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Lower Hybrid Waves Upstream of Comets and Their Implications for the Comet Halley "Bow Wave"

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Observed and theoretical features concerning the nature of so-called cometary "bow shocks" or "bow waves" are discussed. Collective plasma effects associated with the presence of pickup ring ions (protons and water ions) in the vicinity of the supermagnetosonic to submagnetosonic transition region in the quasi-perpendicular limit are considered; the linear and nonlinear evolution of instabilities around the lower hybrid frequency is emphasized. It is shown that lower hybrid waves can lead to heating and produce distributions with magnitudes in reasonable agreement with Giotto data. The implications to the existence and structure of cometary bow shocks are discussed.

1. INTRODUCTION

A most critical and controversial issue arising from the recent encounters of the VEGA, Giotto, Sakigake, and Suisei spacecraft with comet Halley, as well as that of the International Cometary Explorer (ICE) with comet Giacobini-Zinner (G-Z), is the nature of the strong momentum coupling between the solar wind and the cometary plasma often called the bow wave or bow shock [e.g., *Bame et al.*, 1986; *Johnstone et al.*, 1986]. The solar wind flow is decelerated from supermagnetosonic to submagnetosonic at a certain distance upstream from the comet (in the case of Halley, roughly 10^6 km [*Coates et al.*, 1987]). By analogy with the Earth and other planets, one might expect the transition to be accomplished by means of a collisionless bow shock. However, there are reasons for supposing that either a bow shock may not be present at all or, if it does exist, it may look radically different from other planetary bow shocks [*Wallis*, 1972; *Omidi et al.*, 1986], a hypothesis confirmed by the data. One reason for the uncertain structure of cometary "bow shocks" may be found in the unusual nature of a cometary magnetosphere.

As a comet approaches the Sun, material from the nucleus sublimates and moves into the surrounding solar wind plasma. The atoms (principally protons and members of the water group [*Mukai et al.*, 1986]) are ionized by either solar ultraviolet radiation or charge exchange with solar wind ions and can subsequently interact with the ambient solar wind plasma and magnetic fields. The initial phase of this interaction is referred to as solar wind "pickup" of cometary ions. The newly ionized particles are subjected to the solar wind motional electric field and gain energy, whereas the solar wind loses directed energy and is decelerated, a process referred to as solar wind "mass loading." The solar wind is thus continually decelerated as it approaches the comet, in the case of Halley over a distance of many million kilometers.

The exact form of the ion pickup process depends on the spiral angle (ψ) between the solar wind plasma and magnetic field directions, where ψ is defined as $\mathbf{V}_{SW} \cdot \mathbf{B}_{SW} = V_{SW} B_{SW} \cos \psi$. When $\psi \approx 0$, the cometary ions form a field-aligned beam with velocity V_{SW} in the solar wind reference frame, whereas when $\psi \approx \pi/2$, they form a ring distribution function in velocity space. For intermediate angles the distribution is a combination of the two. The resultant distribution functions are unstable to small perturbations, so that collective plasma effects play a significant role in the ultimate energy and momentum coupling of the cometary pickup ions and the solar wind plasma.

Solar wind mass loading can modify the nature of any bow shock or may result in an unshocked supersonic-subsonic flow transition [*Wallis*, 1972]. The continual mass loading implies that the solar wind is steadily decelerated over lengths much longer than a typical shock length scale (which for a low Mach number case will be of the order of a streaming gyroradius ZV_{SW}/Ω_i [*Omidi and Winske*, 1987]). The deceleration is, however, accomplished without an accompanying heating of the solar wind. The kinetic energy lost by the solar wind protons is stored in highly nonthermal velocity ring type distributions of cometary ions. These distributions are unstable to both low-frequency magnetohydrodynamic (MHD) instabilities [*Tsurutani and Smith*, 1986; *Omidi and Winske*, 1987] and electrostatic lower-hybrid waves as discussed in this paper. The possibility arises that these waves which are carried by the solar wind interact near the bow wave with the solar wind protons thereby raising their temperature without any appreciable momentum coupling. A supermagnetosonic to submagnetosonic transition can thus be accomplished through the equivalent of nonlocal viscosity without the mediation of a self-consistently created shock.

