Indian Journal of Radio & Space Physics
Vol. 18, June 1989, pp. 90-94

Plasma maser theory of ordinary mode radiation

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Received 23 February 1988; revised received 1 September 1988

The plasma maser theory of ordinary mode radiation, in the presence of a coherent ion cyclotron wave in a magnetized plasma, due to a high frequency nonlinear force is studied. This nonlinear force which arises from the resonant interaction between electrons and the modulated electric fields caused by coupling between the coherent ion cyclotron wave field and the test high frequency ordinary mode wave field, is calculated. The growth rate of the ordinary mode radiation is obtained and application of the results to the Jovian kilometric radiation is stressed.

1 Introduction

There has been a great deal of interest in recent years in a new maser effect in plasma turbulence, viz., induced bremsstrahlung instability, because of its potential application to numerous space and laboratory plasma radiation phenomena[1]. This new process belongs to the lowest order mode-mode coupling effect in weak plasma turbulence. According to the standard weak turbulence theory, the lowest order mode-mode coupling processes are composed of three parts, namely, (i) the resonant three-wave interaction, (ii) the nonlinear Landau damping and (iii) the plasma maser interaction. The induced bremsstrahlung instability occurs for $\omega = kv$ with $\Omega \neq Kv$, where (ω, k) and (Ω, K) are the frequency and the wave number of the low frequency pump field and the high frequency radiation field, respectively and v is the electron velocity. The plasma maser theory has recently been applied successfully to a number of observations in the space and laboratory plasma[2-6].

The plasma maser effect is basically a frequency-upconversion process in which there is a flow of energy from a low frequency mode to a high frequency mode in a plasma. Evidence of frequency upconversion was also recently reported in the laboratory by Mori and Ohya[7] who detected emission of upper hybrid soliton from ion wave in a beam-plasma system. Pottelette *et al.*[8] reported the emission of very intense bursts of broadband electrostatic noise from the plasma frequency with frequencies up to the local electron cyclotron frequency in the source region of the auroral kilometric radiation (AKR) and suggested possible links between these electrostatic bursts and the AKR generation.

Recently, it has been shown that the induced bremsstrahlung instability mechanism can be best understood in terms of a high frequency nonlinear force[3]. This high frequency nonlinear force arises from the resonant interaction between electrons and modulated electric fields caused by the coupling of low frequency fields with the high frequency ES or EM fields present in the system. Thus electrons suffer acceleration (or deceleration) due to this nonlinear force, and the accelerated electrons can radiate ES or EM wave. In contrast to the parametric interaction process where a low frequency nonlinear force is produced by coupling between a high frequency pump wave and a low frequency wave to make a low frequency wave unstable, the nonlinear force in this mechanism is a high frequency one and makes the high frequency wave unstable.

In this paper we consider the emission of ordinary mode radition in a magnetized plasma containing low frequency ion cyclotron wave excited by a current along the applied magnetic field.

2 Calculation of nonlinear force

In order to calculate the nonlinear force, we need an expression for the modulated electric fields and the electron distribution function due to the low frequency fluctuations. For this purpose, we consider a homogeneous magnetized plasma in the presence of a coherent ion cyclotron wave propagating almost perpendicular to the external magnetic field, $\mathbf{B}_0 = \hat{z} B_0$. The basic equations governing the interaction of the low frequency fields with an electron which leads to the ordinary mode radiation are the set of Vlasov-Maxwell's equations,

$$\left(\frac{\partial}{\partial t}+\mathbf{v}\cdot\frac{\partial}{\partial \mathbf{r}}-\frac{e}{m}\left[\mathbf{E}(\mathbf{r},t)+\frac{1}{c}\mathbf{v}\times\mathbf{B}(\mathbf{r},t)\right]\cdot\frac{\partial}{\partial \mathbf{v}}\right)f_e(\mathbf{r},\mathbf{v},t)=0 \quad \ldots(1)$$

$$\nabla\times\mathbf{E}=-\frac{1}{c}\frac{\partial \mathbf{B}}{\partial t} \quad \ldots(2)$$

$$\nabla\times\mathbf{B}=\frac{1}{c}\frac{\partial \mathbf{E}}{\partial t}+\frac{4\pi}{c}\mathbf{J} \quad \ldots(3)$$

$$\mathbf{J}=-en\int \mathbf{v}f_e(\mathbf{r},\mathbf{v},t)d\mathbf{v} \quad \ldots(4)$$

where f_e is the electron distribution function and other notations are standard ones.

