Nonlinear electron magnetohydrodynamics physics. III.
Electron energization

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Wave-particle interactions of low-frequency whistler modes with wave magnetic fields exceeding the ambient field are investigated experimentally. These highly nonlinear modes are excited with magnetic loop antennas in a large magnetized afterglow plasma. While the nonlinear wave properties are described elsewhere, the present paper focuses on the modification of the electron distribution function by the whistler waves. When the electron current flows in regions of magnetic nulls, such as in spheromak and field-reversed configurations (FRCs), strong electron energization is observed. When the whistler modes are created by electron Hall currents, such as in whistler mirrors, no significant energization occurs. The electron temperature can be raised locally by an order of magnitude. Non-Maxwellian distributions with energetic tail electrons are observed. Electron energization to ≳10 eV produces visible light emission whose time and space dependence is mapped. The light source travels with the subthermal speed of whistler spheromaks. When counterpropagating spheromaks collide, the resultant FRC produces strong local heating and light which dissipates its free magnetic energy. © 2008 American Institute of Physics.

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I. INTRODUCTION

The resonant interaction of electrons with whistler waves has been studied for a long time.1,2 Two wave-particle resonances can occur: Whistlers propagating oblique to the background magnetic field $\mathbf{B}_0$ have parallel electric fields which can give rise to Landau resonance when the electron velocity matches the parallel phase velocity; i.e., $v = \omega / k_0$. These resonant electrons can gain energy via Landau damping or give energy to the wave if the distribution function has a positive slope at the phase velocity.

A second resonant interaction arises from cyclotron resonance even though the wave frequency is below the cyclotron frequency $\omega_c$. It occurs when electrons move along $\mathbf{B}_0$ opposite to the wave vector $\mathbf{k}$, so as to experience a Doppler-upshifted frequency; i.e., $\omega - k \cdot v = \omega_c$. These resonant electrons gain perpendicular energy from the wave or transfer energy to the wave if the electron distribution has excess perpendicular energy.

In large amplitude waves, stochastic rather than resonant wave-particle interactions are dominant.3–5 In the present case, where the wave magnetic field exceeds the ambient field, one has to start from first principles and analyze local heating processes within the nonuniform field. In general, electron heating can be expected in regions of large electric fields along magnetic fields or magnetic null lines. Consequently, these will be regions of large current densities. Hall electric fields ($\mathbf{E} \perp \mathbf{B}$) or currents will not energize electrons. Not every magnetic null point topology produces current sheets, reconnection and dissipation. For the present field topologies of whistler spheromaks and mirrors,6,7,8 the dominant heating occurs in a toroidal $O$-type null line with induced toroidal electric field. Electrons are directly accelerated to high energies, scattered by collisions and fluctuations, and perform nonadiabatic motions adjacent to the null line. Comparatively little electron heating is observed in large-amplitude whistlers without null lines.

II. EXPERIMENTAL SETUP

The experimental device and the method for measuring magnetic fields have been presented in Part I.8 Here we only describe the diagnostic techniques for measuring the properties discussed in the present manuscript; i.e., the electron energization by strong whistler modes and related parameters.

A. Fast Langmuir probe diagnostics

Langmuir probes are a well-established diagnostic tool in plasma physics.10–12 Although the theory has been formulated for steady-state conditions,13 probes are often used in time-varying plasmas or with rapidly varying probe voltages. Two limitations arise under these conditions: First, the external probe circuit has to be designed to provide the fast time resolution. Second, an inherent problem arises when the time variations are faster than the establishment of the dc sheath which is of order of the ion plasma period. The current-voltage characteristics on such fast time scales differs from the dc current-voltage ($I$-$V$) curves, resulting in an incorrect evaluation of the plasma parameters. Well-known examples are the broadening of the $I$-$V$ characteristics in rf plasmas14 and current overshoots on rapidly swept or pulsed probes.15,16

The approach taken in rf plasmas is to let the probe potential oscillate with the plasma potential while varying the dc potential so as to extract only the dc plasma properties. In the present case, however, we are interested in the instantaneous plasma potential and other plasma parameters