The present paper addresses one aspect relevant to the above, admittedly speculative, picture, namely, of the generation of electrostatic lower hybrid waves observed by the VEGA spacecraft [*Galeev et al.*, 1986] which appear first just ahead of the strong coupling region. The thesis of the paper is that the observations associated with these waves are consistent with waves expected by cometary ring distributions in a parameter region where the local magnetosonic Mach number becomes of the order of 1.5–2. Furthermore, these waves, contrary to lower hybrid waves excited by ion reflection in shocks [*Vaisberg et al.*, 1983], can accelerate a fraction of the cometary ions to energies several times higher than their nominal pickup energy. The above conclusions rely on a combination of nonlinear theory of instabilities around the lower hybrid frequency and the VEGA observations. In section 2 we discuss the basic theory (linear and quasi-linear) of such instabilities and compare our results with the plasma wave data obtained from VEGA 1. In section 3 we discuss the nonlinear relaxation of ring distributions and compare our results with the VEGA

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and Giotto observations. The paper concludes with some more speculative remarks on what the results imply for the nature of the supermagnetosonic to submagnetosonic transitions in cometary interactions.

2. Waves Near the Lower Hybrid Frequency

2.1. *Derivation of Threshold Conditions*

In this paper, only the quasi-perpendicular portion of the cometary magnetosphere is considered, so that the angle of relative motion of the comet with respect to the solar wind magnetic field is $>\pi/4$. The results are then applicable to the encounter of VEGA 1, but not of VEGA 2, with comet Halley [*Galeev et al.*, 1986]. In the quasi-perpendicular region the cometary pickup ions form a ring distribution in the solar wind frame with a small field-aligned beam component. The ring ions are assumed to be cold and to have velocity $u = V_{SW} \sin\psi$, where V_{SW} is the solar wind speed and ψ is defined in the introduction. Such ring distributions are susceptible to instabilities around the lower hybrid frequency. We will not be concerned with instabilities driven by the field-aligned beams which give rise to lower-frequency predominantly electromagnetic waves [*Winske and Gary*, 1986]. Other studies of lower hybrid waves of relevance to comets have been carried out by *Buti and Lakhina* [1987], *Sagdeev et al.* [1987], and *Coroniti et al.* [1986]. Buti and Lakhina examined the production of high-energy ions by stochastic acceleration with lower hybrid waves, and Coroniti et al. examined the data at comet G-Z and concluded that the magnetic fluctuations were in agreement with those theoretically expected.

The theory of lower hybrid waves in the presence of ring distributions has been investigated by *Akimoto et al.* [1985]. Under certain restrictions (discussed below) the dispersion relation in the lower hybrid range, $\Omega_i \ll \omega \ll \Omega_e$, evaluated in the solar wind reference frame is

$$1 - \frac{\omega_i^2}{\omega^2} - \frac{\sum_j \alpha_j \omega_i^2 \omega}{(\omega^2 - k_\perp^2 u_j^2)^{3/2}} + \frac{\omega_e^2}{\Omega_e^2}\left(1 + \frac{\omega_e^2}{k^2c^2(1+\beta_e)}\right) - \frac{\omega_e^2}{\omega^2}\frac{\cos^2\theta}{1 + (\omega_e^2/k^2c^2)} = 0 \qquad (1)$$

where ω_i (ω_e) and ω_i (Ω_e) are the proton (electron) plasma and cyclotron frequencies, respectively, and the wave vector $\mathbf{k}$ is assumed to be of the form $\mathbf{k} = (k_\perp, 0, k_z)$, where the ambient magnetic field lies in the positive z direction. The parameter α_j is defined as $\alpha_j = m_p n_j / m_j n_p$, where m_j (n_j) and m_p (n_p) are the masses (number densities) of the jth particle species in the ring and ambient solar wind plasma, respectively; θ is the angle between $\mathbf{k}$ and the ambient magnetic field $\mathbf{B}_0$, and the solar wind is assumed to be a pure proton-electron plasma. Equation (1) assumes that the ring and solar wind ions are cold and unmagnetized, whereas the electrons are cold and strongly magnetized. The cold electron assumption is the most crucial, implying that $\omega/k_z \gg v_e$ and results in considering only flute modes ($k_z \ll |k|$). In deriving (1) the displacement current has been neglected ($\omega/k \ll c$), and the ratio ω_e/Ω_e is assumed large.

The cold ring ion approximation also merits some discussion. Since the solar wind is continually decelerating as it approaches the comet, the velocity of the pickup ions is not constant but also decreases. An idealized distribution function of such a type is shown by *Sagdeev et al.* [1987, Figure 1]. It would thus appear that the ring should be treated as warm, i.e., of finite thermal spread, with regard to the instability. This is untrue. Since the lowest-velocity ring ions are those that have just been picked up, they are approximately cold with $\partial f/\partial v_\perp$ locally large. Hence waves resonating with these ions act as if they were a cold ring. The remainder of the ring acts as a weak damping term, which can generally be neglected.

