ISSN 0915-633X

NATIONAL INSTITUTE FOR FUSION SCIENCE

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(Received – Apr. 20, 1993)

NIFS-225 May 1993

RESEARCH REPORT
NIFS Series

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NAGOYA, JAPAN

Bounce Resonance Heating and Transport in a Magnetic Mirror

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KEYWORDS:

bounce resonance heating, RF-induced transport, mirror plasma,
gyrokinetic equation, power absorption, radial particle flux

ABSTRACT

Bounce resonance heating and associated radial transport in a magnetic mirror are studied. A general expression for the absorbed power and radial particle flux due to bounce resonance is obtained by use of a gyrokinetic equation with the eikonal representation for waves. It is shown that the power absorption yields the outward (inward) particle flux with positive (negative) azimuthal mode number of RF wave fields and the power emission meaning instabilities yields the reverse particle flux with the same mode number for the ion. The particle flux for the electron is opposite to that for the ion.

Electron heating by wave fields in the ion cyclotron range of frequencies (ICRF) has been observed in mirror experiments.[1-3] In these experiments, the heating mechanism of the electron is considered not to be electron drag from the ion, but to be direct heating. Bounce resonance heating is the dominant candidate, because the electron bounce frequency is usually comparable to the ion cyclotron frequency in mirrors and therefore ICRF waves can resonate to the electron bounce motion. We also see that radial transports generally arise associated with RF heating.[4-7]

In this Letter, we study bounce resonance heating and associated radial transport in a magnetic mirror. We derive a general expression of the absorbed power due to the bounce resonance by using the gyrokinetic equation with the eikonal approximation for wave fields applicable to both the electron and ion. By comparing the expression of the absorbed power with the expression of radial particle flux induced by RF wave fields via the bounce resonance obtained previously in ref.6, we find a simple but useful exact relation between the absorbed power and radial particle flux. The relation shows that the power absorption for the ion yields the outward (inward) particle flux with positive (negative) azimuthal mode number of RF fields and the power emission meaning instabilities yields the reverse particle flux with the same mode number. The particle flux for the electron is opposite to that for the ion.

We now derive the expression of the absorbed power due to the bounce resonance in an axisymmetric magnetic mirror. The magnetic field is expressed as $\boldsymbol{B} = B\boldsymbol{b} = \nabla\psi \times \nabla\theta$ 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 consider RF fields described by the eikonal representation for simplicity and then a perturbed quantity $\tilde{X}(\psi,\theta,s,t)$ is expressed as $\tilde{X}(s)\exp[\mathrm{i}S(\psi,\theta) - \mathrm{i}\,\omega t]$, where the wave frequency ω is assumed to be much smaller than the cyclotron frequency. The RF electric and magnetic fields, $\tilde{\boldsymbol{E}}$ and $\tilde{\boldsymbol{B}}$, are expressed in terms of the scalar and vector potentials $\tilde{\phi}$ and $\tilde{\boldsymbol{A}} = \tilde{A}_{\parallel}\boldsymbol{b} + \nabla \times \tilde{\mathcal{A}}\boldsymbol{b}$ as

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

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

where c is the light speed and $\boldsymbol{k}_{\perp} = \nabla S = S_{\psi}\nabla\psi + S_{\theta}\nabla\theta$, $S_{\theta} = m$ being the azimuthal mode number of RF fields. We assumed $|k_{\perp}| >> |\tilde{X}^{-1}\partial\tilde{X} / \partial s|$ in eq.(1).

The starting point is the linearized gyrokinetic equation, since the wave frequency is much smaller than the cyclotron frequency. The perturbed distribution function $\tilde{f}$ for a given species is then given by[8,9]

$$\tilde{f} = \frac{q}{M}\frac{\partial f_0}{\partial \varepsilon}\tilde{\phi} + \frac{q}{MB}\frac{\partial f_0}{\partial \mu}\{(\tilde{\phi} - \frac{v_{\parallel}}{c}\tilde{A}_{\parallel})[1 - J_0(z)\exp(-\mathrm{i}L)] - \frac{v_{\perp}}{k_{\perp}c}\tilde{B}_{\parallel}J_1(z)\exp(-\mathrm{i}L)\} + \tilde{g}\exp(-\mathrm{i}L)\ , \tag{2}$$