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and have taken the following approach: A single planar Langmuir probe of area 1.5 mm$^2$ is mounted on the same shaft as the magnetic probe so as to measure simultaneously electromagnetic fields and plasma parameters in three-dimensional (3-D) space and time from repeated experiments. The radial probe shaft is perpendicular to the toroidal inductive electric field so that no voltages are induced on the transmission line to the probe tip. The probe is biased with a dc voltage. However, the voltage is applied via a transistor switch just prior to the excitation of whistler modes. This minimizes perturbations to the afterglow plasma density when drawing electron saturation current.\(^{17}\) The probe current is measured versus time on a properly terminated transmission line between the probe and a digital oscilloscope with as little as 1 ns time resolution. Simultaneously, the probe voltage is also recorded since it is slightly loaded by different probe currents. At a given probe position, the voltage is incremented from ion to electron saturation currents in typically 300 steps ($\Delta V = 0.5$ V). At each voltage, the experiment is repeated ten times and the data are averaged. Subsequently, the current-voltage characteristics are constructed at each time sample, and the plasma parameters (potential $\phi_{\text{pl}}$, electron temperature $kT_e$, and electron density $n_e$) are evaluated. Incrementing the probe voltage adds effectively a fourth dimension to the data set, making 3-D measurements of plasma parameters extremely time consuming and two-dimensional radial and axial scans more practical.

Although the probe voltage is held constant, rapid variations of the plasma potential produce the same effects as rapidly varying probe voltages in a dc plasma. In the latter case, a positive voltage ramp creates a current overshoot near the plasma potential where the sheath changes from ion-rich to electron-rich. Due to ion inertia, there is temporarily an ion excess above the plasma potential that allows larger electron currents to be collected. The same occurs for a probe with fixed bias when the local plasma potential rapidly decreases. Likewise, when the probe voltage sweep is reversed or the plasma potential rises rapidly, the collected current is reduced near the plasma potential compared to steady-state conditions. Experimentally, pronounced overshoot phenomena are observed on time scales of $t \approx 100$ ns $\approx 2\pi/\Omega_p$.

A fundamental limitation of Langmuir probes is their inability to resolve anisotropic distribution functions that can occur in collisionless plasmas in the presence of strong fields and currents. Although directional velocity analyzers do exist,\(^ {18}\) their use in complex magnetic and electric fields is limited. Non-Maxwellian distributions can be inferred from a non-constant slope of the $\log_{10} I$-$V$ characteristic, but one cannot obtain the angular dependence of, e.g., energetic electron tails.

**B. Light emission diagnostics**

Visible light emission from energetic electrons ($1/2m_e v^2 > 10$ eV) colliding with neutrals and ions\(^ {19}\) is measured with a photomultiplier setup schematically shown in Fig. 1(a). Simple collimating optics provides for spatial resolution transverse to the line of sight. The width at half-intensity is measured at $\Delta x \approx 2$ cm [Fig. 1(b)] using a movable point light source (light-emitting diode). The temporal resolution is obtained by pulsing the light source with a fast rising current. The observed rise and fall times of the optical system are $\Delta t \approx 50$ ns. The optical setup is mounted on a movable platform so as to scan radially or axially across the heated plasma.

The photomultiplier output is proportional to the light intensity, which has been verified with calibrated absorption filters. The measured light intensity depends on various factors such as the electron energy, velocity distribution function, density, dimensions along the line of sight for transparent plasmas, and emission spectrum relative to the spectral response of the detector. Thus, quantitative measurements of the electron properties are difficult to obtain except in special circumstances; for example, in a Maxwellian afterglow plasma. There, the light intensity $I_{\text{light}}$ has been plotted vs electron temperature measured with a Langmuir probe, which yields a nonlinear relation $I_{\text{light}}(kT_e)$. However, such a calibration may not be applicable to wave heating, which can produce runaway electrons. Thus, the main usefulness of the optical diagnostics is to verify the existence of energetic electrons and to determine the light source in space and time.

