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STABILITY OF TRAPPED-PARTICLE MODES
IN LOWER-HYBRID PUMP

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ABSTRACT. The stability of dissipative trapped-particle modes in the presence of a self-consistent lower-hybrid pump is analyzed theoretically. The pump can either be a single coherent wave or a spectrum of turbulent waves. Corresponding dispersion relations modified by the pump are derived and conditions for excitation and suppression of the modes are also discussed.

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1. INTRODUCTION

Recently, there has been active interest in using rf waves near the lower-hybrid frequency as a supplementary heating scheme for tokamak plasmas. Since the dissipative trapped-particle instabilities [1] are expected to significantly influence the energy transport in current and future generations of tokamaks, it is important to study the coupling between the rf waves and these trapped-particle modes, and to determine the effects of such processes on the overall efficiency of the heating scheme. To consider this coupling, we investigate the modifications in the dispersion relations of the dissipative trapped-ion and trapped-electron modes due to the presence of a self-consistent lower-hybrid pump. The pump can either be a coherent single wave corresponding to that excited by an external source or a spectrum of turbulent waves which correspond to waves excited by the incoming wave through the various parametric processes.

The theoretical model and basic approaches are discussed in Sec. 2, and the modified theory for the dissipative trapped-ion mode is presented in Sec. 3. Using the results obtained in Sec. 3, we also consider the dissipative trapped-electron mode in Sec. 4. Our main conclusions and a discussion of the results are presented in Sec. 5.

2. THEORETICAL MODEL AND APPROACH

In our basic theoretical model we assume that at equilibrium there exists in the tokamak plasma a self-consistent electrostatic lower-hybrid field, $\phi_0(\underline{x}, t)$; where

$$\phi_0(\underline{x}, t) = \sum_{\underline{k}} \phi_{\underline{k}} \exp[-i(\omega t + \underline{k} \cdot \underline{x})] + \text{c.c.} \quad (1)$$

Here, $(\omega, \underline{k})$ satisfies the following lower-hybrid dispersion relation

$$\omega^2(\underline{k}) = \omega_{lh}^2 (1 + k_{||}^2 m_i / k_{\perp}^2 m_e) , \quad (2)$$

where $\omega_{lh}^2 = \omega_{pi}^2 / (1 + \omega_{pe}^2 / \omega_{ce}^2)$, $k_{||} = \underline{k} \cdot \underline{B} / B$, $|k_{||} / k_{\perp}| \ll 1$, ω_{pe} and ω_{pi} are, respectively, the electron and ion plasma frequencies, and ω_{ce} is the electron cyclotron frequency. We note that in Eq. (1), $\phi_0(\underline{x}, t)$ is represented as a summation of a spectrum of lower-hybrid waves. The theory presented in this work, however, is also valid if ϕ_0 is a single coherent lower-hybrid wave; i.e., without the summation sign in Eq. (1).

In dealing with the perturbations of this equilibrium, we assume the pump field, $|\phi_{\underline{k}}|$, to be weak, so that only interactions up to the order of $|\phi_{\underline{k}}|^2$ need to be considered. More specifically, for each $\underline{k}$ we will only consider the interactions among the low-frequency $(\Omega, \underline{q})$, the lower-sideband $(\Omega - \omega, \underline{q} - \underline{k})$, the upper-sideband $(\Omega + \omega, \underline{q} + \underline{k})$ and the pump

$(\pm\omega, \pm k)$ modes. Here, the (Ω, q) mode corresponds to the dissipative trapped-particle modes of interest here with $|\omega| \gg |\Omega|$. Since the lower-hybrid modes have wavelengths typically much shorter than the dissipative trapped-particle modes, we can make the adiabatic approximation, $|q_\perp| \ll |k_\perp|$, $|q_\parallel| \ll |k_\parallel|$. Furthermore, for simplicity we shall adopt the slab model and work in the local approximation; i.e., $q = q_\perp \underline{e}_y + q_\parallel \underline{e}_z$ with the nonuniformity in the x-direction.

