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Effect of Anomalous Plasma Transport on Radial Electric Field in Torsatron/Heliotron

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Abstract

Anomalous cross field plasma fluxes induced by the electric field fluctuations has been evaluated in a rotating plasma with shear flow in a helical system. The plasma rotation frequency due to the radial electric field makes the Doppler frequency shift which does not explicity affect the cross field flux. The anomalous ion flux is evaluated by the ion curvature drift resonance continuum in the test particle model. The curvature drift rersonance induces a new force term <B>'/<B> which did not make large influence in the ion flux. The shear flow term in the anomalous flux combined with the electric field in neoclassical flux reduces to a first order differential equation which governs the radial profile of the electric field. A general exact analytical solution for the differential equation is derived and a simple approximate solution for the radial electric field is also given. Numerical results indicate that the shear flow effect is important for the anomalous cross field flux and for determination of the radial electric field particularly in the peripheral region.

Keywords: Radial electric field, poloidal plasma rotation, curvature drift resonance, gyrokinetic solution, test particle model, anomalous cross field plasma flux, electric field fluctuations, magnetic field line curvature effect, shear flow effect.

§ 1. Introduction

In toroidal high temperature plasma confinement devices such as tokamak and Starallator/Heliotron, strong radial electric field has been experimentally observed(1)(2)(3). This electric field is believed to play important roles in various plasma stability, classical and anomalous plasma transports and the L-H transition in tokamaks(4)(5). This radial electric field may be produced by the unbalance of electron and ion transports, and also by the momentum injection associated with auxiliary heatings.

When the radial electric field exists, the plasma suffers rotation due to the ExB drift motion mainly in the poloidal direction with the velocity $v_E=cE_r/B$. In this case, the Larmor gyromotion may be modified, i.e., the perpendicular velocity is shifted by v_E: $w_\perp=(v_x^2+u_y^2)^{1/2}$ with $u_y=v_y-v_E$. The perpendicular constant of particle motion becomes $\mu=w_\perp^2/2B$. This may suggest that the plasma rotation plays an important role through the finite Larmor radius effect in a complicated manner.

The total energy related to the perpendicular motion $\alpha=(v_x^2+u_y^2)/2+e\Phi/m$ and the parallel energy $\beta=v_\parallel^2/2$ are also constants of motion, respectively, where $E_r=-\partial\Phi/\partial x$. The equilibrium distribution function f_0 may be written in the form

$$f_0(\alpha,\beta,X)=N(X)\left(\frac{m}{2\pi T(X)}\right)^{\frac{3}{2}}\exp\left(-\frac{m}{2T(X)}\left(w_\perp^2+v_\parallel^2\right)\right) \tag{1}$$

where the scalar potential Φ has been transformed to the shifted Maxwellian through the invariance condition of particle guiding center $X=x+u_y/\Omega$ and the normalization condition $\int f_0 d\,v=N(x)$ (6) neglecting the trapped particle effect.

In the rotating plasma, the oscillation frequency ω of perturbations may suffer the Doppler frequency shift by the rotation frequency $\omega_E=k_y v_E$. We will show that this Doppler shift has no direct influence on anomalous plasma transport as in the kinetic plasma stability theory. Although the uniform electric field E_r may have no influence on the anomalous transport, whether the shear flow effect due to the non-uniformity of E_r plays an important role or not remains to be investigated. Our purpose of the present paper is, therefore, to examine whether the shear flow effect plays an important role for the anomalous cross field ion flux.

When the radial electric field has the radial dependence, the radial derivative of the equilibrium distribution function (1) has the velocity shear term induced from v_E in $w_\perp$:

$$\frac{\partial f_0}{\partial X}=\left\{\frac{N'}{N}+\frac{T'}{T}(\bar{E}-\frac{3}{2})+\frac{m}{T}\hat{u}_y\hat{v}_E{}'\right\}f_0 \tag{2}$$

where $\hat{v}=v/v_{th}$ and $\bar{E}=(\hat{w}_\perp{}^2+\hat{v}_E{}^2)/2$ is the normalized kinetic energy.

To investigate the effect of plasma transports on the radial electric field E_r, we should consider both neoclassical and anomalous transport processes. When we introduce the shear flow effect, it also should have some influence on the neoclassical transport process. In the neoclassical theory, however, the zeroth order electric field effect does not vanishes, and play a pricipal role, i.e., the flow shear effect due to the nonuniformity of E_r may be higher order, and may be negligible as compared with the zeroth order of E_r. We will take into account this shear flow effect only for anomalous ion flux and neglect in the neoclassical ion flux.

§ 2. Anomalous Cross Field Flux

We consider the anomalous cross field plasma flux Γ_{AN} induced by the electric scalar potential fluctuations $\tilde{\phi}$, defined by

$$\Gamma_{AN}=\int\left\langle\tilde{v}_x\ \tilde{f}\right\rangle d^3v \tag{3}$$

where the angular brackets means the ensemble average for purturbations, $\hat{v}_x$ is the radial velocity perturbation and $\tilde{f}$ is the perturbed distribution function. We apply the gyrokinetic solution given for ions

$$\tilde{f}_i=-\frac{e}{T}\left\{1+J_0^2(\alpha_L)\frac{\omega+\omega_E-\omega*(1+\eta(\bar{E}-\frac{3}{2})-\hat{u}_y\hat{v}_E{}'L_n)}{\omega+\omega_E-\omega_D}\right\}f_0\ \phi(x) \tag{4}$$