**III. EXPERIMENTAL RESULTS**

**A. Probe measurements of $kT_e$, $\phi_{\text{pl}}$, and $n_e$**

We first show basic Langmuir probe $\log_{10} I$-$V$ characteristics and their modifications by strong whistler modes. Prior to the excitation of waves ($t = -5$ $\mu$s in Fig. 2(a)), current is applied to the loop antenna at $t = 0$, the afterglow electron temperature is low (2 eV), uniform and almost constant on the whistler time scale. At $t = 1.2$ $\mu$s, the whistler mirror arrives, which increases the electron temperature to 5 eV and raises the plasma potential to +35 V on axis. It is followed at $t \approx 3.2$ $\mu$s by a whistler spheromak, which heats the electrons dramatically to 20 eV and lowers the plasma potential to $-25$ V. Even stronger heating is observed closer to the antenna where the whistler spheromak fields are more intense.
Figure 2(b) shows three $I$-$V$ traces at different radial positions cutting through the toroidal current layer of a spheromak at $z \approx 7\,\text{cm}$. Within the current layer ($y = 10\,\text{cm}$), one finds almost monoenergetic electrons with $kT_e \approx 30\,\text{eV}$. Adjacent to the current layer ($y = 14\,\text{cm}$) and near the center ($y = -2\,\text{cm}$), the bulk temperature drops to 5 eV but energetic electron tails still remain. Note that the Larmor radius of 30 eV electron in a field of 10 G is only 1.8 cm; hence, the energetic electrons are still highly magnetized and do not simply spiral out of the toroidal current layer.

Next we show in Fig. 3 the plasma parameters, evaluated from the Langmuir probe traces, versus radial position $y$ and time $t$ at $x=0$, $z=23\,\text{cm}$. For reference, the antenna current is shown in Fig. 3(a). The electron temperature [Fig. 3(b)] has been evaluated in the electron retardation region, i.e., just below the knee of the $I$-$V$ characteristics; hence, refers to the bulk electrons rather than the tail. The largest electron heating occurs in whistler spheromaks excited when $I_{\text{coil}} = 0$ with $dI_{\text{coil}}/dt > 0$. Since there is little heat confinement, the electron temperature is raised only in the presence of the spheromak and decays rapidly after it has propagated way. The first spheromak is stronger than the second one; hence, it heats more and extends over a larger radius.

Figure 3(b) shows that the change in plasma potential relative to the undisturbed plasma is positive in whistler mirrors and highly negative in whistler spheromaks. The sign of the space charge is produced by $E_{\text{rad}} \times B$ drifts, which carry electrons out of mirror fields and into spheromak fields. Alternatively, the pinched magnetic field lines of a mirror would like to straighten out and carry the frozen-in electrons axially outward, leaving an excess positive charge or positive plasma potential inside the mirror. Conversely, the bulged field lines in a spheromak push the electrons inward and create a negative charge/potential. In detail, the force balance is probably more complicated since the potential minimum is slightly delayed with respect to the heating maximum; i.e., toroidal current layer. There is also an inductive poloidal electric field due to the time-varying toroidal magnetic field. For a whistler mirror, the radially outward pointing space charge electric field, i.e., $E_{\text{rad}} = -d(\phi_{\text{pol}})/dy$, has the correct direction to produce the toroidal Hall current. The axial space charge electric force balances the outward magnetic force $-\mu \nabla B$.

The alternating sign of the space charge is not a nonlinear feature as it is also present in the linear regime. As the antenna current decays, whistler modes without null points are excited and they must also exhibit positive potentials when $B_{\text{z, wave}} \parallel B_0$ and negative potentials when
The density enhancement propagates radially outward with decreasing speed and amplitude. The former is due to the probe moving oblique to the wavefront that originates at the antenna. The latter is due to the geometric expansion of a density perturbation starting at the antenna. It has earlier been shown that pulsed currents in an insulated antenna produce a density depletion around the antenna wire.\(^{20}\) It is caused by a magnetic \(J \times B\) force on the electrons, which in electron magnetohydrodynamics (EMHD) produces a radial space charge electric field and accelerates ions together with electrons away from the antenna. The antenna wire is located at \(|y| = 7.5\) cm, \(z = 0\). The main perturbation arrives at \(t = 7\) \(\mu\)s after the start of the antenna current and must have traveled axially at a supersonic speed of \(v = \Delta z / \Delta t = 10^6\) cm/s; i.e., the ions initially travel ballistically with about 20 eV of energy. In time, this perturbation steepens into an ion acoustic shock with asymptotic velocity \(v = \Delta y / \Delta t = 2.9 \times 10^5\) cm/s, slightly larger than the ion sound speed \(c_s = (kT_e / m_i)^{1/2} = 2.2 \times 10^5\) cm/s in argon at \(kT_e \approx 2\) eV.

