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THE CIV PROCESSES IN THE CRIT EXPERIMENTS

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Abstract. Recent in situ observations during Critical Ionization Velocity (CIV) experiments under the CRIT project produced puzzling results, apparently at odds with current CIV theories. It is shown that incorporation in the theory of lower hybrid wave (LHW) instabilities of strong turbulence and finite size effects results in reconciliation of the theory with the observations.

Introduction

The critical ionization velocity (CIV) effect is thought to occur when neutral gas streams through a magnetized plasma, with cross field drift energy exceeding the ionization energy of the gas [Alfven, 1954]. Laboratory experiments [Fahleson, 1961; Angerth et al., 1962; Danielson, 1970; Eselevich and Fainshtein, 1986] and theoretical analysis [Galeev, 1981; Formisano et al., 1982; Papadopoulos, 1983, 1984; Raadu, 1978; Abe and Machida, 1985; Goertz et al., 1985; Machida and Goertz, 1986, 1988; Mobius et al., 1987] indicate that a plasma instability driven by the streaming newly ionized neutrals, heats the plasma electrons to energies in excess of the ionization potential on collisionless timescales. The energetic electrons colliding with the neutrals further ionize the gas, inducing a self supporting discharge. The instability responsible for the efficient energy transfer from the streaming ionized gas molecules to electrons has been identified as the lower hybrid instability.

While laboratory experiments and theory seem to be in relatively good agreement with the above physical description, this is not the case for space tests of the CIV hypothesis [Haerendel, 1982; Newell and Torbert, 1985; Wescott et al., 1986a,b; Kelley et al., 1986]. A review of the observations and the discrepancies with the conventional picture can be found in Newell [1985] and Torbert [1988]. Outstanding issues concern the primary instability resulting in electron energization, the ionization rates, the scaling with the ambient electron density, and the CIV triggering process. The recent CRIT I and II sounding rocket experiment experiments [Swenson et al., 1990; Stenbaek-Nielsen et al., 1990; Swenson et al., 1991; Brenning et al., 1991a,b; Kelley et al., 1991] have been unique in that they provided a good complement of in situ diagnostics, including complete vector electric field measurements for frequencies up to 12 kHz [Swenson et al., 1990, 1991]. The purpose of this letter is to reconcile the theoretical foundations of the CIV hypothesis to the space experiments. This letter is a qualitative preliminary presentation of the concepts. Comprehensive quantitative analysis is ongoing and will be presented elsewhere.

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CRIT II Experimental Evidence and Issues

A comprehensive presentation of the CRIT experiments and data can be found in the papers referenced above. We review below selective CRIT II observations taken at a distance approximately 1.7 km from the burst [Swenson et al., 1990], which highlight the issues relevant to this letter

1. The electrostatic wave measurements indicate first a weak activity at .095 secs after the release with a frequency in the vicinity 6–7 kHz. This was followed by the appearance of an oscillatory electric field growing with time. The frequency of the growing oscillations started at 1.6 kHz and gradually dropped to 300 Hz at about .12 secs. The electric field peaked at .12 secs with amplitude of about 400–600 mV/m. The wave structure was very bursty and the polarization almost isotropic. Following .13 secs the wave pattern changed with the waves having lower amplitude and less structure.
2. The ambient electron density was seen to rise from an ambient value of 5.4×10^5 cm^{-3} to a peak value of 2.8×10^6 cm^{-3}. The rise occurred .11 secs following the burst, while the maximum occurred at about .18 secs. The peak production rate was estimated at 5×10^7 cm^{-3} s^{-1}.
3. The interaction was seen to behave as an ionization front riding on the neutral barium with a spatial size of 600–700 m in the stream.

Important data concerning the DC electric and magnetic structures have been also presented in the CRIT I papers mentioned above. Of importance here is that in CRIT I there was no evidence for increase of the electron density past its ambient value which was only 4×10^4 #/cm^3. The wave spectra were similar to CRIT II, but the maximum amplitude was only 150 mV/m.