From the equations of continuity and motion, it is easy to show that the dominant contribution to the sideband nonlinear current comes from the electrons through the coupling between the low-frequency density perturbation, $n_e(\Omega, q)$, and the $\underline{E} \times \underline{B}$ drift induced by the pump wave. Since $|\omega| \gg \omega_{bi}$, ω_{be} and $|k_\parallel| \gg \tilde{q}/R$, all particles may be treated as untrapped, as far as motion in the lower-hybrid waves is concerned. Here, ω_{bi} and ω_{be} are, respectively, the typical bounce frequencies of trapped ions and electrons, $\tilde{q}$ is the safety factor and R is the major toroidal radius. Thus, the potentials of the sidebands are given by

$$\phi_\pm = i\beta \omega_{k_x} q_\perp \frac{n_s}{n_0} \begin{bmatrix} \phi_{k/\epsilon^+} \\ \phi_{k/\epsilon^-}^* \end{bmatrix}, \quad (3)$$

where n_s consists of contributions from both trapped and untrapped particles,

$$n_s(\Omega, q) = n_e^t(\Omega, q) + n_e^u(\Omega, q), \quad (4)$$

$$\beta = - \frac{4\pi n_o e c}{B(1 + \omega_{pe}^2 / \omega_{ce}^2)} , \quad (5)$$

and

$$\epsilon_{\pm} = k_{\pm}^2 [\omega_{\pm}^2 - \omega^2(k_{\pm})] , \quad (6)$$

with $\omega_{\pm} = \Omega \pm \omega$, $k_{\pm} = q \pm k$ and $\omega^2(k_{\pm})$ is given in Eq. (2) .

To describe the low-frequency dynamics, we use the equation of motion and the continuity equation with a model collisional term for the trapped particles as given by Kadomtsev and Pogutse [1] . The nonlinearities due to the couplings between the pump and the sidebands then come in through the ponderomotive force term, $-m_j [(\tilde{v} \cdot \nabla) \tilde{v}]_j$, in the equation of motion and the $\nabla \cdot [n \tilde{v}]_j$ term in the continuity equation with $j = e, i$. The details of the nonlinearities are different not only between the two modes considered here but also between the trapped and the untrapped particles. We will begin this analysis by first treating the dissipative trapped-ion mode in the next section.

3. DISSIPATIVE TRAPPED-ION MODE

Following Kadomtsev and Pogutse [1], the trapped-particle dynamics is described by the following continuity equation

$$\begin{aligned} \Omega n_j^t + i v_{jx}^t (dN/dx)_j^t - q_{\perp} \left[N v_Y + \sum_{\tilde{k}} (n_+ v_{\tilde{k}y}^* + n_{\tilde{k}}^* v_{+y} + n_- v_{\tilde{k}y} + n_{\tilde{k}} v_{y-}) \right]_j^t = -i(v/\epsilon)_j \left[n^t - \sqrt{\epsilon} (n^t + n^u) \right]_j ; \\ j = e, i . \end{aligned} \quad (7)$$

Here, N^t is the equilibrium density of trapped particles, the v is the self-collision frequency, $\epsilon = r/R$ is the small aspect ratio and $n_{k,+,-}$ and $v_{k,+,-}$ are density and velocity perturbations associated with the pump and the two sidebands. For the high-frequency velocity perturbations, we have $(v_{\perp})_{k,+,-}^e = -e_{\perp} \times (cE/B)_{k,+,-}$ and $(v)_{k,+,-}^i = (-eE/m_i \omega)_{k,+,-}$; whereas the low-frequency velocity perturbation v_j is composed of the usual $E \times B$ drift plus that due to the perpendicular ponderomotive force; i.e.,

$$v_j^t(\Omega, q) = \frac{c}{B} (-iq_{\perp} \phi + F_{j\perp}/e_j) \times e_{\perp} , \quad (8)$$

where

$$\begin{aligned} F_j = -im_j \sum_{\tilde{k}} \left[(v_{\tilde{k}} \cdot k_{\tilde{k}}) v_{\tilde{k}-} + (v_{\tilde{k}-} \cdot k) v_{\tilde{k}} + (v_{\tilde{k}}^* \cdot k_{\tilde{k}}) v_{\tilde{k}+} - (v_{\tilde{k}+} \cdot k) v_{\tilde{k}}^* \right]_j , \quad j = e, i . \end{aligned} \quad (9)$$

We note here that since the trapped particles execute bounce motion along the field line and $|\Omega| \ll \omega_{bj}$, perturbations along $\tilde{B}$ are averaged out and, therefore, only perpendicular dynamics needs to be considered.

