ISSN 0915-633X

NATIONAL INSTITUTE FOR FUSION SCIENCE

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(Received – Aug. 17, 1992)

NIFS-172 Sep. 1992

RESEARCH REPORT
NIFS Series

NAGOYA, JAPAN

Physical mechanism of E_ψ-driven current in asymmetric toroidal systems

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Abstract

Physical mechanism of the E_ψ-driven neoclassical current which exists only in asymmetric toroidal systems is clarified. It is found that symmetry-breaking, violating the total angular momentum conservation, changes the direction of flow and heat flux to be damped by the parallel viscosities. The change in the direction depends on the collisionality through the geometric factor which reflects symmetry-breaking. As a result, in contrast to symmetric toroidal systems, the flow and heat flux due to the radial electric field E_ψ can be damped according to the collisionality in asymmetric toroidal systems. If dominant species such as electrons and primary ions have different collisionalities, there remains a difference between the flows due to E_ψ, hence a parallel current proportional directly to the radial electric field, i.e., the E_ψ-driven current, is generated in asymmetric toroidal systems in spite of the charge neutrality condition.

Keywords: E_ψ-driven neoclassical current, the radial electric field, neoclassical theory, asymmetric system, geometric factor, bootstrap current, viscosity

§1. Introduction

As is well known, in axisymmetric toroidal systems, there is no neoclassical parallel currents generated by the radial electric field E_ψ.[1,2] The same result has been obtained in nonaxisymmetric toroidal systems assuming that all the particle species are in the same collisionality regime of the $1/\nu$ regime[3] or the plateau regime[4].

Recently the possibility of the E_ψ-driven current, a non-vanishing neoclassical current generated by the radial electric field in asymmetric toroidal systems, has been indicated in Ref.5. This newly found parallel current can flow only when dominant particle species such as electrons and primary ions have different collisionalities. In symmetric systems, this type of current does not exist under any condition because symmetry-breaking is essential for the E_ψ-driven current to occur.

In this paper, we will show the physical mechanism of this E_ψ-driven neoclassical parallel current in asymmetric systems. Parallel flows are determined by the parallel force balances, i.e., the balance between the parallel viscosity (heat viscosity) and parallel friction (heat friction). The parallel viscosity for a particle species depends directly on the particle orbits, which significantly suffer from the influence of the magnetic field configuration according to the collisionality of the particle species. Thus, the parallel viscosity for a particle species damps the flow and heat flux of the species in such a direction that the magnetic field strength varies according to the collisionality of the species and the magnetic field configuration. It will be shown that the magnitude of the parallel viscosity is mainly determined by the viscosity coefficients, while the direction of the flow and heat flux to be damped is essentially determined by the geometric factor reflecting symmetry-breaking of the magnetic field configuration. Due to the collisionality dependence of the geometric factor in asymmetric systems, the direction of the flow and heat flux of a species to be damped change according to the collisionality of the species. If dominant species such as electrons and primary ions have different collisionalities, contrary to symmetric toroidal systems, the flow and heat flux due to E_ψ are damped according to the collisionality in spite of the total momentum conservation due to Coulomb collision. As a result of it, the E_ψ-driven parallel flow depends on the particle species through its own collisionality, hence, a net parallel current, the E_ψ-driven current, can exist in spite of the charge neutrality.

In marked contrast to it, in symmetric systems, the directions of flow and heat flux to be damped become the same for all the species regardless of their collisionalities since the geometric factor is changed in such a manner that the parallel viscosities damp the flow and heat flux in the direction without symmetry. From this fact and the total momentum conservation of Coulomb collision, the flow and heat flux due to E_ψ are not damped and are independent of particle species. This makes the E_ψ-driven parallel current completely vanish due to the charge neutrality.

The organization of this paper is as follows. The lowest order flow and heat flux are given in §2. The reason is also discussed why the E_ψ-driven current does not exist in both the perpendicular current and Pfirsh-Schlüter current regardless of symmetry in the systems. In §3, the dependence of the flux-surface-averaged parallel viscosity and heat viscosity on the flow and heat flux is considered. The physical meanings of the parallel viscosity, the viscosity coefficients and the geometric factor, and the effect of symmetry are also investigated. The physical mechanism of the E_ψ-driven current is clarified in §4 by placing emphasis on symmetry or symmetry-breaking. Section 5 is devoted to discussion and conclusion.

§2. Lowest order flow and heat flux

In this section we will consider the lowest order flow and heat flux, and especially consider the perpendicular current and Pfirsh-Schlüter current from the viewpoint of the influence of symmetry-breaking in the systems on the role of the radial electric field E_ψ. As well as the standard neoclassical theories,[1-5] the neoclassical ordering[1] and the moment approach method[2] are used. The Boozer coordinates[6] (ψ, θ, ζ) are employed where ψ is the toroidal flux divided by 2π, and θ and ζ are poloidal and toroidal angles, respectively. The typical system size L and frequency ω are defined as $L \simeq |\nabla \ln P|^{-1}$ and $\omega = v_T/L$, respectively, where v_T is a thermal velocity. By using them, the expansion parameter ε is defined as $\varepsilon \equiv \rho/L = \omega/\Omega \ll 1$, where $\rho = v_T/\Omega$ and Ω are the Larmor radius and cyclotron frequency, respectively. The transport ordering $\partial/\partial t \sim \varepsilon^2\omega$ and the drift ordering $e\phi/T \sim 1$ are assumed where e, T, and ϕ are the charge, temperature, and electrostatic potential, respectively.