For small α_j the ring is subjected to two classes of instabilities: resonant with $\omega \approx k_\perp u_j$ and nonresonant with $\omega \neq k_\perp u_j$. The nonresonant modes have small growth and are not of interest for the values of α_j present at comets. For resonant modes the cold electron restriction implies

$$M_A^2 \sin^2\psi \gg \frac{m_i}{m_e}\beta_e \cot^2\theta \qquad (2)$$

where $M_A = V_{SW}/V_A$ is the solar wind Alfvén Mach number and β_e is the electron plasma beta.

Solving (1) for the real and imaginary parts of the frequency ($\omega = \omega_r + i\gamma$) and assuming that all particle species have the same pickup velocity u_j, we find

$$\frac{\omega_r^2}{\omega_H^2} = \left[1 + \frac{(m_i/m_e)\cot^2\theta}{1 + (\omega_e^2/k^2c^2)}\right]\left[1 + \frac{\omega_e^2}{k^2c^2(1+\beta_e)}\right]^{-1} \qquad (3)$$

$$\gamma = \sum_j \frac{\omega_H \alpha_j^{2/5}}{2}\left[1 + \frac{(m_i/m_e)\cot^2\theta}{1 + (\omega_e^2/k^2c^2)}\right] \qquad (4)$$

k is determined from the resonance condition $\omega = k_\perp u_j$ and ω_H is the lower hybrid frequency, defined as $(\Omega_e\Omega_i)^{1/2}$ in the present parameter regime. The instability is of the familiar hydrodynamic type with strong coupling between the lower hybrid and ring waves. Equation (4) and the requirement that $\omega > \Omega_i$ can be used to derive a limit on the value of α needed to drive instabilities.

Consider first oblique modes where $(m_i/m_e)\cot^2\theta \gg 1$. In that case it can be shown that resonant instabilities require that

$$M_A^2 \sin^2\psi < \frac{m_i}{m_e}\frac{\cot^2\theta}{4} \qquad (5)$$

From (2) and (5) we find that electron damping can be neglected if $\beta_e < 1/4$; this condition gives undamped resonant oblique instabilities. In fact, *Akimoto et al.* [1985, Figures 3*a* and 3*b*] showed that electron Landau damping does not suppress the instability completely but makes $\gamma < \Omega_i$, necessitating the inclusion of magnetic field effects on the ion dynamics. In the cometary case the solar wind β_e is either $>\frac{1}{2}$ or close to it; we will therefore ignore these waves in the rest of the paper.

As θ approaches $\pi/2$, the nature of the ring instabilities undergoes a sharp change. The coupling is now of the ion-ion nature so that electron Landau damping is unimportant ($k_z \approx 0$). (A third domain also exists where $(m_i/m_e)\cos^2\theta \approx 1$, the modified two-stream instability [*Lampe et al.*, 1972]; however, in view of the small domain of $M_A \sin\psi$ when this instability will operate and the flute mode will not, we do not discuss it any further.) The flute instability is triggered when

$$\sum_j \alpha_j > 2^{5/2}\left(\frac{m_e}{m_i}\right)^{5/4}\left(\frac{1+\beta_e}{1+\beta_e - M_A^2\sin^2\psi}\right)^{5/4} \qquad (6)$$

and instabilities can only arise if

$$M_A^2 \leq \frac{1+\beta_e}{\sin^2\psi} \qquad (7)$$

Further information on the nature of this instability can be gained by studying the ratio of longitudinal to transverse electric field:

$$\frac{|\mathbf{k} \times \mathbf{E}|}{|\mathbf{k} \cdot \mathbf{E}|} = \left(\frac{m_e}{m_i}\right) \frac{M_A{}^2 \sin^2 \psi}{(1 + \beta_e - M_A{}^2 \sin^2 \psi)^{1/2}} \tag{8}$$

and the ratio

$$\frac{|\mathbf{E}|}{|\mathbf{B}|} = \frac{\Omega_e}{\omega_e} \frac{(1 + \beta_e - M_A{}^2 \sin^2 \psi)^{1/2}}{M_A \sin \psi} \tag{9}$$

Results (7) and (9) will be used in the next section.

2.2. *Experimental Evidence*

We now discuss the processes outlined above in light of the currently available data from the comet Halley encounters. Far from the comet, the solar wind is highly super-Alfvénic and $\alpha \ll 1$ so that the condition for driving flute modes (equation (7)) is not satisfied. In this region where $M_A \gg 1$, we expect that the dominant modes will be the low-frequency ($\omega < \Omega_i$) MHD waves which are triggered by the field-aligned component of the pickup ions as observed by *Tsurutani and Smith* [1986]. As the solar wind approaches the comet, the value of M_A decreases while α increases.