Figure 4(b) shows a contour plot of the absolute density vs radius and time. Both outward and inward propagating density perturbations are visible. Geometric effects enhance the inward propagating perturbation and decrease the outward perturbation. No comparable perturbations are seen at larger axial distances \((\Delta z = 23\) cm) from the antenna; hence, they are not produced by propagating whistler spheromaks.

**B. Plasma parameters in collisions of whistler spheromaks**

As shown in Fig. 8 of the companion paper Part II,\(^9\) the collision of two counterpropagating whistler spheromaks leads to a stationary field-reversed configuration (FRC), whose magnetic field energy dissipates within approximately half an rf period. The magnetic energy is converted into electron thermal energy, which leads to strong heating as confirmed by probe measurements. Any directional acceleration of electrons can be excluded since it would produce currents and measurable magnetic fields. Figure 5(a) shows a contour plot of the bulk electron temperature versus radial position and time approximately in the midplane between the antennas, where the FRC is formed and decays. The bulk electron temperature rises to more than 15 eV over a larger radius and longer time than observed for a single propagating spheromak [see Fig. 3(b)]. Two temperature peaks are observed on axis at different times. The first one is due to the arrival of the two spheromaks at the point of measurement. Given that both arrive almost simultaneously, both spheromaks merge leading to the canceling of their opposing poloidal fields. The second occurs when the resulting toroidal current ring decreases in radius and heats a smaller volume to a higher temperature. It is shown in Part IV that this leads to high frequency whistler instabilities.\(^{21}\) The second peak is definitely not due to the whistler mirror since it arrives later and drives much less heat than the FRC. The second FRC \((8 < t < 10\) \(\mu\)s) is weaker than the first one since the antenna current decays.

Figure 5(b) shows the plasma potential change relative
and the magnetic field is pinched in the center and reverses axial field-aligned currents are driven. Prior to the collision, negative inside the FRC.

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to the unperturbed plasma, i.e., \( \Delta \phi_{pl}(0, y, -23 \text{ cm}, t) \), for the above conditions. Positive potentials are observed in colliding whistler mirrors, negative potentials in colliding spheromaks. As in propagating spheromaks, the temperature maxi- mum and potential minima in the stationary FRC do not exactly coincide. The second potential minimum coincides with the collapse of the toroidal null line on axis forming a minimum \( B \). On axis, the inward pointing magnetic force \(-\mu \nabla B\) is balanced by an outward force \( e \nabla \phi_{pl}\) such that no axial field-aligned currents are driven. Prior to the collision, the magnetic field is pinched in the center and reverses \( \mu \nabla B \) and \( \nabla \phi_{pl}\). Radial gradients in pressure and potential produce cross-field drifts or currents consistent with the reversed magnetic fields. Thus, space charge fields develop self-consistently with magnetic topologies. Potential wells or hills are not the source for closed, field-aligned currents. The latter are driven by rotational electric fields that also determine the dissipation of magnetic energy. The former are necessary to satisfy the force equilibrium or generalized Ohm’s law.

Figure 5(c) extends the time scale of Fig. 5(b) to a time when current no longer flows on the antenna and, therefore, heating and null points are no longer created. The current waveform is the same as the one displayed on Fig. 4(a). The sign of the plasma potential is the same as in the nonlinear regime; i.e., positive when \( B_{wave} \parallel B_0 \) and negative when \( B_{wave} \parallel -B_0 \) relative to the ambient potential. Thus, the net field reversal by the wave field does not change the sign of the space charge.

Density perturbations by colliding nonlinear whistler modes are again negligible in the midplane on short time scales [see Fig. 3(d)], but are noticeable on longer times and near the antenna as shown in Fig. 4(b). The lack of initial density perturbations confirms that the nonlinear properties of the first few spheromaks are not due to modulational instabilities but to wave magnetic field nonlinearities.