where all notations are standard: $\omega_E=k_\theta v_E$, $v_E=cE_r/B$, $\omega^*=(cT/eB)/L_n$, $\eta=d\ln T/d\ln N$, $L_n{}^{-1}=d\ln N/dr$. The precessional drift frequency defined by $\omega_D=\nabla S.\mathbf{b}x(w_\perp{}^2/2\ \nabla\ln B+\kappa(v^2+v_E{}^2/\varepsilon))/\Omega$ is approximated by $\omega_D=2\varepsilon_n\omega^*(\hat{v}_\parallel{}^2+\hat{v}_E{}^2/\varepsilon+\hat{w}_\perp{}^2/2)$ in the low-β plasmas in which $\kappa=\mathbf{b}.\nabla\mathbf{b}=\nabla\ln B$, where $\varepsilon=r/R$, and $\varepsilon_n=R/L_n$. The term $v_E{}^2/\varepsilon$ in ω_D is induced by the centrifugal force due to the plasma rotation in the poloidal direction. In the gyrokinetic solution (4), the E_r-effect is in the Doppler shift $\omega+\omega_E$. It is also involved in the perpendicular velocity $w_\perp$ in E, ω_D, the Bessel function argument $\alpha_L = k_\theta w_\perp / \Omega$ and f_0. The shear flow effect is in the fourth term in the numerator.

The frequency ordering: $\omega > \omega^* \cong \omega_D > \omega_{ti} = k_\parallel v_i$, has been assumed. The ion transit frequency ω_{ti}, therefore, has been dropped in eq.(4). The drift frequency (ω_g) due to the centrifugal force is also neglected in usual theory. It may becomes important when the rotation speed v_E is comparable to ion thermal velocity v_i, or the Mach number is close to unity.

For electrons, the electron transit frequency $\omega_{te} = k_\parallel v_e$ may be much higher than ω: $\omega_{te} >> \omega = \omega^*$, the second nonadiabatic term in the curly braces in eq.(4) may be neglected, and $\tilde{f}$ may be approximated by the first adiabatic term:

$$\tilde{f}_e = -\frac{e\tilde{\phi}}{T_e} f_0 \qquad (5)$$

Assuming that $\tilde{v}_x$ is given by the perturbed ExB drift velocity: $\tilde{v}_x = -ik_\theta \tilde{\phi}/B$, and introducing eq.(4) into eq.(3), we have the cross field ion flux

$$\Gamma_{AN} = N \sum_{k,\omega} \frac{cT}{eB} k_\theta \left| \frac{e\tilde{\phi}}{T} \right|^2 g(k,\omega) \qquad (6)$$

where g ia a function of the Doppler shifted frequency $\bar{\omega} = \omega + \omega_E$:

$$g(k,\omega) = \mathrm{Im} \int J_0^2(\alpha_L) \frac{\bar{\omega} - \omega^* (1 + \eta(\bar{E} - \frac{3}{2}) - \hat{u}_y \hat{v}_E{}' L_n)}{\bar{\omega} - \omega_D} f_0 d^3v \qquad (7)$$

For electrons, on the other hand, applying eq.(3), we have no real flux: $\Gamma_{AN}{}^e = 0$.
It is worthy to note that the integral in eq.(7) is the same one in the electrostatic dispersion relation derived from the neutrality condition(7):

$$\Lambda(k,\omega) = 1 + \frac{1}{\tau} - \int J_0^2(\alpha_L) \frac{\bar{\omega} - \omega^* (1 + \eta(\bar{E} - \frac{3}{2}) - \hat{u}_y \hat{v}_E{}' L_n)}{\bar{\omega} - \omega_D} f_0 d^3v = 0 \qquad (8)$$

The discrete time eigenvalue ω_0 for the η_i-mode is determined by solving eq.(8) for ω. The growth rate γ of an instability is given by $\gamma = \mathrm{Im}\omega_0$ which gives the real cross fieid flux Γ_{AN} due to the instability(discrete eigenvalue). Since the dispersion relation (8) is derived from the neutrality condition, the resultant cross field ion flux should be equal to the electron flux: $\Gamma_{AN}{}^{i} = \Gamma_{AN}{}^{e}$. If we apply the dispersion relation (8), it is clear that the imaginary part of the velocity integral is zero, $g(k,\omega)=0$, i.e., we have no real flux $\Gamma_{AN}{}^{i}=0$ for the η_i-mode. This is due to the fact that $\Gamma_{AN}{}^{e}=0$. Although the ion flux induced by the η_i-mode is zero, the heat flux

$$Q_{AN} = \int \left\langle \tilde{v}_x \ \tilde{f} E \right\rangle d^3v \tag{9}$$

does not , in general, vanish even for the discrete mode(7).

Instead of the flux induced by the growth rate of instability, here we evaluate the flux (6) induced by the drift resonance continuum applying the test particle model without the neutrality condition and assuming that the power spectrum of perturbations, $|e\tilde{\phi}/cT|^2$, has been given. The source of fluctuations $\tilde{\phi}$ may be certain instabilities and/or subcritical turbulence or nonlinear chaos.

