Optogalvanic and laser-induced fluorescence spectroscopy of the Zeeman effect in xenon

IEPC-2007-229

Presented at the 30th International Electric Propulsion Conference, Florence, Italy September 17-20, 2007

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We present optogalvanic and LIF measurements of Zeeman effects on the $6s \ {}^{2}[1/2]_{1}^{0} \rightarrow 6p \ {}^{2}[3/2]_{2}$ absorption line of neutral xenon at 834.682 nm. Excitation was provided by a narrow-linewidth ($\Delta \nu < 10$ MHz) TUI Optics TA-100/830 tapered-amplifier diode laser. Optogalvanic spectra were provided by a Hamamatsu L2783 xenon-neon galvatron placed in the gap of a C-coil electromagnet, while LIF spectra were taken in the plume of a hollow cathode mounted along the axis of a cylindrical solenoid. A half-wave plate and polarizing beamsplitter permitted easy rotation of the beam polarization vector E with respect to the magnetic field B. Zeeman splitting of the π -polarized (E || B) spectrum is negligible at magnetic field intensities up to 300 gauss. The σ -polarized (E \perp B)spectrum, however, shows strong Zeeman splitting, with a 796-MHz gap between the main peaks of the spectrum at 300 gauss.

I. Introduction

Laser-induced fluorescence (LIF) provides a powerful technique for investigating the properties of both ionic and neutral species in the interior and near-field plume of electrostatic thrusters. Though high-speed Langmuir and emissive probe measurement of Hall thruster discharge channels^{1–4} and ion engine discharge chambers^{5–7} has been shown to be useful for measuring electron temperatures, ion densities, and plasma potential contours, it can not return any information about ion velocities. Indeed, electrostatic probes are fundamentally incapable of measuring neutral flow properties.

LIF spectra, on the other hand, can readily return a large amount of information about both ionic and neutral species. The interrogation volume for LIF is defined by the intersection of the laser beam and the collection optics axis, ensuring that LIF measurements can be made with high spatial resolution (usually on the order of a millimeter). Doppler shifting of the LIF spectrum provides high-resolution bulk velocity measurements, even at low signal-to-noise ratios (SNRs) or in the absence of an effective absorption line structure model. Though there are gaps in the spectrum of wavelengths covered by commercially-available extended-cavity diode lasers (ECDLs), usable absorption lines can be found for many species of interest to the electric propulsion community. For example, four neutral xenon (Xe I) transitions (at 823.164 nm, 826.652 nm, 828.012 nm, and 834.682 nm) and one singly-ionized xenon ion (Xe II) transition (at 834.724 nm) are readily probed by ECDLs centered around 830 nm. Published hyperfine structure (HFS) constants are available for many of these transitions, permitting curve-fit extraction of the species temperature for Maxwellian plasmas⁸⁻¹⁰ and deconvolution of the velocity distribution for nonequilibrium plasmas.^{11, 12}

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At high SNRs, Zeeman splitting of the LIF spectrum can also yield information about the local magnetic field. At moderate magnetic field intensities, when the laser beam polarization \mathbf{E} is parallel to the magnetic field \mathbf{B} (π -polarization), Zeeman splitting is negligible and the velocity distribution can be extracted from the LIF spectrum. When the laser beam is perpendicular to the magnetic field (σ -polarization), Zeeman splitting is appreciable, and the magnetic field intensity can be estimated by fitting the LIF spectrum to a computational model of the Zeeman effect on hyperfine energy levels.¹³

In a series of Xe I optogalvanic spectrocopy (OGS) experiments carried out with a Hamamatsu L2783 xenon-neon galvatron placed in the air gap of a custom electromagnet, we have measured Zeeman splitting of the $6s \ {}^{2}[1/2]_{1}^{0} \rightarrow 6p \ {}^{2}[3/2]_{2}$ transition of neutral xenon in both π and σ polarizations. As expected, Zeeman splitting of the π -polarized spectrum is negligible at magnetic field intensities up to 300 gauss. The σ -polarized spectrum, however, shows strong Zeeman splitting. Subsequent LIF experiments carried out in the plume of a 10-A hollow cathode also show strong π -polarized Zeeman splitting in a magnetic field intensity of 173 Gauss.