From (7) it is easy to see that if condition (6) is satisfied, the flute ion-ion instability will always be triggered ahead of the supersonic to subsonic transition. For example, for the Giotto transition $\psi \approx 56°$ [*Johnstone et al.*, 1986], and the condition given by (7) will be satisfied everywhere inside the region where the magnetosonic speed $M_S \equiv M_A/(1 + \beta_e)^{1/2} \lesssim 1/\sin \psi \approx 1.2$. Since in this case $\beta_e \approx 0.5$, the above condition is equivalent to $M_A \leq 1.5$. Similar analysis for the VEGA 1 gives a value of $M_A \lesssim 1.4$. The condition given by (6) can be approximated as $\alpha > 5 \times 10^{-4}$. This is satisfied further ahead of the bow interaction than condition (7). For example, *Balsinger et al.* [1986], using Giotto data, estimate the concentration of cometary protons to be 0.2% at a distance of 2.2×10^6 km from the nucleus increasing to 0.7% at 1.2×10^6 km near the bow interaction. Note that although the number density of pickup water ions is comparable to the pickup protons, they play a secondary role in initially triggering the instability. This is due to the fact that $\alpha_{H_2O^+}$ is smaller than α_p by a factor equal to their mass ratio. Hence although the water ions are largely responsible for the mass loading of the solar wind, they do not drive the initial linear phase of the instability. However, their velocity is in phase with the waves, and they can interact coherently with them.

We examine next the evidence for the existence of the lower hybrid waves as recorded by the plasma wave instruments MISCHA and APV-V on VEGA 1 [*Galeev et al.*, 1986; *Klimov et al.*, 1986]. We are interested in plasma wave data at the lower hybrid frequency which for the VEGA 1 magnetic field strength of 10 γ [*Galeev et al.*, 1986] at the bow wave is around 4–6 Hz. *Galeev et al.* [1986], *Riedler et al.* [1986], and *Klimov et al.* [1986] have shown data in the frequency range 1–4 Hz. With each of the respective instruments (APV-N and MISCHA), a significant increase in the emission at these frequencies was seen at the bow wave, with an intensity of a few mV/m [*Klimov et al.*, 1986; Figures 1*b* and 1*c*]. Further, by combining the results of the two instruments, *Galeev et al.* [1986] provided an approximate estimate of the ratio given in (9): this is shown in their Figure 5. Near 5 Hz the ratio of E/B is approximately 50. To evaluate (9), we use the VEGA 1 density and magnetic field strengths [*Galeev*, 1986; *Galeev et al.*, 1986] of 12 cm^{-3} and 10 γ and find that $E/B \approx 10^{-2}$, in reasonable agreement with the data. Hence both the location of the enhancement of lower hybrid waves and the E/B ratio are consistent with the onset of ring-driven flute modes as M_A approaches unity.

Note that other means of generating hybrid waves (such as the "shock" postulated by *Omidi and Winske* [1987]) would give similar results. This is because, given an appropriate free energy source, the frequency and E/B ratio depend only upon the background plasma parameters. The gyrating upstream ions observed by Omidi and Winske will form a beam with respect to the background plasma in the shock foot rather than a ring. However, this would not affect the properties of the lower hybrid waves. We also point out that U.S. spacecraft cannot observe plasma waves in this frequency range owing to their spin rate; so we cannot discuss this mechanism as it might apply to comet G-Z.

3. HEATING AND ACCELERATION OF PICKUP COMETARY PARTICLES

3.1. *Theoretical Considerations*

Having established the cause of the sharp increase in the amplitude of the electrostatic lower hybrid waves in the vicinity but ahead of any potential shock transition, one needs to examine the resulting anomalous transport. In this paper we confine our analysis to the relaxation of the ring distribution of the picked-up cometary ions. Energy transfer and heating of the solar wind will be examined in a future publication.

The relaxation is computed on the basis of a quasi-linear theory. As noted above, the instability results in a cylindrically symmetric spectrum in $k_\perp$ space about the magnetic field. This fact makes the quasi-linear analysis valid even for narrow values of the spectrum [*Davidson*, 1972; *Kulygin et al.*, 1971]. Since the relaxation time is much shorter ($\lesssim$20 s) than the ionization time scale, the source term is neglected in the diffusion equation. The evolution of a single cometary picked-up species is discussed first, followed by a multispecies analysis.

