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Radial Transport Induced by Rotating RF Fields and Breakdown of Intrinsic Ambipolarity in a Magnetic Mirror

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Synopsis

Radial transport induced by rotating RF fields in a magnetic mirror is studied. The used RF frequency is in the ion cyclotron range of frequnecy and also comparable to the electron bounce frequency. The radial particle fluxes for ions and electrons are obtained by use of the gyrokinetic equation. With a positive azimuthal mode number of the RF field, the radial particle flux is outward for Maxwellian ions and is inward for Maxwellian electrons. This shows that the intrinsic ambipolarity is broken in this RF-induced transport. The present analyses can be applied to the thermal-barrier pumping of a tandem mirror.

Radial transport induced by RF fields is an important research subject in magnetically confined plasmas. It is a critical issue in understanding anomalous transport phenomena associated with instabilities or with RF plasma heating.[1-4] It is also useful in the control of electrostatic potential in a plasma. In tandem mirrors, barrier pumping of trapped ions in the thermal barrier with use of RF fields has been discussed.[5]

In this Letter we study radial transport induced by rotating RF fields in the ion cyclotron range of frequencies(ICRF) in a magnetic mirror, where the electron bounce frequency is usually comparable to the ion cyclotron frequency. We here obtain expressions of the radial particle fluxes induced by RF fields for ions and electrons. The orgin of transport is the cyclotron resonance for ions and the bounce resonance for electrons. We show that with a positive azimuthal mode number of the RF field, the radial particle flux is outward for Maxwellian ions and is inward for Maxwellian electrons, and therefore that the intrinsic ambipolarity is broken in the present RF-induced transport, which is different from that associated with drift instabilities in tokamaks.[6]

We now derive the expression of the RF-induced radial flux in a magnetic mirror. We consider an axisymmetric mirror. The magnetic field is expressed as $\underline{B} = B\underline{b} = \nabla \psi \times \nabla \theta$ in the magnetic flux coordinates (ψ, θ, s), where ψ is the flux coordinate, θ is an anglelike coordinate and s is the distance along a field line. We assume a parabolic profile given by $B(s) = B_0(1 + s^2/L_B{}^2)$ for the magnetic field. We also consider RF fields described by the eikonal representation for simplicity and then a perturbed quantity

$\tilde{X}(\psi, \theta, s, t)$ is expressed as $\tilde{X}(s)\exp[iS(\psi, \theta) - i\omega t]$. RF electric and magnetic fields, $\tilde{\underline{E}}$ and $\tilde{\underline{B}}$, are expressed in terms of the scalar and vector potentials $\tilde{\phi}$ and $\tilde{\underline{A}} = \tilde{A}_{\parallel}\underline{b} + \nabla \times \tilde{A}\,\underline{b}$ as

$$\tilde{\underline{E}} = -\nabla\tilde{\phi} + i(\omega/c)\tilde{\underline{A}} ,$$

$$\tilde{\underline{B}} = \nabla \times \tilde{\underline{A}} = k_{\perp}^2\tilde{A}\,\underline{b} + i\tilde{A}_{\parallel}\underline{k}_{\perp} \times \underline{b} = \tilde{B}_{\parallel}\underline{b} + \tilde{\underline{B}}_{\perp} , \tag{1}$$

where ω is the RF wave frequency, c the light speed and $\underline{k}_{\perp} = \nabla S = k_{\psi}\nabla\psi + k_{\theta}\nabla\theta$. $k_{\theta} \equiv m$ denotes the mode number of rotating RF fields. We assumed $|\underline{k}_{\perp}| \gg |\tilde{X}^{-1}\partial\tilde{X}/\partial s|$ in eq. (1).

