because the electric fields are stronger there as well.

The self-consistent physical picture that emerges is that of two nested hollow concentric cylinders of plasma, a hot electron dominated outer cylinder encircling the inner cylindrical ion beam. As the hot electrons and bulk ions follow the magnetic field lines, they drag the beam ions radially outward and downstream through the intermediary electric field. The magnetic forces on beam ions are inconsequential to their radial expansion or their parallel acceleration. The electric field arising from the field-aligned, ion hole structure determines the rate of radial expansion of the beam ions. Since both the energetic electrons and the ion beam travel from the source to the expansion region, i.e., the energetic electrons are not moving antiparallel to the ions, the observed particle motion is inconsistent with expectations for a DL that stretches across the entire expansion region.

The hypothesis proposed here for the origins of the strong potential difference between the source and expansion region that spontaneously appears in these low pressure, expanding plasmas is consistent with a variety of other phenomena that have been reported in these sources. For example, *Chakraborty Thakur et al.* have demonstrated significant changes in ion beam creation, plasma density profiles, instability growth, and plasma rotation depending on whether or not the inner surface of the expansion chamber is conducting or insulating.¹³²

Those observations reflect the critical role electrons in the plasma edge flowing downstream from the source play in setting up the overall potential structure upstream and downstream of the expansion location. *Charles* and *Boswell* were one of the first to identify the neutral pressure threshold for DL formation¹ and have reported the existence of energetic electrons³⁸ and hollow density profiles³⁶ in their experiments. Other groups have also reported hints of energetic electrons in their expanding helicon source plasmas.^{80,124}

Prior work at WVU found a strong correlation between enhanced upstream density and the appearance of the downstream ion beam. For rf coupling levels that did not result in formation of an ion beam, the upstream plasma density was lower than when a beam formed⁵³. This is consistent with the idea that energetic electrons were passing through the upstream plasma with sufficient energy to enhance ionization of the background neutral gas. While these reports of energetic electrons flowing downstream into the expansion region are inconsistent with a DL, they are consistent with this new paradigm.

The perpendicular IVDFs reported here have introduced a fundamentally different perspective in the study of ion beam formation in expanding helicon plasmas. The upstream ions are not only accelerated along the field, they are accelerated by a complex, multi-dimensional, annular electric field structure that results in an effective radial ion temperature of many eV. The radial electric field strengths are consistent with the measured perpendicular IVDFs. These highly anisotropic ion distributions are likely to drive a variety of plasma instabilities.

As seen in other experiments, LIF measurements of the $v = 8,000$ m/s ion beam fade with downstream distance in a manner consistent with quenching of the initial ion metastable state needed for LIF. The calculated quenching cross section of 2.7×10^{-18} m² is comparable to expectations for inelastic collisions of the ions with electrons. Perhaps somewhat surprising was the resonant effect of the downstream magnetic field strength on the ion beam velocity. The ion

beam velocity increased from 8,000 m/s to 12,000 m/s and the beam amplitude also increased in these experiments for a downstream magnetic field of 31 G. Therefore, for plasma thruster applications, a finite downstream magnetic field may prove beneficial even though additional resources are required to produce the downstream magnetic field. Experiments in pure argon and pure xenon plasmas determined that the accelerating mechanism was the same for similar pressures. For multi-ion species plasmas, the work presented here arrives at a different conclusion than *Biloiu* and *Scime*.⁶³ The reason for this discrepancy is because *Biloiu* and *Scime*⁶³ kept the flow rate equal to 10 SCCM for the duration of the experiment. In this work, a similar case with argon and helium was shown in Fig. 5.3. The effect of each gas is unclear from these measurements alone. Therefore, additional measurements were conducted, as discussed in Sec. 5.2, to isolate the effect of each gas species. The effect of adding a second gas at fixed total pressure is to slow the ion beam regardless of the relative individual densities. However, the ion beam is faster with addition of a light species compared to the baseline single species plasmas. A light species has a smaller atomic radius and will produce less drag which will result in a faster ion beam. The theoretical prediction by *Baalrud et al.*¹⁶ for multi-ion sheaths is not necessarily refuted, rather it is simply not applicable to the accelerating structures that appear in expanding helicon plasmas. Experiments by *Yip et al.* and *Severn et al.* did not isolate the effect of each gas as we have done nor did they account for changes due to collisions. If sheath physics was really responsible for the ion acceleration in expanding helicon sources, it might be tempting to combine mixtures of ion species to maximize both thrust and specific impulse in a helicon based thruster. However, these measurements suggest that the effects on beam velocity are explainable by simple collisional processes.

The experiments in this dissertation have shown that the ion acceleration mechanism is not a true DL. Instead, these measurements suggests a mechanism specific to helicon sources. Future studies of these plasmas will focus on the electrostatic wave activity and its impact on

the complex IVDFs reported here. Particle tracking, through LIF optical tagging and/or computational work, could provide an explanation for the origins of the various groups of particles. Measurements of the neutral particle velocity distribution functions with LIF could provide further insight into the role they play in multi-ion species plasmas.

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