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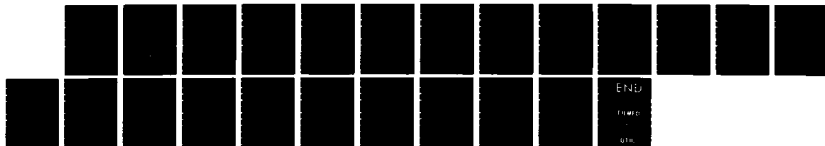
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Production and Control of Ion Cyclotron Instabilities in the High Latitude Ionosphere by High Power Radio Waves

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December 30, 1985

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PRODUCTION AND CONTROL OF ION CYCLOTRON INSTABILITIES IN THE HIGH LATITUDE IONOSPHERE BY HIGH POWER RADIO WAVES

I. INTRODUCTION

It is now recognized that many of the phenomena occurring in the auroral ionosphere may be associated with the presence of the field-aligned current-systems which are a constant feature at these latitudes [see e.g., Burke, 1981]. In addition to being a key component of the ionosphere-magnetosphere coupling circuit, these currents may also be responsible in generating micro-and/or macro-structures (turbulence) in the ambient plasmas at low altitudes. Such structures have been observed by various measurement techniques (ground based ionosonde, VHF radar backscatter, optical measurements, scintillations, in-situ rocket and satellite measurements, etc.) A catalogue of the structure scale sizes spans the domain from a few meters to a few hundreds of kilometers [see e.g., Aarons, 1973; Clark and Raitt, 1976; Dyson, 1969; Dyson and Winningham, 1974; Fremouw et al., 1977; Greenwald, 1974; Hanuise et al., 1981; Hower et al., 1966; Kelley et al., 1980; Ogawa et al., 1976; Olesen et al., 1976; Phelps and Sagalyn, 1976; Sagalyn et al., 1974; Vickrey et al., 1980; Weaver, 1965]. Recent theoretical studies have attempted to interpret these structures in terms of plasma instabilities, neutral turbulence, structured particle precipitation, convection flows, etc. [Keskinen and Ossakow, 1983]. Due to the ever-varying conditions prevalent in the auroral ionosphere, it seems likely that all of these processes may be occurring at one time or another. In some instances, different processes can result in similar observable effects which makes the resolution of individual mechanisms difficult and introduces complications in the modelling process. The use of active experiments in the near-earth space offers an

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attractive complement to the studies of understanding of natural plasma processes occurring in the ionosphere. One technique that has been widely used is the modification of the ionosphere using powerful radio waves [Utlaut, 1970; Fialer, 1974; Thide et al., 1982; Stubbe et al., 1981; Wong et al., 1981]. A considerable amount of theoretical work has been done in this field [Perkins and Kaw, 1971; Perkins et al., 1974; Fejer, 1979; Vaskov and Gurevich, 1977; Das and Fejer, 1979; Stenflo, 1983]. One of the possible consequences of the interaction of a large amplitude EM wave with the ionosphere is the generation of structures in the ambient plasma density. This may come about several ways; one of them being the excitation of plasma instabilities by an electromagnetic pump wave under sub-threshold conditions. By the same token, one could use the same principle of parametric coupling of the external pump and the natural plasma modes to suppress a current-driven or gradient-driven plasma instability in the system. Work on basic parametric instability processes in plasmas is exhaustive [Silin, 1965; DuBois and Goldman, 1965; Nishikawa, 1968]. The application of these ideas to the possible control of ionospheric irregularities was originally made by Lee et al. [1972] for the case of equatorial electrojet; and, has since been applied for the equatorial spread F situation [Bujarbarua and Sen, 1978] and the auroral ionosphere [Keskinen et al., 1983]. In this report, we examine the possibility of external control (excitation or suppression) of the current-driven ion-cyclotron instability in the auroral ionosphere. In addition to being excited in the topside ionosphere [Kindel and Kennel, 1971], the current-driven ion cyclotron instability may also be triggered in the collisional bottomside ionosphere [Chaturvedi, 1976; Satyanarayana et al., 1985; Fejer et al., 1984] by a radio-wave at the local upper-hybrid frequency. In the next section, we outline the theory, and, estimate the

incident power density required to stabilize or destabilize the instability. The final section considers application of these results to the bottomside high latitude F region ionosphere.

