Measurement of ion motion in a shear Alfvén wave

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In this study, the technique of laser-induced fluorescence (LIF) has been used to measure $T_i$ and the $\mathbf{E} \times \mathbf{B}_0$ and polarization drifts of shear Alfvén waves in the Large Plasma Device at UCLA [W. Gekelman, H. Pfister, Z. Lucky, J. Bamber, D. Leneman, and J. Maggs, Rev. Sci. Instrum. 62, 2875 (1991)]. The waves were launched by an antenna located at the end of the device and were observed to propagate along the axis of a 9 m long, 40 cm diameter cylindrical argon plasma in the kinetic regime [$\beta_e \approx 9.5 (m_i/m_e)$], with $f_{\text{wave}}/f_{ci} \approx 0.8$. Care was taken to record the measurements from various diagnostics at the same spatial positions on four cross-sectional planes along the length of the plasma. Two-dimensional LIF measurements of the ion drifts perpendicular to $\mathbf{B}_0$ were undertaken. Ion drifts were observed to be as large as 14% of the ion thermal speed. The ion polarization and $\mathbf{E} \times \mathbf{B}_0$ drifts were distinguished by their phase relation to $\mathbf{B}_\text{wave}$. The measured drifts are compared to kinetic theory, $\mathbf{E}_\perp$ (the transverse component of $\mathbf{E}_\text{wave}$) was computed from the drift velocities, and $E_\parallel$ was estimated from $E_\perp$. © 2005 American Institute of Physics.

I. INTRODUCTION

In space physics, the high resolution measurements of spacecraft such as Freja$^1$ and Fast Auroral Snapshot Explorer (FAST) (Ref. 2) have generated great interest in the study of Alfvén waves in the auroral ionosphere. The Freja mission observed strong low frequency electromagnetic spikes with $\Delta E/\Delta B = v_e/c$ in conjunction with deep ($\delta n/n \approx 0.7$) spatially narrow (width $\sim \delta = c/(\omega_{pe})$) density depressions.$^3$ The authors propose that highly nonlinear Alfvén waves with an electric field parallel to the ambient magnetic field were responsible for the density depletion.

The FAST satellite has observed strong wave-particle interactions between energetic electrons and intense “electromagnetic ion cyclotron waves” (i.e., shear Alfvén waves).$^4$ The interactions include the coherent modulation of field-aligned electron fluxes as well as the acceleration of secondary electrons to form counterstreaming field-aligned fluxes with energies up to 300 eV. In another mission, the Polar spacecraft measured intense electric and magnetic field structures associated with Alfvén waves at the outer boundary of the plasma sheet. The events were observed when the spacecraft was field aligned with intense auroral structures below (detected by an onboard UV imager). The authors conclude that the Poynting flux associated with these waves is of one to two orders of magnitude larger than previously estimated and may be sufficient to power the acceleration of auroral electrons (as well as the energization of upflowing ions and Joule heating of the ionosphere).$^2$ A related study$^6$ correlates the observations of the FAST and the Polar satellites to identify the principal magnetospheric drivers of auroral acceleration mechanisms. The authors identify large-amplitude Alfvén waves (propagating from the magnetotail into the auroral region) as one of three main drivers of auroral acceleration, the other two being field-aligned currents and earthward flows of high-energy ion beams. The GEODESIC sounding rocket made simultaneous observations of two-dimensional ion distribution functions and plasma waves in the high-latitude topside ionosphere.$^7$ The rocket encountered large amplitude, short perpendicular wavelength Alfvén waves, and the corresponding $\mathbf{E} \times \mathbf{B}_0$ drifts were measured.

Until recently Alfvén waves have been difficult to observe in laboratory devices because of their low frequencies and long wavelengths. Early experiments were limited by noisy sources, high collisionality, and physical size. A review of the progress of laboratory experiments on Alfvén waves and their relationship to space observations was given by Gekelman.$^8$

Since the introduction of the Large Plasma Device (LAPD) at UCLA, detailed experimental studies of Alfvén waves have been carried out in tandem with the development of a detailed theoretical description. Morales, Loritsch, and Maggs$^9$ analytically described the structure of shear Alfvén waves with small transverse scale (on the order of $\delta = c/(\omega_{pe})$) for the inertial regime. Their theory was directly compared to experimental observations of shear Alfvén wave propagation in the LAPD (Ref. 10) with excellent agreement. Morales and Maggs extended their theory to describe shear Alfvén waves in the kinetic regime,$^{11}$ and other LAPD Alfvén wave experiments have been conducted$^{12–17}$ with an emphasis on their relevance to space plasmas and auroral processes.