II. Theory



Figure 1. Optogalvanic and LIF schemes using Xe I absorption at 834.682 nm.

Figure 1(a) shows the fine structure transitions for optogalvanic and non-resonant LIF spectroscopy of the 834.682 nm transition of Xe I. Absorption of 834.682 nm laser light causes xenon atoms resting in the lower state $6s \ {}^{2}[1/2]_{1}^{0}$ to jump to the excited upper state $6p \ {}^{2}[3/2]_{2}$. Since this upper state is closer to the first ionization potential for xenon (12.1298 eV) than the lower state, the probability of electron-impact ionization (and thus the galvatron discharge current) increases linearly with the upper state population, providing an optogalvanic signal. Those atoms which are not ionized out of the upper state quickly decay to a lower state via an allowed transition. One of these decay paths is the $6p \ {}^{2}[3/2]_{1} \rightarrow 6s \ {}^{2}[3/2]_{1}^{0}$ fluorescence at 473.415 nm, providing a non-resonant LIF signal.

A. Hyperfine structure

In xenon (and other species with HFS), the LIF spectrum can be modeled as the convolution of the hyperfine splitting function $h(\nu)$ with natural and Doppler broadening functions,

$$i(\nu) = h(\nu) \otimes l(\nu) \otimes d(\nu).$$
(1)

The hyperfine splitting $h(\nu)$ is a series of n Dirac delta functions,

$$h(\nu) = \sum_{j=1}^{n} p_j \delta(\nu - \nu_j).$$
 (2)

where ν_j is the jth hyperfine line center and p_j is the jth hyperfine line intensity. The 6s ${}^2[1/2]_1^0 \rightarrow 6p {}^2[3/2]_2$ transition of Xe I has n = 18 such hyperfine splittings.

Hyperfine structure consists of two components. Isotopic splitting is caused by the mass and nuclear volume differences between isotopes in the sample. Nuclear-spin splitting is caused by the coupling between the nuclear spin \mathbf{I} and the total electronic angular momentum \mathbf{J} . For the sake of brevity, we will henceforth use the term "HFS constants" to denote the combined set of isotopic shifts and nuclear-spin interaction constants.

There are nine stable isotopes of xenon, seven of which have natural abundances greater than one percent. Each of these isotopes causes a slightly different transition energy E_{ij} , resulting in isotopic splitting. Table 1 presents naturally-occurring xenon isotopic abundances and shifts (conventionally defined relative to the center frequency for ¹³²Xe).

Table 1. Isotopic abundance and shift relative to 132 Xe for the Xe I 6s ${}^{2}[1/2]_{1}^{0} \rightarrow 6 {}^{2}[3/2]_{2}$ transition at 834.682 nm (from Suzuki *et al.*²⁵). No isotopic shift values are available for 124 Xe and 126 Xe.

Mass (amu)	124	126	128	129	130	131	132	134	136
Abundance $(\%)$	0.0096	0.009	1.92	26.4	4.1	21.1	26.9	10.4	8.9
Shift (MHz)			84.1	96.5	37.1	72.8	0.0	-38.9	-116.9

Seven of the nine stable isotopes of xenon have even atomic mass, resulting in no nuclear spin **I**. The two isotopes with an odd atomic mass, however, have non-zero nuclear spin quantum numbers I. The lighter isotope, ¹²⁹Xe, has I = 1/2, while ¹³¹Xe has I = 3/2. These non-zero nuclear spins cause nuclear-spin splitting of the atomic energy levels. This nuclear-spin splitting is 1 to 2 orders of magnitude broader than the isotopic splitting, and provides most of the characteristic shape of our absorption line.