where $\varepsilon(=v^2/2 + q\Phi/M)$ is the particle energy per unit mass, $\mu(=v_{\perp}^2/2B)$ the magnetic moment per unit mass, $L = \boldsymbol{v}_{\perp}\cdot(\boldsymbol{k}_{\perp}\times\boldsymbol{b})/\omega_{\mathrm{c}}$, $z = k_{\perp}v_{\perp}/\omega_{\mathrm{c}}$, $\omega_{\mathrm{c}}(=qB/Mc)$ the cyclotron frequency, q the charge, M the mass, $\boldsymbol{v}$ the velocity, Φ the equilibrium electrostatic potential, $J_n(z)$ the Bessel function of order n and f_0 the unperturbed distribution function. The function $\tilde{g}$ is given by

$$
\begin{aligned}
&[v_{\parallel}\frac{\partial}{\partial s}-\mathrm{i}(\omega-\omega_{\mathrm{d}})]\tilde{g} \\
&=i(\omega-\omega_{*})\frac{q}{M}\frac{\partial f_0}{\partial \varepsilon}[(\tilde{\phi}-\frac{v_{\parallel}}{c}\tilde{A}_{\parallel})J_0(z)+\frac{v_{\perp}}{k_{\perp}c}\tilde{B}_{\parallel}J_1(z)] \ ,
\end{aligned} \tag{3}
$$

where $\omega_{\mathrm{d}}=\omega_{\mathrm{E}}+\omega_{\mathrm{b}}+\omega_{\kappa}$ and ω_{*} are drift frequencies defined by, respectively,

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

and $\boldsymbol{\kappa}=(\boldsymbol{b}\cdot\nabla)\boldsymbol{b}$ is the curvature of a magnetic field line. Hereafter we neglect ω_{d} and ω_{*} in eq.(3) by assuming that the wave frequency ω is much larger than ω_{d} and ω_{*}.

If we introduce a timelike variable defined by $\tau=\int \mathrm{d}s/|v_{\parallel}|$, the function $\tilde{g}$ can be expressed by an expansion in the harmonics of the bounce motion as

$$
\tilde{g}=\sum_{\lambda=-\infty}^{\infty}\tilde{g}(\lambda)\exp(\mathrm{i}\lambda\omega_{\mathrm{B}}\tau) \ , \tag{5}
$$

$$
\tilde{g}(\lambda)=-\omega\frac{q}{M}\frac{\partial f_0}{\partial\varepsilon}\frac{\tilde{H}(\lambda)}{\omega-\lambda\omega_{\mathrm{B}}} \ , \tag{6}
$$

where ω_{B} is the bounce frequency and $\tilde{H}(\lambda)$ is given by

$$
\tilde{H}(\lambda)=\int\mathrm{d}\tau[\tilde{\phi}-\frac{v_{\parallel}}{c}\tilde{A}_{\parallel}+\frac{M\mu}{q}\tilde{B}_{\parallel}]\exp(-\mathrm{i}\lambda\omega_{\mathrm{B}}\tau)/\int\mathrm{d}\tau \ , \tag{7}
$$

with use of the approximation of $J_0(z) \cong 1$ and $J_1(z) \cong z/2$.

We now calculate the absorbed power due to the bounce resonance. Since the perturbed current density $\boldsymbol{j}$ is expressed as $\boldsymbol{j} = q\int \boldsymbol{v}\tilde{f}d\boldsymbol{v}$, a line-integrated absorbed power can be given by

$$P = \frac{1}{2}\mathrm{Re}\{q\int \frac{\mathrm{d}s}{B} \sum_{\pm|v_{\parallel}|} \int \frac{B\,\mathrm{d}\varepsilon\,\mathrm{d}\mu\,\mathrm{d}\zeta}{|v_{\parallel}|} < \tilde{\boldsymbol{E}}^* \cdot \boldsymbol{v}\tilde{f} >\} \ , \tag{8}$$

where ζ is the gyrophase angle between $\boldsymbol{v}_{\perp}$ and $\boldsymbol{k}_{\perp}$, $<\cdots>$ denotes time averaging over the wave oscillation period and $\mathrm{Re}\{\cdots\}$ represents the real part of a complex quantity. Here $\tilde{\boldsymbol{E}} \cdot \boldsymbol{v}$ is expressed as

$$\tilde{\boldsymbol{E}} \cdot \boldsymbol{v} = E_{\parallel}v_{\parallel} + \frac{(\boldsymbol{k}_{\perp} \cdot \boldsymbol{v}_{\perp})(\boldsymbol{k}_{\perp} \cdot \tilde{\boldsymbol{E}}_{\perp})}{k_{\perp}^2} + \frac{\boldsymbol{b} \cdot (\boldsymbol{k}_{\perp} \times \boldsymbol{v}_{\perp})\boldsymbol{b} \cdot (\boldsymbol{k}_{\perp} \times \tilde{\boldsymbol{E}}_{\perp})}{k_{\perp}^2} \ . \tag{9}$$