Let f_a be the distribution function of particle species a in a phase space $(\psi, \theta, \zeta, E, \mu, \varphi)$ where E, μ, and φ are the total energy, magnetic moment, and gyrophase, respectively. In the neoclassical theory the 0th order distribution function is assumed to be a local Maxwellian:

$$f_{aM}(\psi, E) = \frac{n_a}{\pi^{3/2} v_{Ta}^3} \exp\left[-\frac{E - e_a\phi}{T_a}\right]. \tag{1}$$

The expansion of the distribution function around the local Maxwellian distribution gives

$$f_a = f_{aM} + \tilde{f}_{a1} + \bar{f}_{a1} + O(\varepsilon^2) \tag{2}$$

where

$$\tilde{f}_{a1} = -\vec{\rho}_a \cdot \nabla\psi \frac{\partial f_{aM}}{\partial \psi}, \tag{3}$$

$$\bar{f}_{a1} = -\vec{\delta}_a \cdot \nabla\psi \frac{\partial f_{aM}}{\partial \psi}. \tag{4}$$

The 1st order distribution function f_{a1} has been divided into two parts. One is the gyro-varying part $\tilde{f}_{a1}$, which describes particle gyration motions and $\vec{\rho}_a$ is the gyro-radius for species a of the order of ε. The other is the gyro-averaged part $\bar{f}_{a1}$, which describes drift motions of guiding centers and $\vec{\delta}_a$ is the drift-radius for species a of the order of ε. From Eq.(1) we see

$$\frac{\partial f_{aM}}{\partial \psi} = \left[A_{a1}(\psi) L_0^{(3/2)}(x_a) - A_{a2}(\psi) L_1^{(3/2)}(x_a)\right] \frac{e_a}{T_a} f_{aM}, \tag{5}$$

where $L_0^{(\alpha)}(x_a) = 1$ and $L_1^{(\alpha)}(x_a) = \alpha + 1 - x_a$ are the Laguerre polynomial functions with $x_a = (v/v_{Ta})^2$ and

$$A_{a1}(\psi) = \frac{1}{e_a n_a}\frac{dP_a}{d\psi} + \frac{d\phi}{d\psi}, \quad A_{a2}(\psi) = \frac{1}{e_a}\frac{dT_a}{d\psi}. \tag{6}$$

$A_{a1}(\psi)$ and $A_{a2}(\psi)$ are called the thermodynamic force, which are driving sources for particle fluxes (flows), heat fluxes, and currents as will be shown later. From Eqs.(2) to (4) and the moment approach method given in Ref.2, we see the lowest order flow $\vec{u}_a$ and heat flux $\vec{q}_a$ are of the order of ε:[1-5)]

$$\vec{u}_{a1} = u_{\parallel a1}\hat{n} + \vec{u}_{\perp a1}, \quad \vec{q}_{a1} = q_{\parallel a1}\hat{n} + \vec{q}_{\perp a1}, \tag{7}$$

and the perpendicular (parallel) flows and heat fluxes come from the gyro-varing (gyro-averaged) distribution functions. It follows that both the perpendicular and parallel flows and heat fluxes include the thermodynamic forces $A_{a1}(\psi)$ and/or $A_{a2}(\psi)$. The properties of the 1st order flow and heat flux are given by[1–5)]

$$\nabla\cdot\vec{u}_{a1}=0,\qquad \vec{u}_{\perp a1}=A_{a1}(\psi)\frac{\vec{B}\times\nabla\psi}{B^2},\qquad \vec{u}_{a1}\cdot\nabla\psi=0, \tag{8}$$

$$\nabla\cdot\vec{q}_{a1}=0,\ \ \vec{q}_{\perp a1}=\frac{5P_a}{2}A_{a2}(\psi)\frac{\vec{B}\times\nabla\psi}{B^2},\ \ \vec{q}_{a1}\cdot\nabla\psi=0. \tag{9}$$

The gyration motion is the motion around a magnetic field line, so that it is insensitive to whether the system has a symmetry or not. Hence, the lowest order perpendicular particle flow and heat flux obtained by the gyro-varying distribution function lead to the same qualitative results between symmetric systems and asymmetric systems. Since the $\vec{E}\times\vec{B}$ drift is independent of particle species in both systems, the perpendicular current becomes independent of the radial electric field $E_\psi=-d\phi/d\psi$ due to the charge neutrality $\sum_a e_a n_a=0$:

$$\vec{J}_\perp=\sum_a e_a n_a\vec{u}_{\perp a1}=\frac{dP}{d\psi}\frac{\vec{B}\times\nabla\psi}{B^2} \tag{10}$$

where $P=\sum_a P_a$.