The integral (7) can also be expressed in terms of force terms as in the neoclassial theory

$$g(k,\omega) = c_1\omega + c_2\frac{N'}{N} + c_3\frac{T'}{T} + c_2\frac{eE_r}{T} + c_4\frac{eE_r{}'}{T} \tag{10}$$

where the coefficients c_j have been defined by

$$c_1 = \mathrm{Im}\int J_0^2(\alpha_L)\,\frac{f_0}{\bar{\omega}-\omega_D}\,d^3v \tag{11}$$

$$c_2 = \frac{cT}{eB}k_\theta\, c_1 \tag{12}$$

$$c_3 = \mathrm{Im}\int J_0^2(\alpha_L)\,\frac{(\bar{E}-3/2)\,f_0}{\bar{\omega}-\omega_D}\,d^3v\,\frac{cT}{eB}k_\theta \tag{13}$$

$$c_4 = \mathrm{Im}\int J_0^2(\alpha_L)\,\frac{u_y f_0}{\bar{\omega}-\omega_D}\,d^3v\,\frac{k_\theta}{\Omega} \tag{14}$$

The fourth term in eq.(10) comes from the rotation frequency ω_E in eq.(4). If the first term in eq.(10) is combined with the fourth term, they give the Doppler shifted frequency: $c_1(\omega+\omega_E)$. The fourth term with electric field Φ' is always transfored to the

Doppler shift, which is proved in a more general formalism in Appendix A. Since the frequency power spectrum $|\tilde{\phi}|^2$ in the rotating system suffers the same Doppler shift when observed in the Laboratory system, the cross field flux has no electric field effect when integrated over frequency ω in eq.(4).

Since $\alpha_L = k_\theta w_\perp/\Omega$ and ω_D depend on both $w_\perp$ and $v_\parallel$, and the shear flow term involves u_y the velocity integrals in eq.(7) and eqs.(11) – (14) become triple integral with respect to v_x, u_y and $v_\parallel$. In this case, eq.(7) can be written in the form

$$g(k,\omega) = \mathrm{Im}\,\pi^{-3/2}\int_{-\infty}^{\infty} dx e^{-x^2}\int_{-\infty}^{\infty} dy e^{-y^2} \times\int_{-\infty}^{\infty} dz e^{-z^2} J_0^2(\alpha_L)\frac{\bar{\omega}-\omega*(1+\eta(x^2+y^2+z^2-3/2)+L_n y\hat{v}_E{}')}{\bar{\omega}-\hat{\omega}_D((x^2+y^2)/2+z^2)} \tag{15}$$

The triple integrals may only be possible by numerical calculations making use of high speed computer. The major complexity in eq.(15) is due to the shear flow term.

The simplest approximation for the shear flow term may be to assume the ion velocity in the poloidal direction is thermal velocity: $u_y=v_y-v_E=v_i$, or y=1. In this case, the shear flow effect may approximately be evaluated as $(v_E/v_i)L_n/L_E$, where L_E is the scale length of electric field variation. Since $v_E/v_i \ll 1$, the shear flow effect may be negligible unless $L_n/L_E \gg 1$, i.e., the variation of the electric field E_r should be much larger than that of ion density N. This may not happen in usual situations. In the Torsatron/Heliotron, however, the density shows hollow profile(2)(3) and $L_n/L_E \gg 1$ may occur near the peak point where $L_n \to \infty$. If this shear flow term is assumed to be a constant, or u_y is approximated by $w_\perp$, and the curvature drift frequency is approximated by $\omega_D = \hat{\omega}_D E$, the triple integral can be reduced to a double integral.

$$g(k,\omega) = \frac{2}{\sqrt{\pi}}\mathrm{Im}\int_0^1 \frac{d\lambda}{\sqrt{1-\lambda}}\int_0^{\infty} J_0^2(\alpha_L)\frac{\bar{\omega}-\omega*(1+\eta(E-3/2)+L_n\hat{v}_E{}'\sqrt{\lambda E})}{\bar{\omega}-\hat{\omega}_D E}e^{-E}\sqrt{E}dE \tag{16}$$

where the transformation of variables: $E=w_\perp^2+v_\parallel^2$, $\lambda=2\mu B/E=w_\perp^2/(w_\perp^2+v_\parallel^2)$ and $w_\perp = (\lambda E)^{1/2}$ have been employed.

If we neglect the finite Larmor radius effect, $\alpha_L=0$, by carrying out the intagration with respect to λ, eq.(16) is reduced to a single integral form:

$$g(k,\omega) = \frac{4}{\sqrt{\pi}} \mathrm{Im} \int_0^\infty \frac{\bar{\omega} - \omega_* (1 + \eta (E - 3/2) + L_n \hat{v}_E' \sqrt{E\pi/4})}{\bar{\omega} - \hat{\omega}_D E} e^{-E} \sqrt{E} dE \qquad (17)$$

By the same manner, all velocity integrals in eqs.(11)~(14) can be reduced to the single integral with respect to E as follows

$$c_1 = \frac{4}{\sqrt{\pi}} \mathrm{Im} \int_0^\infty \frac{\sqrt{E} e^{-E} dE}{\bar{\omega} - \hat{\omega}_D E} \qquad (18)$$

$$c_3 = -\frac{4}{\sqrt{\pi}} \mathrm{Im} \frac{cT}{eB} k_\theta \int_0^\infty \frac{E - 3/2}{\bar{\omega} - \hat{\omega}_D E} \sqrt{E} e^{-E} dE \qquad (19)$$

$$c_4 = \sqrt{\pi} \frac{v_E'}{v_i} \mathrm{Im} \frac{cT}{eB} k_\theta \int_0^\infty \frac{E e^{-E}}{\bar{\omega} - \hat{\omega}_D E} dE \qquad (20)$$

This single integral may be convenient for numerical integration.