The total angular momentum $\mathbf{F}=\mathbf{I}+\mathbf{J}$ takes quantum number values

$$F = I + J, I + J - 1, \dots, |I - J|$$
(3)

where J is the total electronic angular momentum quantum number.²⁹ The extra term energy due to nuclear-spin splitting is given by²⁸

$$E_{\rm nss} = A\frac{C}{2} + BD,\tag{4}$$

where A is the nuclear magnetic dipole interaction constant, B is the nuclear electric quadrupole interaction constant, and the terms

$$C = F(F+1) - I(I+1) - J(J+1)$$
(5)

and

$$D = \frac{(3C/4)(C+1) - I(I+1)J(J+1)}{2I(2I-1)J(2J-1)}$$
(6)

contain the nuclear spin-orbit interactions.

Table 2 presents energy levels and nuclear-spin structure constants for the Xe I absorption line at 834.682 nm. Constants for the lower state $(6s\ ^2[1/2]_1^0)$ are taken from Suzuki *et al.*,²⁵ while Jackson and Coulombe²⁶ provide constants for the upper state $(6s\ ^2[3/2]_2)$.

Table 2. Energy levels and nuclear-spin structure constants for the Xe I $6s \ {}^{2}[1/2]_{1}^{0} \rightarrow 6 \ {}^{2}[3/2]_{2}$ transition at 834.682 nm.

	E_i	A_{129}	A_{131}	B_{131}
Level	(eV)	(MHz)	(MHz)	(MHz)
$6s\ ^2[1/2]_1^0$	9.56972	-5801.1	1713.7	24.5
$6s^{\ 2}[3/2]_2$	11.0547	-2892	831.0	26.8

The transition rule for nuclear-spin splitting is $\Delta F \equiv F - F' = [0, \pm 1]$, where F is the upper and F' is the lower state's total angular momentum quantum number. (The zero-zero transition is forbidden,

 $F = 0 \not\rightarrow F' = 0.$) Figure 1 shows the fine structure and nuclear-spin splitting for Xe I 6s ${}^{2}[1/2]_{1}^{0} \rightarrow 6 {}^{2}[3/2]_{2}$ transition at 834.682 nm.

The relative intensity of each nuclear-spin split component is given for a $J \rightarrow J - 1$ transition by³⁰

$$I(F \to F - 1) \propto \frac{P(F)P(F - 1)}{F}$$
 (7)

$$I(F \to F) \propto \frac{(2F+1)}{F(F+1)} P(F)Q(F)$$
(8)

$$I(F-1 \to F) \propto \frac{Q(F)Q(F-1)}{F}$$
(9)

where P(F) = (F + J)(F + J + 1) - I(I + 1) and Q(F) = I(I + 1) - (F - J)(F - J + 1).

For the isotopes with even mass numbers (i.e., without nuclear-spin splitting), the line intensity p_j is linearly proportional to the naturally-occurring abundance for each isotope. The line intensities of ¹²⁹Xe and ¹³¹Xe are linearly proportional to the product of the isotopic abundance and the relative intensity of the nuclear-spin split components.

B. Zeeman splitting

In species (or isotopes) without nuclear spin splitting, anomalous Zeeman splitting results from precession of the total electronic angular momentum \mathbf{J} around the the magnetic field \mathbf{B} . The resulting magnetic quantum number M_J takes the values

$$M_J = -J, -J + 1, \dots, J.$$
(10)

The transition rule for Zeeman splitting is $\Delta M_J \equiv M_J - M'_J = [0, \pm 1]$ where M_J is the upper and M'_J is is the lower state's magnetic quantum number. Light emitted (or absorbed) by the $\Delta M_J = 0$ transitions is polarized parallel to the magnetic field ($\mathbf{E} \parallel \mathbf{B}$) in what is called π -polarization. The remaining $\Delta M_J = \pm 1$ transitions emit (or absorb) light that is polarized at right angles to the magnetic field ($\mathbf{E} \perp \mathbf{B}$) in what is called σ -polarization.

