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QUASILINEAR THEORY OF THE ORDINARY-MODE ELECTRON-
CYCLOTRON RESONANCE IN PLASMAS

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ABSTRACT

A coupled set of equations, one describing the time evolution of the ordinary-mode wave energy and the other describing the time evolution of the electron distribution function is presented. The wave damping is mainly determined by $T_{\parallel}$ while the radiative equilibrium is mainly an equipartition with $T_{\perp}$. The time rate of change of $T_{\perp}$, $T_{\parallel}$, particle (N_0), and current ($J_{\parallel}$) densities are examined for finite $k_{\parallel}$ electron-cyclotron-resonance heating of plasmas.

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Recently, there has been renewed interest in electron cyclotron resonance heating (ECRH) of plasmas both in connection with the need for supplementary heating in magnetically confined devices and with ionospheric heating experiments. In the ordinary (0) mode ECRH, the electric field $\vec{E}$ of the heating wave is parallel to the confining magnetic field $\vec{B} = B\hat{z}$, and the heating is a consequence of the finite size of the electron Larmor orbits (ρ). Hence intuitive expectations would suggest that the O-mode ECRH would only lead to a heating in the parallel direction (i.e., to an increase in the parallel electron temperature $T_{\parallel}$ since $\vec{E} \parallel \vec{B}$) with a heating rate that is determined mainly by $T_{\perp}$ (since this heating is due to ρ). However, in this paper we will show that the predictions of the quasilinear theory are in sharp contrast to the intuitive expectations. In particular, we will show that the perpendicular heating rate ($\partial T_{\perp}/\partial t$) is very much larger than the parallel heating rate ($\partial T_{\parallel}/\partial t$) and the wave damping is determined mainly by $T_{\parallel}$, while the radiative equilibrium is an equipartition mainly with $T_{\perp}$. Finally, we will obtain closed form expressions for the time rate of change of $T_{\perp}$, $T_{\parallel}$, particle (N_0), and current ($J_{\parallel}$) densities for finite $k_{\parallel}$ O-mode ECRH.

We will begin by assuming Landau's quantized particle motion of the electrons in a uniform magnetic field^{1,2} and we shall use the notations and the results found in Ref. 1. For 0-mode the transitions probabilities of absorption j_A and emission j_E of Eq. (8) of Ref. 1 become

$$[j_A/N_{ks}] = [j_E/(N_{ks} + 1)] = \begin{Bmatrix} n+1 \\ n \end{Bmatrix} \left(\frac{2\pi^2 q^2 k_{\perp}^2}{DL^3 \mu \omega \omega_b} \right) (v_z + \hbar k_z/2\mu)^2$$

$$\delta_{v_z}, v_z + \hbar k_z/\mu \quad \delta[\omega - \omega_b - k_z(v_z + \hbar k_z/2\mu)] \quad , \quad (1)$$

where the term $\left\{ \begin{smallmatrix} n+1 \\ n \end{smallmatrix} \right\}$ means that $(n+1)$ belongs to transitions between the quantum states $|n, v_z\rangle$ and $|n+1, v'_z\rangle$, and n belongs to transitions between the states $|n, v_z\rangle$ and $|n-1, v'_z\rangle$. Here $D \approx (1 - \omega_p^2/\omega^2)$, $\omega_p^2 = (4\pi N_0 q^2/\mu)$, $\omega_b = (qB/\mu c)$, q and μ are the charge and mass of the electron, respectively, and the electrons are contained in a box of volume L^3 . In Eq. (1), the Kronecker δ indicates the conservation of z -momentum and the Dirac δ function indicates the conservation of energy. On making use of Eq. (1), and Eq. (29) of Ref. 1, the classical limit of Eq. (13) of Ref. 1 becomes

$$\left(\frac{\partial \epsilon_k}{\partial t}\right) = \int_{-\infty}^{\infty} dv_z \int_0^{\infty} dE_{\perp} \left(\frac{\pi k_{\perp}^2 \omega^2}{2D \omega_b^2} \right) (v_z^2 E_{\perp}) \left\{ \left(\frac{\omega_b}{\omega} \epsilon_k \right) [\delta(\omega - \omega_b - k_z v_z)] \right. \\ \left. Q F(E_{\perp}, v_z) \right\} + \left\{ \delta(\omega - \omega_b - k_z v_z) F(E_{\perp}, v_z) \right\} \left\{ \right\} , \quad (2)$$