If we introduce a new variable $\tilde{\chi}$ defined by

$$\tilde{A}_{\parallel} = -\mathrm{i}\frac{c}{\omega}\frac{\partial \tilde{\chi}}{\partial s} \ , \tag{10}$$

and use eq.(1) and $\tilde{B}_{\parallel} = (c/\omega)\boldsymbol{b} \cdot (\boldsymbol{k}_{\perp} \times \tilde{\boldsymbol{E}}_{\perp})$ obtainable from Maxwell's equation, eq.(9) is reduced to

$$\tilde{\boldsymbol{E}} \cdot \boldsymbol{v} = -\frac{\partial}{\partial s}(\tilde{\phi} - \tilde{\chi}) + \frac{(\boldsymbol{k}_{\perp} \cdot \boldsymbol{v}_{\perp})(\boldsymbol{k}_{\perp} \cdot \tilde{\boldsymbol{E}}_{\perp})}{k_{\perp}^2} - \frac{\omega \boldsymbol{v}_{\perp} \cdot (\boldsymbol{k}_{\perp} \times \boldsymbol{b})\tilde{B}_{\parallel}}{ck_{\perp}^2} \ . \tag{11}$$

Substituting eqs.(2) and (11) into (8) and performing the integral in the gyrophase angle ζ, we can obtain

$$P = \mathrm{Re}\{q\int \mathrm{d}\varepsilon\,\mathrm{d}\mu\int \mathrm{d}\,\tau[-J_0\frac{\partial}{\partial s}(\tilde{\phi}-\tilde{\chi})^* + \mathrm{i}\frac{\omega v_\perp}{ck_\perp}\tilde{B}_\parallel^{\,*}J_1]\tilde{g}\}$$
$$\cong \pi\mathrm{Re}\{q\int \mathrm{d}\varepsilon\,\mathrm{d}\mu\int \mathrm{d}\,\tau[-\frac{\partial}{\partial s}(\tilde{\phi}-\tilde{\chi})^* + \mathrm{i}\,\omega\frac{M\mu}{q}\tilde{B}_\parallel^{\,*}]\tilde{g}\}\ , \tag{12}$$

where we employed

$$\int \mathrm{d}\,\zeta\exp(-\mathrm{i}L)/2\pi = J_0(z)\ ,$$
$$\int \mathrm{d}\,\zeta \boldsymbol{v}_\perp \exp(-\mathrm{i}L)/2\pi = -\mathrm{i}(v_\perp/k_\perp)J_1(z)\boldsymbol{k}_\perp\times\boldsymbol{b}\ , \tag{13}$$

and used that the first and second terms of $\tilde{f}$ given by eq.(2) do not contribute the Joule heating of $\tilde{\boldsymbol{E}}^*\cdot\boldsymbol{j}$ since they are reactive parts of the distribution function.

When we perform the integral in τ after the expansion of the perturbed quantities in the bounce harmonics, we obtain

$$P = \pi\mathrm{Re}\{\mathrm{i}\,q\int \mathrm{d}\varepsilon\,\mathrm{d}\mu\tau_\mathrm{B}\sum_{\lambda=-\infty}^{\infty}[\lambda\omega_\mathrm{B}(\tilde{\phi}-\tilde{\chi})+\omega\frac{M\mu}{q}\tilde{B}_\parallel]^*\tilde{g}(\lambda)\}$$
$$= \pi\frac{2\pi^3q^2}{M}\omega\int \mathrm{d}\varepsilon\,\mathrm{d}\mu(-\frac{\partial f_0}{\partial\varepsilon})\sum_{\lambda=1}^{\infty}\lambda\left|\tilde{H}(\lambda)\right|^2\delta[\omega-\lambda\omega_\mathrm{B}]\ , \tag{14}$$

where $\tau_\mathrm{B} = 2\pi/\omega_\mathrm{B}$, $\delta[x]$ the δ-function and $\tilde{H}(\lambda)$ is rewritten as

$$\tilde{H}(\lambda) = \int \mathrm{d}\,\tau[\tilde{\phi}-\tilde{\chi}+\frac{M\mu}{q}\tilde{B}_\parallel]\exp(-\mathrm{i}\lambda\omega_\mathrm{B}\tau)/\int \mathrm{d}\,\tau\ , \tag{7'}$$

by use of the resonance condition $\omega = \lambda\omega_\mathrm{B}$. As seen from eqs.(9) and (11), $\tilde{\phi}-\tilde{\chi}$ in $\tilde{H}(\lambda)$ is related to the perturbed parallel electric field $\tilde{E}_\parallel$ by