By the transformation of variable : $E = x^2$, and making use of the relation

$$\frac{1}{x^2 - \varsigma^2} = \frac{1}{2\varsigma} \left(\frac{1}{x - \varsigma} - \frac{1}{x + \varsigma} \right)$$

the above integrals with respect to E can be transformed to the usual plasma dispersion function as follows

$$\int_0^\infty \frac{\sqrt{E} e^{-E} dE}{\bar{\omega} - \hat{\omega}_D E} = -\frac{\sqrt{\pi}}{\hat{\omega}_D} Z_1(\varsigma)$$

where $\zeta = (\bar{\omega}/\hat{\omega}_D)^{1/2}$, and the plasma dispersion function $Z_j(\zeta)$ has been defined by

$$Z_j(\varsigma) = \frac{1}{\sqrt{\pi}} \int_{-\infty}^{\infty} \frac{x^j}{x - \varsigma} e^{-x^2} dx \qquad (21)$$

$Z_0(\zeta)$ is the usual plasma dispersion function. The higher moments are expressed by the lower moments: $Z_1(\zeta) = 1 + \zeta Z_0(\zeta)$, $Z_2(\zeta) = \zeta Z_1(\zeta)$, and so forth(8).

The coefficients c_j are expressed in term of the plasma dispersion function:

$$c_1 = -\frac{4}{\hat{\omega}_D}\mathrm{Im}Z_1(\varsigma) \tag{22}$$

$$c_3 = -\frac{4}{\hat{\omega}_D}\mathrm{Im}\left(Z_3(\varsigma) - \frac{3}{2}Z_1(\varsigma)\right)\frac{cT}{eB}k_\theta \tag{23}$$

The integral in c_4 , however, can not be expressed by the Z-function. The argument $\zeta=(\bar{\omega}/\hat{\omega}_D)^{1/2}$ indicates that the coefficients have the branch points at the origin $\bar{\omega}=0$ and at infinity $\bar{\omega}=-\infty$, and the branch cut is the whole negative real axis in the complex $\bar{\omega}$-plane [9]. This branch cut is the Vlasov continuum due to the curvature drift resonance.

Since the frequency integration is real in eq.(6), the coefficients c_j have no real compornent except the resonance condition: $\bar{\omega}-\hat{\omega}_D E=0$, i.e., without the wave-particle resonance all coefficints c_j vanish, and we have no anomalous particle flux, $\Gamma_{AN}=0$. At the drift resonance condition, $\bar{\omega}-\hat{\omega}_D E=0$ $(0<E<\infty)$, applying the formula

$$\frac{1}{\omega-\omega_D} = P\frac{1}{\omega-\omega_D} - i\pi\delta(\omega-\omega_D) \tag{25}$$

where P means the principal value, we have the real component of the flux from the second term in eq.(25), and eq.(17) can be written in the form

$$g(k,\omega) = -\frac{4}{\sqrt{\pi}\hat{\omega}_D}\varsigma\left\{\bar{\omega}-\omega^*\left(1+\eta(\varsigma^2-\frac{3}{2})+\frac{\pi}{4}L_n\hat{v}_E{}'\varsigma\right)\right\}e^{-\varsigma^2} \tag{26}$$

which can also be derived by using the formula $\mathrm{Im}Z_0(\zeta)=\pi^{1/2}\exp(-\zeta^2)$ for the coefficients in eqs.(22)~(24).

If we assume the separation of variables in the power spectrum

$$\left|\frac{e\tilde{\phi}}{T}\right|^2_{k,\omega} = \left|\frac{e\tilde{\phi}}{T}\right|^2_{k} S(\omega) \tag{27}$$

introducing eqs.(26) and (27) into eq.(6), and transforming the sum over the frequency into an integral with respect to $\bar{\omega}/\hat{\omega}_D$, we have

$$\Gamma_{AN}=\sum_k \frac{cT}{eB}k_\theta\left|\frac{e\phi}{T}\right|_k^2 N\left\{(\hat{\omega}_D-\omega*\eta)I_1-\omega*\left(1-\frac{3}{2}\eta\right)I_0-\frac{\pi}{4}L_n\hat{V}_E\omega*I_{\frac{1}{2}}\right\} \tag{28}$$

where the moment integral I_j has been defined by

$$I_j=\frac{4}{\sqrt{\pi}}\int_0^\infty \varsigma^{(j+1)}e^{-\varsigma^2}S\,d\varsigma \tag{29}$$

The flux (28) can be rewritten in the neoclassical expression

$$\Gamma_{AN}=\sum_k\left|\frac{cT}{eB}k_\theta\right|^2\left|\frac{e\phi}{T}\right|_k^2 N\left\{-I_0\frac{N'}{N}-\left(\frac{3}{2}I_0-I_1\right)\frac{T'}{T}-2I_1\frac{\langle B\rangle'}{\langle B\rangle}-\frac{\pi}{4}\frac{v_E}{v_i}I_{\frac{1}{2}}\frac{E_r'}{E_r}\right\} \tag{30}$$

Introducing the E_r effect, we have seen that the frequency suffers the Doppler frequency shift, i.e., ω is replaced by $\bar{\omega}=\omega+\omega_E$. When observed in the Laboratory system, the frequency spectrum $S(\omega)$ should also be a function of the Doppler shift : $\bar{\omega}=\omega+\omega_E$. The moment integral with respect to ω is transformed to the integral about ω, and therefore, we have no ω_E-effect.

If we assume a simple cut-off frequency spectrum with width $2\omega_s$ of the form $S(\omega)=1/2\omega_s$ for $|\omega|<\omega_s$, and the width ω_s is wide enough $\omega_s>>|\omega_D|$, the moment integral (29) can be approximated by the white noise spectrum: $I_0=1/\omega_s$, $I_{1/2}=2/\pi^{1/2}\omega_s$ and $I_1=3/2\omega_s$(10). In this case, the temperature gradient term in eq.(30) cancels out. Bearing in mind the relations, $\omega^*=(cT/eB)N'/N$, $\hat{\omega}_D=2\varepsilon_n\omega^*=2(cT/eB)B'/B$, eq.(30) can also be rewritten in the neoclassical from:

$$\Gamma_{AN}=\sum_k\left|\frac{e\phi}{T}\right|_k^2\left|\frac{cT}{eB}k_\theta\right|^2\frac{N}{\omega_s}\left\{-\frac{N'}{N}-3\frac{\langle B\rangle'}{\langle B\rangle}-\frac{\sqrt{\pi}}{2}\frac{v_E}{v_i}\frac{E_r'}{E_r}\right\} \tag{31}$$

The second force term B'/B in the curly braces in eq.(31), which comes from the curvature drift frequency ω_D, is a new term. The third term due to the shear flow is also a new term.