The starting point is the linearized gyrokinetic equation applicable to arbitrary frequency.[7,8] In this theory, the perturbed distribution function $\tilde{f}$ for a given species is given in terms of $\tilde{\phi}$, $\tilde{A}_{\parallel}$ and $\tilde{B}_{\parallel}$ as

$$\tilde{f} = \frac{q}{M}\frac{\partial f_0}{\partial \varepsilon}\tilde{\phi} - 2\psi\frac{\partial}{\partial \psi}\left(\frac{\tilde{B}_{\parallel}}{k_{\perp}^2}\right)\frac{\partial f_0}{\partial \psi} + \sum_{n=-\infty}^{n=\infty}\tilde{g}_n\exp(-iL_n) , \tag{2}$$

where $\varepsilon = v^2/2 + q\Phi/M$ is the particle energy per unit mass, $L_n = \underline{v}_{\perp}\cdot(\underline{k}_{\perp}\times\underline{b})/\omega_c + n\zeta$, ζ the gyrophase angle between $\underline{v}_{\perp}$ and $\underline{k}_{\perp}$, $\omega_c (= qB/Mc)$ the cyclotron frequency, q the charge, M the mass, Φ the equilibrium electrostatic potential and f_0 the unperturbed distribution function, which is here assumed to be a function of ε only. The function $\tilde{g}_n$ is given by

$$\left[v_{\parallel}\frac{\partial}{\partial s} - i(\omega - \omega_d - n\omega_c)\right]\tilde{g}_n$$

$$= i\frac{q}{M}(\omega - \omega_*)\frac{\partial f_0}{\partial \varepsilon}\left[\left(\tilde{\phi} - \frac{v_{\parallel}}{c}\tilde{A}_{\parallel}\right)J_n - \frac{v_{\perp}}{k_{\perp}c}\tilde{B}_{\parallel}J_n'\right] , \tag{3}$$

where $J_n(z)$ is the Bessel function of order n, $z=k_\perp v_\perp/\omega_c$, $J_n'=dJ_n/dz$, $\omega_d=\omega_E+\omega_b+\omega_\kappa$ and ω_* are drift frequencies defined by, respectively,

$$\begin{aligned}
&\omega_E=(c/B)(\underline{k}_\perp\times\underline{b})\cdot\nabla\Phi,\\
&\omega_b=(v_\perp^2/2\omega_c)(\underline{k}_\perp\times\underline{b})\cdot\nabla\ln B,\\
&\omega_\kappa=(v_\parallel^2/\omega_c)(\underline{k}_\perp\times\underline{b})\cdot\underline{\kappa},\\
&\omega_*=-(\underline{k}_\perp\times\underline{b})\cdot\nabla f_0/(\omega_c\partial f_0/\partial\varepsilon),
\end{aligned} \tag{4}$$

and $\underline{\kappa}=(\underline{b}\cdot\nabla)\underline{b}$ is the curvature of a magnetic field line. In eq. (3), we here neglect ω_d and ω_* since we consider the case that the RF frequency ω is comparable to ω_{ci} and much larger than ω_d and ω_*. We retain only n=1 for ions and also n=0 for electrons in the summation over n in eq. (2).

The cross-field particle drift $\tilde{V}_\perp$ driven by RF fields is dominantly given by

$$\begin{aligned}
\tilde{V}_\perp &= c\frac{\tilde{E}\times b}{B}+v_\parallel\frac{\tilde{B}_\perp}{B}-\underline{v}_\perp\frac{\tilde{B}_\parallel}{B}\\
&=-i\frac{c}{B}[(\tilde{\phi}-\frac{v_\parallel}{c}\tilde{A}_\parallel)\underline{k}_\perp\times\underline{b}-i\frac{\underline{v}_\perp}{c}\tilde{B}_\parallel],
\end{aligned} \tag{5}$$

where eq. (1) is used and $|k_\theta\nabla\theta|\gg|k_\phi\nabla\psi|$ is assumed. By using eq. (5) and the perturbed distribution function $\tilde{f}$, we can obtain a cross-field particle flux induced by RF fields as

$$\underline{\Gamma}_\perp=(1/2)\mathrm{Re}\{\int d^3v<\underline{\tilde{V}}_\perp{}^*\tilde{f}>\}, \tag{6}$$

where $<\cdots>$ denotes time averaging over the RF oscillation period and $\mathrm{Re}\{\cdots\}$ represents the real part of a complex quantity.