The diffusion equation for the newly born ions (protons or water ions) is given by *Davidson* [1972] as

$$\frac{\partial f_j}{\partial \tau_j} = \frac{4}{25} \frac{\partial}{\partial w} \left(\frac{1}{w^{1/2}} \frac{\partial f}{\partial w}\right) \qquad w > 0 \tag{10}$$

where $w = u_\perp{}^2 - v_0{}^2$, $v_0 = \omega_0/k$, $w_0 = v_R{}^2 - v_0{}^2$ with v_R being the speed of a newly born ion ring (common for both species), and j represents the jth species. The distribution functions f_j are normalized to unity, i.e., $\int_0^\infty dw f_j = 1$, and the effective time variable $\tau_j(t)$ is defined as follows:

$$\tau_j(t) = 25(16\pi^2 e^2 v_0{}^2/m_j{}^2)\, \Delta_0 \int_0^t dt' T_{k0}(t') \tag{11}$$

Note that τ_j can be written in terms of τ_p as $\tau_j = (m_p/m_j)^2 \tau_p$, where τ_p is the proton diffusion time. This highlights the self-similar nature of the diffusion process when more than one species is present. The initial condition is chosen as a mono-energetic ring:

$$f_j(\tau = 0) = \delta(w - w_0) \tag{12}$$

This assumption is not critical to the evolution nor to the final state of the system (*Davidson* [1972] and *Kulygin et al.* [1971]; see also section 2).

The excited spectrum is assumed to be a narrow ring in k space, $k_1 < k_\perp < k_2$ centered at $k_\perp = k_0$ and with a width

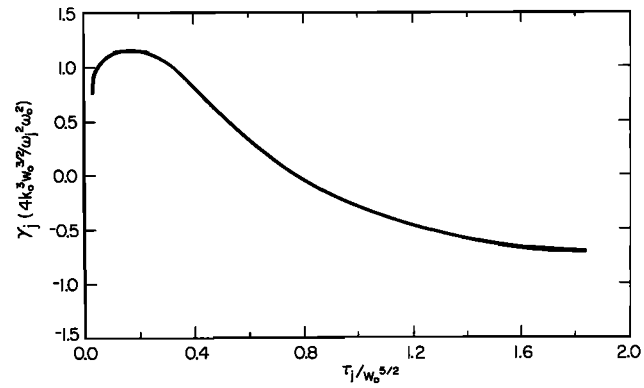

Fig. 1. The normalized growth rate $\gamma_j/(\omega_j{}^2\omega_0{}^2/4k_0{}^3w_0{}^{3/2})$ as a function of $\tau_j/w_0{}^{5/2}$ if only the jth species is considered.

$\Delta_0 = k_2 - k_1 \ll k_0$. The spectral energy density T_{k_0} is related to the electric field energy density by [*Kulygin et al.*, 1971]

$$T_{k_0}k_0\Delta_0 = \frac{E^2}{8\pi} \tag{13}$$

We note that since the resonant region in velocity space ($\mathbf{k} \cdot \mathbf{v}_\perp = \omega_0$) occupies the infinite domain $v_\perp > \omega_0/k_\perp = v_0$, the phase speed v_0 can be very small in comparison with v_R and yet the narrow spectrum is able to affect protons at $v_\perp = v_R$ upon suitable choice of the angle between $\mathbf{k}$ and $\mathbf{v}_\perp$. This is a unique feature of two-dimensional diffusion compared with one-dimensional diffusion. Hence the initial hydrodynamic instability evolves into a kinetic instability.

The growth rate and the spectral energy density are simultaneously calculated as functions of τ (or t) from the standard equations, namely,

$$\gamma = \sum_j \gamma_j = \sum_j \frac{\omega_0{}^2\omega_j{}^2}{2k_0{}^3} \int_0^\infty dw \, \frac{1}{w^{1/2}} \frac{\partial f_j(w, \tau_j)}{\partial w} \tag{14}$$

and

$$\frac{\partial T_{k_0}}{\partial \tau_j} = 2\gamma_j T_{k_0} \tag{15}$$

where ω_j is the plasma frequency of the jth species. The last equation can be replaced by the following one:

$$T_{k_0}(\tau) - T_{k_0}(0) = \sum_j \frac{2}{25(16\pi^2e^2v_0{}^2/m_j{}^2)\Delta_0} \int_0^{\tau_j} d\tau' \gamma_j(\tau') \tag{16}$$

after making use of relation (11).

The solution of (10) is

$$f_j(w, \tau_j) = \frac{5}{2\tau_j}(ww_0)^{3/4} \cdot \exp\left(-\frac{w^{5/2} + w_0{}^{5/2}}{\tau_j}\right) I_{-3/5}\left[\frac{2(ww_0)^{5/4}}{\tau_j}\right] \tag{17}$$

where $I_{-3/5}$ is the modified Bessel function of order $-3/5$.

