TITLE ION-ACOUSTIC WAVES AND DRIFT WAVES IN NEGATIVE ION SOURCES

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Ion-acoustic waves and drift waves in negative ion sources

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ABSTRACT

Attention is directed towards instabilities induced in Hydrogen plasmas due to the presence of the desired H^- ions. Preliminary investigations indicate that when the electron temperature exceeds the temperature of the background ions (H^+), the H^- beam causes collective instabilities even at very low (even vanishing) beam velocities, provided that the non-ideal properties of the background plasma are taken into consideration. The most dangerous instabilities involve oscillations transverse to the direction of beam extraction, and may thereby degrade the beam emittance.

1. INTRODUCTION

Coherent ion-acoustic oscillations have been observed to be spontaneously generated in the plasma of a positive ion source device, probably by the presence of fast electrons accelerated through cathode sheaths.¹ The authors of this work, K. Schoenberg and W. Kunkel, state that the fluctuation properties of the ion source plasma set the lower limits of beam divergence, because the phase-space volume of extracted ions is incompressible (Liouville's theorem). For the same basic reason, it is important to investigate the macroscopic fluctuations that may reside in negative ion source plasmas, especially in view of the severe requirements on H^- beam divergence in the NPB system.

Attention will be directed towards instabilities induced in (H^+, e) plasmas by the presence of the H^- species. Preliminary investigations by this writer indicate that, in a condition of electron temperature exceeding the ion temperature, the H^- beam can induce instabilities even at very low (or vanishing) velocity relative to the background plasma.

It is important to remark that the boundary condition at the H^- extraction surface in negative ion sources is such as to reflect H^+ ions, unlike the boundary condition associated with the Bohm-like sheaths in positive ion sources. Thus, the argument made in Ref. 1 that Bohm sheath boundaries tend to absorb ion acoustic waves, thereby limiting the size of fluctuation amplitudes in the plasma volume, is not applicable to H^- sources. Therefore, H^- sources might be more susceptible than positive ion sources to the build up of macroscopic oscillations, both because of the different boundary conditions, and because of the presence of H^- ions.

Nezlin² has reviewed the experimental and theoretical work on instabilities induced by the presence of ion beams traversing a background plasma. His work is oriented towards fast beams traversing an ideal (non-dissipative) plasma. This work, or modifications of it, may be applied to surface sources (SPS), in which H^- ions are generated by secondary emission at Cesium-coated cathodes, and accelerated to high energies (~ 100 eV) along the ambient magnetic field. However, Nezlin's model is not applicable to the case of very small H^- beam velocities, less than ion thermal velocities, because he has assumed that the background ions are cold. Slow H^- beams would be relevant to Volume Sources, and to cross-field drift of H^- ions towards the lateral anodes in Surface Sources. Briggs³ has included the temperature of the background ions in his monograph, and finds that there are no instabilities excited in the ideal plasma for beam velocities smaller than ion thermal velocities. Briggs also points out that the low-velocity stable regime can be dissipatively destabilized if the non-ideal properties of the background plasma are taken into account. The present writer has found, in a preliminary investigation, that the dissipative destabilization can persist even down to vanishing beam velocity, provided that the plasma electrons are hotter than the plasma ions.

2. SURVEY OF MACROSCOPIC OSCILLATIONS INDUCED BY H^- BEAM IN IDEAL PLASMA

2.1 The dispersion relation

The linearized equations governing electrostatic perturbations of various plasma and field quantities q of the form $(\delta q) \exp[ik_x x + i(t - \lambda t)]$ in an infinite uniform H^- beam plasma system (without a magnetic field) are as follows:

$$k^2 \delta \phi = -k_z (\delta n_e + \delta n_i + \delta n_{H^-}) \quad (1)$$

$$\delta \vec{E} = -ik\delta\phi \quad (2)$$

$$\frac{\delta n_e}{n_e} = \frac{\vec{k} \cdot \delta \vec{V}_e}{\omega} \quad (3)$$

$$\delta P_e = n_e e \delta \phi \quad (4)$$

$$\delta P_e = \gamma_e T_e \delta n_e \quad (5)$$

$$\frac{\delta n_i}{n_i} = \frac{\vec{k} \cdot \delta \vec{V}_i}{\omega} \quad (6)$$