The radial flux can be obtained by multiplying $\underline{\Gamma}_\perp$ by $\nabla\psi$ as

$$\underline{\Gamma}_\perp \cdot \nabla\psi = (1/2)\mathrm{Re}\{\sum_{v_\parallel \gtrless 0} \int \frac{B d\varepsilon\, d\mu}{|v_\parallel|} \int d\zeta < \underline{\widetilde{V}}_\perp{}^* \cdot \nabla\psi\, \widetilde{f} > \} . \tag{7}$$

Substituting eq. (5) into eq. (7) and performing the integral in the gyrophase angle ζ, we can obtain

$$\underline{\Gamma}_\perp \cdot \nabla\psi = \pi\, c m \mathrm{Re}\{ i \sum_{v_\parallel \gtrless 0} \int \frac{B dE d\mu}{|v_\parallel|}\, \widetilde{H}_n{}^* \widetilde{g}_n \} , \tag{8}$$

$$\widetilde{H}_n(s) = (\widetilde{\phi} - \frac{v_\parallel}{c}\widetilde{A}_\parallel) J_n(z) - \frac{v_\perp}{k_\perp c}\widetilde{B}_\parallel J_n{}'(z) , \tag{9}$$

where $n = 1$ for ions and $n = 0$ for electrons and we employed

$$\int d\zeta / 2\pi \exp(-iL_n) = J_n(z) ,$$
$$\int d\zeta / 2\pi\, \underline{v}_\perp \exp(-iL_n) = i(v_\perp / k_\perp) J_n{}'(z) \underline{k}_\perp \times \underline{b} , \tag{10}$$

and $(k_\perp \times b) \cdot \nabla\psi = k_\theta B = mB$. We note that the first and second terms in eq. (1) do not contribute the transport since they are reactive parts of the ditribution function. As the magnetic field B(s) is varying along the field line, we define a line-integrated radial flux by

$$\Gamma_\psi = \int \frac{ds}{B} \underline{\Gamma}_\perp \cdot \nabla\psi = \pi\, c m \mathrm{Re}\{ i \int d\varepsilon\, d\mu \sum_{v_\parallel \gtrless 0} \int \frac{ds}{|v_\parallel|} \widetilde{H}_n{}^* \widetilde{g}_n \} . \tag{11}$$

If we define a line-integrated density as $N = \int ds n_0 / B$ where n_0 is the density, N and Γ_ψ satisfy $\partial N / \partial t + \partial \Gamma_\psi / \partial \psi = 0$.[3)]

We first consider the ion radial flux induced by RF fields. The ion bounce frequency is less than the ion cyclotron frequency and then the function $\tilde{g}_{i1}$ is given by the local approximation. Setting $\partial \tilde{X}/\partial s = ik_{\parallel}\tilde{X}$ in eq. (3), we obtain

$$\tilde{g}_{i1} = -\frac{e}{M_i}\omega\frac{\partial f_{0i}}{\partial \varepsilon}\frac{\tilde{H}_1}{\omega - n\omega_{ci} - k_{\parallel}v_{\parallel}} . \quad (12)$$