Gekelman et al.$^{18}$ made measurements of the propagation of large-amplitude ($B_{\text{wave}}/B_0 \sim 10^{-3}$) Alfvén waves launched by an inductive loop antenna in an argon plasma. The ion temperatures and drifts of the waves were measured using laser-induced fluorescence (LIF) at a limited number of spatial locations. LIF has previously been used by others to measure the properties of electrostatic waves in plasmas.$^{19,20}$ In this work we present the first detailed LIF measurements of ion motion in a shear Alfvén wave.
TABLE I. Plasma parameters, derived quantities, and discharge characteristics.

<table>
<thead>
<tr>
<th>Description</th>
<th>Symbol</th>
<th>Value</th>
</tr>
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<tbody>
<tr>
<td>Gas fill pressure (Ar)</td>
<td>$P_{Ar}$</td>
<td>$5 \times 10^{-3}$ torr</td>
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<tr>
<td>Discharge current</td>
<td>$I_{dis}$</td>
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<td>Discharge voltage</td>
<td>$V_{dis}$</td>
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<tr>
<td>Discharge duration</td>
<td>$\Delta t_{dis}$</td>
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<tr>
<td>Magnetic field</td>
<td>$B_o$</td>
<td>1200 G</td>
</tr>
<tr>
<td>Plasma density</td>
<td>$n$</td>
<td>$\approx 1.3 \times 10^{12}$ cm$^{-3}$</td>
</tr>
<tr>
<td>Electron temperature</td>
<td>$T_e$</td>
<td>3.6 eV</td>
</tr>
<tr>
<td>Ion temperature</td>
<td>$T_i$</td>
<td>$\approx 0.8$ eV</td>
</tr>
<tr>
<td>Ion-cyclotron frequency</td>
<td>$f_{ci}$</td>
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<tr>
<td>Wave frequency</td>
<td>$f_{wave}$</td>
<td>36.0 kHz</td>
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<tr>
<td>Wave-to-cyclotron-frequency ratio</td>
<td>$\tilde{\omega}$</td>
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<td>Ion thermal speed</td>
<td>$v_{Ti}$</td>
<td>$2.0 \times 10^5$ cm/s</td>
</tr>
<tr>
<td>Ion sound speed</td>
<td>$c_i$</td>
<td>$2.9 \times 10^5$ cm/s</td>
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<tr>
<td>Electron thermal speed</td>
<td>$v_{Te}$</td>
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<tr>
<td>Alfvén speed</td>
<td>$v_A$</td>
<td>$3.6 \times 10^7$ cm/s</td>
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<td>$\beta_i(m_i/m_e)$</td>
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<td>$\lambda_D$</td>
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<td>Ion sound gyroradius</td>
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<td>Perpendicular wavelength</td>
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<td>Parallel wavelength</td>
<td>$\lambda_l$</td>
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<td>Perpendicular wave number</td>
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<tr>
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<tr>
<td>Collision-to-angular-frequency ratio</td>
<td>$\Gamma = v_{ei}/\omega$</td>
<td>11</td>
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II. SHEAR ALFVÉN WAVES

The shear Alfvén wave propagates along the ambient magnetic field and (to first order) causes magnetic field fluctuations orthogonal to the direction of the background magnetic field. When the frequency of the wave is a substantial fraction of the ion cyclotron frequency and the parallel phase velocity of the wave falls within the electron velocity distribution so that Landau damping is important, a kinetic description of the wave dispersion is necessary. This experiment was performed in the kinetic regime for shear waves where $v_A < v_{Te}$ and $\beta_i \geq m_i/m_e$, with $v_A = B_0/\sqrt{4\pi nM_p}$, $\beta_i = 8\pi nT_i/B_0^2$, and $v_{Te} = \sqrt{2T_e/m_e}$. In this case the kinetic dispersion relation for shear Alfvén waves (neglecting collisions) reduces to

$$\omega^2 = \frac{\omega_p^2}{k^2} (1 - \bar{\omega}^2 + k^2 \rho_i^2),$$

where $\rho_i = c_i/\omega_i$ is the ion sound gyroradius, with $c_i = \sqrt{T_i/m_i}$ and $\bar{\omega} = \omega/\omega_{ci}$. Note that in the limits $k \rightarrow 0$ and $\bar{\omega} \rightarrow 0$, Eq. (1) reduces to the pure magnetohydrodynamics relation. In previous experiments, wave propagation in both the kinetic and inertial regimes has been observed. The relation between the parallel and perpendicular components of the wave electric field in the kinetic regime is [see Ref. 12, Eq. (9)]