These measurements pose a series of theoretical questions, such as:

1. Why was the frequency of the observed waves variable and low? For low hybrid instabilities, the frequency should be close to either the barium or the oxygen lower hybrid frequencies, which for the CRIT II conditions were 2.3 and 6.9 kHz correspondingly.
2. Why was the wave polarization almost isotropic? According to linear theory the wave polarization in lower hybrid instabilities is strongly anisotropic, with the field aligned electric field smaller by more than an order of magnitude than the transverse component.
3. Can the electron energization be reconciled with the observed wave spectrum? If the wave spectrum is isotropic the phase velocity of the waves in the magnetic field direction would be too low to energize electrons to ionizing energies, at least within the context of quasi-

linear and resonance broadening theories [Galeev, 1981; Papadopoulos, 1983; Goertz et al., 1985].

4. What is the role of the ambient electron density in determining the CIV triggering conditions?

The Plasma Physics of the Interaction

An essential ingredient of almost all of the CIV theories [Galeev, 1981; Formisano et al., 1982; Papadopoulos, 1983, 1984; Abe and Machida, 1985; Goertz et al., 1985; Machida and Goertz, 1986, 1988; Mobius et al., 1987] is the lower hybrid instability driven by the streaming of an ion beam generated by the interaction of the neutral cloud with the ambient ionosphere. During the initial ignition stage, when the cloud density is relatively high, the ion beam can be generated either by charge exchange and charge stripping of the expanding barium cloud interacting with the ionospheric gas [Stenbaek-Nielson et al., 1990], or by collisional reflection and snowplowing of the ambient plasma [Papadopoulos, 1991]. In the first case the ion beam will be a barium ion beam moving with the neutral front speed, while in the second it will be an oxygen beam moving with speed faster than the neutral front. Following ignition, the interaction occurs between a barium ion beam and the ambient magnetoplasma.

The dispersion relation given by for such a system has been examined by several authors [Papadopoulos, 1983, 1984; Galeev, 1981; Mobius et al., 1987]. It incorporates two types of instabilities. For a monoenergetic beam moving with velocity U there is a coherent hydrodynamic instability. Of more interest here is the beam kinetic instability which has frequency and temporal growth given by

$$\omega_k^2 = \omega_{LH}^2 \left(1 + \frac{M_o}{m}\frac{k_z^2}{k^2}\right) \tag{1}$$

$$\gamma_k \approx \frac{\alpha\, \omega_{LH}}{\left(1 + \frac{k_z^2}{k^2}\frac{M_o}{m}\right)^{1/2}} \left(\frac{U}{\Delta v}\right)^2 \tag{2}$$

where ω_{LH} is the oxygen lower hybrid frequency, α the ratio fo the beam to ion plasma frequency, M_o the mass of O_2, and Δv is the velocity spread of the ion beam and the magnetic field $\underline{B} = B_o \hat{e}_z$. For an infinite, homogeneous plasma maximum growth occurs for wavenumbers satisfying

$$\frac{\omega_k}{k} = U \qquad \frac{M_o}{m}\frac{k_z^2}{k^2} \lesssim 1 \tag{3}$$

The above instability has been simulated in one and two dimensions using particle codes with periodic boundary conditions and sometimes including ionization and some form of collisional effects [Abe and Machida, 1985; Goertz et al., 1985; Machida and Goertz, 1986, 1988; Tanaka and Papadopoulos, 1984]. It has been demonstrated that the electrons gain energy in the direction parallel to magnetic field. In addition to particles simulations, the equations for the waves and the particles have been solved in the quasi-linear homogeneous infinite medium approximation [Galeev, 1981; Papadopoulos, 1983, 1984]. It is apparent that the wave properties expected on the basis of eqs. (1–3) and seen in the simulations and analytic solutions are at odds with the CRIT data. At this stage we can either discard the lower hybrid hypothesis, which has been experimentally verified in laboratory experiments and seems a natural one for CIV settings, or examine linear and nonlinear effects not included in the analysis or the simulations but present in space releases. We take the latter approach.