We assume the saturation level of the power spectrum

$$\left|\frac{e\phi}{T}\right|_k^2=\frac{1}{k_\perp^2L^2}$$

and the perpendicular wave number is of the range $k_\perp\rho_i=\lambda^{-1}$ with λ is close to unity. In this case, the flux may be expressed in the simple form[10]

$$\Gamma_{AN}=\lambda\frac{v\rho_i^2}{L}N\left\{-\frac{N'}{N}-3\frac{\langle B\rangle'}{\langle B\rangle}-\frac{\pi}{2}\frac{v_E}{v_i}I\frac{E_r'}{E_r}\right\} \tag{32}$$

This flux involves no radial electric field E_r. Notice that the disappearance of T'/T term in eq.(31) is due to the cancellation of the moment integrals for the particular spectral function. This cancellation may not occur in general, and we have T'/T term in the flux.

So far, we have neglected the centrifugal force effect due to plasma rotation. When the rotation speed is comparable to ion thermal velocity or the Mach number is close to unity, the drift frequency ω_g may become important for the flux Γ. If the centrifugal force effect is included, the moment integral defined in eq.(10) should be modified by mutiplying a nonlinear factor $\exp(-\omega_g/\hat{\omega}_D)$.

§3. Radial Electric Field

In this section we derive the radial electric field E_r from the ambipolar condition taking into account neoclassical and anomalous transport coefficients neglecting the second order E_r-effect in ω_g. The neoclassical cross field flux has been expressed in the form

$$\Gamma_{NC}=-a_{NC}\left(\frac{N'}{N}-\frac{eE_r}{T}\right)-b_{NC}\frac{T'}{T} \tag{33}$$

where coefficients a_{NC} and b_{NC} consist of two parts, the axisymmetric part associated with transit or trapped particles, and the asymmetric part associated with locally trapped particles arising through violation of the system's axial symmetry :

$$a_{NC}=\frac{1.09a_i}{1+0.87\alpha_i}+\frac{0.67b_i^{(1)}}{1+1.02\beta_i^{(1)}} \tag{34}$$

$$b_{NC}=\frac{3}{2}\frac{1.46a_i}{1+0.39\alpha_i}-\frac{1.09a_i}{1+0.87\alpha_i}+\frac{1.9b_i^{(1)}}{1+0.59\beta_i^{(1)}}-\frac{3}{2}\frac{0.67b_i^{(1)}}{1+1.02\beta_i^{(1)}} \tag{35}$$

where coefficients a_i b_i α_i and β_i for ions have been given in Ref.(11). These coefficients for the electron flux have also been given in Ref.(11). Combing eqs.(31) and (33), the total ion flux $\Gamma=\Gamma_{AN}+\Gamma_{NC}$ may be written in the form

$$\Gamma = -(a_{NC}+a_{AN})\frac{N'}{N} + a_{NC}\frac{eE_r}{T} - b_{NC}\frac{T'}{T} - c_{AN}\frac{\langle B\rangle'}{\langle B\rangle} - d_{AN}\frac{eE_r'}{T} \tag{36}$$

where the neoclassical coefficients a_{NC} and b_{NC} have been given in eqs.(34) and (35), and the anomalous coefficients may be approximated by $a_{AN}=\lambda v_i \rho_i^2/L$, $c_{AN}=3a_{AN}$ and $d_{AN}=\pi^{1/2}\rho_i a_{AN}/2$. In order to adjust to the neoclassical electric field term and make a differential equation with respect to E_r, the shear flow term in eq.(31) has been changed by the relation $v_E E_r'/E_r = (\rho_i/2)eE_r'/T$.

Although the anomalous electron flux has been neglected in §2, if we take into account trapped electron effect, strong trapped electron modes induced by temperature gradient and curvature drift effects may be excited. These trapped electron modes may induce the anomalous cross field flux(12). Even these TEM are absent, the trapped electron resonance may make an anomalous electron flux as long as electric field fluctuations exist as in the case of ion resonace treated in §2. We assume the anomalous electron flux in the same form as eq.(31)

$$\Gamma^e_{AN} = -a^e_{AN}\frac{N'}{N} - b^e_{AN}\frac{T_e'}{T_e} - c^e_{AN}\frac{\langle B\rangle'}{\langle B\rangle}$$

We assume the anomalous electron flux $\Gamma_{AN}{}^e$ is much larger than the neoclassical electron flux $\Gamma_{NC}{}^e$ and neglect $\Gamma_{NC}{}^e$. In this case,. the ambipolar condition or neutrality condition,

$$\Gamma^i = \Gamma^e \tag{37}$$

determines the radial electric field E_r. Since the coefficient $\beta^{(1)}$ in eqs. (34) and (35) depends on $E_r=-\Phi'$, the neoclassical flux, in general, depends nonlinearly on E_r. In this case, E_r may be determined iteratively solving eq.(37) for each r(13). When the coefficient $\beta^{(1)}$ is negligible, which may be valid for low collisional regimes, the ambipolar condition (37) depends linearly on E_r and E_r'. In this case, eq.(37) reduces to a first order differential equation for E_r:

$$E_r - \frac{d_{AN}}{a_{NC}}E_r' = E_{r0}(r) \tag{38}$$

where E_{r0} is the electric field in the case without shear flow, $d_{AN}=0$, and given by

$$E_{r0}(r) = \frac{T}{e}\left\{\left(1+\frac{a_{AN}-a^e_{AN}}{a_{NC}}\right)\frac{N'}{N}+\frac{b_{NC}}{a_{NC}}\frac{T'}{T}-\frac{a^e_{AN}}{a_{NC}}\frac{T_e'}{T_e}+\frac{c^e_{AN}-c_{AN}}{a_{NC}}\frac{\langle B\rangle'}{\langle B\rangle}\right\} \quad (39)$$