$$E_{||} = \frac{ikk_p \rho_i^2}{(1 - \bar{\omega}^2)} E_\perp.$$

Collisions affect shear wave propagation in two important ways. First, they cause additional dissipation (damping) as the wave propagates axially. Second, large perpendicular wave numbers are more heavily damped, which tends to smooth out and broaden the transverse structure of the waves. Because the LAPD plasma is highly ionized, electron-ion Coulomb collisions ($\nu_{ei} = 2.5 \times 10^8$ s$^{-1}$) are the dominant collision mechanism (see Table I). The experiments reported in this work were done in the kinetic regime ($v_A < v_{Te}$) and the plasma was made as collisionless as possible. Since the LIF diagnostic required the use of an argon plasma to make the parallel wavelength as short as possible without experiencing cyclotron damping, the wave frequency was set to 80% of $f_{ci}$. The experimental parameters are summarized in Table I.

III. EXPERIMENTAL SETUP

The experiments for this work were carried out in the original LAPD at UCLA. The plasma column was 9 m long and 40 cm in diameter, with $n_e \approx 10^{12}$ cm$^{-3}$ and $T_e$ of several eV. Since this work was completed, the LAPD has been upgraded to be even larger and more flexible. The plasma is created by a pulsed dc electrical discharge between an oxide coated cathode and nearby anode.

The experimental setup used to launch Alfvén waves and record diagnostic data is summarized in Fig. 1. As shown in the right side of the figure, Alfvén waves were launched into the plasma with an inductive loop antenna. The antenna is...
made of electrically insulated copper tubing bent into a long loop, 46 cm along \( B_0 \) and 6.5 cm across. When current is driven through the loop, the antenna currents along \( B_0 \) and the azimuthal antenna currents inductively couple into the plasma to drive shear Alfvén waves. The rf input signal was produced by an arbitrary function generator, and the antenna current was driven by a rf amplifier through an impedance-matched, frequency-resonant LC network. The wave driving signal for these experiments was a tone burst of ten sinusoidal cycles at 36 kHz with a maximum amplitude of 600 A peak-to-peak through the antenna.

A three-axis induction coil probe was used to measure the magnetic fields of the propagating shear Alfvén waves. The probe utilized shielded back-to-back coils to minimize electrostatic pickup noise. The magnetic field signals were recorded over a period of 1024 \( \mu \)s at a sample rate of 1 MHz with a 30 shot average by a rf amplifier through an impedance-matched, frequency-resonant \( LC \) network. The data window for the magnetic field signals begins 100 \( \mu \)s before the rf tone burst. A Langmuir probe was used to measure the density and temperature of the plasma. The magnetic field probe, Langmuir probe, and ion saturation current measurements were all taken at multiple positions on four cross-sectional planes along the plasma column, as shown in Fig. 2. The distances of the four \( x-y \) planes from the antenna (labeled as \( \Delta z \) in the figure) were 66 cm, 223 cm, 349 cm, and 475 cm, respectively. (Their positions were chosen to provide measurements at regular intervals along the \( z \) axis.) All data planes were centered on the same \( x-y \) position. Magnetic field measurements were taken on 20 \( \times \) 20 cm\(^2\) planes with 1 cm spacing between points.

To make reliable spatial correlations between wave and plasma properties that were measured with different probes a pulsed electron beam was used to reproducibly align the various diagnostics to the same magnetic field line (the same \( x-y \) coordinates) on each measurement plane. The electron beam was a tightly wound tungsten coil mounted 443 cm from the antenna. It emitted a field-aligned beam of electrons \( \approx 5 \) mm in diameter. The pulsed beam produced a readily observable signal on all planes. For each data plane each diagnostic probe was moved with stepper motors and the data plane aligned with the beam. The magnetic field data planes were aligned with an accuracy of \( \pm 3 \) mm and the LIF and Langmuir data planes with an accuracy of \( \pm 2 \) mm.