The extra term energy due to anomalous Zeeman splitting is given by 28

$$E_Z = g_J M_J \mu_B B \tag{11}$$

where μ_B is the Bohr magneton, B is the magnetic field intensity, and g_J is the Landé factor

$$g_J = 1 + \frac{J(J+1) - L(L+1) + S(S+1)}{2J(J+1)}.$$
(12)

The relative intensity of each Zeeman-split component is given for a $J \rightarrow J + 1$ transition by

$$I(M_J \to M_J) \propto 4(J^2 - M_J^2) \tag{13}$$

$$I(M_J \to M_J - 1) \propto (J + M_J)(J + M_J - 1)$$
 (14)

$$I(M_J \to M_J + 1) \propto (J - M_J)(J - M_J - 1).$$
 (15)

In species (and isotopes) with nuclear spin splitting, the above theory for anomalous Zeeman splitting is no longer valid. We refer interested readers to the companion paper by Ngom *at al*,¹³ which describes a nonlinear Zeeman-hyperfine model that separates the time-independent Schrodinger wave-equation into independent wavefunctions for the nucleus and a single orbiting electron. The resulting eigenvalue problem can be solved for all allowable energy levels, and the intensities for allowed transitions between quantum states can be calculated from the resulting wavefunctions for each unique set of quantum numbers (J, F, M_J, M_I) .

III. Apparatus and Procedure

A. Laser and optics

The laser system used in this experiment is a TUI Optics TA-100/830 tapered-amplifier diode laser. This system has a nominal linewidth of ~ 10 MHz, a tuning range of 820 nm to 836 nm, and a mode-hop-free range of over 10 GHz. Though capable of output powers in excess of 500 mW, off-peak efficiency drops and

Faraday isolation of the tapered amplifier reduce the available output power at the upper end of the system wavelength range $(835 \pm 1 \text{ nm})$ to approximately 250 mW.

The master oscillator for this system is an ECDL. Coarse wavelength control is provided by Peltier cooling of the ECDL and micrometer adjustment of the ECDL feedback grating angle. PC control of the fine tuning is provided by a LabVIEW VI, which uses the auxiliary output of a Stanford SR810 lock-in amplifier to send a control voltage to a piezoelectric transducer (PZT) driving the ECDL feedback grating. Feedback circuitry in the ECDL power supply decreases the master oscillator diode current during a PZT sweep to maximize the mode-hop-free range. In order to include the effect of changing laser output power, we normalize our reported LIF and OGS spectra by a quadratic curve-fit to the master oscillator power. It is important to note that this normalization carries an implicit assumption that the absorption line is not saturated. Steering mirrors inside the TA-100 enclosure couple the ECDL output into the tapered amplifier diode, which boosts the ECDL output power by a factor of more than 20 while maintaining the master oscillator wavelength and linewidth.

Figure 2 shows a schematic of the diode laser table contents. A borosilicate glass window splits off approximately 5% of the TA-100 output beam into a beamlet (beam A). The remaining 95% goes to a 50T/50R beam splitter; the reflected ~ 50% is further divided by a 70T/30R beam splitter into a transmitted 70% beamlet (beam B) and a reflected 30% beamlet (beam C).



Figure 2. Optical schematic for Zeeman-splitting experiment.

Beam A goes through an iris and is directed by two mirror into a half-wave plate that rotates the polarization vector to a user-selectable angle. The beam then passes through a Stanford SR540 chopper running at ~ 1.6 kHz into a polarizing beamsplitter cube; by changing the orientation of the polarizing cube, the user can ensure that the transmitted beam polarization is entirely parallel (or entirely perpendicular) to the surrounding magnetic field. The transmitted portion of the beam continues on into the Hamamatsu L2783 galvatron (placed in the electromagnet air gap) and terminates at a Thorlab DET110 photodiode. During saturation studies, additional neutral density filters can be added between the cube and the optogalvanic cell.