For the untrapped particles, however, the dominant dynamics is parallel to $\tilde{B}$. Since $|\Omega| \ll |q_n v_{tj}|$, particles maintain the Boltzman distributions; i.e.,

$$n_j^u = -(1 - \sqrt{\epsilon}) N_j e_j (\phi + \psi_j) / T_j, \quad j = e, i, \quad (10)$$

where ψ_j is the effective potential of the parallel ponderomotive force; i.e.,

$$\psi_j = i F_{j\parallel} / q_n e_j, \quad j = e, i. \quad (11)$$

Using the above relations, we then can express n_j^t and n_j^u in terms of ϕ_+ , ϕ_- and ϕ_k . ϕ_+ and ϕ_- are in turn related to ϕ_k and $n_e = n_e^t + n_e^u$ through Eq. (3). Combining these relations and the quasi-neutrality condition

$$n_e^t + n_e^u = n_i^t + n_i^u, \quad (12)$$

it is straightforward to derive the dispersion relation of the dissipative trapped-ion mode modified by the lower-hybrid pump. Let us note that for the parameters of interest here, the dominant nonlinear contributions are due to those of the trapped ions and the untrapped electrons. The corresponding dispersion

relation is then given by

$$\epsilon_s = -\chi_i G - i\chi_e H / \Omega, \quad (13)$$

with

$$\epsilon_s = \chi_e + \chi_i, \quad (14)$$

$$\chi_e = 1 - \sqrt{\epsilon} (\Omega - \omega_e^*) / (\Omega + i\nu_e / \epsilon), \quad (15)$$

$$\chi_i = T [1 - \sqrt{\epsilon} (\Omega + \omega_i^*) / (\Omega + i\nu_i / \epsilon)], \quad (16)$$

where $\omega_j^* = q_{\perp} c \kappa T_j / eB$, $\kappa = -d \ln N / dx$ and $T = T_e / T_i$. The factors G and H are terms corresponding to the modification of the dispersion relation by the pump and are given by

$$G = \frac{q_{\perp}^2 \omega_{pi}^2}{(1 + \omega_{pe}^2 / \omega_{ce}^2)} \sum_{\tilde{k}} \left| \frac{c k_x \phi_{\tilde{k}}}{BC_s} \right|^2 \left(\frac{1}{\epsilon_+} + \frac{1}{\epsilon_-} \right), \quad (17)$$

and

$$H = \frac{\sqrt{\epsilon} q_{\perp}^2 \omega_{pi}^2}{(1 + \omega_{pe}^2 / \omega_{ce}^2)} \sum_{\tilde{k}} \left| \frac{c k \phi_{\tilde{k}}}{BC_s} \right|^2 \left(\frac{2 k_x k_y c_s^2}{\omega^2} \omega_{ci} + \omega_e^* \frac{\omega_{ci}}{\omega} \frac{k_x}{q_{\perp}} \right) \left(\frac{1}{\epsilon_+} + \frac{1}{\epsilon_-} \right). \quad (18)$$

Since the pump wave is a self-consistent lower-hybrid wave and $|\omega| \gg |\Omega|$, $|\mathbf{k}| \gg |\mathbf{q}|$; we can expand $\epsilon_{\pm} = \epsilon(\Omega \pm \omega, \mathbf{q} \pm \mathbf{k})$ about $(\omega, \mathbf{k})$ and obtain

$$\frac{1}{\epsilon_{+}} + \frac{1}{\epsilon_{-}} \approx \frac{(1 - \omega_{lh}^2/\omega^2)(q_{\parallel}^2 k_{\perp}^2/k_{\parallel}^2 - q_{\perp}^2) - (k_{\perp} \Omega/\omega)^2}{2 k_{\perp}^4 (\Omega - \mathbf{q} \cdot \mathbf{v}_{gk})^2}, \quad (19)$$

and

$$\mathbf{q} \cdot \mathbf{v}_{gk} = \omega \left(1 - \frac{\omega_{lh}^2}{\omega^2} \right) (q_{\parallel}/k_{\parallel} - q_{\perp} k_y/k_{\perp}^2). \quad (20)$$