IV. USING LIF TO MEASURE VELOCITY DISTRIBUTIONS

A. Basic principle

LIF is a technique which enables highly localized and time-resolved measurements of the ion velocity distribution in a plasma. A tunable narrow-band laser is directed through a selected volume in the plasma and scanned over a frequency range that spans the spectral linewidth of a suit-
able ion transition. Ions will be excited by the laser and fluoresce if their velocity satisfies the Doppler relation

\[ k_i \cdot v_i = 2 \pi (n_i - v_p), \]

where \( k_i \) and \( v_i \) are the laser wave number vector and frequency, \( v_i \) is the ion velocity, and \( v_p \) is the frequency of the ion transition in the rest frame. The fluorescence at each frequency of the laser is proportional to the number of ions in the detection volume that have a specific velocity component along \( k_i \). Scanning the laser over the entire spectral line and recording the fluorescent intensity at each frequency step thus maps out the distribution function \( f_i(r, v, t) \) at position \( r \). Under ordinary conditions \( f_i \) will be Maxwellian-like the measured fluorescent light profiles shown in Fig. 3. The ion temperature can be computed from the width of the distribution, while the bulk ion drift velocity is derived from the center offset of the distribution relative to \( v_p \). Two- and three-dimensional velocity distributions can be obtained by redirecting the laser to pass through the detection volume along orthogonal axes.

B. LIF in argon

Input laser light at 611.492 nm excites argon ions in the metastable \( 3d^2^2G_{9/2} \) state to the \( 4p^2^2F_{5/2} \) state. Within a few nanoseconds \(^{29}\) they spontaneously decay to the \( 4s^2^2D_{5/2} \) state emitting fluorescent light at 460.957 nm (wavelengths measured in air). The spectral profile can thus be interpreted solely as a Doppler broadened Gaussian as long as other possible sources of spectral broadening are accounted for. In this experiment, the natural linewidth \(^{29}\) pressure broadening, \(^{30,31}\) Stark broadening \(^{30,31}\) Stark splitting, \(^{32,33}\) laser bandwidth, and instrumental bandwidth of the detection system were determined to be negligible. To minimize the effect of Zeeman splitting the input laser was polarized parallel to \( B_0 \) to exclude the \( \Delta m_j = \pm 1 \) transitions. When the composite profile includes only the \( \Delta m_j = 0 \) peaks at \( B = 1200 \) G, the measured full width at half maximum of the profile differs from the purely Doppler-broadened width by less than 0.3%.

Power broadening is an artificial broadening of the spectral line shape that occurs when the laser intensity becomes high enough to saturate the atomic transition, i.e., when the optical pumping rate and stimulated emission dominate over spontaneous emission. \(^{34-36}\) Goeckner, Goree, and Sheridan \(^{36}\) measured the spectral line broadening as a function of laser intensity for the same Ar\(^+\) transitions used in these experiments. Their results indicate that the laser intensity used in these experiments was close to power broadening and the apparent ion temperature may be inflated by a factor of 2. No obvious evidence of saturation was observed during the acquisition of LIF data and the ion temperatures measured in these experiments was \( T_i = 0.8 \) eV. As no evidence of power broadening was observed when the laser power was decreased, we will assume the ion temperature is the measured value. In these experiments the maximum laser intensity was needed to optimize the ratio of LIF to spontaneous light. The measurement of ion temperature in these experiments was secondary to the measurement of ion drifts which are unaffected by slight power broadening.

Figure 4 shows a schematic of the laser setup. The ion excitation light was produced by a Coherent 899-21 tuneable dye laser and Rhodamine 6G dye, with a typical output of 500 mW. The 899-21 was optically pumped with 6–8 W from an argon ion laser. The other pictured components served to calibrate, modulate, or monitor the laser output, and will be described subsequently.

The first attempts at detecting laser-induced fluorescent light in the LAPD plasma followed the conventional approach implemented by other labs—i.e., completely noninvasive external optics. \(^{27,37}\) This approach works well for plasmas in which the intensity of the laser-induced fluorescence is comparable to the spontaneously emitted ion light, or the LIF signal can be extracted through extensive signal averaging. The plasma used in this experiment satisfied neither criterion. Because the plasma is pulsed approximately once per second, signal averaging over more than ~100 shots is time consuming and impractical. The LAPD plasma (produced by fast electrons) also has a bright background of spontaneously emitted light at the same wavelength as the fluorescent light. This problem is exacerbated by the large
diameter of the plasma (≈40 cm). To image a fluorescent spot near the center of the plasma, the detection optics must peer through ≈20 cm of bright background light. Even when the best available spatial-filtering optical techniques were employed, the laser-induced fluorescent light was only a small fraction (≈7%) of the total light collected. The solution to this problem was to move the detection optics inside the plasma, as close as possible to the fluorescent spot. This led to the development of a robust optical probe which could measure two components of the ion velocity vector in the plane perpendicular to the axis of the background magnetic field. By mounting a single input beam at an angle of 45° with respect to the detection axis (see Fig. 5) measurements could be made along two orthogonal axes simply by rotating the entire probe shaft 180°.