$$M\omega(\vec{k} \cdot \delta \vec{V}_i + k^2 \frac{\delta P_i}{n_i}) = -ek^2\delta\phi \quad (7)$$

$$\delta P_i = \gamma_i T_i \delta n_i \quad (8)$$

$$\frac{\delta n_b}{n_b} = \frac{\vec{k} \cdot \delta \vec{V}_b}{\omega} \quad (9)$$

$$M\omega(\vec{k} \cdot \delta \vec{V}_b) = ek^2\delta\phi \quad (10)$$

The subscripts e, i, b, refer respectively to plasma electrons, plasma ions (H^+), and beam ions (H^-). Here, electron inertia is ignored, a cold, coasting H^- beam is assumed, and ideal energy equations, with adiabatic indexes γ_e and γ_i , are taken to govern the pressure perturbations in the warm e and H^+ species. Moreover, H^+ and H^- ions are taken to have mass M , and $\omega(\omega - \vec{k} \cdot \vec{V}_b)$ represents the Doppler shifted frequency seen by the beam. Lastly, number densities are represented by n , temperatures by T , velocities by $\vec{V}$, pressures by P , the electrostatic potential by ϕ , and the zero-order beam velocity by $\vec{V}_b$.

These equations yield the following dispersion relation

$$1 - \frac{\omega_p^2}{\omega^2 - k^2 v_{Te}^2} - \frac{\omega_p^2}{(\omega - \vec{k} \cdot \vec{V}_b)^2 - k^2 v_{Ti}^2} - \frac{\omega_p^2 n_e}{k^2 v_{Te}^2 n_i} = 0 \quad (11)$$

where the characteristic thermal velocities and plasma frequencies are

$$v_{Te} = (T_e/L_e/M)^{1/2} \quad (12)$$

$$v_{Ti} = (T_i/L_i/M)^{1/2} \quad (13)$$

$$\omega_{pi} = (4\pi n_i e^2 / M)^{1/2}, \quad j = e, b. \quad (14)$$

Moreover, zero-order neutrality holds, namely

$$n_e + n_b = n_i. \quad (15)$$

In the case of frequencies far below the ion plasma frequency (and wavelengths far above the Debye length), we may take the quasi-neutral limit of Eq. (11), equivalent to setting the rhs of Eq. (1) to zero. The result is

$$\frac{\omega_{pi}^2}{k^2 \epsilon^2} \left(1 - \frac{n_b}{n_i} \right) = \omega^2 - \frac{\omega_{pb}^2}{k^2 v_{th}^2} + (\omega_c)^2. \quad (16)$$

This dispersion relation can be rewritten in terms of a reference frequency ω_p , defined by

$$\omega_p^2 = k^2 \epsilon^2 \left(1 - \frac{n_b}{n_i} \right) \quad (17)$$

and a coupling (beam strength) parameter ϵ , defined by

$$\epsilon = n_b / n_i. \quad (18)$$

The rewritten dispersion relation is

$$1 = \frac{1}{\Omega^2 - \sigma^2} + \frac{\epsilon}{(\Omega - U)^2} \quad (19)$$

where the normalized frequency Ω is

$$\Omega = \omega / \omega_p, \quad (20)$$

the normalized drift parameter U is

$$U = k V_{th} / \omega_p, \quad (21)$$

and the warm-ion parameter σ is defined by

$$\sigma^2 = \frac{v_{th}^2}{\epsilon^2} \left(1 - \frac{n_b}{n_i} \right). \quad (22)$$

(Note that a slight redefinition of the symbols serves to remove the factor $(1 - n_b/n_i)$.)

2.2. Solutions of the dispersion relation, and interpretation

Equation (19) can be further re-written as

$$(\Omega - U)^2(\Omega^2 - \sigma^2 + 1) = \epsilon(\Omega^2 - \sigma^2). \quad (23)$$

We now examine the complex frequency roots, Ω , of this dispersion relation. For negligible beam density, $\epsilon \rightarrow 0$, there are two identical drift wave solutions, $\Omega = U$, or $\omega = k \cdot V_b$, and there are forward and backward acoustic waves, $\Omega = \pm(\sigma^2 + 1)^{1/2}$, or $\omega = \pm k(c_{th}^2 + s^2)^{1/2} = \pm kc_s$, where c_s is the speed of ion-acoustic waves. The drift wave solutions represent slight ripples fixed on the beam passing through the background plasma un-noticed because of the low beam density. The acoustic waves represent quasi-neutral longitudinal waves propagating through the background plasma in the absence of the beam.