Substituting eq. (12) into eq. (11) with n= 1, we can obtain

$$\Gamma_{\phi i} = -\pi^2 m\frac{ce\omega}{M_i}\int d\varepsilon\, d\mu \frac{\partial f_{0i}}{\partial \varepsilon}$$

$$\times \sum_{v_{\parallel} \gtrless 0}\int \frac{ds}{|v_{\parallel}|}|\tilde{H}_1(s)|^2 \delta\left[\omega - \omega_{ci}(s) - k_{\parallel}v_{\parallel}(s)\right] , \quad (13)$$

where $v_{\parallel}(s) = \pm\left[2(\varepsilon - e\Phi(s)/M_i - \mu B(s))\right]^{1/2}$. We here assume that the equilibrium electrostatic potential $\Phi(s)$ is constant as $\Phi(s) = \Phi_0$. Then the solution of $\omega - \omega_{ci}(s) - k_{\parallel}v_{\parallel}(s) = 0$ in eq. (13) is determined from the equation, for $v_{\parallel} \gtrless 0$,

$$\bar{s}^2 - \{\omega - \omega_{c0} \pm k_{\parallel}[2(\varepsilon - e\Phi_0/M_i - \mu B_0) - 2\mu B_0\bar{s}^2]^{1/2}\}/\omega_{c0} = 0, \quad (14)$$

where $\bar{s} = s/L_B$ and $\omega_{c0} = eB_0/M_i c$. For small $k_{\parallel}$, we can approximately drop the s-dependence in $v_{\parallel}$ and then obtain, for $v_{\parallel} \gtrless 0$,

$$s_{\pm} = \pm L_B[(\omega - \omega_{c0} \pm k_{\parallel}[2(\varepsilon - e\Phi_0/M_i - \mu B_0)]^{1/2})/\omega_{c0}]^{1/2} . \quad (15)$$

In this case if we perform the integral in s, we can obtain, for $\omega - \omega_{c0} - k_{\parallel}\bar{v}_{\parallel} > 0$,

$$\Gamma_{\phi i} = -\frac{\pi}{2^{3/2}}m\frac{ce\omega L_B}{M_i}\int d\varepsilon\, d\mu\frac{\partial f_{0i}}{\partial \varepsilon} \quad (16)$$

$$\times \sum_{v_{\parallel} \gtrless 0}\sum_{s=s_{\pm}}\frac{|\tilde{H}_1(s)|^2}{\{[\varepsilon - \mu B_0(\omega - k_{\parallel}\bar{v}_{\parallel})/\omega_{c0}][\omega_{c0}(\omega - \omega_{c0} - k_{\parallel}\bar{v}_{\parallel})]\}^{1/2}} ,$$

where $\overline{v}_{\parallel} = v_{\parallel}(s=0)$. The ion radial flux $\Gamma_{\psi i}$ induced by RF fields is proportional to the mode number m of the rotating RF fields and the direction of $\Gamma_{\psi i}$ is outward with positive m for Maxwellian ions, $\partial f_{0i}/\partial \varepsilon < 0$.

Next we discuss the electron radial flux. The electron bounce frequency is assumed to be comparable to the ion cyclotron frequency and then the function $\tilde{g}_{e0}$ can be expressed by an expansion in harmonics of the bounce motion as

$$\tilde{g}_{e0} = \sum_{\lambda=-\infty}^{\infty} \tilde{g}_{e0}(\lambda) \exp(i\lambda \omega_B \tau) , \tag{17}$$

$$\tilde{g}_{e0}(\lambda) = \frac{e}{M_e} \omega \frac{\partial f_{0e}}{\partial \varepsilon} \frac{\tilde{H}_0(\lambda)}{\omega - \lambda \omega_{Be}} , \tag{18}$$

where $\tau = \int ds/|v_{\parallel}|$ and ω_{Be} is the electron bounce frequency, given by $\omega_{Be} \equiv 2\pi/\tau_{Be} = (2\mu B_0)^{1/2}/L_B$. If we use the approximation of $J_0 \simeq 1$ and $J_0' \simeq k_\perp v_\perp / 2|\omega_{ce}|$, $\tilde{H}_0(\lambda)$ is given by

$$\tilde{H}_0(\lambda) = \int d\tau \left[\tilde{\phi} - \frac{v_{\parallel}}{c}\tilde{A}_{\parallel} - \frac{M_e \mu}{e}\tilde{B}_{\parallel}\right] \exp(-i\lambda \omega_{Be} \tau) / \int d\tau . \tag{19}$$