Equations (14)–(16) can be solved analytically by using the series representation of the Bessel function in (17) [*Gradshteyn and Ryzhik*, 1980]. The resulting series can then be summed numerically. The expression for the growth rate γ and the spectral energy density T_{k_0} are functions of the variable τ. Conversion back to the time variable t can be achieved through (12) or its equivalent differential form. In Figures 1 and 2, $\gamma_j/(\omega_j{}^2\omega_0{}^2/4k_0{}^3w_0{}^{3/2})$ and $[T_{k_0} - T_{k_0}(t = 0)]_j/(n_jm_jw_0/200\pi\Delta_0k_0)$ are plotted as functions of $\tau_j/w_0{}^{5/2}$. ($[T_{k_0} - T_{k_0}(t = 0)]_j$ is the contribution from jth species in (16).) Solutions have been computed using both MACSYMA and standard numerical packages in the SLATEC library.

Assuming that the oscillations grow from a seed noise level such as $T_{k_0}(t = 0)/T_{k_0\max} \approx 8.0 \times 10^{-6}$, or, equivalently, $E(t = 0)/E_{\max} \approx 2.8 \times 10^{-3}$, we are able to evaluate the maximum electric field amplitude reached, $E_{\max}$. It is given, in volts per meter, by the following expression:

$$E_{\max} \simeq 5.5 \times 10^{-4}(m_i/m_p)^{1/2}\, n_p{}^{1/2}(v_R{}^2 - \lambda^2 f^2)^{1/2}(V/m) \tag{18}$$

where v_R, $\lambda = 2\pi/k_0$, $f = \omega_0/2\pi$ are measured in kilometers per second, kilometers, and hertz, respectively, and n_p is the proton ring number density. The value of the seed noise level, if it is chosen properly small, does not affect the result in (18). The maximum growth rate is also found to be

$$\gamma_{\max} = 8 \times 10^4\left(\frac{m_i}{m_p}\right)n_p f^{-1}\left(\frac{u_R{}^2}{\lambda^2 f^2} - 1\right)^{-3/2} (\mathrm{s}^{-1}) \tag{19}$$

Without replenishment the oscillations eventually subside back to noise level. If the final noise level is set equal to zero, then the corresponding asymptotic γ is negative and approximately 60% of the maximum growth rate (Figure 1). On the other hand, the first moment of the distribution function $f(w, \tau)$, namely, $\langle w \rangle = \int_0^\infty dw w f(w, \tau)$, becomes asymptotically equal to the respective moment of the initial distribution function (equation (11)), namely, w_0. Therefore all the field energy goes back to the particles, as is expected in two-dimensional quasi-linear theory. Solutions to (14)–(17) can also be easily computed for constant wave energy level. In this case, t scales linearly with τ, and only (11) and (17) are required.

We digress briefly to discuss further the importance of the two-dimensional diffusion process. Previous workers [e.g., *Lee and Birdsall*, 1979; *Roth and Hudson*, 1985] have studied the diffusion of a cold ring (with $\omega_e/\Omega_e \ll 1$). Using one-dimensional particle or hybrid codes, they found that trapping occurred rapidly. This is to be expected in one-dimension, since the autocorrelation time $\tau_{ac} \sim 1/k\Delta v$, where Δv is the ring's thermal spread. When the ring is cold, $\Delta v \ll v$, $\tau_{ac} \sim (v/\Delta v)(1/\omega_H)$, which is large. However, we do not believe that this is the correct approach. In two dimensions the resonance condition is $\omega \approx \mathbf{k} \cdot \mathbf{v} \neq kv$ so that $\tau_{ac} \sim 1/kv \sim 1/\omega_H$, which is the shortest time scale. The different resonance condition

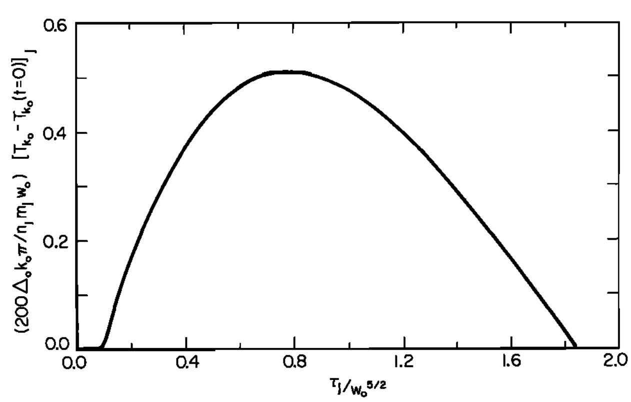

Fig. 2. The normalized spectral energy density $[T_{k_0} - T_{k_0}(t = 0)]/(n_jm_jw_0/200\,\pi\Delta_0k_0)$ as a function of $\tau_j/w_0{}^{5/2}$ if only the jth species is considered.