where the wave energy $\epsilon_k = N_{ks} \hbar \omega$, $E_{\perp} = \mu v_{\perp}^2/2$, and the linear differential operator $Q = \{(\partial/\partial E_{\perp}) + (k_z/\mu \omega_b) (\partial/\partial v_z)\}$. Similarly, by applying the principle of detailed balance for the transition probabilities per unit volume of emission and absorption between the pairs of quantum states, $(|n, v_z\rangle; |n+1, v'_z\rangle)$ and $(|n, v_z\rangle; |n-1, v'_z\rangle)$, and then taking the classical limit, after a certain amount of algebra one can show³ that the time evolution of the electron distribution function $F(E_{\perp}, v_z)$ may be written

$$\partial F(E_{\perp}, v_z)/\partial t = \sum_k \left(2\pi^2 q^2 k_{\perp}^2 v_z^2 / DL^3 \mu \omega_b \right) \left(\omega_b \epsilon_k / \omega \right)$$

$$\left[(1 + E_1 Q) \{ \delta(\omega - \omega_b - k_z v_z) Q F(E_1, v_z) \} \right] + \left[(1 + E_1 Q) \{ \delta(\omega - \omega_b - k_z v_z) F(E_1, v_z) \} \right] \quad (3)$$

The set of Eqs. (2) and (3) is the coupled pair of classical quasilinear equations, one describing the time evolution of the ordinary mode wave energy and the other describing the time evolution of the electron distribution function. In Eq. (2), the first term which is proportional to ϵ_k yields the growth or damping rate of the wave and is a consequence of a balance between induced emission and absorption; while the second term which is independent of ϵ_k is a consequence of spontaneous emission. It is seen that Eq. (3) is a Fokker-Planck equation whose first and second terms account for "diffusion" and "dynamical friction," respectively. The velocity-space diffusion term is proportional to the wave energy ϵ_k and is again a consequence of a balance between stimulated emission and absorption. On the other hand, the dynamical friction term [i.e., the term which is independent of ϵ_k on the right-hand side of Eq. (3)] is a consequence of spontaneous emission.

If we now assume that $F(E_1, v_z)$ is a Maxwell-Boltzmann distribution function, and $\epsilon_k = 0$ at time $t = 0$, then the solution of Eq. (2) is $\epsilon_k = \epsilon_k^{(0)} [1 - \exp(-2\gamma_k t)]$ where $\epsilon_k^{(0)} = (\omega \kappa T_1 / \omega_b) / [1 + (T_1 / T_1)] (\omega - \omega_b) / \omega_b] \approx \kappa T_1$ for $\omega \approx \omega_b$, the damping rate

$$2\gamma_k \approx (\pi^{1/2} \omega \frac{\omega_p^2}{p R} \frac{2}{2c^2} \omega_b k_z v_z) [1 + (T_1 / T_1)] \{ (\omega - \omega_b) / \omega_b \}$$

$$\exp[-v_R^2/v_{\parallel}^2] , \quad (4)$$

$v_R = (\omega - \omega_b)/k_z$, $v_{\parallel}^2 = (2\kappa T_{\parallel}/\mu)$, and $D = (c^2 k_{\perp}^2/\omega^2)$ for $k_{\perp} \gg k_{\parallel}$. In the limit $k_z \rightarrow 0$, $2\gamma_k = (\kappa T_{\parallel}/\mu) (\pi \omega_p^2/2c^2) \delta(\omega - \omega_b)$. It is therefore interesting to find that although the wave damping $2\gamma_k$ is determined mainly by $T_{\parallel}$, the radiative equilibrium between the electrons and the O-mode radiation field is an equipartition mainly with $T_{\perp}$. At equilibrium the energy per mode $\epsilon_k^{(0)} = \kappa T_{\perp}$ as a result of the equipartition theorem.⁴ That is, the balance between induced emission and absorption is determined mainly by $T_{\parallel}$; while the global balance between spontaneous emission, induced emission, and absorption is determined mainly by $T_{\perp}$. This result is somewhat in contrast to intuitive expectations.