Taking into account the shear flow effect and solving the differential equation (39), we have the solution in the form

$$E_r(r) = \exp\left(\int_0^r \frac{a_{NC}}{d_{AN}}dr''\right)\left\{-\int_0^r dr'\frac{a_{NC}}{d_{AN}}E_{r0}(r')\exp\left(-\int_0^{r'}\frac{a_{NC}}{d_{AN}}dr''\right)+E_{r0}(0)\right\} \quad (40)$$

If d_{AN} is small as in usual situations, by repeating the partial integration for eq.(40), we have an approximate expansion

$$E_r = E_{r0}+\frac{d_{AN}}{a_{NC}}E_{r0}'-\left(\frac{d_{AN}}{a_{NC}}\right)^2 E_{r0}'' \ \ldots \quad (41)$$

Equation (41) indicates that, to the first order of $d_{AN}/a_{AN}(>0)$, the decreasing electric field ($E_r'<0$) makes E_r more smaller.

§4. Numerical Result

First we examine the gradB-effect on the radial electric field E_{ro} given by eq.(39) without the shear flow effect. If we consider the frequency range $\omega_{it}<<\omega<<\omega_{te}$ and neglect trapped electrons as mentioned in §2, the anmalous electron flux Γ_{AN}^e may be neglected as compared with ion flux Γ_{AN}^i. We neglect the trapped particle effect and assume $\Gamma_{AN}^e=0$ for the sake of simplicity. To compare theoretical results with experimental observations, the anomalous as well as neoclassical electron fluxes should be taken into account. We leave this problem in future study, and concentrate here on investigation of the field line curvature and shear flow effects on the radial electric field.

The inverse scale length of the helically averaged helical magnetic field, $L_B^{-1} =\langle B\rangle'/\langle B\rangle$, is approximately given by eq.(B.9) in Appendix B. For CHS with R=100cm, m=8, l=2 and a=20cm, the radial variation of L_B^{-1} calculated by eq.(B.9) is presented in Fig.1. On the average, this quantity is essentially the same order of the curvature of field lines, R^{-1}. With this curvature effect, the radial variation of E_{ro} given by eq.(39) without electron transport effect is shown in Fig.2, where the coefficient ratio has been approximated by $(a_{NC}+a_{AN})/a_{NC}=a_o+a_1 x^2$ with a_o and a_1 being constants. This model simulate the experimental observations[2] for the ratio of ion thermal diffusivity χ_{exp}/χ_{NC} which increases up to 10-100 as radius x increases from the center to the periphery.

In calculation of eq.(39), temperature and density profiles have been given by[13]

$$T(x)=(T_o-T_b)(1-x^{\alpha_t})^{\beta_t}+T_b \tag{42}$$

$$n(x) = (n_0-n_b)\frac{1-x^{\alpha_n}}{g}\left\{1-(1-g)\left(1-x^{\alpha_n}\right)^{\beta_n-1}\right\}+n_b \tag{43}$$

The profile parameters α_t, α_n, β_t, β_n, n_o, n_b, T_o, T_b and g are chosen to simulate experimentally observed profiles. These parameters are tabulated in Ref.(13) for two typical low and high density discharges in CHS.

We also calculated E_{ro} without <B>' term in eq.(39) for the same high density case. The result was almost the same as shown in Fig.2, i.e., the curvature effect may be negligible in evaluation of the radial electric field E_{ro}.

We now examine the shear flow effect on the radial electric field. Since $d_{AN}/(a\, a_{AN}) = \pi^{1/2}\rho_i/2a \ll 1$, the expansion in eq.(41) up to the first order term may be good enough. The radial derivative Ero' is numerically calculated making use of eqs.(39), (42), and (43) for the radial difference Δx=0.001. Result is presented in Fig.3 for the high density case. Although analytical limit of Ero' tends to zero as x->0 for the profiles given by eqs.(42) and (43), numerical result indicates a singularity at x=0 even for Δx=0.001. This singular behavior may be induced by the numerical error due to the difference of two very close quantities. Therefore, numerical result at the center x=0 may not be reliable. Although $d_{AN}/(a\, a_{AN})$ is small, the derivative dE_{ro}/dx is large as seen in Fig.3, the shear flow effect may be important particularly in the peripheral region.

Making use of the numerically calculated E_{ro}', the electric field is calculated by eq.(41). Result is presented in Fig. 2, and comparison is also made with the zeroth order electric field E_{ro}. As seen Fig.2, the shear flow effect makes important contribution to the determination of electric field profile. The difference between E_{ro} and E_r is so large in the peripheral region, the higher order term in eq.(41) or more precise solution (38) may be necessary for more complete theory.

§5. Summary

Anomalous cross field ion flux induced by the curvature drift resonance has been evaluated by the test particle model in a rotating CHS plasma , assuming electric field back ground fluctuations exist. The radial electric field E_r only makes the Doppler frequency shift and makes no influence on the anomalous flux as in the case of transit resonance model.