The unperturbed electron distribution function and the electric fields are

$$F_{0e}=f_{0e}+\varepsilon f_{1e}+\varepsilon^2 f_{2e} \quad \ldots(5)$$

$$\mathbf{E}_{0l}=\varepsilon\mathbf{E}_l \quad \ldots(6)$$

where f_{0e} is the space and time averaged part of the distribution function, f_{1e} the fluctuating part of the electron distribution function due to the low frequency field and f_{2e} its higher order fluctuating part. Here ε is a small parameter and is associated with the amplitude of the low frequency ion cyclotron fluctuations. Eq. $\mathbf{E}_l=\mathbf{E}_l\cos(k_\parallel z+k_\perp x-\omega t)$ is the electrostatic ion cyclotron wave field which is assumed to be in the xz-plane with $k_\perp \gg k_\parallel$. Parameters ω and $\mathbf{k}$ are the frequency and the wave vector and ($\perp$, $\parallel$) refer to directions perpendicular and parallel to the applied magnetic field, respectively.

To the order of ε, we obtain from Eq. (1)

$$\left(\frac{\partial}{\partial t}+\mathbf{v}\cdot\frac{\partial}{\partial \mathbf{r}}-\frac{e}{mc}\mathbf{v}\times\mathbf{B}_0\cdot\frac{\partial}{\partial \mathbf{v}}\right)f_{1e}=\frac{e}{m}\mathbf{E}_l\cdot\frac{\partial}{\partial \mathbf{v}}f_{0e} \quad \ldots(7)$$

For low frequency ion cyclotron wave, electron motion along the magnetic field is important. The Fourier component of the corresponding electron distribution function is obtained from Eq. (7) as

$$f_{1e}(\mathbf{k},\omega)=\frac{(\mathrm{i}\,e/m)\,E_{l\parallel}(\mathbf{k},\omega)\,(\partial/\partial v_\parallel)\,f_{0e}}{2\,(\omega-k_\parallel v_\parallel+\mathrm{i}0)} \quad \ldots(8)$$

Here i0 represents the small imaginary part of ω.

We now perturb the steady state by the high frequency ordinary mode test wave fields $\mu\delta\mathbf{E}_h$ and $\mu\delta\mathbf{B}_h$. Here μ represents the ordering of the high frequency perturbation and we have assumed that $\mu \ll \varepsilon$. The total perturbed electric and magnetic fields and the electron distribution function according to the linear response theory of a turbulent plasma[9] can be written as

$$\delta\mathbf{E}=\mu\delta\mathbf{E}_h+\mu\varepsilon\delta\mathbf{E}_{lh}+\mu\varepsilon^2\Delta\mathbf{E} \quad \ldots(9)$$

$$\delta\mathbf{B}=\mu\delta\mathbf{B}_h+\mu\varepsilon\delta\mathbf{B}_{lh}+\mu\varepsilon^2\Delta\mathbf{B} \quad \ldots(10)$$

$$\delta f=\mu\delta f_h+\mu\varepsilon\delta f_{lh}+\mu\varepsilon^2\Delta f \quad \ldots(11)$$

where $\delta\mathbf{E}_{lh}$ and $\delta\mathbf{B}_{lh}$ are the electric and magnetic fields, respectively, of the mixed mode, $\Delta\mathbf{E}$ and $\Delta\mathbf{B}$ the next higher order perturbed electric and magnetic fields, respectively, δf_{lh} is the perturbed electron distribution function due to the mixed mode and Δf the next higher order perturbed electron distribution function. The ordinary mode is assumed to propagate along x direction with wave vector $\mathbf{K}=(K,0,0)$.