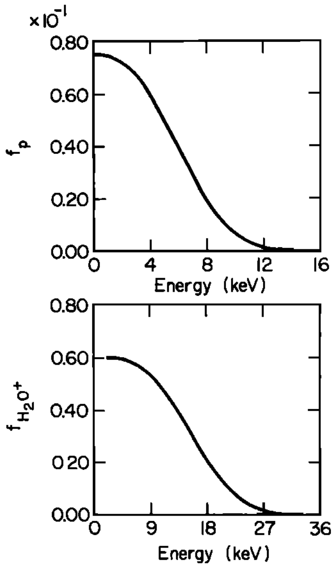

Fig. 3. The proton and H_2O^+ distribution after 8 s of heating. A constant electric field $E = 3.4 \times 10^{-2}$ V/m is assumed. The density and velocity are 0.05 cm^{-3} and 300 km s^{-1}, respectively, for both species.

also implies that heating can easily occur, which will not be the one-dimensional result.

3.2. *Heating of Cometary Species*

We turn next to the applications of the theory to ring distributions at cometary bow shocks. Owing to the different masses of protons and H_2O^+ ions, the relaxation of each of the ring distributions occurs on very different time scales. In the first phase the protons give rise to fluctuations which grow rapidly and damp their energy mostly on the protons. The water group ions initially play no role in this process, since they diffuse in energy about 300 times slower. Thus the H_2O^+ heating is accomplished on a time scale 324 times longer than that of protons. At first sight it would appear that almost all the energy must go to the protons, since the diffusion time $\tau \propto m_j^2$. However, this does not arise for the following reasons. First, as the protons are heated, their distribution function is flattened so that the energy transfer to them is slowed down substantially. Second, the region of the distribution that the waves can damp on is at the high-energy end of the spectrum. Since the characteristic diffusion time is $\propto v_\perp^5$ (see (10)), it is very difficult to produce further heating. The combination of these two effects makes it possible to heat both water ions and protons. We now proceed to make these statements more quantitative.

A difficulty arises if one wishes to solve (14)–(17) for both protons and H_2O^+ ions. Since very disparate length scales arise in both velocity space and time, it is all but impossible (as we have found) to solve the complete problem. We instead solve the following problem which, while simpler, retains the relevant physics:

1. We solve (13)–(17) for the initial relaxation of the protons as an initial value problem.

2. We solve the same equations for the H_2O^+ distribution, only ignoring the protons, to calculate the maximum electric field.

3. We then assume that an approximately steady state level of electric field fluctuations exists and study the response of both protons and H_2O^+ ions to this field. The level of fluctuations is chosen to be of the order of the maximum electric field due to an H_2O^+ ring (to within a factor of 2).

This method of solution allows us to make use of some simple scalings of the results. The growth rates and maximum electric fields scale very simply with the ion mass as shown in (18) and (19). Also, for a constant electric field fluctuation level the solution (17) is self-similar, since the solution for protons at time t_1 corresponds to the solution for H_2O^+ at a time $t_1(m_i/m_p)^2$.

We present results for a "standard" set of parameters. It is assumed that the density of cometary protons and H_2O^+ is 0.05 cm^{-3}. The velocity of the ring ions is assumed to be 300 km s^{-1} so that the protons have an energy of 470 eV and the water ions have an energy of roughly 8.5 keV, as seen by an observer in the solar wind frame. This compares with energies of 1.5 keV and 20 keV, respectively, evaluated in spacecraft frame by Giotto [*Wilken et al.*, 1986], and the minimum phase velocity of the waves is taken as 150 km s^{-1}. The numerical solution of (13)–(17) separately for the protons and H_2O^+ gives values for the maximum turbulent electric field in the lower hybrid region as $E_{max} = 3.2 \times 10^{-2}$ and 1.4×10^{-1} V/m correspondingly. These values are consistent with the theoretical values given by (18) and are of the same order of magnitude as the ones shown in the work of *Galeev et al.* [1986]. The subsequent heating of both protons and H_2O^+ was studied by assuming a stationary spectrum of lower hybrid waves at a saturated value of $E^2/8\pi \approx 5 \times 10^{-15}$ J/m³ which corresponds to an electric field $E \sim 3.4 \times 10^{-2}$ V/m. The resultant distribution for the protons and H_2O^+ after 8 s of heating which corresponds to the time scale for the decay of the waves is shown in Figure 3. The H_2O^+ particles have thermal spread of the order of 14 keV, while the proton spread is about 800 eV. This is consistent with the self-similar nature of the solution. The results are consistent with the observations reported by *Wilken et al.* [1986], although the heating time scales are faster than the time resolution of the current published data.