From Eq. (19) it is noted that when $\Omega \approx \mathbf{q} \cdot \mathbf{v}_{gk}$, a new modulational type instability may be excited by the pump lower-hybrid waves [2]. Physically, one may interpret the excitation as a Landau type wave-"particle" interaction where the short-wavelength lower-hybrid waves behave as "particles" traveling with the appropriate group velocities and "resonate" with the phase velocity of propagating trapped-particle modes. Now, since $\omega \sim \omega_{pi} \gg |\Omega| \sim \sqrt{\epsilon} \omega_e^*$, one can see that the resonance condition $\Omega - \mathbf{q} \cdot \mathbf{v}_{gk} \approx 0$ is satisfied only when ω is very close to local ω_{lh} and/or $\mathbf{q}$ and $\mathbf{k}$ satisfy a rather special condition. We, therefore, concentrate on the non-resonant interaction and assume $|q_{\parallel}/k_{\parallel}|$, $|q_{\perp}/k_{\perp}| \gg |\Omega/\omega| (1 - \omega_{lh}^2/\omega^2)^{-1}$, which is approximately

$$(1 - \omega_{lh}^2/\omega^2) \gg \sqrt{\epsilon} (k p_1) (k_{\perp} \lambda_{Di}). \quad (21)$$

Note that $\sqrt{\epsilon} (k \rho_i) (k_{\perp} \lambda_{Di})$ is very small; usually $O(10^{-3})$ or less. With these assumptions, Eq. (19) reduces to

$$\left(\frac{1}{\epsilon_+} + \frac{1}{\epsilon_-} \right) \approx \frac{1}{2} \left(\frac{q_{\parallel}^2}{k_{\parallel}^2} - \frac{q_{\perp}^2}{k_{\perp}^2} \right) (k_{\perp} \omega)^{-2} \left(1 - \frac{\omega_{1h}^2}{\omega^2} \right)^{-1} \left(\frac{q_{\parallel}}{k_{\parallel}} - \frac{q_{\perp} k_y}{k_{\perp}^2} \right)^{-2}, \quad (22)$$

which is in general of order $(k_{\perp} \omega)^{-2} (1 - \omega_{1h}^2/\omega^2)^{-1}$.

Let us now examine the dispersion relation, Eq. (13), in more detail. Equation (13) can be written as

$$\chi_i (1 + G) + \chi_e (1 + iH/\Omega) = 0. \quad (23)$$

With $\Omega = \Omega_r + i\Omega_i$ and $|\Omega_r| \gg |\Omega_i|$, we have from the real part of Eq. (23)

$$\Omega_r = \sqrt{\epsilon} \omega_e^* (1 + G - \epsilon H/v_e) / (1 + T + TG), \quad (24)$$

and

$$\Omega_i = \frac{\epsilon^2 \omega_e^{*2}}{v_e (1+G)} \frac{(1+G-\epsilon H/v_e)^2}{(1+T+TG)^2} - H \frac{1+G-\epsilon H/v_e}{(1+G)(1+T+TG)} - \frac{v_i}{\epsilon}. \quad (25)$$

From Eq. (25), one can obtain the conditions for parametric excitation or suppression of the dissipative trapped-ion mode.