For a small relative density of beam ions, $\epsilon \ll 1$, one can iteratively determine how these four roots of the dispersion relation are distorted by the beam. In the parameter range where $\sigma^2 = U^2 < (\sigma^2 + 1)$, one easily finds that the acoustic waves and one drift wave remain stable, whereas the remaining drift wave is destabilized. It possesses the following root

$$\Omega = U + \epsilon^{1/2} \frac{U(U^2 - \sigma^2 + 1)^{1/2}}{(\sigma^2 + 1 - U^2)}. \quad (24)$$

This expression is valid only for the imaginary part representing a small correction (denominator not too close to zero). Since U is $(V_b/s)\cos\theta$, where θ is the angle between V_b and k , it transpires that beam aligned waves ($\theta = 0$) are the most rapidly growing. The range of beam velocities governing this behavior is then given by $c_s < V_b < 2c_s$.

It can easily be shown that there are no instabilities for $U^2 = \sigma^2$, nor for $U^2 = (\sigma^2 + 1)$, that is, for $V_b \cos\theta = c_s$, and for $V_b \cos\theta = 2c_s$. Thus, the range of unstable behavior lies between the ion thermal velocity and the ion acoustic velocity, except for a special case discussed below. Eq. (24) represents a non-resonant drift instability (sometimes called the Anomalous Doppler Instability, or the Reactive Medium Instability). It can be interpreted as arising from a negative dielectric property of the plasma in this parameter range.

A special case arises when the denominator vanishes in Eq. (24). This condition, $U^2 = \sigma^2 + 1$, can also be written as

$$V_b \cos\theta = c_s. \quad (25)$$

The angle of wave propagation relative to the beam direction given by Eq. (25) is called the Mach angle, or the angle of Cerenkov Resonance. It has the property that a plane wave structure of this particular wave vector orientation $\vec{k}$, but rigidly drifting axially with the beam, exactly keeps up with oblique waves in the background plasma propagating along $\vec{k}$ with phase speed c_s . In this manner, a special perturbation drifting with the beam can resonantly drive a natural mode of the plasma. Note that for very large beam velocities V_b relative to c_s , according to Eq. (25) the resonant waves are propagating almost at right angles to the beam direction.

In the $\theta = 0$ resonant case, the one unstable root of Eq. (23) reads, in the small ϵ limit

$$\Omega = U + \epsilon^{1/3} \left(\frac{U}{\sigma} \right)^{1/2} \frac{1}{(\sigma^2 + 1)^{1/2}}. \quad (26)$$

Note that the growth rate depends only upon the cube root of the beam density ratio and is therefore a more robust growth rate (for weak beams) than is the non-resonant growth rate. Note also that the resonant condition, Eq. (25), can occur only for beam velocities that equal or exceed the ion acoustic velocity. The corresponding instability, Eq. (26), is associated with the vanishing of the plasma dielectric property.

For negative ion sources of the Penning geometry (Surface Sources), in which H^- ions generated at opposing cathodes are accelerated through the cathode sheaths to form fast beams parallel and anti-parallel to the ambient magnetic field, an electrostatic model of macroscopic oscillations has to be augmented by an electromagnetic model. Here, one can expect that the beam velocity compared to the Alfvén velocity (especially at low beta, $v_A \ll V_A$) will constitute an important criterion to distinguish non-resonant and resonant instabilities, and to determine the obliqueness of the latter. A build-up of such waves in the plasma chamber may be a concern as regards the quality of the internally extracted H^- beam. For a small amount of obliqueness, the period of such standing waves would vary in proportion to the cathode spacing and the square root of the ion mass, and inversely as the magnetic field strength.

Electrostatic ion acoustic waves can also be driven along the ambient magnetic field by a distribution of fast electrons originally generated at the cathodes, provided that the resulting electron velocity distribution effectively has $T_e \gg T_i$. (The electrons generated at the cathodes fill in their energy distribution by making various inelastic collisions.) The period of such waves would have the dependences mentioned above, except that the magnetic field dependence would be absent. Preliminary theoretical work done on this mechanism has achieved some qualitative experimental confirmation.⁴ The period of these aligned electrostatic waves would be longer than those of the H^- beam-driven magnetic waves in the ratio (V_A/v_A) . There is still some uncertainty about the boundary conditions at the cathodes that are needed in theory to support the build-up of such beam-aligned ion acoustic waves.