We now linearize Eq. (1) for $\mathbf{E}_l \gg \delta\mathbf{E}$ to obtain

$$\left(\frac{\partial}{\partial t}+\mathbf{v}\cdot\frac{\partial}{\partial \mathbf{r}}-\frac{e}{m}\left[\varepsilon\mathbf{E}_l+\frac{1}{c}\mathbf{v}\times\mathbf{B}_0\right]\cdot\frac{\partial}{\partial \mathbf{v}}\right)\delta f$$
$$-\frac{e}{m}\left(\delta\mathbf{E}+\frac{1}{c}\mathbf{v}\times\delta\mathbf{B}\right)\cdot\frac{\partial}{\partial \mathbf{v}}(f_{0e}+\varepsilon f_{1e}+\varepsilon^2 f_{2e})=0 \quad \ldots(12)$$

To the order of μ, $\mu\varepsilon$ and $\mu\varepsilon^2$, we obtain from Eq. (12)

$$P\delta f_h-\frac{e}{m}\left(\delta\mathbf{E}_h+\frac{1}{c}\mathbf{v}\times\delta\mathbf{B}_h\right)\cdot\frac{\partial}{\partial \mathbf{v}}f_{0e}=0 \quad \ldots(13)$$

$$P\delta f_{lh}-\frac{e}{m}\mathbf{E}_l\cdot\frac{\partial}{\partial \mathbf{v}}\delta f_h-\frac{e}{m}\left(\delta\mathbf{E}_h+\frac{1}{c}\mathbf{v}\times\delta\mathbf{B}_h\right)\cdot\frac{\partial}{\partial \mathbf{v}}f_{1e}$$
$$-\frac{e}{m}\left(\delta\mathbf{E}_{lh}+\frac{1}{c}\mathbf{v}\times\delta\mathbf{B}_{lh}\right)\cdot\frac{\partial}{\partial \mathbf{v}}f_{0e}=0 \quad \ldots(14)$$

$$P\Delta f-\frac{e}{m}\left(\mathbf{E}_l\cdot\frac{\partial}{\partial \mathbf{v}}\delta f_{lh}+\left[\delta\mathbf{E}_{lh}+\frac{1}{c}\mathbf{v}\times\delta\mathbf{B}_{lh}\right]\cdot\frac{\partial}{\partial \mathbf{v}}f_{1e}\right)$$
$$-\frac{e}{m}\left(\Delta\mathbf{E}+\frac{1}{c}\mathbf{v}\times\Delta\mathbf{B}\right)\cdot\frac{\partial}{\partial \mathbf{v}}f_{0e}=0 \quad \ldots(15)$$

where the operator

$$P\equiv\frac{\partial}{\partial t}+\mathbf{v}\cdot\frac{\partial}{\partial \mathbf{r}}-\frac{e}{mc}\mathbf{v}\times\mathbf{B}_0\cdot\frac{\partial}{\partial \mathbf{v}}$$

Introducing transforms in space and time for the various quantities according as

$$A(\mathbf{r},\mathbf{v},t)=\sum_{\mathbf{K},\Omega}A(\mathbf{v},\mathbf{K},\Omega)\exp[i(\mathbf{K}\cdot\mathbf{r}-\Omega t]\qquad\ldots(16)$$

and integrating along the orbits of the particles in the unperturbed fields[10], we solve Eqs (13)-(15) to obtain the mixed mode electric field[11] as

$$\delta E_{lh}(\mathbf{K},\Omega)=-D_0^{-1}(\mathbf{K},\Omega)\frac{\omega_{pe}^2}{\Omega}\int 2\pi v_\perp dv_\perp dv_\| v_\|$$