Due to the curvature drift resonance, $\omega=\omega_D$, the anomalous flux involves a new term contributed from B'/B.This curvature effect was, however, not significant in a high density CHS discharge case.

When the shear flow effect is taken into account, the velocity integral in the cross field flux becomes complicated triple integral form. Evaluation of this integral may only be possible by numerical calculationbs. By employing a simple model, neglecting the finite Larmor radius effect, the triple velocity integral has been reduced to a single integral with respect to energy, which made it possible to analytically evaluate the flux.

Combing neoclassical and anomalous fluxes, from the ambipolar condition, a first order differential equation for the radial electric field E_r has been derived. The exact analytical solution for the differential equation has been derived in general form. By partial integrations of the solution, an approximate solution has been derived for a weak shear flow case, which was applied for numerical evaluations. In a high density CHS discharge, numerical results indicate that the shear flow effect is impotant particularly in the periphery region as in stability theories.

In our calculation, we have assumed a simple white noise like frequency spectrum. In this particular case, the anomalous coefficient for the term T'/T in the flux disappeared. In general form of the spectra, this cancellation may not occur, and we have some anomalous transport contribution to T'/T term which may modify the electric field profile.

Acknowledgement

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Appendix A Radial Electric Field in Shaing's Flux

Assuming the existence of electric field fluctuations in a toroidal system, Shaing has derived the cross filed particle flux induced by the transit resonance, $\omega=k_{\parallel}v_{\parallel}$, in the form(eq.(46) in Ref.(14))

$$\Gamma_{AN}=\frac{3c}{e\chi\psi'}\sum_{mn\omega}\frac{\sqrt{\pi}}{6}B\langle P\rangle\left|\frac{e\Phi_{mn\omega}}{\langle T\rangle}\right|^2\frac{m^2}{v_i|m-nq|}\exp\left(-\left(\frac{\omega_{mn}^E}{\omega_t}\right)^2\right)$$
$$\times\left\{-\left\langle\frac{\omega}{m}\right\rangle-U_p+\frac{U_t}{q}+\frac{2}{5}\left(-\frac{3}{2}+\left(\frac{\omega_{mn}^E}{\omega_t}\right)^2\right)\left(-\frac{q_p}{\langle P\rangle}+\frac{1}{q}\frac{q_t}{\langle P\rangle}\right)\right\} \qquad \text{(A.1)}$$

where Φ_{mn} is the Fourier coefficient of scalar potential fluctuations with m and n being the poloidal and toroidal mode numbers, respectively, $\omega_{mn}{}^E=\omega+\omega_E$, q_p and q_t are heat fluxes, and all other notaions are standard as defined in Ref.(14). Let us confirm that the term with the radial electric field Φ' in the mass flow velocities U_p and U_t in eq.(A.1) can be tranformed to the Doppler shift in the first frequncy term.

From the definitions of U_q and U_t as given by eq.(18) in Ref.(14) near the rational surface, q=m/n, we have

$$-U_p+\frac{U_t}{q}=-\frac{e\langle T\rangle}{e\varphi'}\frac{\langle P'\rangle}{\langle P\rangle}-\frac{\omega_E}{m}$$

That is, the electric field involved in U_q and U_t can be expressed by the rotation frequency ω_E. Since the frequency average has been defined by

$$\left\langle\frac{\omega}{m}\right\rangle\equiv\sum_{mn\omega}\frac{m\omega}{|m-nq|}\Phi^2_{mn\omega}\exp\left(-\left(\frac{\omega_{mn}^E}{\omega_t}\right)^2\right)\left\{\sum_{mn\omega}\frac{m^2}{|m-nq|}\Phi^2_{mn\omega}\exp\left(-\left(\frac{\omega_{mn}^E}{\omega_t}\right)^2\right)\right\}^{-1}$$

we have

$$\left\langle\frac{\omega}{m}\right\rangle+\frac{\omega_E}{m}=\left\langle\frac{\omega_{mn}^E}{m}\right\rangle$$

By introducing this relation into eq.(A.1), the flux can be written in term of the Doppler shifted frequency in the form

$$\Gamma_{AN}=\frac{3c}{e\chi\psi'}\sum_{mn\omega}\frac{\sqrt{\pi}}{6}B\langle P\rangle\left|\frac{e\Phi_{mn\omega}}{\langle T\rangle}\right|^2\frac{m^2}{v_i|m-nq|}\exp\left(-\left(\frac{\omega_{mn}^E}{\omega_t}\right)^2\right)$$
$$\times\left\{-\left\langle\frac{\omega_{mn}^E}{m}\right\rangle-\frac{e\langle T\rangle}{e\varphi'}\frac{\langle P'\rangle}{\langle P\rangle}+\frac{2}{5}\left(-\frac{3}{2}+\left(\frac{\omega_{mn}^E}{\omega_t}\right)^2\right)\frac{e\langle T\rangle}{e\varphi'}\frac{\langle T'\rangle}{\langle T\rangle}\right\} \qquad \text{(B.2)}$$

Since the frequency power spectrum $|e\Phi_{mn\omega}|^2$ in the rotating system may suffer the same Doppler shift as observed in Labolatory system, integrating over ω in eq.(A.2), the effect of ω_E may be eliminated in the flux Γ_{AN}.