II. THEORY

We follow the theoretical procedure outlined in Lee et al. [1972] and Keskinen et al. [1983] for the calculation of the dynamic plasma response to a large amplitude electromagnetic pump wave. The equilibrium consists of a field-aligned current $J_0 \hat{z} (= -n_0 e V_0 \hat{z})$ and an oscillating electric field $\underline{E}_0 = \underline{E}_p \cos \omega_0 t$. We have equated the zero-order current to an equilibrium electron drift $V_0 \hat{z}$ and have used the so-called dipole approximation for the pump wave which is valid when the wavelength of the perturbations ($\sim k^{-1}$) is considered much smaller than the scale size of the variation of the pump-field (k_0^{-1}), i.e., $k_0 \ll k$. We shall use the two-fluid equations in describing the electron - and ion - fluid dynamics,

$$\frac{\partial}{\partial t} N_\alpha + \nabla \cdot (N_\alpha \underline{V}_\alpha) = 0 \quad (1)$$

$$N_\alpha M_\alpha \left(\frac{\partial}{\partial t} + \underline{V}_\alpha \cdot \nabla \right) \underline{V}_\alpha = - \nabla P_\alpha \pm e N_\alpha \left(\underline{E} + \frac{\underline{V}_\alpha \times \underline{B}_0}{c} \right) - \nu_\alpha M_\alpha N_\alpha \underline{V}_\alpha \quad (2)$$

$$P_\alpha = N_\alpha T_\alpha \quad (3)$$

$$\nabla \cdot \underline{E} = 4\pi e (N_i - N_e) \quad (4)$$

We initially assume cold ions ($T_i = 0$). Later, we will present results that include the ion-temperature effects ($T_i \neq 0$). In Eq. (1)-(4), M, N, V, P, T denote the mass, number density, velocity, pressure, and temperature, respectively, of the species α ; ν_α is the collision frequency of

species α with the neutrals, and e, c, E , and B represent electric charge, velocity of light, electric field and magnetic field, respectively. Since we restrict this study to bottom-side altitudes we take $\nu_{ii} \leq \nu_{in}$ where ν_{ii} is the ion-ion collision frequency and ν_{in} the ion-neutral collision frequency. In equilibrium, we have

$$\nabla \cdot N_{\alpha} \left(\underline{V}_{\alpha} + \frac{\underline{V}_{\alpha p}}{\omega_{\alpha}} \right) = 0 \quad (5)$$

$$0 = \pm \frac{e}{m_{\alpha}} \left(\underline{E}_0 + \frac{\underline{V}_{\alpha 0} \times \underline{B}}{c} \right) - \nu_{\alpha} \underline{V}_{\alpha 0} - \nabla p_{\alpha 0} / m_{\alpha} N_{\alpha} \quad (6)$$

$$\frac{d\underline{V}_{\alpha p}}{dt} = \pm \frac{e}{m_{\alpha}} \left(\underline{E}_p + \frac{\underline{V}_{\alpha p} \times \underline{B}}{c} \right) - \nu_{\alpha} \underline{V}_{\alpha p} \quad (7)$$

where $\alpha = e, i$, and $\underline{V}_0$ and $\underline{V}_p$ refer to drifts induced by the ambient and pump fields, respectively. Note that eq. (7) describes the approximate equilibrium oscillatory motion of plasma particles under the action of the RF electromagnetic field. Linearizing (1) - (4) for perturbations, $f_{\alpha} = f_{\alpha 0} + \delta f_{\alpha}$, and assuming $\delta f_{\alpha} \ll f_{\alpha 0}$, we obtain

$$\frac{\partial}{\partial t} \delta n_{\alpha} + (\underline{V}_{\alpha 0} + \underline{V}_{\alpha p}) \cdot \nabla \delta n_{\alpha} + \delta \underline{V}_{\alpha} \cdot \nabla N_0 + N_0 \nabla \cdot \delta \underline{V}_{\alpha} = 0 \quad (8)$$

$$\frac{\partial \delta \underline{V}_{\alpha}}{\partial t} + (\underline{V}_{\alpha 0} + \underline{V}_{\alpha p}) \cdot \nabla \delta \underline{V}_{\alpha} = \pm \frac{e}{m_{\alpha}} \left[\delta \underline{E} + \frac{\delta \underline{V}_{\alpha} \times \underline{B}}{c} \right] - \frac{\nabla \delta p_{\alpha}}{m_{\alpha} N_0} - \nu_{\alpha} \delta \underline{V}_{\alpha} \quad (9)$$

$$\nabla \cdot \delta \underline{E} = 4\pi e (\delta n_i - \delta n_e) \quad (10)$$