3. SURVEY OF MACROSCOPIC OSCILLATIONS INDUCED BY H^- BEAM IN DISSIPATIVE PLASMA

Real plasmas are dissipative. In this section, we include the processes of ion viscosity, ion and electron thermal conduction, and electron-ion friction (resistivity) due to momentum transfer (Coulomb collisions with plasmas ions eH^+) only. The general dispersion relation that governs quasi-neutral electrostatic perturbations is derived, but approximate solutions to it are so far obtained only by neglecting resistivity. The latter assumption is examined, and the significance of the dispersion relation roots is discussed in the context of negative ion sources.

3.1. The dispersion relation

The electron momentum equation (Eq. (4)) is replaced by

$$k^2 \delta P_e = -n_e e k^2 \delta \phi - \frac{1}{4} m_e n_e v_{Te} (k \cdot \nabla V - \nabla^2 V), \quad (27)$$

where v_{Te} is the electron-ion collision frequency representing the cumulative effect of many small angle scattering.⁵

The ion momentum equation (Eq. (7)) is replaced by

$$\frac{1}{2} M (k \cdot \nabla V) + k^2 \frac{\delta P_i}{n_i} + \frac{1}{4} \mu k^2 (k \cdot \nabla V) = -k^2 \delta \phi - \frac{1}{4} m_i \frac{n_i}{n_i} v_{Ti} (k \cdot \nabla V - \nabla^2 V), \quad (28)$$

The ion viscosity coefficient μ shall be written as

$$\mu = M n_i D_V, \quad (29)$$

where D_V represents the viscous diffusivity (kinematic viscosity) of the ion velocity field.

With thermal conduction included, the electron and ion linearized energy equations retain their earlier forms (Eqs. (5) and (8)), but the indexes ω_e and ω_i now become complex quantities. To avoid the following prescription

$$\frac{1}{1 - \omega_e} = \frac{1}{1 - \omega_e} + \frac{v_{Te}^2 k^2 D_e}{1 - \omega_e}, \quad (30)$$

where D_e represents the thermal diffusivity of the electrons (or ions). Note that heat sources have been ignored for the energy equations because such processes are second-order in fluctuation amplitudes. It is also important to note that, in the process of forming the dispersion relation, the characteristic thermal velocities can conveniently be introduced as complex quantities according to

$$v_{Te} = (T_e/T_i)^{1/2} M_i^{-1/2} = (T_e/T_i)^{1/2} M_i^{-1/2} \quad (31)$$

The remaining equations are as they appear in section (2.1). From all of these equations, the general dispersion relation (quasi-neutrality not yet invoked) is found to be as follows:

$$1 - \frac{\omega_{pe}^2}{\omega^2} \frac{F^2 - F^2_e}{\text{Denom}} = (1 - \epsilon) \omega_{pi}^2 \frac{I^2 + F^2_e}{\text{Denom}} + \frac{\omega_{pi}^2}{(\omega - i)^2} \quad (32)$$

Here, the following symbols have been employed:

$$I^2 = k^2 c^2 \quad (33)$$

$$I^2 = \omega_{pi}^2 + \omega^2 + i\omega k^2 D_A \quad (34)$$

$$F^2 = \frac{m_e}{M} (\omega - i)^2 \quad (35)$$

$$\text{Denom} = I^2(\omega^2 - F^2)(I^2 + F^2_e) + F^2 I^2_e \quad (36)$$

If resistivity is neglected ($F = 0$), it is easy to show that the dispersion relation (Eq. (32)) reduces to the form of Eq. (11), but generalized to include ion viscous damping (D_A) and ion and electron thermal conduction (complex ω_{pi}^2 and complex c^2). That is:

$$1 - \frac{\omega_{pe}^2}{\omega^2 + i\omega k^2 D_A} = (1 - \epsilon) \frac{\omega_{pi}^2}{k^2 c^2} + \frac{\omega_{pi}^2}{(\omega - i)^2} \quad (37)$$