The expanding barium cloud in the CRIT and other experiments has a finite size. As a result the ion beam that drives the instability has a finite size transversely to the direction of propagation. Referring to eq. (3) we note that, linearly, the finite size of the beam is most important in the field aligned direction. To facilitate the discussion we assume that the beam propagates in the x-direction. If we denote by L_{zo} the size of the beam in the z-direction the infinite homogeneous analysis and the simulations using periodic boundary conditions require $k_z L_{zo} >> 1$. This condition is clearly not satisfied in the early time of the expansion where the CRIT measurements take place. For the CRIT II observations near 1.7 km, the field aligned length is of the order of 200 m and thus maximum growth occurs not at $\frac{k_z^2}{k^2}\frac{M_o}{m} << 1$, but at the first confined mode, i.e.

$$\pi / k_{zo} = L_{zo}. \tag{4}$$

Of profound significance is that, since the k_{zo} is fixed by the finite size the group velocity V_g of the excited waves should be taken subject to a constant value of k_z. It is given by

$$V_g = -U \frac{\omega_k^2 - \omega_{LH}^2}{\omega_k^2} \tag{5}$$

For frequencies above the lower hybrid frequency, the instability is a backward wave instability (i.e. the group and phase velocities are in opposite directions). The single wave field aligned structure and the backward wavepacket nature of the unstable waves play a critical role in the interpretation of the observations as well as in the evolution of the interaction. We will discuss them after addressing first the role of the electron nonlinearity, which was mostly neglected in previous CIV studies.

Quasi-linear theories neglect electron nonlinearities in all directions. Two dimensional simulations, while maintaining electron nonlinearities along the magnetic field, seriously underestimate the electron nonlinearity across the magnetic field by ignoring the induced electron drift in the third dimension (note that two dimensionality implies $k_y = 0$). This is seen by considering the coupled nonlinear system of high frequency ($\omega \gtrsim \omega_{LH}$), and low frequency ($\omega << \omega_{LH}$) perturbations for lower hybrid waves (LHW).

The equations for the complex amplitude of the lower hybrid frequency potential ϕ and the low frequency density fluctuations δn are [Shapiro and Shevchenko, 1984]

$$\begin{aligned} i\frac{\partial}{\partial t}\nabla^2\phi + \frac{3}{4}\omega_{LH} r_e^2 \nabla^4 \phi - \frac{M_o}{2m}\omega_{LH}\nabla_z^2\phi = \\ i\frac{M_o}{m}\frac{\omega_{LH}^2}{\Omega n_o}[\nabla n \times \nabla\phi]_z \end{aligned} \tag{6}$$

$$\frac{\partial^2}{\partial t^2}\delta n - \frac{T}{M_o}\nabla^2 \delta n = -i\frac{ecn_o}{\omega_{LH}}\frac{\nabla^2}{M_o B}[\nabla\phi \times \nabla\phi^*]_z \tag{7}$$

where r_e is the electron gyroradius T is the plasma temperature and ϕ^* is the complex conjugate of ϕ. In eq. (6), for the high frequency mode, the term causing the wave dispersion has the same form as in the linear theory. There is, however, an additional term associated with the slow density modulation on the right hand side. This modulation requires inclusion of the first order ExB drift of the electrons. It was noted by Musher and Sturman [1975] that the coupling of LHW with low frequency perturbations, becomes extremely important if $\delta n(\underline{r})$ is two dimensional in the plane perpendicular to the magnetic field (i.e. the x-y plane). This is because the electron drift $\underline{V}_D = c\frac{E\times B}{B^2}$ causes the high frequency electron density perturbation $\frac{\partial n_e}{\partial t} \sim c\frac{E_y}{B}\frac{\partial \delta n}{\partial x} + c\frac{E_x}{B}\frac{\partial \delta n}{\partial y}$. For a simple geometry with $k_y = 0$ (i.e. homogeneous in the y-direction) the only high frequency density perturbation is due to E_z, which is by a factor $\frac{\omega_k}{\Omega}$ smaller.

A detailed analysis of the nonlinear consequences of eqs. (6) and (7) is presently underway. We note, here, the similarity of eqs. (6) and (7) to those describing collapse of oblique Langmuir waves [Zakharov, 1984]. Similar considerations indicate that for values of the electric field E such that the cross field electron drift V_D is of the order of the sound speed C_s, i.e. $E \geq E_T \approx \lambda B\frac{C_s}{c}$ where λ is a factor of order (but smaller) unity, the ponderomotive force dominates leading to the formation of density cavities in which the LHW are trapped. This, of course, leads to short scalelengths and large localized electric fields predominantly in the field aligned direction. For $C_s \approx 10^5 cm/sec$ the threshold field electric E_T is about 10–20 mV/m.