It is convenient to transform to an oscillating frame, defined by

$$\underline{r}'_j = \underline{r} - \underline{R}_j(t) \quad (11)$$

$$\frac{d\underline{R}_j(t)}{dt} = \underline{v}_{\alpha 0}$$

The plasma quantities in the oscillating frame are

$\delta \tilde{f}_\alpha(\tilde{\underline{r}}'_\alpha, t) = \delta f_\alpha(\underline{r}_\alpha, t)$ where the prime denotes the oscillating frame quantities. Equations (3) - (9) are similar in the oscillating frame, with substitutions $(V_{\alpha 0} + V_{\alpha p}) \rightarrow V_{\alpha 0}$, $\delta n_\alpha \rightarrow \delta \tilde{n}_\alpha$, $\delta \underline{v}_\alpha \rightarrow \delta \tilde{\underline{v}}_\alpha$, $\delta \underline{E} \rightarrow \delta \tilde{\underline{E}}$ and $\underline{r} \rightarrow \tilde{\underline{r}}_\alpha$. Assuming a harmonic perturbation, $\exp[i(\underline{k} \cdot \tilde{\underline{r}}_\alpha - \omega t)]$ with $\underline{k} = (k_x x + k_y y + k_z z)$ and $\omega = \omega_r + iY$, we may combine (3) and (9) and write

$$\delta \tilde{n}_e = \frac{ik\chi_e}{e} \delta \tilde{E}, \quad (12a)$$

$$\delta \tilde{n}_i = -i \frac{k\chi_i}{e} \delta \tilde{E} \quad (12b)$$

where

$$4\pi\chi_e = \frac{(\omega_{pe}^2/k^2)}{\left(\omega_e - \frac{v_e^2 k^2}{\omega} + \frac{v_e^2 k^2}{\omega_e^2 - \omega^2}\right)} \left(-\frac{k_z^2}{\omega} + \frac{k^2}{\omega} \frac{1}{\omega_e^2 - \omega^2}\right) \quad (13)$$

and

$$4\pi\chi_i = -\frac{\omega_{pi}^2(1 + i\frac{v_i}{\omega})}{[(\omega + iv_i)^2 - \omega_i^2]} \quad (14)$$

and

$$\omega_{p\alpha}^2 = \frac{4\pi n_0 e^2}{m_\alpha}, \quad \omega = \tilde{\omega}_e + iv_e, \quad \tilde{\omega}_\alpha = \omega - \underline{k} \cdot \underline{v}_0, \quad q_\alpha = \pm e, \quad \omega_\alpha = \frac{q_\alpha B_0}{m_\alpha c}$$

We have used $\nabla N_0 = 0$, $T_i = 0$ in deriving (13) - (14). One now transforms back to the laboratory frame and writes down the dispersion relation for the modes by using the Poisson's equation. This transformation has been discussed in Lee et al. [1972], and the dispersion relation, thus obtained, is

$$[1 + 4\pi(\chi_i + \chi_e)] = -J_1^2(\zeta) 4\pi \chi_i (1 + 4\pi \chi_e) \left[\frac{1}{H_e(\omega_+)} + \frac{1}{H_e(\omega_-)} \right] \quad (15)$$

where $\omega_{\pm} = \omega \pm \omega_0$ and,

$$J_n(\zeta) = \frac{i^{-n}}{\pi} \int_0^\pi e^{i\zeta \cos\theta} \cos n\theta d\theta \quad (16)$$

and ζ is defined by $\underline{k} \cdot \underline{R} = \zeta \sin(\omega_0 t + \beta)$, and, $H_e = - \frac{1}{(4\pi \chi_e)} (1 + 4\pi \chi_e)$. Note that we have ignored the effect of the pump wave on ions and have used the approximation $\zeta \ll 1$ in deriving (15) which states that the electron excursion length in the pump field is smaller than the perturbation wavelength. In absence of the pump field, the left hand side of (15) leads to the dispersion relation for ion-cyclotron waves [Chaturvedi, 1976],

$$(\omega + i\nu_i)^2 - \omega_i^2 - [C_s^2 k^2 - i \frac{m_e}{m_i} \frac{k^2}{k_\perp^2} v_e (\omega - k_z V_0)] (1 + i \frac{\nu_i}{\omega}) = 0 \quad (17)$$

which describes a resistive ion-cyclotron instability for modes with real frequency,