$$\times\sum_{a=-\infty}^{\infty}\frac{J_a^2(Kv_\perp/\Omega_e)}{\Omega-a\Omega_e}\Big[\delta E_h(\mathbf{K}-\mathbf{k},\Omega-\omega)\frac{\partial}{\partial v_\|}f_{1e}(\mathbf{k},\omega)$$

$$+\delta E_h(\mathbf{K}+\mathbf{k},\Omega+\omega)\frac{\partial}{\partial v_\|}f_{1e}(-\mathbf{k},-\omega)\Big]\qquad\ldots(17)$$

where $D_0(\mathbf{K},\Omega)$ is the linear dielectric constant of the ordinary mode given by

$$D_0(\mathbf{K},\Omega)=1-\left(\frac{cK}{\Omega}\right)^2+\frac{\omega_{pe}^2}{\Omega}\int 2\pi v_\perp dv_\perp dv_\| v_\|$$

$$\times\sum_{a=-\infty}^{\infty}\frac{J_a^2(Kv_\perp/\Omega_e)}{\Omega-a\Omega_e}\left[\frac{a\Omega_e}{\Omega v_\perp}\left(v_\|\frac{\partial}{\partial v_\perp}-v_\perp\frac{\partial}{\partial v_\|}\right)\right.$$

$$\left.\times f_{0e}+\frac{\partial}{\partial v_\|}f_{0e}\right]\qquad\ldots(18)$$

Here ω_{pe} and Ω_e are, respectively, the electron plasma and cyclotron frequencies, and $J_a\ (Kv_\perp/\Omega_e)$ is the Bessel function. In deriving Eq. (7) we have neglected the contribution of the terms $\delta\mathbf{B}_h$ and $\delta\mathbf{B}_{lh}$, because the induced bremsstrahlung arises from electron acceleration through modulation electric fields.

From Eq. (15), retaining the most dominant nonlinear term in the high frequency perturbation[12], the high frequency nonlinear force $\mathbf{F}_N$ can be written in the Fourier space as follows[3].

$$\mathbf{F}_N(\mathbf{K},\Omega)=en_0\int\delta\mathbf{E}_{lh}(\mathbf{K}-\mathbf{k},\Omega-\omega)\cdot\frac{\partial}{\partial\mathbf{v}}f_{1e}(\mathbf{k},\omega)\mathbf{v}d\mathbf{v}\qquad\ldots(19)$$

Substituting Eqs (17) and (8) for $\delta E_{lh}(\mathbf{K}-\mathbf{k},\Omega-\omega)$ and $f_{1e}(\mathbf{k},\omega)$ respectively, we obtain from Eq. (19) the z-component of the nonlinear force as

$$F_{NZ}(\mathbf{K},\Omega)=f\delta E_h(\mathbf{K},\Omega)\qquad\ldots(20)$$

where

$$f=-en_0D_0^{-1}(\mathbf{K}-\mathbf{k},\Omega-\omega)$$

$$\times\frac{|E_{l\|}(\mathbf{k},\omega)|^2\omega_{pe}^2\pi^2v_e^2}{2\Omega^3k_\|}\left(\frac{m}{2\pi T}\right)^3\left\{1+\eta Z(\eta)\right\}\qquad\ldots(21)$$

Here $v_e=(2T/m)^{1/2}$ is the electron thermal velocity, T the electron temperature and $Z(\eta)$ the plasma dispersion function, where $\eta=(\omega-k_\|v_0)/k_\|v_e$. In deriving Eq. (20) we have assumed the electron distribution function f_{0e} to be a shifted Maxwellian f_{0e} which is given by

$$f_{0e}=(m/2\pi T)^{3/2}\exp[-m(v_\|-v_0)^2/2T]\exp[-mv_\perp^2/2T]\qquad\ldots(22)$$

where v_0 is the drift velocity of the electron.