Substituting eqs.(17) and (18) to eq.(11) with n=0, we obtain

$$\Gamma_{\psi e} = \pi^2 m \frac{ce\omega}{M_e} \int d\varepsilon\, d\mu \frac{\partial f_{0e}}{\partial \varepsilon}$$

$$\times \sum_{\lambda=-\infty}^{\infty} \int d\tau\, \tilde{H}_0^* \tilde{H}_0(\lambda) \exp(i\lambda \omega_{Be} \tau)\, \delta[\omega - \lambda \omega_{Be}] . \tag{20}$$

And if we expand $\tilde{H}_0$ as $\tilde{H}_0 = \sum_{\lambda=-\infty}^{\infty} \tilde{H}_0(\lambda) \exp(i\lambda \omega_{Be} \tau)$ and perform the integral in τ, we obtain

$$\Gamma_{\phi e} = \pi^2 m \frac{ce\omega}{M_e} \int d\varepsilon \, d\mu \frac{\partial f_{0e}}{\partial \varepsilon} \tau_{Be} \sum_{\lambda=1}^{\infty} |\tilde{H}_0(\lambda)|^2 \delta[\omega - \lambda \omega_{Be}] . \quad (21)$$

We here introduce a new variable $\tilde{\chi}$ defined by

$$\tilde{A}_{\parallel} = -i(c/\omega) \partial \tilde{\chi} / \partial s , \quad (22)$$

in place of $A_{\parallel}$. If we express the electron bounce motion as $s(\tau) = a(E, \mu)\sin(\omega_{Be}\tau)$ with $a = L_B[(\varepsilon + e\Phi_0/M_e - \mu B_0)/\mu B_0]^{1/2}$ and use $\tilde{X}(s) = \bar{X}\exp(ik_{\parallel}s)$, substitution of eq. (22) into eq. (19) then yields

$$\tilde{H}_0(\lambda) = J_\lambda(k_{\parallel}a)[\tilde{\phi} - \frac{\lambda \omega_{Be}}{\omega}\tilde{\chi} - \frac{M_e \mu}{e}\tilde{B}_{\parallel}] . \quad (23)$$

The electron radial flux is also proportional to the mode number m and the direction of $\Gamma_{\phi e}$ is inward with positive m for Maxwellian electrons, $\partial f_{0e}/\partial \varepsilon < 0$, contrary to that for the ion.

We have shown here that in mirrors with RF fields of $\omega \sim \omega_{ci} \sim \omega_{Be}$ and the positive mode number m, the ion radial flux due to the cyclotron resonansce is directed outward for Maxwellian ions, on the other hand, the electron radial flux due to the bounce resonance becomes inward for Maxwellian electrons. This means the breakdown of the intrinsic ambipolarity in the present RF-induced transport. The intrisic ambipolarity is satisfied for anomalous transport associated with drift wave instabilities in tokamaks.[5) This intrisic ambipolarity is satisfied only when there exsits the θ-symmetry for the unperturbed state, the wave fields are produced by microscopic instabilities and there is no external momentum input in the θ- direction. It breaks down even if one of them is not satisfied. In the present case, the rotating RF fields yield the external momentum input into a plasma in the θ-direction via wave-particle resonances. Therefore, the intrinsic ambipolarity in the present RF-induced transport breaks down and then the associated selective radial flux results.

In conclusion, we have studied the RF-induced radial transport in mirrors and have obtained the expressions of the radial paricle fluxes for ions and electrons. We believe that the present analyses are useful in studying transport phenomena associated with ICRF heating in mirrors. The present analyses can be applied to the barrier pumping of trapped ions in the thermal barrier of a tandem mirror. With use of an appropriate RF mode, the RF fields pump out the trapped ions from the thermal barrier. Since the electron radial flux is then directed inward, therefore the RF-induced transport tends to enhance the well depth of the thermal barrier potential.

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