$$\omega_r^2 = \omega_i^2 + \langle C_s^2 \rangle^2 \quad (18a)$$

and growth rate

(18b)

$$\gamma \approx -\frac{1}{2} \frac{m_e}{m_i} \frac{k^2}{k_z^2} v_e \left(1 - \frac{k_z V_0}{\omega_r} \right) - \nu_i$$

where $C_s^2 = T_e/m_i$. An application of this instability for the auroral ionosphere has recently been discussed by Satyanarayana et al. [1985]. Physically, the instability is related to the dissipative effect on electrons in their parallel motion which impedes the "instantaneous" redistribution of the electron fluid to a Boltzmann-like distribution in the wave potential. In the collisionless limit, the dissipation is due to the wave-particle (Landau) resonance [Drummond and Rosenbluth, 1962], while in the collisional case, it is caused by the electron collisions with ions (or neutrals) [Chaturvedi and Kaw, 1975]. When the electron drift velocity exceeds the parallel wave phase velocity, the electrostatic ion cyclotron (EIC) wave becomes unstable.

In the presence of the pump-wave, the dispersion relation (17) is modified to

$$\left[(\omega + i\nu_i)^2 - \Omega_i^2 - \{k^2 C_s^2 + (m_e/m_i)(\hat{\omega}^2 - (k_z^2/k^2)\Omega_e^2)^{-1}(\hat{\omega} - \tilde{\omega}_e)(\Omega_e^2 - \hat{\omega}^2)\} \right. \\ \left. \left(1 + i \frac{\nu_i}{\omega} \right) \right] = A \quad (19)$$

which gives

$$\omega_r^2 = \Omega_i^2 + k^2 C_s^2 + A \quad (20)$$

and

$$\gamma = \frac{m_e}{2m_i} \frac{k^2}{k_z^2} \frac{v_e}{\Omega_i} (k_z v_0 - \bar{\Omega}_i - \frac{A}{2\bar{\Omega}_i}) - v_i \quad (21)$$

where

$$A = 2J_1^2(\zeta) (\omega_{pe}^2 \omega_{pi}^2) \frac{(\delta + \frac{v_e}{\omega_0} \beta)}{(\delta^2 + \beta^2)} \quad (22)$$

$$\bar{\Omega}_i^2 = \Omega_i^2 + c_s^2 k^2,$$

$$\delta = \omega_0^2 - \omega_{UH}^2, \quad \omega_{UH}^2 = \Omega_e^2 + \omega_{pe}^2 + v_e^2$$

$$\omega_{pe}^2 = \omega_{pe}^2 + v_e^2 k^2, \quad \beta = v_e \omega_0 (2 - \frac{\omega_{pe}^2}{\omega_0^2})$$

and,

$$\begin{aligned} \zeta^2 = & \frac{\omega_e^2}{\pi_e^2 (\omega_0^2 - \Omega_e^2)^2} \{ k_x^2 E_{px}^2 + k_y^2 E_{py}^2 + \frac{\Omega_e^2}{\omega_0^2} (k_y^2 E_{px}^2 + k_x^2 E_{py}^2) + \\ & 2E_{px} E_{py} [k_x k_y \cos \psi (1 - \frac{\Omega_e^2}{\omega_0^2}) - \frac{\Omega_e}{\omega_0} k_{\perp}^2 \sin \psi] \\ & + \frac{(\omega_0^2 - \Omega_e^2)}{\omega_0^2} k_z E_{pz} [k_x E_{px} - \frac{\Omega_e}{\omega_0} k_x E_{px} \sin \psi + k_y E_{py} \cos \psi] \} \end{aligned} \quad (23)$$

Here, the pump wave is taken as [Lee et al., 1972]

$$\underline{E}_p = \frac{E_{px}}{-2i} \exp(-i\omega_0 t) \hat{x} + \frac{E_{py}}{-2i} \exp(-i\omega_0 t - i\psi) \hat{y} + \frac{E_{pz}}{-2i} \exp(-i\omega_0 t) \hat{z} + c.c. \quad (24)$$

where, E_{pz} , is included as we are considering modes with $k_z \neq 0$. We have

used $\Omega_e/v_e > k/k_z$, $\omega_p > \gamma$, v_{in} , and, $1 > A/\bar{\Omega}_i^2$, in deriving (20) - (21). We can rewrite the modified growth rate of the modes as