Let us assume that the pump is sufficiently weak so that

$|G|, |\epsilon H/v_e| < 1$, Eq. (25) then reduces to

$$\Omega_i \approx \Omega_{iL} - [H + (T-1)G\Omega_i^e]/(1+T) . \quad (26)$$

Here, with $\Omega_i^e = \epsilon^2 \omega_e^{*2}/v_e(1+T)^2$, $\Omega_{iL} = \Omega_i^e - v_i/\epsilon$ is the linear growth (damping) rate without the pump. Note that in obtaining Eq. (26) we use the condition $\epsilon^{3/2} \omega_e^* \ll v_e$. Equation (26), thus, indicates parametric suppression (excitation) if $[H+(T-1)G\Omega_i^e]$ is positive (negative). Note that in general one has $T = T_e/T_i > 1$. Furthermore, we see from Eqs. (17) and (18) that G and H have the same sign as that of $(1/\epsilon_+ + 1/\epsilon_-)$ given by Eq. (19). Hence, the lower-hybrid pump wave has a stabilizing effect for dissipative trapped-ion modes with

$$(q_{\parallel}/q_{\perp})^2 > (k_{\parallel}/k_{\perp})^2 \sim O(m_e/m_i) . \quad (27)$$

Modes with $(q_{\parallel}/q_{\perp})^2 < O(m_e/m_i)$ tend to be further destabilized by the pump wave.

For modes far away from the linear marginal stability condition, $\Omega_{iL} \approx \Omega_i^e$ and, because $|G| < 1$, Ω_i of Eq. (26) reduces to

$$\Omega_i \approx \Omega_i^e - H/(1+T) . \quad (28)$$

Thus, in this case parametric stabilization is primarily provided through the trapped ions. Complete stabilization occurs when $H > (1+T)\Omega_i^e$, or, from Eq. (18), we have approximately the following condition

$$\sum_{\mathbf{k}} \left| \frac{c \tilde{E}_{\mathbf{k}}}{B c_s} \right|^2 \left(2 \frac{k_{\perp}^2 c_s^2}{\omega^2} + \frac{\omega_e^*}{\omega} \frac{k_{\perp}}{q_{\perp}} \right) \frac{q_{\perp}^2 / k_{\perp}^2}{(1 - \omega_{lh}^2 / \omega^2)} \gtrsim \frac{\epsilon_e^{3/2} \omega_e^*}{v_{e ci}} \quad (29)$$

4. DISSIPATIVE TRAPPED-ELECTRON MODE

It is well known that for this mode ion dissipative effects and ion trapping effects are negligible. Because $|\Omega| \gg |q_{\perp} v_i|$, ion dynamics is primarily in the perpendicular (to $\mathbf{B}$) direction and, therefore, is similar to that of trapped ions discussed in the preceding section. Similar to the trapped-ion mode, electron parametric coupling to the pump is mainly through the parallel ponderomotive force acting on the untrapped electrons. Using the results obtained in the preceding section and Ref. [1], we have the following dispersion relation

$$\tilde{\epsilon}_s = -\tilde{\chi}_i G - i \tilde{\chi}_e \tilde{H} / \Omega, \quad (30)$$

where

$$\tilde{\epsilon}_s = \tilde{\chi}_e + \tilde{\chi}_i, \quad (31)$$

$$\tilde{\chi}_i = -\omega_e^* / \Omega, \quad (32)$$

$$\tilde{\chi}_e = 1 - \sqrt{\epsilon} \left\langle \frac{\omega - \omega_e^*}{\omega + i v_e / \epsilon} \right\rangle, \quad (33)$$

and

$$\tilde{\omega}_e^* = \omega_e^* [1 + \eta (m_e v^2 / 2T_e - 3/2)] . \quad (34)$$

Here, $\eta = d \ln T_e / d \ln N$ and $\tilde{\omega}_e^*$ is the electron drift frequency due to both density and temperature gradient. The angular bracket in Eq. (33) indicates averaging over the electron Maxwellian distribution f_{oe} [1]. G is given in Eq. (17) and

$$\tilde{H} = H / \sqrt{\epsilon} , \quad (35)$$

with H defined by Eq. (18).

Again, we are interested in the non-resonant interaction.