Restricting consideration to quasi-neutral dynamics, we shall ignore the $(1 - \epsilon)$ of Eq. (32) (i.e. $\epsilon \rightarrow 0$) and refer to the remaining terms:

3.2. Solutions of the dispersion relation and interpretation

First, we consider what is required to neglect electron ion friction. From the quasi-neutral form of Eq. (32), one finds that essentially two conditions are necessary to leading order (assuming small ϵ and small dissipative rates). These are:

$$\epsilon \frac{m_e}{M} \frac{v_{Te}}{c} \ll 1 \quad (38)$$

$$\epsilon \frac{m_e}{M} \frac{v_{Te}}{c} \ll \frac{\omega}{\omega_{pe}} \quad (39)$$

We shall assume that these conditions are satisfied in the limit that the electron ion mass ratio (i.e. M) is taken to be arbitrarily small. However, they ought to be checked for each specific case, with realistic parameter values (e.g. $\epsilon = 1$). In special cases, they may be violated. For instance, if $(\omega_{pe}^2)^{-1}$ happens to be of order unity, then condition (38) is always satisfied for small ϵ . Also, condition (38) is satisfied at a resonance (i.e. where $\omega = \omega_{pe}$), however, (39) is not then satisfied for the non-resonant drift instability because $\omega < \omega_{pe}$ (i.e. $\omega < \omega_{pe}$) that scenario. Moreover, higher order effects may be important for obtaining accurate growth rates of instabilities, which will be considered numerically.

Ignoring the “1” in Eq. (37) by quasi-neutrality, we can re-write (37) in the canonical form of Eq. (19),

$$1 = \frac{k^2 s^2}{\omega^2 + i\omega k^2 D_V - k^2 v_{th}^2} + \epsilon \frac{k^2 s^2}{(\omega')^2} \quad (39)$$

Here, the (complex) s^2 has been slightly re-defined to absorb $(1 - n_b/n_i)$, that is, $s^2(1 - \epsilon)^{-1}$ is our new s^2 . Correspondingly, the electron temperature is redefined so that $T_e(1 - \epsilon)^{-1}$ is our new T_e .

Finally, for further analysis based upon the smallness of the coupling parameter, ϵ , we re-arrange Eq. (39) as

$$(\omega')^2 (\omega^2 + i\omega k^2 D_V - k^2 c_s^2) = \epsilon k^2 s^2 (\omega^2 + i\omega k^2 D_V - k^2 c_s^2 + k^2 s^2), \quad (40)$$

where the complex-valued acoustic speed is defined by

$$c_s^2 = v_{th}^2 + s^2 \quad (41)$$

Equation (40) constitutes the dispersion relation that governs quasi-neutral electrostatic oscillations for the system of a cold H^- beam penetrating a warm dissipative plasma.

In the elucidation of the various solutions of Eq. (40), it proves helpful to factor the expression describing the propagation of damped acoustic waves in the plasma. Thus

$$\omega^2 + i\omega k^2 D_V - k^2 c_s^2 = (\omega - \omega_+)(\omega - \omega_-), \quad (42)$$

where

$$\omega_{\pm} = \pm \omega_0 + i\gamma, \quad (43)$$

When all dissipation rates are small compared to the ion-acoustic frequency, so that only first-order corrections are kept, the real frequency carried by the plasma is

$$\omega_0 = k v_{ph} = k (v_b I + s_e T_e) M_i^{-1/2} \quad (44)$$

where v_{ph} is the real part of v_+ and the damping rate is

$$\gamma = k^2 D_{eff} \quad (45)$$

where D_{eff} is an effective diffusivity resulting from a combination of electron thermal conduction, ion thermal conduction, and ion viscosity.

$$D_{eff} = \frac{1}{2} (D_V + s_e^{-1} D_{ee}) + (1 - s_e) \frac{T_e}{M_i v_b^2} + (s_e - 1)^2 D_{ee} \frac{T_e}{M_i v_b^2} \quad (46)$$

Use of Eqs. (42)–(46) in Eq. (40) allows one to verify the existence of resonant and non-resonant instabilities similar to those discussed in Section 4.2, and in essentially the same range of beam velocities. The Cerenkov resonance is somewhat more complicated (in fact, there are two close but distinct resonances), and the addition of dissipation can, in some circumstances, contribute to destabilization rather than damping. However, our interest in this paper will be confined to the range in which $v_b < V_{ph} = v_{ph}$. Without dissipation, this range of V_b is stable.