3 Growth rate of ordinary mode

As has already been mentioned, the plasma maser effect can be explained in terms of a high frequency nonlinear force. We now use this nonlinear force [Eq. (19)] in the equation of motion for electron to obtain[3]

$$mn\left[\frac{\partial\mathbf{v}}{\partial t}+(\mathbf{v}\cdot\nabla)\mathbf{v}\right]=-en\left[\mathbf{E}+\frac{1}{c}\mathbf{v}\times\mathbf{B}\right]+\mathbf{F}_N\qquad\ldots(23)$$

Maxwell's equations are:

$$\nabla\times\mathbf{E}=-\frac{1}{c}\frac{\partial\mathbf{B}}{\partial t}$$

$$\nabla\times\mathbf{B}=\frac{1}{c}\frac{\partial\mathbf{E}}{\partial t}+\frac{4\pi}{c}\mathbf{J}$$

$$\mathbf{J}=-en\mathbf{v}\qquad\ldots(24)$$

where the symbols have their usual meanings. After linearization and Fourier analyses we get from Eqs (23) and (24),

$$\mathbf{v}=-\frac{ie}{\Omega m}\delta\mathbf{E}_h(\mathbf{K},\Omega)+\frac{i\mathbf{F}_N(K,\Omega)}{mn_0\Omega}\qquad\ldots(25)$$

$$(\Omega^2-c^2K^2)\delta\mathbf{E}_h(\mathbf{K},\Omega)=4\pi i\,en_0\Omega\mathbf{v}\qquad\ldots(26)$$

Using Eqs (20) and (25), we finally obtain from Eq. (26) the dispersion relation for the ordinary mode as

$$\Omega^2=c^2K^2+\omega_{pe}^2\left(1-\frac{f}{en_0}\right)\qquad\ldots(27)$$

The ordinary mode instability can be obtained by puting $\Omega=\Omega_r - i\gamma$ in Eq. (27), where Ω_r is the real frequency, and γ the growth rate of the ordinary mode. The real frequency and growth rate are obtained, neglecting the nonlinear frequency shift as

$$\Omega_r^2 = c^2K^2 + \omega_{pe}^2 \quad \ldots (28)$$

$$\gamma = -\,\omega_{pe}^2\, I_m\, f/2en_0\Omega \quad \ldots (29)$$

Here I_m stands for the imaginary part.

Case I

For $\Omega \approx \omega_{pe}$ and $\omega \ll \Omega$, $k_\parallel \ll k_\perp$, the linear dielectric function can be written as

$$D_0(\mathbf{K}-\mathbf{k}, \Omega-\omega) \approx c^2(2Kk_\perp - k_\perp^2)/\Omega^2 \quad \ldots (30)$$

In the long wavelength limit ($K \to 0$)

$$D_0(\mathbf{K}-\mathbf{k}, \Omega-\omega) \approx -\,c^2k_\perp^2/\Omega^2 \quad \ldots (31)$$

For $|(\omega - k_\parallel v_0)|/k_\parallel v_e \ll 1$,

$$I_m Z(\eta) \approx i(\pi/2)^{1/2} \quad \ldots (32)$$

Substituting Eqs (31) and (32) in Eq. (29), we obtain the growth rate of the ordinary mode in the long wavelength limit for $\Omega \approx \omega_{pe}$ as

$$\gamma/\omega_{pe} = \left(\frac{2}{\pi}\right)^{1/2}\left(\frac{v_0}{v_e}\right)\left(\frac{\Omega}{ck_\perp}\right)^2\left(\frac{k_e}{k_\parallel}\right)\frac{|E_{l\parallel}(\mathbf{k},\omega)|^2}{8\pi n_0 T} \quad \ldots (33)$$

where k_e is the electron Debye wave number.