We now wish to calculate the time rate of change of $T_{\perp}$, $T_{\parallel}$, particle (N_O), and current ($J_{\parallel}$) densities for finite $k_{\parallel}$ ECRH. These are given by

$$\begin{aligned} & \left[\partial \langle E_{\perp} \rangle / \partial t = \partial(\kappa T_{\perp}) / \partial t, \quad \partial \langle E_{\parallel} \rangle / \partial t = \partial(\kappa T_{\parallel}/2) / \partial t, \quad \partial N_O / \partial t, \quad \partial J_{\parallel} / \partial t \right] \\ & = \int_{-\infty}^{\infty} dv_z \int_0^{\infty} dE_{\perp} [E_{\perp}, (\mu v_z^2/2), N_O, N_O q v_z] \{ \partial F(E_{\perp}, v_z) / \partial t \} , \end{aligned} \quad (5)$$

where each term in the square bracket on the left-hand side of Eq. (5) goes with the corresponding term in the square bracket under the integral on the

right-hand side in chronological order. Here the angular brackets refer to an average, and we have written $\langle E_{\perp} \rangle = \langle \mu v_{\perp}^2/2 \rangle = \langle \mu(v_x^2 + v_y^2)/2 \rangle = \kappa T_{\perp}$ and $\langle E_{\parallel} \rangle = \langle \mu v_z^2/2 \rangle = \kappa T_{\parallel}/2$ since the perpendicular motion has two degrees of freedom, while the parallel motion has only one; the average energy per degree of freedom is thus $\kappa T/2$ as a result of the equipartition theorem. On making use of Eq. (3), Eq. (5) yields

$$\left[\partial(\kappa T_{\perp})/\partial t; \partial(\kappa T_{\parallel}/2)/\partial t; \partial N_0/\partial t; \partial J_{\parallel}/\partial t \right] = \sum_k \left[\left(v_{\parallel}^2/\omega_b N_0 \right) \left\{ \left(\omega_b v_R/k_z v_{\parallel}^2 \right) + \left(2T_{\perp}/T_{\parallel} \right) \right\} ; \left(2v_{\parallel}^2/\omega_b N_0 \right) \left\{ v_R^2/v_{\parallel}^2 \right\} ; \right. \\ \left. 2/\mu \omega_b ; \left(3q v_{\parallel}/\mu \omega_b \right) \left\{ v_R/v_{\parallel} \right\} \right] \left(v_k/L^3 \right) , \quad (6)$$

where the terms in the square brackets on both sides of Eq. (6) have the similar one-to-one correspondence as in Eq. (5),

$$y_k = \left(\frac{\pi k_{\perp}^2 \omega_p^2 v_R}{2D \omega^2} \right) \left[\epsilon_k \left\{ F_{\parallel}(v_z) - \frac{\kappa T_{\perp} k_z}{\mu \omega_b} \frac{\partial F_{\parallel}(v_z)}{\partial v_z} \right\} - \frac{\omega \kappa T_{\perp}}{\omega_b} F_{\parallel}(v_z) \right]_{v_z = v_R} , \quad (7)$$

and $F_{\parallel}(v_z) = \int_0^{\infty} dE_{\perp} F(E_{\perp}, v_z)$. The somewhat lengthy and tedious algebra leading to Eqs. (6) and (7) only makes use of integration by parts over $dE_{\perp}$ and over dv_z ; and the constant terms resulting from these integrations vanish at $E_{\perp} = 0$ and $E_{\perp} = \infty$ and at $v_z = \pm \infty$. Similarly Eq. (2) can be rewritten as

$$\partial \epsilon_k / \partial t = - (v_{\parallel}^2 / \omega_b) \{ \omega v_R / k_z v_{\parallel}^2 \} Y_k. \quad (8)$$

If we now assume that $F_{\parallel}(v_z)$ is a Maxwell-Boltzmann distribution function so that $F_{\parallel}(v_z) = (1/\pi^{1/2} v_{\parallel}) \exp(-v_z^2/v_{\parallel}^2)$, then $\partial F_{\parallel}/\partial v_z = -(2v_z/v_{\parallel}) F_{\parallel}$ and Y_k of Eq. (7) become

$$Y_k = (\pi \omega_p^2 v_R / 2c^2) \{ \epsilon_k \{ 1 + (T_{\perp}/T_{\parallel}) (\omega - \omega_b)/\omega_b \} - (\omega k T_{\perp} / \omega_b) \} F_{\parallel}(v_z = v_R) \quad (9)$$

where we have written $D \approx (c^2 k_{\perp}^2 / \omega^2)$ for $k_{\perp} \gg k_{\parallel}$. In Eqs. (7) and (9) the term that is proportional to ϵ_k is a consequence of a balance between induced emission and absorption, while the term that is independent of ϵ_k is a consequence of spontaneous emission. It may be noted that $Y_k = 0$ of Eq. (9) yields the previously discussed radiative equilibrium solution $\epsilon_k^{(0)} = kT_{\perp}$ for $\omega \approx \omega_b$.