Assuming $\Omega = \Omega_r + i\Omega_i$ and $|\Omega_r| \gg |\Omega_i|$, we obtain from Eq. (30)

$$\Omega_r \approx \omega_e^* \left[1 + G + (\epsilon H / v_e) \langle 1 - \tilde{\omega}_e^* / \omega_e^* \rangle \right] , \quad (36)$$

and

$$\Omega_i = \tilde{\Omega}_{iL} - \epsilon H - \epsilon^{3/2} (\omega_e^*)^2 G / v_e . \quad (37)$$

Here, $\tilde{\Omega}_{iL} = \epsilon^{3/2} \omega_e^* (\tilde{\omega}_e^* - \omega_e^*) / v_e$ is the linear growth (damping)

rate. In deriving Eqs. (36) and (37), we have assumed

$|v_e| > |\epsilon \omega_e^*|$ and, as in the preceding section, the pump is weak so that $|G|, |\epsilon H / v_e| < 1$. Equation (37), again, indicates that

the lower-hybrid pump is stabilizing for modes with

$(q_{||} / q_{\perp})^2 > (k_{||} / k_{\perp})^2 \sim O(m_e / m_i)$. Otherwise, it has a further

destablizing effect. A condition similar to Eq. (29) for total stablization of this instability can also be derived from Eq. (37).

5. CONCLUSIONS

We have examined the influence of lower-hybrid waves (used for auxiliary heating of fusion plasmas) on dissipative trapped-particle modes (believed to be responsible for anomalous energy and particle transport). The primary mechanism of influence is the parametric coupling of the pump (or decay) lower-hybrid waves to trapped-particle modes through ponderomotive forces on ions and electrons. It is concluded that the influence may be either stabilizing or destabilizing. Roughly speaking, dissipative trapped-particle modes with $q_{\parallel}/q_{\perp} < (m_e/m_i)^{1/2}$ are further destabilized, whereas modes with $(q_{\parallel}/q_{\perp}) > (m_e/m_i)^{1/2}$ are stabilized. Now it is well-known that for trapped-electron modes $\omega \sim \omega^* > q_{\parallel} v_i$ i.e., $q_{\parallel}/q_{\perp} < (\rho_i \kappa)$, whereas for trapped-ion modes $q_{\parallel}/q_{\perp} > \sqrt{\epsilon} (\rho_i \kappa)$. Since $\rho_i \kappa \leq (m_e/m_i)^{1/2}$ for typical toroidal systems, it is concluded that a lower-hybrid pump wave will further destabilize all the trapped-electron modes and some fraction of trapped-ion modes, (those with $q_{\parallel}/q_{\perp}$ between $\sqrt{\epsilon} \rho_i \kappa$ and $(m_e/m_i)^{1/2}$); the majority of trapped-ion modes will, however, experience stabilization. Thus, in the relatively low temperature operation of toroidal systems where trapped-electron modes are expected to be significant, lower-hybrid waves may lead to an enhancement of the growth rate γ for dissipative low-frequency modes (and an associated increase in

anomalous transport), whereas in the higher-temperature regions, one expects beneficial effects of reduced γ (and anomalous transport) for the trapped-ion mode.

To really get a quantitative appreciation of the effects involved, one must undoubtedly do the calculation in more realistic geometry including such details as geometrical effects, mode structure effects etc. Here, we give an order of magnitude calculation based on the above slab model to show that for typical lower-hybrid pump fields, the effects are significant. From Eq. (29) we note that the relevant condition for order of magnitude estimates is

$$\left| \frac{CE}{BC_s} \right|^2 > \frac{\epsilon^{3/2} \omega_e^* \omega}{v_e \omega_{ci}} \frac{k_{\perp}}{q_{\perp}},$$

where we assume $\omega \gtrsim (2 \text{ to } 3\omega_{lh})$ everywhere in the plasma. Since the right side $\sim \epsilon^{3/2} (\omega_e^*/v_e) (\omega/\omega_{ci}) (k_{\perp}/q_{\perp})$ can readily be $\lesssim 0(1)$ and $(CE/B) \gtrsim 0(1)$ inside the lower-hybrid heated plasma, we see that conditions are appropriate for a significant modification of dissipative modes by the lower-hybrid pump waves.

Finally, it should be mentioned that the above calculations assume that the lower hybrid fields uniformly permeate the entire plasma. If the lower hybrid fields are strongly localized in space due to severe absorption effects, the above threshold calculations may have to be modified.

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