We now examine the drift-wave solutions of Eq. (40), assuming that $\epsilon \ll 1$, i.e., on weak H^- beams. Assuming beam-aligned waves for simplicity, Eq. (40) is therefore re-written as

$$(\omega')^2 (\omega^2 + i\omega k^2 D_V - k^2 c_s^2) = \epsilon k^2 s^2 (\omega^2 + i\omega k^2 D_V - k^2 c_s^2 + k^2 s^2) \quad (47)$$

in which $\omega = kV_b$ is substituted on the rhs, valid when the rhs constitutes a small correction. Also, on the rhs, an expansion is made to first order in all dissipative processes. Hence, expressions for $\omega_{\pm}$ in Eq's. (43)–(45) are invoked. Thus, it is assumed that for any dissipation rate $\nu_{D_{\text{eff}}}$,

$$\frac{\nu_{D_{\text{eff}}}}{Re(\omega)} \ll 1, \quad \frac{\nu_{D_{\text{eff}}}}{kV_b} \ll 1, \quad \frac{\nu_{D_{\text{eff}}}}{ks} \ll 1. \quad (48)$$

Here, s_0^2 is the leading order part of $s^2 = (I_p I_e / M)$, obtained from the leading order part of I_p , namely $s_0 = cS$ (see Eq. (30)). We shall also need to expand complex s^2 to first order in the dissipation rate associated with electron thermal transport, in Eq. (47). In so doing we shall write

$$s^2 = s_0^2 + i(s^2)_1 \quad (49)$$

in which the first order dissipative part of I_p yields

$$(s^2)_1 = \frac{i\nu_{Te}}{Re(\omega)} cS = 10^2 \frac{I_e}{M} \quad (50)$$

where ν_{Te} is a dissipation rate associated with electron thermal diffusivity

$$\nu_{Te} = k^2 D_{Te}. \quad (51)$$

Here, it is important to note that the sign of $(s^2)_1$ is the same as the sign of $Re(\omega)$.

The result of the above described treatment of the rhs of Eq. (47) finally reads as follows

$$\omega = kV_b \pm \omega^{(2)} k \left[1 + \frac{\nu_{Te}^2}{s_0^2} \left(\frac{V_b^2}{V_e^2} \right)^{\frac{1}{2}} \right] \pm \frac{1}{2} \left[1 + \frac{1}{2} \frac{(s^2)_1}{s_0^2} \right] \left(\frac{\nu_{Te}^2}{V_e^2} + \frac{V_b^2}{V_e^2} \right) \frac{s_0^2}{V_e^2} \pm \frac{kV_b}{2} \left(\frac{\nu_{Te}^2}{V_e^2} + \frac{V_b^2}{V_e^2} \right)^{\frac{1}{2}}. \quad (52)$$

Here, ν_{Te}^2 is the leading order part of the complex ion thermal velocity squared, (I_p / M) , which is cS / M .

At this point, we can observe several features of this expression for the drift wave roots of Eq. (49). First note that the assumed ordering breaks down when V_b is too close to ν_{Te} in the same sense that Eq. (24) breaks down near Cerenkov resonance and has to be re-done to give Eq. (26). Thus, we assume that V_b lies sufficiently below ν_{Te} that the iterative correction employed in Eq. (52) is valid.

Second, notice that if $I_p = I_e$ (assuming $s_0 = cS$), then the term $(s^2)_1 = \nu_{Te}^2$ vanishes in (52). In this equal temperature case, one sees that $I\nu_{Te}$ vanishes in the limit that the beam velocity vanishes. $I\nu_{Te}$ also vanishes in the absence of all dissipative effects in the plasma. That is, the equal temperature plasma is stable unless the H^- beam is in motion through the non-ideal plasma.

Third, notice that when $I_p = I_e$, one of the two drift waves is always dissipatively unstable in the range $0 < V_b < \nu_{Te}$. In the example of an unmagnetized plasma, one would expect D_{Te} to be the dominant dissipative process. The upper sign in Eq. (52) then proves to be the unstable root, with a growth rate (for $I_p = I_e$) in the order of $I\nu_{Te} \pm \omega^{(2)} k V_b \pm \nu_{Te} \pm k^2 D_{Te}$.