The output frequency of the dye laser was tracked and measured two ways—with a wave meter (a Fizeau interferometer, model 7711 from New Focus) and by comparison with the absorption spectrum from an iodine cell. A beam splitter was used to sample ≈100 mW from the dye laser output beam which was then steered through the iodine cell and into the wave meter, as shown in Fig. 4. The wave meter was useful for coarse frequency tuning because it provides nearly instantaneous readings of the wavelength with a precision of 0.001 nm (Δf ≈ 0.8 GHz) or a measured velocity resolution which is ≈20% of v_{Ti} for 1 eV ions.

Higher precision frequency calibration was achieved using the iodine cell. Using tabulated wavelengths of iodine absorption peaks, the laser frequency was calibrated with an accuracy of ≈0.1 GHz. This corresponds to a resolution in velocity space of 0.06×10^5 cm/s or ≈3% of v_{Ti} for 1 eV ions.

Because of the large number of shots required for signal averaging and the long duration (≈1 h) of each frequency scan, LIF data could only be acquired at a fraction of the spatial positions where other diagnostics were recorded. On each of the four planes where data were taken (see Fig. 2), LIF data were taken with both the 45° downward and 45° upward probe orientations at 9 positions on a 3×3 cm^2 grid at the center of the plane (except for plane 4, where only three positions were recorded). Each LIF data run consisted of one full frequency scan at a single position with the probe in either orientation. The data recorded included the LIF signal, the laser power, the Alfvén antenna current, and the wavelength reading from the wave meter.

V. WAVE MAGNETIC FIELD DATA

Since the signals from the magnetic field coils are proportional to the time derivative of the magnetic field an integration technique was used to obtain the magnetic field. Its absolute value was obtained by using the probe to measure the field close to a long straight wire through which a known oscillating current close to the experimental frequency flowed. A rotation corrected the probe data for both the mounted angle of the probe coils and the angle of the probe shaft at each position.

Figure 6 shows a sample of B_{wave} data from measurement plane 2. Successive frames in the figure represent time increments of 4.0 μs. The full wave period is 28 μs (f_{wave} ≈ 36 kHz). The basic pattern of the magnetic fields is that of two counterstreaming axial currents. The B_{wave} vectors add constructively in the center of the pattern and destructively towards the edges. This wave field structure is very similar to that of previous studies in which shear Alfvén waves were launched by axial currents from two small disk antennas driven 180° out of phase.12,13 The wave fields have a broad transverse spatial pattern that is largely independent of small (~1 cm) inhomogeneities in plasma density. An interesting feature is that although the entire pattern oscillates predomi-
nately in a linear fashion along a nearly 45° diagonal, it also rotates as the wave fields pass through their minima and the axial currents twist around one another. For each of the four planes and for the temporal interval in which the wave was present a two-dimensional Fourier transform was used to determine the perpendicular wave number, which was $k_\perp = 0.628 \pm 0.30$ cm$^{-1}$. The large error reflects the large perpendicular wavelength with respect to the grid over which data were acquired. This error made it impossible to measure small changes in wave number from plane to plane.

To illustrate the propagation of the wave along $B_0$, Fig. 7 shows a snapshot in time of the $B_{wave}$ field on the four cross-sectional measurement planes. The parallel phase velocity was measured by tracking the axial propagation of a fixed point in phase along a field line. The average parallel phase velocity for these experiments was $v_{\parallel} = 3 \times 10^7$ cm/s with an uncertainty of $\approx 10\%$. Table II lists the amplitude of the wave magnetic field on each of the four measurement planes.

Using the measured components of $B_{wave}$ the axial current density can be computed as follows:

$$ j_z = \frac{c}{4\pi} (\nabla \times B_{wave}) \cdot \hat{z}. $$

For these low frequency waves, the displacement current is negligible. Figure 8 shows a sample of the computed axial current density along with $B_{wave}$ at a time when the fields are maximum on measurement plane 2.
VI. LIF DATA

For each position the digital LIF signals were acquired as a time series at each laser frequency step. A time series with the laser off was also acquired as a background reference. To reduce extraneous noise the data were smoothed over three time steps (out of 1024) and digitally filtered (a low-pass filter with cutoff at $f_c$). The laser-off signal was subtracted from the laser-on signals and normalized to the measured instantaneous laser power. Finally, a nonlinear least-squares fitting algorithm was used to fit a Gaussian curve to the profile of LIF versus ion velocity at each time step. Examples of such fits are shown in Fig. 3. The parameters of the fitted curves then yield the total ion drift velocity $v_{\text{drift}}$ and ion temperature $T_i$ at each time step. The average ion temperature measured on each plane is shown in Table III. There was no observable ion heating due to the Alfvén waves.