C. Light emission from whistler spheromaks and mirrors

In the present parameter regime, light emission arises from inelastic collisions of energetic (>10 eV) electrons with neutrals (density \( n_0 \approx 9 \times 10^{12} \text{ cm}^{-3} \)) and ions (density \( n_i \approx 10^{12} \text{ cm}^{-3} \)).25 There is little light emission in the afterglow plasma at 150 \( \mu \text{s} \) after the discharge voltage is turned off. The excitation of linear whistler modes produces no light. The excitation of nonlinear whistler modes with \( B_{z, wave} \parallel B_0 \) and \( B_{z, wave} \parallel -B_0 \) produces some light near the antenna, but copious light emission is seen within the plasma volume for the opposite polarity; i.e., \( B_{z, wave} \parallel -B_0 \).

When the antenna axis is rotated by an angle \( \theta \) with respect to \( B_0 \), the light emission decreases. At \( \theta = 90^\circ \), no light is observed in the plasma volume (\( \Delta z = 25 \text{ cm} \)) for either polarity, although weak light emission occurs directly at the antenna at every half-cycle of the rf current.

Figure 6 shows the temporal behavior of light emission at a fixed axial distance (\( \Delta z = 25 \text{ cm} \)) from the antenna with a radial line of sight through the center of the plasma where
whistler modes propagate. Light pulses are excited from whistler spheromaks, which are launched during coil current reversal [top trace in Fig. 6(b)] with $dI_{\text{coil}}/dt>0$ and propagate at $v_i \approx 20 \, \text{cm}/\mu\text{s}$. Growth and decay of the light pulse are comparable and much faster than the light decay of the afterglow plasma ($\tau = \tau_{\text{light}}/[dI_{\text{light}}/dt] \approx 180 \, \mu\text{s}$, not shown here). On a logarithmic scale [Fig. 6(b)], the light decay shows a second time scale of $\tau = 12 \, \mu\text{s}$ after the passage of several spheromaks. This slower light decay reflects the decay of heat deposited by the spheromaks in the background plasma. The duration of the light pulses is comparable to the transit time of the spheromak through the line of sight. Thus, electrons are energized only in the spheromak and “lose” their energy rapidly thereafter. This could occur due to several reasons: The energetic electrons could stay within the toroidal current ring and travel axially with it, an unlikely scenario since the electrons would have to be trapped in the spheromak, which travels slowly compared to the electron thermal velocity. Another possibility could be that the electrons are energized unidirectionally so as to carry the toroidal current. As they travel out of the current layer, they could transfer their energy back to the fields. This is also unlikely since the peak current density [see Fig. 8(c) in Part I\textsuperscript{8}] is too small for a bulk drift at $>10 \, \text{eV}$. It would also imply conservation of magnetic energy which is not the case. The most likely scenario is that the energized electrons simply convect out of the current layer and are replaced by cold electrons. This lowers the light intensity for two reasons: (i) The density of hot electrons decreases as they expand along poloidal field lines into a larger volume (except toward the center of the spheromak where the temperature often peaks) and (ii) the electrons transfer energy by collisions with the mass of colder surrounding electrons as well as by inelastic collisions with neutrals and ions; e.g., by excitation of light. In this situation, the electrons provide for a true dissipation channel for the wave magnetic energy which is observed to decay. The wave damping is much stronger than by elastic electron-neutral collisions which dominate in whistler mirrors or linear whistler waves. In whistler spheromaks or FRCs, the direct acceleration of electrons effectively produces an “inertial” resistivity that is not present in other topologies.

Figure 7 shows the light intensity versus time at different axial distances from the antenna ($z=0$). Light emission from the antenna region is the largest. It also shows a peak when whistler mirrors are excited ($I_{\text{coil}}\approx 0$ with $dI_{\text{coil}}/dt>0$). Since the light from mirrors disappears at $z>25 \, \text{cm}$ from the antenna, it is not the result of wave currents but by electron acceleration in the toroidal magnetic null line at $z=0$ just outside the coil [see Fig. 6(c) in Part I\textsuperscript{8}]. With increasing axial distance the main light peak is delayed and decays similar to the whistler spheromak intensity. The axial propagation speed $v_i \approx 30 \, \text{cm}/\mu\text{s}$ is much smaller than the thermal speed of the $10 \, \text{eV}$ electrons producing light (180 cm/µs). Thus, the source of the light is the propagating whistler spheromak, not the thermal expansion of an initial heat pulse at the antenna. Heat transport causes the light pulse to spread out with distance.