Self-similar solutions indicate that during collapse [Sotnikov et al., 1979] the wave energy density W scales as

$$k_z(t) \sim \frac{W}{n_o T}, \quad k_\perp(t) \sim \sqrt{\frac{W}{n_o T}}$$
$$\frac{k_z(t)}{k_\perp^2(t)} \approx constant, \quad k_z^{-1}(t) \sim \frac{1}{t - t_o} \tag{8}$$

As a result the frequency $\omega_k(t)$ of a collapsing wavepacket in a reference frame where $V_g = 0$ will be given by

$$\omega_k(t) = \omega_{LH}\left[1 + \frac{M_o k_{zo}}{m k_o^2} k_z(t)\right]^{1/2} \tag{9}$$

where k_{zo} and k_o are the values corresponding to the linear mode, i.e.

$$k_o \approx \frac{\omega_k}{U}, \quad k_{zo} \approx \frac{\pi}{L_{zo}} \tag{10}$$

The energization of the electrons can be described by a Fokker-Planck equation similar to the one derived by Morales and Lee (1974) for Langmuir turbulence. For LHW the electron energization can be approximately described by

$$\frac{\partial f}{\partial t} = \frac{\partial}{\partial v_z} D(v_z) \frac{\partial f}{\partial v_z} - \mathcal{L}(f) \tag{11a}$$

$$D(v_z) = \begin{cases} 2\frac{e^2 E_z^2}{m^2 v_z^2}\frac{L_z^2}{\tau} & v_z > L_z\omega_k \\ 0 & v_z < L_z\omega_k \end{cases} \tag{11b}$$

where $\mathcal{L}(f)$ represents the inelastic losses and τ is the mean collision time of the electrons with the localized electric fields. The form of $D(v_z)$ has a simple physical interpretation. Slow electrons, with $v_z < L_z\omega_k$, simply quiver in the oscillating electric fields and gain no energy. Fast electrons, with $v_z > L_z\omega_k$, transit through the localized fields before the electric field changes sign. As a result they see an essentially static field, have almost straight orbits and, depending on the phase of the wave, are accelerated or deccelerated. On the average they gain energy from the wave due to a second order Fermi process [Morales and Lee, (1974); Manheimer and Papadopoulos (1975)]. As a result, wave spectra with spectral peak in the low phase velocity region can produce energetic tails.

Taking for the mean collision time $\tau \approx 1/N_s v_z$, where $N_s \approx k_{zo}$ is the linear density of the localized wavepackets, we find from eq. (11) that energization to an energy of ϵ_o requires a heating time

$$\tau_H(\epsilon_o) \approx \frac{1}{32\pi^2}\left(\frac{\epsilon_o}{\tilde{\epsilon}}\right)\left(\frac{M_o}{m}\right)\left(\frac{k_z}{k_{zo}}\right)^2\left(\frac{k_{zo}v_z}{\omega_{LH}}\right)\frac{1}{\omega_{LH}} \tag{12}$$

$$\tilde{\epsilon} \equiv \frac{1}{2}\frac{e^2 E^2}{m\,\Omega_e^2} \approx 1.6 \times 10^{-3}\left(\frac{E}{V/m}\right)^2 eV \tag{13}$$

For Ba the peak of the ionization rate is near 10 eV. The optimum ionization rate ν_i will be approximately twice the heating time to $\epsilon_o \approx 10eV$.

$$\nu_i = 6.5 \times 10^2\left(\frac{E}{V/m}\right)^2\left(\frac{k_z}{k_{zo}}\right)^2\left(\frac{\omega_{LH}}{k_{zo}v_z}\right)\omega_{LH}. \tag{14}$$