$$\gamma = \frac{-k^2}{2k_z^2} \frac{v_e}{\Omega_e} (k_z v_0 - \bar{\Omega}_i) - v_i] - \frac{k^2}{2k_z^2} \frac{v_e}{\Omega_e} \frac{(z^2/16)}{\bar{\Omega}_i} \frac{(\omega_{pe}^2 \omega_{pi}^2) \delta}{(\delta^2 + \beta^2)} = \gamma^L - \gamma_s \quad (25)$$

where γ^L denotes the linear growth contribution in absence of the pump-field while γ_s represents the modification introduced by the pump. We have again used $z \ll 1$ in going from (21) to (25). We see from inspection of (25) that the presence of the pump can be stabilizing (destabilizing) if $\delta > 0$ (< 0), or, when $\omega_0 > \omega_{UH}$ ($\omega_0 < \omega_{UH}$). We note that similar criteria were also obtained by earlier workers [Lee et al., 1972, Bujarbarua and Sen, 1978, Keskinen et al., 1983].

We can treat the finite ion temperature effects in a similar fashion. Since the ion-cyclotron instability has maximum growth rate in the ionosphere near $k_{\perp} \rho_i \sim 1$, [Satyanarayana et al., 1985; Kindel and Kennel, 1971] one needs to use the kinetic (collisionless Vlasov-Boltzmann) equation for ions [Dum and Dupree, 1970]. In this case, the modified dispersion relation becomes,

$$\omega = \Omega_i \left[1 + \Gamma_{1i} \tau \left\{ \left(1 - i \frac{v_e \bar{\omega}_e}{k_z^2 v_e^2} \right) + 2J_1^2(z) \frac{\omega_{pe}^4}{k^2 v_e^2} \frac{(\delta + \frac{v_e}{\omega_0})}{(\delta^2 + \beta^2)} \right\} \left\{ 1 - i\sqrt{\pi} \frac{\omega t}{\Omega_i} e^{-t^2} \right\} - i(v_i/\Omega_i) \right] \quad (26)$$

where,

$$\tau = \frac{\Gamma_e}{\Gamma_i}; \quad t = \left(\frac{\omega - \Omega_i}{k_z v_i} \right); \quad \Gamma_n(\mu) = I_n(\mu) e^{-\mu^2}; \quad \mu = \frac{k_{\perp} v}{\Omega_n} \quad (27)$$

In the absence of the pump wave, this expression yields the ion-cyclotron wave dispersion relation,

$$\omega = \Omega_i [1 + \Gamma_{1i} \tau (1 - i \frac{v_e \bar{\omega}_e}{k_z^2 v_e^2} - i \sqrt{\pi} \frac{\omega}{\Omega_i} \bar{e}^{-t^2}) - i(v_i/\Omega_i)] \quad (28)$$

which yields the expressions for the real frequency and the growth rate

$$\begin{aligned} \omega_r &= \Omega_i [1 + \Gamma_{1i} \tau], \\ \gamma &= -\Omega_i \Gamma_{1i} \tau \left[\frac{v_e \bar{\omega}_e}{k_z^2 v_e^2} + \sqrt{\pi} \frac{\omega_r}{k_z v_i} \Gamma_{1i} \bar{e}^{-t^2} \right] - \nu_i - \nu_{ii} k_{\perp}^2 \rho_i^2 \end{aligned} \quad (29)$$

where the second term in (29) in the expression for γ results from the ion-Landau damping and we have added a contribution (last term) due to ion-viscous damping [Dum and Dupree, 1970] for completeness. The collisional ion-cyclotron instability described by (29) may have relevance for the short wavelength observations (such that $k_{\perp} \rho_i \sim 1$) of ion-cyclotron waves at low ionospheric altitudes [Bering, 1984; Satyanarayana et al., 1985].

The inclusion of pump wave effects in equation (28) leads to modification of (29) in the form given by

$$\begin{aligned} \omega_r &= \Omega_i \left[1 + \Gamma_{1i} \left(\tau + \frac{A}{k_z^2 v_i^2} \right) \right], \\ \gamma &= \gamma_{\perp} - (\Gamma_{1i})^2 \frac{\tau}{\mu_i} \frac{v_e}{k_z^2 v_e^2} A \end{aligned} \quad (30)$$

where A and $\gamma_{\perp}$ are given by (22) and (29) respectively. We see that the pump wave is stabilizing (destabilizing) for $\delta > 0$ (< 0), i.e.,

when $\omega_0 > \omega_{UH}$ ($\omega_0 < \omega_{UH}$); as in case of fluid ions (eq. (25)). One can readily extend this analysis to the case of the kinetic EIC instability [Drummond and Rosenbluth, 1962]. However, in the ionosphere, the kinetic instability is likely to occur at higher altitudes [Kindel and Kennel, 1971] such as above the F peak; whereas, the modification experiments have access only to the lower collisional altitudes.