Case II

For $\Omega \approx \Omega_e$, the linear dielectric function can be expanded as

$$D_0(\mathbf{K}-\mathbf{k},\Omega-\omega) = -\left(\frac{\omega_{pe}}{\Omega_e}\right)^2 \frac{k_\perp^2 v_e^2}{\Omega^2 - \Omega_e^2} \quad \ldots (34)$$

Substituting Eqs (32) and (34) in Eq. (29), we obtain the growth rate of the ordinary mode in the long wavelength limit for $\Omega \approx \Omega_e$ and $\Omega_e \gg \omega_{pe}$ as

$$\gamma/\Omega_e = \left(\frac{2}{\pi}\right)^{1/2}\left(\frac{v_0}{v_e}\right)\left(\frac{k_e}{k_\perp}\right)^2 (k_e\rho_e)\,\frac{|E_l(\mathbf{k},\omega)|^2}{8\pi n_0 T} \quad \ldots (35)$$

where we assumed $K\rho_e \approx 1$, ρ_e being the electron gyroradius.

4 Discussion

We have thus obtained the growth rates of the ordinary mode radiation for $\Omega \approx \omega_{pe}$ and $\Omega \approx \Omega_e$. We have shown that the high frequency nonlinear force which arises from the resonant interaction between the modulation fields and the electrons drives the instability. The emission mechanism of the high frequency ordinary mode radiation is explained on the basis of this nonlinear force.

Recent detailed observations of Voyager data[13,14] have shown that the broadband Jovian kilometric radiation, the bKOM, and the narrowband Jovian kilometric radiation, the nKOM (Refs 15-17) are in the extraordinary and ordinary modes, respectively, and their source region is probably the plasma torus of the Jupiter's satellite Io. Further, strong plasma wave activity is also reported from this region[17].

We now apply the results of our investigation to estimate the growth rate of the Jupiter's narrowband kilometric radiation, which is emitted with the peak frequency around 100 kHz, having the probable source region at the outer edge of the Io plasma torus where the plasma frequency is about 100 kHz (Refs 14,18). We take the following probable parameters: $\lambda = 10$ km, $K = 2\pi \times 10^{-6}$, $k_\perp = 4K$, $k_\perp = 20k_\parallel$, $n = 120\ \text{cm}^{-3}$, $v_0 = 0.5v_e$, $T = 50$ eV and $E_{l\parallel} = 20$ mV/m. With these parameters, the growth rate of the ordinary mode radiation [Eq. (33)] becomes $\gamma/\omega_{pe} \approx 10^{-3}$.

The generation of ordinary mode radiation in a magnetoplasma in the presence of a coherent ion cyclotron wave was previously studied by Nambu[11]. However, the result was erroneous because of the mistake in calculating the imaginary part contribution in the nonlinear dielectric function[19]. In our study, the contribution from the polarization term has been correctly computed. Further, the growth rate of the ordianry mode radiation has been obtained from the consideration of the high frequency nonlinear force which drives the instability.

5 Conclusions

As the mechanism discussed in this paper requires the presence of a low frequency wave for the amplification of a high frequency wave, and as this essential condition is satisfied in the probable source region of the Jupiter's kilometric radiation as reported from observation, we believe that our mechanism may be an efective process for the generation of the Jupiter's kilometric radiation. More specifically, the theory presented in this paper may be applicable to the nKOM which is in the ordinary mode.

In the present study we have assumed a stationary turbulent state for the low frequency fluctuations. This assumption remains valid as long as the time of interaction is small compared to the characteristic time scale of the low frequency fluctuations.

Finally, we comment on the paper by DuBois and Pesme[20]. For unmagnetized plasma, the most dominant polarization contribution vanishes because of the cancellation of the polarization terms for electrostatic wave. However, the cancellation does not occur for electromagnetic wave for an unmagnetized plasma as well as for a magnetized plasma and, therefore, the most dominant contribution for plasma maser effect comes from the polarization term[12].

Acknowledgement

The authors are grateful to the referees pointing out a couple of references. This study was partially supported by a research grant from the Science, Technology and Environment Department, Government of Assam.

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