$\tilde{E}_{\parallel} = -\mathrm{i}k_{\parallel}(\tilde{\phi} - \tilde{\chi})$. Equation (14) is a general expression of the absorbed power due to the bounce resonance. We can obtain a more explicit expression for the absorbed power if we give an explicit model for the axial profiles of the magnetic field and electrostatic potential to find the bounce frequency and also the particle trajectory, which will be reported in a separate paper.

Hereafter, we discuss a relation between the absorbed power P due to the bounce resonance and the radial particle flux induced by RF fields which we have obtained in ref.6. The line-integrated radial particle flux in ref.6 is defined by

$$\Gamma_{\psi} = \int \frac{\mathrm{d}s}{B} \Gamma_{\perp} \cdot \nabla \psi = \frac{1}{2} \mathrm{Re}\{\int \frac{\mathrm{d}s}{B} \sum_{\pm|v_{\parallel}|} \int \frac{B\,\mathrm{d}\varepsilon\,\mathrm{d}\mu\,\mathrm{d}\zeta}{|v_{\parallel}|} < \tilde{V}_{\perp}^{*} \cdot \nabla \psi \tilde{f} >\} , \qquad (15)$$

where $\tilde{V}_{\perp}$ is the cross-field particle drift driven by RF fields. Though the radial particle flux of eq.(15) is obtained for the electron in ref.6, we can easily obtain that for the ion in the same manner. Then the expression of the radial particle flux induced by RF fields for a given species is given by

$$\Gamma_{\psi} = m \frac{2\pi^{3} q c}{M} \int \mathrm{d}\varepsilon\,\mathrm{d}\mu (-\frac{\partial f_{0}}{\partial \varepsilon}) \sum_{\lambda=1}^{\infty} \lambda |\tilde{H}(\lambda)|^{2} \delta[\omega - \lambda \omega_{\mathrm{B}}] . \qquad (16)$$

Equation (16) just coincides with that obtained in ref.6 for the electron. By comparing eq.(14) with eq.(16), we can obtain a simple but useful exact relation between P and Γ_{ψ} as

$$\Gamma_{\psi} = \frac{mc}{q\omega} P . \qquad (17)$$

We briefly discuss about the physics of the equality (17). The power absorption (or, emission) due to wave-particle resonances causes the change in a wave momentum. From the momentum conservation of a plasma, this change in the wave momentum is transferred to plasma particles to change the particle momentum. For RF wave fields with azimuthal wave number $k_\theta = m / r$, the force yielding the momentum input in a spatially uniform plasma is a azimuthal force and given by[10,11]

$$F_\theta = \frac{k_\theta}{\omega} P \ .$$

We have found that the above relation holds for the bounce resonance absorption in a non-uniform mirror plasma. Then the $\boldsymbol{F} \times \boldsymbol{B}$ drift is a radial drift and yields the radial particle flux given by eq.(17). That is,

$$\Gamma_\psi = rB\Gamma_r = rB\frac{cF_\theta}{qB} = \frac{mc}{q\omega} P \ ,$$

since the flux Ψ is given by $\psi = r^2 B / 2$ for an axisymmetric plasma.

In conclusion, we have studied bounce resonance heating and the associated radial transport in mirror plasmas. We have derived a general expression of the absorbed power and have found that it is closely related to the radial particle flux induced by RF wave fields. We see from eq.(17) that the power absorption (P>0) for the ion yields the outward (inward) particle flux with a positive (negative) azimuthal mode number of RF wave fields and the direction of the particle flux is reverse for the electron. On the other hand, the direction of the particle flux for the power emission (P<0) meaning instabilities is opposite to that for the power absorption with the same mode number. We believe that the present analyses are also useful in studying the heating process and the associated radial transport of trapped particles in toroidal plasmas.

Finally, we are grateful to members of the Plasma Research Center, University of Tsukuba, for useful discussions. This work was partly supported by a Grant-in-Aid for Scientific Research from the Ministry of Education, Science and Culture.

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