In the usual formulation of the classical quasilinear plasma kinetic theory (CQPKT)^{5,6} based on the Vlasov-Maxwell equations, it is extremely difficult to obtain the dynamical friction term of Eq. (3). That is, in CQPKT one assumes that the wave energy ϵ_k is much greater than $kT_{\perp}$; and consequently, one can neglect the dynamical friction term in Eq. (3) and the spontaneous emission terms of Eqs. (5) to (9). Then Y_k of Eq. (9) becomes

$$Y_k' = A_k (v_R/v_{\parallel} \pi^{1/2}) [1 + (T_{\perp}/T_{\parallel}) (\omega - \omega_b)/\omega_b] \exp\{-v_R^2/v_{\parallel}^2\}, \quad (10)$$

where $A_k = (\epsilon_k \pi \omega_p^2 / 2c^2)$. In Fig. 1 we have plotted the dimensionless functions $\{(\omega_b v_R / k_z v_{\parallel}^2) + (2T_{\perp}/T_{\parallel})\} (Y_k'/A_k) = Z_k(T_{\perp})$, $\{v_R^2/v_{\parallel}^2\} (Y_k'/A_k) = Z_k(T_{\parallel})$, $(Y_k'/A_k) = Z_k(N_0)$, $\{v_R/v_{\parallel}\} (Y_k'/A_k) = Z_k(J_{\parallel})$, and $\{\omega v_R / k_z v_{\parallel}^2\} (Y_k'/A_k) = Z_k(\epsilon_k)$ for two different angles of propagation. In this figure these functions $Z_k(T_{\perp})$, $Z_k(T_{\parallel})$, $Z_k(N_0)$, $Z_k(J_{\parallel})$, and $Z_k(\epsilon_k)$ illustrate the behavior of $\partial T_{\perp}/\partial t$, $\partial T_{\parallel}/\partial t$, $\partial N_0/\partial t$, $\partial J_{\parallel}/\partial t$, and $\partial \epsilon_k/\partial t$, respectively, as a function of ω/ω_b for a given value of $\vec{k}$ [see Eqs. (6) and (8)]. It is interesting to note from this figure that the Doppler effect due to the finite value of k_z does not yield the usually expected Gaussian broadening as found in the literature.⁴ However, it leads to a somewhat symmetric Doppler splitting for $\partial T_{\perp}/\partial t$, $\partial J_{\parallel}/\partial t$, and $\partial \epsilon_k/\partial t$ as a function of ω/ω_b , and to an antisymmetric splitting for $\partial T_{\parallel}/\partial t$ and $\partial N_0/\partial t$ as a function of ω/ω_b . This again is somewhat in contrast to intuitive expectations.

In Eq. (6) the sum over $\vec{k}$ goes over essentially³ as an integral over $d\omega$. Thus for radiative equilibrium the antisymmetric functions in Eq. (6) will integrate to zero. That is, for the antisymmetric terms of Eq. (6) whatever the particles gained from or lost to the waves for $\omega < \omega_b$ is again lost to or gained from the waves, respectively, for $\omega > \omega_b$. Consequently, from Eqs. (6) and (8) one can show that

$$\sum_{\vec{k}} (\partial \epsilon_k / \partial t) = - [\partial (L^3 N_0 \epsilon_{T_{\perp}}) / \partial t],$$

$$\sum_k (k_z/\omega) (\partial \epsilon_k / \partial t) = - [(L^3 \mu/q) (\partial J_{\parallel} / \partial t) - (L^3 \mu v_R) (\partial N_0 / \partial t)],$$

$$(\partial N_0 / \partial t) = 0,$$

$$\text{and } [\partial (L^3 N_0 k T_{\parallel}) / \partial t] = 0. \quad (11)$$

The set of relations in Eq. (11) represents the conservation laws of energy, momentum, particle number density, and the fact that an equilibrium distribution of waves does not spend any of its energy to increase the particle's parallel temperature, respectively. This clearly shows that the coupled set of Eqs. (2) and (3) is self-consistent, and the heating is primarily in the perpendicular direction. Indeed from Eq. (6) it is relatively easy to see⁷ that $(\partial T_{\perp} / \partial t) / (\partial T_{\parallel} / \partial t) \ll [\omega_p / (\omega - \omega_p)] \gg 1$. That is, the 0-mode ECRH would lead to a heating mainly in the perpendicular direction. This result is again somewhat contrary to physical intuition since for 0-mode ECRH the electric field $\vec{E} \parallel \vec{B}$.