To distinguish the ion drifts due to the waves from random fluctuations, correlation techniques were employed. First the cross correlation between the filtered $v_{d,\text{wave}}$ signal and the antenna signal was used to synchronize those two signals in time. Then, using a sliding window four wave periods in width, a cross correlation coefficient was computed at each time step for the synchronized $v_{d,\text{wave}}$ and antenna signals. The correlation coefficients from each time step were then turned into a weighting function with values ranging from zero to one, indicating the probability that the two signals were related. Finally, the filtered $v_{d,\text{wave}}$ signal was multiplied by the weighting function to attenuate the uncorrelated noise. A series of plots of the resulting measured ion drifts due to the waves on plane 1 is shown in Fig. 9.

The drifts measured on the other planes were similar to those of plane 1 but smaller in amplitude and with more random scatter in the orientation of the vectors as the waves decayed.

To quantify the uncertainty in the drift measurements the
value of $v_{d,\text{wave}}$ in proportion to $B_{\text{wave}}$ was examined at one wave period during the most steady-state period of wave propagation. The average and standard deviation of $v_{d,\text{wave}}$ for those time steps was computed at each position. These, in turn, were averaged over the four positions in the LIF measurement plane which best represented the center of the wave field. The values thus computed for each measurement plane are summarized in Table IV.

Figure 10 shows the comparison of the experimentally measured value of $v_\perp$ with the theoretical value for all four measurement planes. The error bars on $v_\perp$ are based on uncertainties in $B_{\text{wave}}$ and $k_\perp$, and represent the largest experimental uncertainty. The predicted values are derived from the plasma dispersion relation in the kinetic limit and Ampère’s law.

The ion drift formulas for finite frequency ($\omega \sim \omega_i$) are

$$v_{\perp} = \frac{c}{B_0} \frac{E_\perp \times B_0}{(1 - \omega^2)}$$  \hspace{1cm} (4)

$$v_p = \frac{c}{\omega_c B_0 (1 - \omega^2)} \frac{\partial E_\perp}{\partial t}.$$  \hspace{1cm} (5)

Table III. The average ion temperature measured with LIF on each plane. $\Delta z$ is measured from the antenna (at the far end of the machine); the warmer temperatures at larger $\Delta z$ values are closer to the source region.

<table>
<thead>
<tr>
<th>$\Delta z$ (cm)</th>
<th>$T_i$ (eV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>66</td>
<td>0.77</td>
</tr>
<tr>
<td>223</td>
<td>0.89</td>
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<td>349</td>
<td>1.04</td>
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<tr>
<td>475</td>
<td>1.10</td>
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</table>

The two drifts have comparable magnitude (the polarization drift is smaller by a factor of $\alpha$) but can be distinguished by their phase relationship to $B_{\text{wave}}$. The two drifts have a $90^\circ$ phase difference---$v_{E \times B_0}$ is maximum when $E_\perp$ and $B_{\text{wave}}$ are maximum, while $v_p$ is maximum when $E_\perp$ and $B_{\text{wave}}$ are passing through zero. At the axial position where $v_{E \times B_0}$ is maximum, it is antiparallel to $B_{\text{wave}}$ at the midpoint between the axial current channels. From this point at a distance of $\lambda/4$ farther along $z$, $v_p$ is maximum and carries polarization current across $B_0$ from one axial current channel to the other. Thus, the orientation of the $v_{d,\text{wave}}$ vectors also helps to distinguish the two drifts, in conjunction with their temporal relationship to $B_{\text{wave}}$.

Figure 11 shows $B_{\text{wave}}$ and $v_{d,\text{wave}}$ measured at two different moments in time on plane 1. In the upper frame at $t = 265$ $\mu$s, $B_{\text{wave}}$ has just passed through a minimum and $v_{d,\text{wave}}$ is directed from one current channel to the other—an example of polarization drift. In the lower frame, approximately one-quarter wave period later at $t = 271$ $\mu$s, $B_{\text{wave}}$ is maximum and $v_{d,\text{wave}}$ is antiparallel to $B_{\text{wave}}$. The drift in that case is $E \times B_0$.