As a word of caution, we recall that Eq. (48) (small dissipation rates) has been assumed in our approximate analysis. In the present example, this condition requires that $k^2 D_{Te} \ll kV_b$. Thus, the size of the plasma, l , needs to be sufficiently large to allow the considered instabilities to have large wavelengths that entail only a small amount of dissipation. For shorter wavelengths, one must resort to a numerical treatment of the original dispersion relation. The extent to which the dissipatively unstable waves can exist at shorter wavelengths thus remains to be determined.

A singularly interesting case occurs when the H^- beam is practically at rest in the plasma, $V_b \approx 0$. This situation may approximately correspond to the slow mass field diffusion of H^- towards the extraction region in both surface Sources and Volume sources. The root of the dispersion relation, Eq. (52), then reduces to

$$\omega \approx \pm \epsilon^{1/2} (k\epsilon_0) \left(\frac{v_{th0}}{c\epsilon_0} \right) \left\{ 1 \pm \frac{1}{2} \frac{(s^2)_1}{s_0^2} \frac{s_0^2}{v_{th0}^2} \right\} \quad (53a)$$

From this expression, we see that because the sign of $(s^2)_1$ is the same as the sign of $Re(\omega)$, the stability properties of both roots are the same. This makes physical sense, because, if there is no beam velocity, then the sign of $Re(\omega)$ has no real significance. Equation (53a) can also be written as (for $\gamma_e = \gamma_i = \gamma$),

$$\omega \approx \pm \epsilon^{1/2} k \left(\frac{\gamma - T_e T_i}{M T_e + T_i} \right)^{1/2} \left\{ 1 \pm \frac{1}{2} \frac{v_{th0}}{Re(\omega)} \frac{(\gamma - 1)^2 T_e - T_i}{T_i} \right\} \quad (53b)$$

Thus, the presence of the H^- at rest in the (H^+, e) plasma can induce an instability driven by electron thermal conduction, provided that $T_e > T_i$. This non-equilibrium thermal state of the plasma apparently constitutes a source of free energy that can be tapped by the impurity fraction (H^-) in the presence of electron thermal conduction.

It is important to remark that electron thermal conduction is largest along $\hat{B}$, the ambient magnetic field. In principle, we can imagine our wave vector k to be oriented in this direction, since the H^- beam is nominally at rest in this special example. Then, for electrostatic waves along $\hat{B}$, the entire derivation given above should be valid. From (53b), the growth rate of this instability is (for $\gamma = 5/3$),

$$Im(\omega) \approx \frac{2}{15} \frac{T_e - T_i}{T_i} k^2 D_{th0} \quad (54)$$

Again, we emphasize the caution that the dissipative rates were assumed small in our approximate analysis. In the case of Eq. (53b), this means that $k D_{th0} < \epsilon^{1/2} c/c_0$. Shorter wavelengths require numerical investigation.

4. SUMMARY AND CONCLUSIONS

A preliminary exploration was carried out of the microscopic instability properties of a uniform warm, dissipative Hydrogen plasma penetrated by a uniform cold, weak H^- beam. The possibility, in principle, of the generation of unstable waves at very low beam velocities was pointed out, and explicit expressions for the frequencies and growth rates of these waves were derived for the unmagnetized plasma. The possible importance of electron-ion temperature differences was mentioned.

Because the plasma transport coefficients (especially electron thermal conduction) are large for unmagnetized transport, only large wavelengths are available to the approximate analysis of this paper, for typical negative ion source plasma parameters. A more realistic treatment including shorter wavelengths and concomitant larger dissipation rates will require numerical analysis of the dispersion relation.

In this connection, it is important to note that transport processes across the ambient magnetic field of negative ion sources will be greatly reduced, especially electron thermal conduction. Thus, extending the theory to explicitly include waves across an ambient magnetic field will allow shorter wavelengths to exist with reduced dissipation rates.

The above-mentioned extensions of the theory are beginning to be implemented, together with other effects (beam-driven H^- beam, inclusion of resistivity, nonlinear simulation).

5. ACKNOWLEDGEMENT

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6. REFERENCES

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