Beam B is directed by two mirrors through a variable neutral density filter, and then into a Burleigh WA-1000 wavemeter. The wavemeter can determine laser wavelength with a 1 pm resolution, which is sufficient for coarse alignment and troubleshooting of the laser. Oversampling of the PZT control voltage permits greater wavelength resolution by a quadratic curve-fit to the wavemeter output, but does not improve the absolute bias of the laser wavelength measurement.

Beam C is directed by two mirrors through an iris into a Burleigh SA-91 etalon assembly with a free spectral range of 2 GHz and a finese of 256. This provides a series of sharply-defined frequency markers, allowing high-resolution measurement of the frequency interval swept during a laser scan.

B. Optogalvanic cell and electromagnet

Figure 3 shows a more detailed view of the optogalvanic cell, electromagnet, and Hall-effect probe used in this experiment. The optogalvanic cell is a Hamamatsu L2783-42 XeNe-Mo galvatron. The core of the galvatron is a $\phi 6.25$ mm pair of cylindrical Mo electrodes, centered in a $\phi 25 \times 120$ mm glass cylinder filled with



Figure 3. Detailed view of the optogalvanic cell and Hall-effect probe mounted in the electomagnet air gap.

approximately 3 Torr of xenon and and 4 Torr of neon. The ends of the cylinder are angled approximately 10 degrees from a perpendicular to the electrode axis; though this is less than Brewsters angle for a glassair interface, it reduces the likelihood of retroreflected beams entering the discharge. This design permits unobstructed passage of a laser beam along the axis shared by the cylinder and electrodes.

A 250 V discharge across the galvatron electrodes produces a reasonably cool, dense, stationary plasma. A ballast resistor prevents runaway current growth after the discharge is struck, and provides a simple method of measuring the discharge current. The voltage drop across the ballast resistor is connected to the Stanford SR-810 lock-in amplifier through a 3.33 μ F, 450 V capacitor. The capacitor passes the AC component of the voltage drop, while isolating the lock-in from high DC voltages.

The electromagnet frame consists of two ϕ 76 x 74 mm iron cores on a iron C-frame. A Sorensen DCS 55-55 energizes the 14-gauge copper windings, which are hooked up in series for this experiment. Air cooling of resistive losses in the windings is provided by a large fan. Simple tests with the Hall-effect probe show that within a region extending about 3 cm from the optogalvanic cell interrogation volume, non-axial components of the magnetic field constitute less than half a percent of total magnetic field intensity. Two aluminum support struts place the electomagnet winding cores on the same horizontal plane as the optogalvanic cell axis.

Based on the geometry and the material of the optogalvanic cell, we assumed that the optogalvanic cell produces negligible distortion to the magnetic field. This assumption was verified by subsequent Hall-effect probe measurements, which showed that the magnetic field magnitude and direction in the immediate vicinity of the optogalvanic cell did not change appreciably when the cell was dormant, operating with a 250 V discharge, or removed entirely from the setup.

The Hall-effect probe is a model MMZ-2508-UH 3-axis probe connected to a LakeShore model 460 gaussmeter. A vertical translation stage places the probe axis parallel to the optogalvanic cell axis, with the probe tip exactly 2 cm above the intersection of the electromagnet and optogalvanic cell axes. This placement lets us use the Hall-effect to monitor drifting of the magnetic field over the course of a laser scan. Calibration runs to correlate this off-center magnetic field with the magnetic field in the interrogation volume showed that the two field strengths are proportional to within 1%. Although the probe in question can measure magnetic field strength to within 0.25%, the uncertainty is limited by the use of off-axis measurement to predict the centerline value.

C. Vacuum facility

LIF experiments were performed in the Large Vacuum Test Facility (LVTF) at the University of Michigan. The LVTF is a $\phi 6 \text{ m} \times 9$ stainless steel-clad cylindrical chamber. Pumping is provided by seven single-stage CVI model TM-1200 cryopumps and liquid nitrogen shrouds, with a nominal pumping speed of 240,000 l/s. Facility pressure is monitored by two hot-cathode ionization gauges, and base pressure is approximately

1.2×10^{-7} Torr.