Appendix B Averaged Helical Magnetic Field

Each compornent od helical magnetic files may be written by

$$B_r=b_l\, l\, I_l(z) \sin\phi \tag{B.1}$$

$$B_\theta=(lb_l/\alpha r)I_l(z) \cos\phi \tag{B.2}$$

$$B_z=B_0(1-\varepsilon_h \cos\phi) \tag{B.3}$$

where $z=l\alpha r$, $\phi=\lambda\theta-l\alpha z$, $\alpha=m/lR$, $\varepsilon_h=c_l\dot{I}_l(z)$, $c_l=lb_l/Bo$ and the dot over the modified Bessel function I_l means the differental with respect to the argument.The equation of magnetic field lines is

$$\frac{dz}{B_z}=\frac{dr}{B_r}=\frac{rd\theta}{B_\theta} \tag{B.4}$$

The radial variation of magnetic field line is given by integration of eq.(B.4):

$$r = ro + \Delta r \cos\phi \tag{B.5}$$

where the hehical amplitude is given by $\Delta r=c_l R I_l(z)$.

If we define the averaged magnetic field <B> over the helically deformed magnetic surface by

$$\langle B\rangle = \frac{1}{2\pi}\oint B(r(\phi),\phi)\, d\phi \tag{B.6}$$

introducing eq.(B.5) into eq.(B.3), and applying the Taylor expansion, we have

$$\langle B\rangle = B_o\left(1-\frac{1}{2}\Delta r\varepsilon_h'\right) \tag{B.7}$$

The radial derivative of the averaged field becomes

$$\langle B\rangle' = -B_o c_l^2 (\alpha l)^2 R\dot{I}_l(z)\ddot{I}_l(z) \tag{B.8}$$

From eqs.(B.7) and (B.8), we have

$$\frac{\langle B\rangle'}{\langle B\rangle}=-\frac{m^2}{8R}c_l^2\,(I_{l-1}+I_{l+1})\,(I_{l-2}+2I_l+I_{l+2}) \tag{B.9}$$

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Figures Captions

Fig.1: Variation of charateristic scale length of helically averaged helical magnetic field, $\langle B\rangle'/\langle B\rangle$, versus normalized radius x=r/a for m=8, l=2, and c_l=0.382.

Fig.2: Variations of electric fields Ero and Er calculated by eqs.(40) and (41), respectively, versus x for "high density" case with b_{NC}/a_{NC}=0.15 and a_1=5.

Fig.3: Variation of shear flow effect normalized by E_{TO}, $R_E = (d_{AN}/a_{NC})E_{TO}'/E_{TO}$, versus x for the same case as in Fig.2.

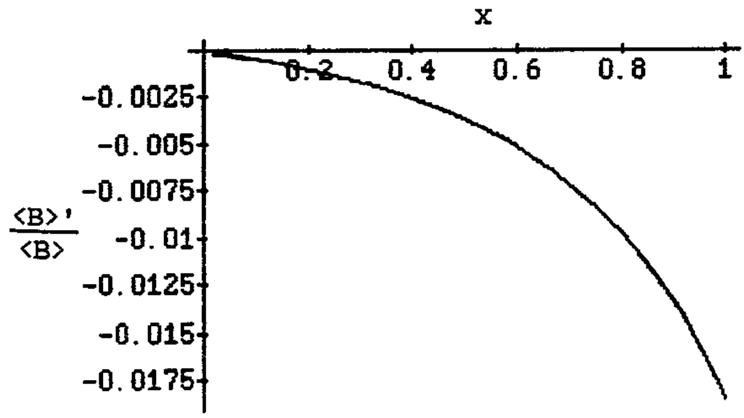

Fig.1

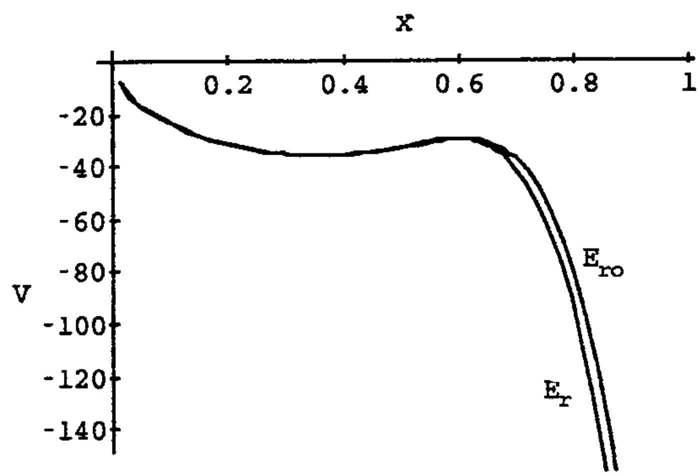

Fig.2

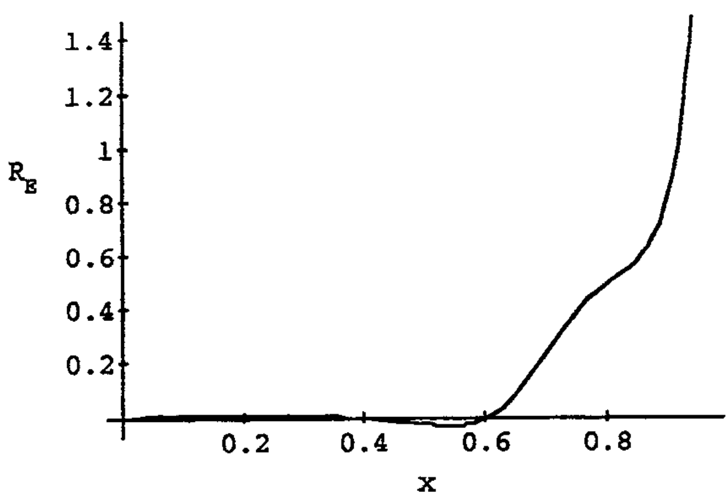

Fig.3

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