The radial temperature profile can be inferred from axial light measurements that provide radial spatial resolution but are axially averaged. Figure 8(a) shows antenna current and light emission from the center of the plasma ($r=0$) and from the coil wire ($r_{\text{coil}}\approx 7.5 \, \text{cm}$) versus time. Figure 8(b) shows the light emission versus radius at different times within the light pulse. Initially, the light peaks near the coil, where the toroidal current rings of the two spheromaks originate. The current layer peaks in the toroidal magnetic null region, whose radius decreases as the total current decays. Thus, the light peaks move toward the axis and the hollow light/temperature profile disappears. Note that the light pulse is much shorter than the long rf period chosen here. The relation of spheromak duration and rf period has been briefly discussed in Part II\textsuperscript{9} but is expanded below.

Simultaneous probe and light measurements confirm the coincidence between light, heating, and the presence of the spheromak. For the current waveform shown in Fig. 8(a), the time dependence of the light emission, viewed radially at $\Delta z=25 \, \text{cm}$ from the antenna, is shown in Fig. 9(a) on an
after leaving the wave fields. The plasma potential change only energized within the spheromak but lose their energy light and temperature profiles indicate that the electrons are expanded time scale. Figure 9(b) shows the radial profile of the electron temperature from Langmuir probe measurements, i.e., \( kT_e(x=0,y,z=25 \text{ cm},t) \); Fig. 9(c) displays the toroidal current density \( J_t(x=0,y,z=25 \text{ cm},t) \), indicating the location of the spheromak. (d) Plasma potential profile \( \phi_p(x=0,y,z=25 \text{ cm},t) \).

FIG. 9. (Color online) Temporal coincidence of light and plasma parameters in whistler spheromaks. Antenna current as in Fig. 8(a). (a) Light intensity vs time at \( \Delta z=25 \text{ cm} \) from the antenna along a radial line of sight. (b) Radial bulk electron temperature profile \( kT_e(x=0,y,z=25 \text{ cm},t) \). (c) Toroidal current density \( J_t(x=0,y,z=25 \text{ cm},t) \), indicating the location of the spheromak. (d) Plasma potential profile \( \phi_p(x=0,y,z=25 \text{ cm},t) \).

...toroidal current density \( J_t(x=0,y,z=25 \text{ cm},t) \), but must be forced into the spheromak by magnetic force. Although an ideal spheromak has a force-free field \( \mathbf{J} \times \mathbf{B}=0 \), the actual whistler spheromak must exert a small inward magnetic force to balance the electron pressure gradient and to create a small excess electron charge. The \( \mathbf{J} \times \mathbf{B} \) force points radially inward irrespective of propagation direction and wave intensity. The axial wave propagation in the presence of a radial electron drift may cause the lag between the potential minimum relative to the spheromak center.

D. Parameter dependence of light emission

By varying the parameters of the plasma and of the applied field, we have investigated what determines the spheromak strength and size. The spheromak current is driven by the applied electric field near the antenna; i.e., \( E_y = V_{\text{coil}}/(2\pi r_{\text{coil}}) \times dI_{\text{coil}}/dt \). For a sinusoidal current, \( (dI_{\text{coil}}/dt)_{\text{max}} = \omega (I_{\text{coil}})_{\text{max}} \), such that for \( \omega = \text{const} \), one expects \( I_{\text{plasma}} \propto I_{\text{coil}} \). Figure 10(a) shows the intensity and half-width of the light pulse observed radially at 25 cm from the antenna versus peak coil current. The current has to exceed a threshold before field reversal and spheromak formation are possible. Thereafter, the light intensity increases rapidly with \( I_{\text{coil}} \). The width of the light pulses also increases with \( I_{\text{coil}} \). As shown elsewhere (Fig. 8, Part II), the radius of the current ring increases with increasing plasma current while its axial propagation speed decreases (Fig. 3, Part II), both of which contribute to increasing the spheromak light intensity and duration.
versus $B_0$ and time at two different axial distances from the antenna. For purpose of comparison, the current waveform is also shown [Fig. 11(a)]. At $\Delta z=25$ cm [Fig. 11(b)], the light emission arises from propagating whistler spheromaks. For very low background magnetic fields ($B_0<3$ G), the total field is essentially that of the antenna, which does not have a spheromak topology. For large background fields ($B_0>20$ G) and a given coil/plasma current, the whistler mode cannot reverse the external field to form a spheromak. Thus, there is an optimum field ($B_0=10$ G) for the most efficient heating and light emission. The graph also shows that with decreasing $B_0$ or increasing $B_{\text{wave}}/B_0$, the propagation speed decreases and the half-width increases, as earlier shown in Figs. 10(a) and (3), Part II. However, as $B_0 \to 0$, the delay is better described by field diffusion than convection by a wave.