From the dispersion relation and Ampère’s law the ratio of the shear wave perpendicular electric to magnetic field in the kinetic limit is

$$\frac{E_\perp}{B_{\text{wave}}} = \frac{v_A}{c} \frac{1 - \omega^2}{(1 - \omega^2 + k_\perp^2 \rho_i^2)^{1/2}}.$$  \hspace{1cm} (6)

Comparison of the theoretical and experimental values of $E_\perp$ is shown in Table V.

By substituting $E_\perp$ in Eq. (4) the $E \times B_0$ drift may be written as

FIG. 8. (Color). The axial wave current density $j_z$ computed from $B_{\text{wave}}$ (at a moment in time when $B_{\text{wave}}$ was maximum) on measurement plane 2.
FIG. 9. (Color). A time series showing $v_{d,\text{wave}}$ vectors (pink) overlaid on $B_{\text{wave}}$ at 4 $\mu$s intervals on measurement plane 1. The largest drift measured on plane 1 was $v_{d,\text{wave}} = 2.1 \times 10^4$ cm/s, which is $\approx 11\%$ of $v_T$ for 0.77 eV ions.
The expression for $v_{E \times B_0}$ considers only the dominant $k_\perp$ mode and strictly speaking is the Fourier component of the velocity for this $k_\perp$. The wave number spectra does have a single mode, broad enough to yield a 50% uncertainty in $k_\perp$. The sensitive dependence on $k_\perp$ in the denominator of Eq. (7) results in a 50% uncertainty in the theoretical prediction when the experimental values are inserted. With this in mind the curves in Fig. 10 are in agreement. The measured perpendicular ion drift velocity on plane 1 is compared to the theoretical estimate of the polarization and $E \times B_0$ drift velocity in Fig. 12. The ion drift should oscillate in time between the $E \times B_0$ and polarization drifts. In Fig. 12 the maximum $E \times B_0$ agrees well with the prediction but not with the polarization drift. The polarization drift for an $m=0$ mode is expected to be largest at the moment that the wave magnetic field goes through zero. In the data the magnetic field never has a null between the current channels. Figure 11, top panel is a snapshot at a time where the wave field should be zero, yet this is clearly not the case. The polarization drift closest to the midpoint of the current channels at that time points in the correct direction, i.e., polarization current flows from one to the other. The pattern should not rotate for an $m=0$ mode. This implies the existence of higher order $m$ numbers not included in the theory. The amplitude as a function of time of the $v_{E \times B_0}$ is reasonable.

To make a precise prediction of the measured quantities in this experiment one must use a kinetic theory with the inclusion of collisions in conjunction with a model of the antenna including the near field. This is beyond the scope of

### Table IV.

<table>
<thead>
<tr>
<th>$\Delta z$ (cm)</th>
<th>Average $v_{d,\text{wave}}$ (cm/s)</th>
<th>Standard dev $\sigma_{vd}$ (cm/s)</th>
<th>Ratio $\sigma_{vd}/v_{d,\text{wave}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>66</td>
<td>$1.33 \times 10^4$</td>
<td>$1.88 \times 10^3$</td>
<td>14%</td>
</tr>
<tr>
<td>223</td>
<td>$8.24 \times 10^3$</td>
<td>$1.84 \times 10^3$</td>
<td>22%</td>
</tr>
<tr>
<td>349</td>
<td>$4.91 \times 10^3$</td>
<td>$1.66 \times 10^3$</td>
<td>34%</td>
</tr>
<tr>
<td>475</td>
<td>$6.19 \times 10^3$</td>
<td>$3.43 \times 10^3$</td>
<td>55%</td>
</tr>
</tbody>
</table>

\[
v_{E \times B_0} = \left( \frac{B_{\text{wave}}}{B_0} \right) \frac{v_A}{(1 - \omega^2 + k_\perp^2 \beta_i^2)_{1/2}} (-\hat{\theta}).
\]

FIG. 10. (Color) Axial dependence of $v_\perp$ taken from averages from the center of each plane. The red curve for $v_{l,\text{exp}}$ is the experimentally derived value, and for comparison, in blue for $v_{l,\text{theory}}$, see formula [Eq (7)]. The dashed curves are the error of $v_{l,\text{exp}}$. They are large because of the uncertainty in $\lambda_\perp$.