Xe I LIF spectra were recorded at the exit plane of a hollow cathode operating in diode mode with a xenon flow rate of 10 sccm. During these experiments, the hollow cathode was mounted on the axis of an iron-core solenoid, providing a user-controlled axial magnetic field at the cathode orifice. Post-test calibration of the magnetic field intensity as a function of solenoid current was provided by the same Hall-effect probe and gaussmeter used in the optogalvanic cell experiments. Research-grade xenon propellant (99.999% pure) was supplied to the hollow cathode by a commercially available flow meter and controller, having an accuracy of $\pm 1\%$ of full scale. Calibration of the flow system by the constant volume method⁴ includes the effects of compressibility.



Figure 4. Laser beam delivery and fluorescence collection optics in the LVTF.

Figures 4(a) and 4(b) show the LVTF beam handling setup. The beam enters the LVTF through a sapphire window. Approximately 50 mW of beam power enters the periscope through a replacable antire-flection (AR) coated window. Three near-infrared mirrors within the periscope direct the beam to a focusing telescope parallel to the cathode and solenoid axis, reducing the beam diameter to less than 1 mm.

Two separate AR windows protect the $\phi 100 \text{ mm}$, f/2.5 collection lens. The collimated fluorescence from the thruster plume passes out of the LVTF through another sapphire window and into a $\phi 100 \text{ mm}$, f/5 lens. Two protected aluminum mirrors direct the focusing beam onto a Spex H-10 monochromator (using 500- μ m slits) with a Hamamatsu 928 PMT. A Stanford SR850 lock-in amplifier extracts the LIF signal from the PMT output. The interrogation volume, defined by the intersection of the laser beam and the doubly-magnified monochromator slit image, is a $\phi 1 \times 1$ mm cylinder oriented parallel to the cathode axis.

IV. Results

A. Saturation studies

At low laser intensities, the optogalvanic signal scales linearly with laser power. As the laser intensity increases, the lower-state depopulation rate due to absorption increases. Since the upper-state depopulation rate due to spontaneous emission remains constant, the optogalvanic signal eventually goes asymptotic at high laser intensities. The resulting *saturation* tends to flatten out peaks of the lineshape, causing the optogalvanic spectrum to appear artificially broadened.

Figure 5 shows the results of a series of 834.682 nm scans at varying levels of laser power entering the galvatron. The lock-in amplifier time constant was set to 300 ms, and the total time for each scan was 300 s. Neutral density filters placed between the 0.4 mW input beam and the galvatron reduced input power by regular increments. The saturation curve shown in Figure 5 is highly linear, indicating that the integrated optogalvanic signal intensity is effectively not saturated at 0.4 mW.

B. Optogalvanic spectra

All remaining optogalvanic spectra were taken with a 0.4 mW laser input to the galvatron. The lock-in time constant was 300 ms, and the total time for each scan was 300 s.



Figure 5. Saturation curve for optogalvanic spectrum of Xe I at 834.682 nm.



Figure 6. Zeeman-split optogalvanic spectra of Xe I at 834.682 nm, π -polarization (B || E).

Optogalvanic spectra taken with the laser polarization vector parallel to the magnetic field (π -polarization) are summarized in Figures 6(a) and 6(b). The plots are essentially identical, with no apparent Zeeman splitting as the magnetic field intensity rises from 0 to 300 gauss (0-0.03 T). All eleven scans taken in π -polarization (at approximately 30-gauss intervals) returned the same unsplit optogalvanic spectrum.

As in the π -polarization scans described above, we took σ -polarized scans at 30-gauss intervals between 0 and 300 gauss. Figures 7(a) through 7(f) show selected optogalvanic spectra taken with the laser polarization vector perpendicular to the magnetic field (σ -polarization). Figures 8(a) presents the resulting spectra as a three-dimensional surface, while Figure 8(b) shows optogalvanic signal intensity contours as a function of frequency detuning and magnetic field intensity.