Reconciliation with Observations

We examine next which of the questions raised by the CRIT data can be resolved on the basis of the previous considerations. The frequency ω_o measured by a stationary probe in the ionosphere is given and is given by

$$\omega_o = \omega_k + \underline{k} \cdot \underline{V}_g \tag{15}$$

Using eq. (7) for $\underline{V}_g$ and $\omega_k = \underline{k} \cdot \underline{U}$ we find that

$$\frac{\omega_o(t)}{\omega_{LH}} = \frac{\omega_{LH}}{\omega_k(t)}. \tag{16}$$

This is conjunction with eq. (8) gives

$$\frac{\omega_o(t)}{\omega_{LH}} = \frac{1}{\left[1 + \frac{M_o k_{zo}^2}{m k_o^2}\frac{k_z(t)}{k_{zo}}\right]^{1/2}} \tag{17}$$

Equation (17) determines the observed variation of the frequency with time. At t= .095, $k_{zo} \approx 1.5 \times 10^{-2} m^{-1}$ and $k_o \approx .4\ m^{-1}$, giving $\omega_o(t = 0.95) \approx \omega_{LH}$ consistent with the observations. For later times $\omega_o(t)$ can be approximated by

$$\omega_o(t) \approx \omega_{LH}\left(\frac{k_{zo}}{k_z(t)}\right)^{1/2} \tag{18}$$

As a result a frequency of 1.6 kHz would correspond to $\frac{k_z}{k_{zo}} \approx 16$ or field aligned scale lengths of 12–15 m. The observed asymptotic frequency 300 Hz corresponds to field aligned lengths of $L_z \approx .5m$. In essence the nonlinear backward wavepacket hypothesis requires field aligned collapse to meter like structures. Notice also that electrons with values of $v_z > \frac{\omega_k}{k_z} \approx 10^7 cm/sec$ can be accelerated by the observed turbulence. Thus, the hypothesis of the backward propagating, nonlinearly collapsing wavepacket can resolve the questions the 1–3 posed in the introduction.

We utilize next eq. (14) in conjunction with eq. (18) and the observed values to examine the resultant ionization rates. From (14) and (18) we find

$$\nu_i = 6.5 \times 10^2 \left(\frac{E}{V/m}\right)^2 \left(\frac{\omega_o}{\omega_{LH}}\right)^2 \left(\frac{\omega_o}{k_{zo} v_z}\right) \omega_o \tag{19}$$

Taking $\omega_o \approx 2 \times 10^3 sec^{-1}$, $v_z \approx 10^8 cm/sec$ and $k_{zo} \approx 1.5 \times 10^{-2} m^{-1}$, we find that $\nu_i = 250(E/V/m)^2 sec^{-1}$.For CRIT II, $E \approx .5V/m$ and $n_o \approx 5 \times 10^5 \#/cm^3$, we find $\nu_i \approx 63\, sec^{-1}$ and a production rate $\nu_i n_o \approx 3 \times 10^7 \#/cm^3$ consistent with the experiments. For CRIT I, $E \approx .15V/m$, $n_o \approx 5 \times 10^4 \#/cm^3\, sec$, giving $\nu_i \approx 5\, sec^{-1}$ and $\nu_i n_o \approx 2.5 \times 10^5 \#/cm^3\, sec$. A production rate by two orders of magnitude weaker. For rocket experiments, assuming a particle loss rate n_{Ba}/Ω_{Ba}, where Ω_{Ba^+} the Ba^+ cyclotron frequency, the triggering condition for discharge is $\nu_i > \Omega_{Ba^+}$, which for $\Omega_{Ba} \approx 40\, sec^{-1}$, requires E > 300 mV/m.

The correlation of the ambient density n_o to the value of E, will be discussed elsewhere. We simply comment here that assuming that the wave energy density E^2 scales linearly with the free energy of the seed ions, i.e. $E^2 \sim n_s M_j U^2$,where n_s is the seed ion density, we expect $E \sim \sqrt{n_s}$. As long as n_s is produced by charge exchange or snowplowing $n_s \sim n_o$. As a result $E \sim \sqrt{n_o}$. The CRIT results are consistent with this scaling. In the context of the above theory the differences in the CRIT I and II experiment can be attributed to the lower free energy available for triggering at lower ambient densities.

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