Let us now examine the limiting case of Eq. (6) for exactly perpendicular propagation. In the limit $k_z \rightarrow 0$, $Z_k(T_{\perp}) = \{(\omega_p v_R / k_z v_{\parallel}^2) + (2T_{\perp} / T_{\parallel})\} (Y_k' / A_k) + (\omega_p / 2) \delta(\omega - \omega_p)$, $Z_k(T_{\parallel}) = \{v_R^2 / v_{\parallel}^2\} (Y_k' / A_k) + 0$, $Z_k(N_0) = (Y_k' / A_k) + 0$, $Z_k(J_{\parallel}) = \{v_R / v_{\parallel}\} (Y_k' / A_k) + 0$, and $Z_k(e_k) = \{\omega v_R / k_z v_{\parallel}^2\} (Y_k' / A_k) + (\omega / 2) \delta(\omega - \omega_p)$. Here we have used the fact that the Dirac δ -function satisfies the relations $x \delta(x) = 0$ and $x \delta'(x) = -\delta(x)$. This limiting case illustrates the reason for grouping the terms in the form shown in Eqs. (6) and (8). It may

again be noted that $Z_k(T_\perp) = -Z_k(\epsilon_k)$ is a reflection of the conservation laws of energy when $k_z \rightarrow 0$. In the conventional collisionless hot-plasma quasilinear theory the only mechanism responsible for broadening and/or splitting the resonances is the longitudinal Doppler effect. Since for $k_z \rightarrow 0$, the longitudinal Doppler shift vanishes, it is apparent why this theory predicts a δ -function resonance. A more exact theory should incorporate other mechanisms^B that may be responsible for resonance broadening and/or splitting.

In conclusion we have presented a coupled pair of quasilinear equations for the 0-mode electron cyclotron resonance in plasmas. One of these equations describes the time evolution of the 0-mode wave energy and the other describes the time evolution of the electron distribution function. We have shown that this coupled pair of equations is self-consistent by proving that these equations satisfy the conservation laws of energy, momentum, and particle density. It is found that although the wave damping is determined mainly by $T_\perp$, the radiative equilibrium is an equipartition mainly with $T_\perp$. In 0-mode ECRH the heating is primarily in the perpendicular direction even though $\vec{E} \parallel \vec{B}$. We have also presented closed form expressions for the time rate of change of $T_\perp$, $T_\parallel$, N_0 , $J_\parallel$, and ϵ_k . It is found that the Doppler effect due to the finite value of k_z does not yield the usually expected Gaussian broadening. However, it leads to a symmetric Doppler splitting for $\partial T_\perp / \partial t$, $\partial J_\parallel / \partial t$, and $\partial \epsilon_k / \partial t$, and to an antisymmetric splitting for $\partial T_\parallel / \partial t$ and $\partial N_0 / \partial t$ as a function of ω / ω_b . All these results are somewhat in contrast to intuitive expectations.

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FIGURE CAPTIONS

FIG. 1 Plots of the dimensionless functions $Z_k(T_{\perp})$, $Z_k(T_{\parallel})$, $Z_k(N_O)$, $Z_k(J_{\parallel})$, and $Z_k(\epsilon_k)$ as a function of ω/ω_p for two different angles of propagation: _____ is for $\theta = 2^\circ$, _ _ _ is for $\theta = 10^\circ$. Set (a) is $Z_k(T_{\perp})$, (b) is $Z_k(J_{\parallel})$, (c) is $Z_k(\epsilon_k)$, (d) is $Z_k(T_{\parallel})$, and (e) is $Z_k(N_O)$. Conditions are $T_{\perp} = 3$ keV, $T_{\parallel} = 2$ keV, and $(\omega_p/2\pi) = 9 \times 10^{10}$ cps.

Fig. 1

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