Due to the multiply-peaked hyperfine structure of this transition, Zeeman splitting of the major peak is not immediately apparent until the magnetic field strength is greater than 120 gauss. At 300 gauss, however, the major peak has clearly split into two equally-strong peaks with a 1-GHz separation.

C. Cathode plume LIF

Figures 9(a) and 9(b) show LIF spectra taken at the exit plane of a hollow cathode running on xenon in the Large Vacuum Test Facility. The lock-in time constant was 300 ms for the B = 0 scan and 1 s for the B = 174 gauss scan. The signal-to-noise ratio (SNR) in Figure 9(a) is appreciably lower than the SNR for our optogalvanic spectra. Figure 9(b), in turn, has a SNR that is approximately half that of the unsplit spectrum.

The unsplit LIF spectrum of Figure 9(a) compares favorably to the unsplit optogalvanic spectrum shown in Figure 7(a). Despite the decreased SNR, the LIF spectrum does not look appreciably different from the OGS spectrum. The peak widths are essentially identical, as are the relative depths of the valleys between hyperfine splittings.

However, comparison of Figure 9(b) with the σ -polarized 180-gauss optogalvanic spectrum of Figure 7(d) shows some striking differences. Despite the slightly smaller magnetic field strength, the LIF spectrum appears to have a slightly more pronounced splitting (~ 700 MHz between major peaks, compared to ~ 600 MHz for the OGS spectrum). The main peaks are also more strongly asymmetric, and the valley between the main peaks is more pronounced. All of these differences, though, are readily explained by the lower SNR and coarser frequency resolution of the LIF scan.

V. Conclusions

We have demonstrated that $6s \ {}^{2}[1/2]_{1}^{0} \rightarrow 6p \ {}^{2}[3/2]_{2}$ absorption line of neutral xenon at 834.682 nm is readily-accessible to commercially-available diode lasers. Optogalvanic spectra taken by lock-in amplification of the discharge current signal from a standard galvatron provide a simple, low-noise method of investigating the properties of this transition. Multiple π -polarized ($\mathbf{E} \parallel \mathbf{B}$) spectra show neglible Zeeman splitting at magnetic field intensities up to 300 gauss. Similar σ -polarized ($\mathbf{E} \perp \mathbf{B}$) spectra, however, shows a monotonic increase in Zeeman splitting with magnetic field intensity. At 300 gauss, the gap between main peaks of the σ -polarized spectrum is 796-MHz.

Future applications of this technique will use a pair of LIF spectra taken at perpendicular polarizations to nonintrusively measure the velocity distribution (via deconvolution of the π -polarized spectrum with a computational model of the zero-temperature hyperfine-broadened spectrum), the magnetic field intensity (by least-squares fitting of the σ -polarized spectrum to a computational model of the Zeeman effect on hyperfine energy levels), and the magnetic field direction (from the polarization direction vectors). The resulting diagnostic would be useful over a wide range of near-field and internal flows of interest in electrostatic thruster research.

VI. Acknowledgements

This research was supported by the Air Force Office of Scientific Research grants F49620-02-1-0051 and F49620-01-1-0061 (Dr. Mitat Birkan is the contract monitor for both). The authors are greatly appreciative of this support.

We would like to thank Bryan Reid, who installed, calibrated, and operated the hollow cathode used in the LIF portion of this research. We would also like to thank the graduate students at PEPL for their



Figure 7. Zeeman-split optogalvanic spectra of Xe I at 834.682 nm, σ -polarization (B \perp E).



Figure 8. Three-dimensional surface and contour plots of σ -polarized optogalvanic spectra of Xe I at 834.682 nm.



Figure 9. Zeeman splitting LIF spectra of Xe I at 834.682 nm, taken at hollow cathode exit.

assistance in experimental setup and operation, and the Department's technical support staff for their aid with facility repairs and modifications.

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