4. DISCUSSION AND CONCLUSIONS

In this paper the cause and the implications of the electrostatic lower hybrid turbulence in the bow interaction region of comets were discussed. The turbulence is driven by ions (both protons and H_2O^+ group) of cometary origin in the outer regions of a cometary magnetosphere. It was shown that

1. There is a critical solar wind speed at which such ring distributions are unstable to predominantly electrostatic lower hybrid waves propagating perpendicular to the ambient magnetic field. For cometary parameters the critical velocity is in the neighborhood of the observed bow interaction region and always ahead of the supersonic-subsonic transition.

2. The concentration of protons is large enough to drive these lower hybrid modes.

3. The theoretically expected ratio E/B of 10^{-2} for cometary parameter agrees well with the observations [*Galeev et al.*, 1986].

4. The theoretically expected amplitude and the triggering location agree reasonably well with the observations [*Galeev et al.*, 1986].

5. The waves heat the cometary protons extremely rapidly ($\ll 1$ s) and the H_2O^+ ions more slowly (10 s).

This is faster than current observations are able to resolve. We also note that the low-frequency waves ($\omega < \Omega_i$) can also contribute to heating of the cometary ions, as was found by *Winske et al.* [1985] for the case of the AMPTE lithium releases. The energy density of the low-frequency waves is much bigger than the lower hybrid waves by at least a factor of 100 [*Tsurutani and Smith*, 1986; *Galeev et al.*, 1986]. On the contrary, the heating time scale of the lower hybrid waves is a few seconds as opposed to minutes for the low-frequency waves. Clearly, there is a trade-off between energy density and heating rate, and a more careful analysis, especially of the low-frequency waves for cometary parameters, is needed to establish the precise importance of each mechanism.

This paper is part of an ongoing extensive investigation that examines the energy and momentum transfer from the solar wind to cometary ions and the potential role of a shock transition. Viewed from a global system point of view, we deal with a transition from a cold proton dominated solar wind upstream to a hot plasma downstream composed of hot solar wind protons and energetic suprathermal cometary distributions of H_2O^+ and protons. The present part addressed the physical processes in which energy is transfered from the solar wind protons to cometary ion rings with extremely small energy dispersion to hot energetic populations with large energy dispersion (spread). It did not, however, address the next critical question, namely, is any part of the wave energy, electrostatic or electromagnetic, generated by the cometary ring distributions reabsorbed by the solar wind protons? Such a process would allow the supersonic to subsonic solar wind transition to be accomplished without a shock due to an increase in the magnetosonic speed. It would essentially correspond to a nonlocal viscosity concept, although in the present case, momentum coupling to the solar wind cannot take place. Rather, nonlocal viscous heating is what occurs.

The possibility that resonance absorption of the low-frequency electromagnetic cyclotron waves by the protons in the transition region can lead to solar wind proton heating is currently under study and will be reported elsewhere [*Sharma et al.*, 1988]. It is also possible that an extension of the present theory can produce the same effect by heating the solar wind protons by absorption of the electrostatic lower hybrid waves.

Next, let us turn to the role of lower hybrid instabilities in cometary "bow shocks." We pose the following question "Does the cometary bow shock give rise to the lower hybrid instabilities, or are the consequences of the instability what people identify as the shock?" In the former case the waves would have to be driven by cross-field currents in the transition region or by reflected solar wind protons. It has been argued here and supported by the data that the lower-hybrid waves are consistent with thermalization of the ring and not part of any possible shock transition.

The envisaged scenario is as follows: deceleration of the solar wind results in the ring velocity falling through the critical condition needed for generating flute modes. Strong lower-hybrid turbulence arises (as is observed) giving rise to quasi-linear proton and H_2O^+ heating. Subsequently, nonlinear processes couple in the bulk solar wind, so that the solar wind heating makes the flow locally submagnetosonic. The observed plasma and magnetic field compression are simply an adiabatic response to the ion heating. One does not have a "shock" in the accepted sense of the word, but a turbulent interaction region, the initiation of which is due to instabilities of the cometary ions. We stress that the latter parts of this model are extremely speculative at present but worthy of further investigation.

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