FIG. 11. (Color) Polarization drift on measurement plane 1 at $t=265$ $\mu$s (upper frame) and $E \times B_0$ drift on the same plane at $t=271$ $\mu$s (lower frame), approximately one-quarter wave period later. The pink $v_{d,\text{wave}}$ drift vectors are overlaid on $B_{\text{wave}}$. The wave is traveling into the page. The background magnetic field points into the page.
TABLE V. The average experimental perpendicular electric field \( E_\perp \) and standard deviation of the experimental values on the four measurement planes. The error in the theoretical estimate is 50% due to the uncertainty in \( k_\perp \). The values of electric field are in mV/cm.

<table>
<thead>
<tr>
<th>( \Delta z ) (cm)</th>
<th>( E_\perp ) ( \text{expt} )</th>
<th>( \sigma_\text{expt} )</th>
<th>( E_\perp ) ( \text{theory} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>66</td>
<td>75.9</td>
<td>10.3</td>
<td>90.3</td>
</tr>
<tr>
<td>223</td>
<td>47.1</td>
<td>9.0</td>
<td>67.7</td>
</tr>
<tr>
<td>349</td>
<td>28.0</td>
<td>8.3</td>
<td>56.8</td>
</tr>
<tr>
<td>475</td>
<td>19.6</td>
<td>22.1</td>
<td>52.4</td>
</tr>
</tbody>
</table>

The parallel current density \( j_z \) can be derived from data such as those in Fig. 8. Since the derived parallel electric field was spatially averaged, the axial current was averaged as well. The current \( j_i(r) \) in each of the two channels of Fig. 8 was fitted to a parabola with a base 3 cm in radius, and the average current density in each was evaluated \( (j_i=5.6 \times 10^{-2} \text{ A/cm}^2) \). The current parallel to the magnetic field in a shear Alfvén wave is carried by electrons and the cross-field current by the ion polarization drift. These currents must be continuous, \( j_A \parallel = j_i A \perp \). Here \( j_i \) is the current density in either channel. \( A \parallel \) has a normal along the background field and is the region in Fig. 8 in either of the magnetic O points. \( A \perp \) can be approximated by a rectangle with one field-aligned side, \( \delta z \), and another, \( L \), in the transverse plane. In the case of Fig. 8, \( L \) is a line between the two current channels parallel to the magnetic field vectors. This is the line that the polarization current \( j_\perp \) crosses. We estimate \( L=14 \text{ cm} \). From continuity, we estimate that \( \delta z \approx 1 \text{ m} \) or about 1/8 of the parallel wavelength. The polarization current closes in the range of spatial locations where \( B_{\text{wave}} \approx 0 \) and the ion drift is directed from one current channel to another, as at \( t=265 \mu \text{s} \) in Fig. 11.

VII. SUMMARY AND CONCLUSIONS

This is a basic plasma physics experiment on the ion motion, electric fields, and current systems of a shear Alfvén wave. While not reporting anything unanticipated, we measure, for the first time to the authors’ knowledge, the \( Y_{\text{E} \times B_0} \) and ion polarization drifts of an Alfvén wave using laser-induced fluorescence. The drifts were measured at different times during the wave cycle and the measurements agreed very well with kinetic theory. The cross-field current of a shear Alfvén wave with finite \( k_\parallel \) is carried by the ions and limits the final wave amplitude since the parallel electron mobility is very large in most cases. For example, in this experiment the cross-field ion drift was about 14% of the ion thermal speed. The wave magnetic fields were measured and the parallel wave current derived from them was found to be consistent with the cross-field current. The components of the perpendicular wave electric field were calculated from the drift velocity. The parallel wave electric field was evaluated using the dispersion relation and the value for \( E_\perp \) and was found to be very small \((\approx 0.41 \text{ mV/cm})\). Measurement of a predicted 0.4 V change in electron energy over the 10 m plasma column was not attempted.

Parallel wave electric fields are important for electron acceleration in the auroral ionosphere where parallel wave-
lengths are of order 100 km. Since both density and magnetic field vary along the wave trajectory effects do not cancel over a wave cycle. The wave behavior in nonuniform media is not obvious. For example, a calculation of the propagation of a large transverse scale Alfvén wave with a density amplitude is large enough, ponderomotive effects appear.39 Nonlinear effects were not observed in this experiment but could be expected if the wave fields were ten times larger. Although the Alfvén waves are a fundamental mode in plasmas, their interaction with the plasma they travel through is subtle to say the least. The full implications of their behavior in space, astrophysical and thermonuclear plasmas is just beginning to emerge.

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