Figure 11(c) shows the light emission at the antenna, $z=0$. In addition to the strong spheromak light, one can also see light emission when whistler mirrors are emitted. Since the light is only seen near the antenna, the energetic electrons are not produced in the wave field but in the magnetic null line just outside the antenna [see topology in Fig. 6(c), Part I]. Another interesting observation is that copious light is produced as $B_0 \to 0$ on every half-cycle of the applied current. Electrons are energized in the toroidal null line inside or outside the coil and confined by the coil magnetic field.

The production of fast electrons, measured by the light intensity, has also been investigated as a function of plasma and neutral densities. As expected, the light intensity increases with neutral gas pressure but surprisingly little with plasma density, which is varied with discharge current $I_{\text{dis}}$. The half-widths of the light pulses increase with ion density and gas pressure. The latter also changes the plasma density when $I_{\text{dis}}=\text{const}$.

We have also measured the light emission from colliding spheromaks and mirrors. Figure 12 shows light intensity vs time in the midplane between two identical loop antennas separated axially by $50$ cm. The loops can be displaced radially by a distance $\Delta r$ across $B_0$. The line of sight is through the center of the spheromaks along the direction of radial
displacement. For a head-on collision ($\Delta r=0$), the light intensity is the largest and lasts the longest ($\Delta t=7 \mu s$), similar to the electron temperature measured with probes [Fig. 5(a)]. During this time, the magnetic energy of the FRC is converted into electron thermal energy. For glancing collisions of spheromaks ($\Delta r>0$), the interaction weakens, the light pulses become shorter, and the interaction vanishes for ($\Delta r>2r_{coil}$). The light emission becomes the same as that of a propagating spheromak, which is governed by its transit time $\Delta t=\Delta z/v_z$.

Spheromak collisions are the strongest and most interesting wave-wave interactions. Mirror collisions produce negligible heating and mirror-spheromak collisions are dominated by spheromak heating.

Finally, it should be recalled that EMHD FRCs can also be produced by pulsed rather than oscillatory coil currents (see Fig. 3, Part I). In this case the decaying FRC also produces long lasting light emission as shown in Fig. 13, provided the coil field is opposite to the ambient field.

IV. CONCLUSIONS

This paper continues the experimental investigation of nonlinear whistler modes. Its focus is the energization of electrons in whistlers whose wave magnetic field exceeds the ambient field. The plasma parameters are directly measured with Langmuir probes and the presence of energetic electrons is indirectly inferred from light emission measurements. A newly discovered heating process differs from familiar damping processes such as Landau and cyclotron damping. The observed electron acceleration depends on the topology of the magnetic field, which in turn depends on field amplitude and direction. For a sinusoidal excitation, the heating process greatly varies from one half-cycle to the next. The strongest electron heating is found when the wave produces magnetic null lines such as in spheromak and FRC topologies. Electrons are accelerated along the null lines by an inductive electric field. The wave magnetic energy is converted into electron heat, which is a fundamental process of magnetic reconnection or annihilation. Non-Maxwellian electron distributions are produced. The scaling of light emission with plasma parameters, antenna size and orientation has been explored. Acceleration processes near the antenna have been explored, such as expulsion of ions by space charge electric fields and the acceleration of electrons in null lines of the antenna near zone field. These results are of intrinsic interest in nonlinear wave physics and of practical importance for efficiently exciting whistler modes from magnetic antennas.

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