A rough estimate for the threshold amplitude of the pump-field required to stabilize or destabilize the collisional EIC instability, reveals that (using (30)),

$$\frac{cE_p^2}{3\pi} = \frac{\mu_i}{\Gamma_{ii}} \frac{B_0^2}{ck_{\perp}^2} \frac{\omega_0^2(\omega_0 - \Omega_e)^2}{\Omega_e^2} \frac{k_z^2 v_e^2 \omega_0 \Omega_e}{\omega_{pe}^4} \left[\frac{v_e(k_z V_0 - \omega_r)}{k_z^2 v_e^2} - \pi^{1/2} \frac{\omega_r}{k_z v_i} \Gamma_{ii} e^{-t^2} \right. \\ \left. - (\Gamma_{ii} \tau \Omega_i)^{-1} (v_i + v_{ii} k_{\perp}^2 \rho_i^2) \right] \quad (31)$$

where t is given by (27) and we have considered a circularly-polarized O-mode, so $\psi = -\pi/2$ and $|E_{px}| = |E_{py}|$. For typical high-latitude F-region parameters [Satyanarayana et al., 1985], we have, $B_0 \sim 0.5$ G, $v_e/\Omega_e \sim 10^{-4}$, $\frac{k_z}{k_{\perp}} \sim 10^{-2}$, $k_{\perp} \rho_i \sim 1.5$, $\rho_i \sim 10^3$ cm, $\omega_0 \sim \omega_{pe} \sim 3 \times 10^7$ s $^{-1}$, $T_e = T_i = 0.1$ eV, $\rho_e \approx 3$ cm, $V_0 \sim 2 \times 10^6$ cm s $^{-1}$, $v_i = 0.1$ sec $^{-1}$, $v_{ii} = 0.01$ sec $^{-1}$, giving that $c|E_p|^2/8\pi \approx 5 \times 10^{-4}$ W/m 2 .

III. DISCUSSION

We have studied the effects of a high frequency electromagnetic pump wave on the evolution of ion cyclotron instabilities in a weakly ionized collisional plasma. We find that an electromagnetic pump wave oscillating

with the upper hybrid frequency can lead to the artificial production or control of ion cyclotron instabilities in a collisional plasma. The basic physical mechanism of this process can be expressed as follows. The electromagnetic pump wave with frequency ω_{UH} interacts with low frequency ion cyclotron modes ω and drives sidebands at $\omega \pm \omega_{UH}$. These sidebands can, in turn, couple to the pump fields and nonlinearly affect the ion cyclotron modes. This interaction is ponderomotive in nature. However, the partial pressure force (Fejer, 1979) can also play a role depending upon the $k_{\perp}$ of the sidebands. Fejer [1979] has shown that the partial pressure force exceeds the ponderomotive force roughly when $\lambda_{\perp} > 4\lambda_e$, where λ_e is the electron mean free path. For the collisional EIC, $(k_z/k_{\perp}) \sim 0.06$, $(k_{\perp}\rho_i) \sim 1.5$ and $\lambda_e \sim 1$ km. [Satyanarayana et al., 1985], which implies, $\lambda_{\perp} < \lambda_e$. Thus, for the ion-cyclotron modes, the predominant pump-wave-induced parametric effects are likely to be of ponderomotive force-like in nature. For parameters typical of the high latitude ionosphere, we find that ion cyclotron modes can be stabilized or destabilized with an incident pump electric field on the order of 0.1 V/m. The polarization of the pump wave should be O mode. These power density levels could be accessible using the ionospheric heaters in Alaska (Wong et al., 1981) and Norway (Stubbe et al., 1981).

We have made several simplifying assumptions in the course of this paper. We have assumed the dipole approximation, i.e., taken the wavelength of the pump wave $\sim k_p^{-1}$ to be much larger than the wavelength of the ion cyclotron modes k^{-1} . Since $k_p^{-1} \sim 100$'s m and $k^{-1} \sim \rho_i \sim 10$ m, this assumption is valid. In addition, we have neglected pump electric field swelling near the reflection point (Ginzburg, 1964) and, as a result, our threshold pump electric fields should be considered as lower bounds.

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