The same result is obtained for the Pfirsh-Schlüter current. Using $\nabla\cdot\vec{u}_{a1}=0$ and $\nabla\cdot\vec{q}_{a1}=0$, we can see

$$Bu_{\|a1} = - A_{a1}(\psi)\{g_2-\frac{B^2}{\langle B^2\rangle}\langle g_2\rangle\}+\frac{B^2}{\langle B^2\rangle}\left\langle Bu_{\|a1}\right\rangle, \tag{11}$$

$$Bq_{\|a1} = -\frac{5P_a}{2}A_{a2}(\psi)\{g_2-\frac{B^2}{\langle B^2\rangle}\langle g_2\rangle\}+\frac{B^2}{\langle B^2\rangle}\left\langle Bq_{\|a1}\right\rangle, \tag{12}$$

where the function g_2 is the solution of the following equation:

$$\vec{B}\cdot\nabla\left(\frac{g_2}{B^2}\right)=\vec{B}\times\nabla\psi\cdot\nabla\left(\frac{1}{B^2}\right),\ \ g_2(B=B_{max})=0. \tag{13}$$

Here, $\langle\cdots\rangle$ indicates a flux-surface average and solvability conditions of Eq.(13) are satisfied regardless of symmetry in the systems. From the Ampere's law, the current density must be divergence-free, i.e., $\nabla\cdot\vec{J}=0$. The Pfirsh-Schlüter current, which vanishes in flux-surface averaging, is a current flowing in the plasma to satisfy the divergence-free condition. As the perpendicular current is independent of E_ψ, the resultant Pfirsh-Schlüter current is also independent of E_ψ:

$$BJ_{\|\,PS}=-\sum_a A_{a1}(\psi)\{g_2-\frac{B^2}{\langle B^2\rangle}\langle g_2\rangle\}=-\frac{dP}{d\psi}\{g_2-\frac{B^2}{\langle B^2\rangle}\langle g_2\rangle\}. \tag{14}$$

The parallel current which does not vanish in the flux-surface averaging consists of flux-surface-averaged parallel flows. Since these parallel flows are determined by the parallel force balance, we will consider the parallel force balance in the next section.

§3. Flow and heat flux dependence of the flux-surface-averaged parallel viscosity and heat viscosity

To determine the 1st order flux-surface-averaged parallel flow $\left\langle Bu_{\|a1}\right\rangle$ and heat flux $\left\langle Bq_{\|a1}\right\rangle$ in Eqs.(11) and (12), the flux-surface-averaged parallel momentum and heat flux balance equations are used:[2)]

$$\begin{bmatrix} \left\langle \vec{B}\cdot\nabla\cdot\overleftrightarrow{\Pi}_a \right\rangle \\ -\left\langle \vec{B}\cdot\nabla\cdot\overleftrightarrow{\Theta}_a \right\rangle \end{bmatrix} = \begin{bmatrix} \left\langle \vec{B}\cdot\vec{F}_{a1} \right\rangle \\ -\left\langle \vec{B}\cdot\vec{F}_{a2} \right\rangle \end{bmatrix} + \begin{bmatrix} n_a e_a \left\langle \vec{B}\cdot\vec{E}^{(A)} \right\rangle \\ 0 \end{bmatrix}, \tag{15}$$

where $\overleftrightarrow{\Pi}_a$ and $\overleftrightarrow{\Theta}_a$ are the viscosity and heat viscosity tensors for particle species a, respectively. $\vec{F}_{a1}$ and $\vec{F}_{a2}$ are the friction and heat friction of particle species a and an external inductive electric field $\vec{E}^{(A)}$ is included.

The 1st order parallel friction $\left\langle \vec{B}\cdot\vec{F}_{a1}\right\rangle$ and heat friction $\left\langle \vec{B}\cdot\vec{F}_{a2}\right\rangle$ are expressed in terms of the 1st order parallel flow $\left\langle Bu_{\|a1}\right\rangle$ and heat flux $\left\langle Bq_{\|a1}\right\rangle$ for all particle species:[2)]

$$\begin{bmatrix} \left\langle \vec{B}\cdot\vec{F}_{a1} \right\rangle \\ -\left\langle \vec{B}\cdot\vec{F}_{a2} \right\rangle \end{bmatrix} = \sum_b \begin{bmatrix} l_{11}^{ab} & l_{12}^{ab} \\ l_{21}^{ab} & l_{22}^{ab} \end{bmatrix} \begin{bmatrix} \left\langle Bu_{\|b1} \right\rangle \\ -\dfrac{2}{5P_b}\left\langle Bq_{\|b1} \right\rangle \end{bmatrix}, \tag{16}$$

where l_{ij}^{ab} $(i,j=1,2)$ are the friction coefficients defined in Ref.2, which are independent of the magnetic field $\vec{B}$ and collisionality, and are valid in an arbitrary toroidal system. The relevant physical roles of the frictions are to transfer the momentum between unlike particles and to conserve the total momentum: $\sum_a \vec{F}_{a1} = 0$. Thus, the friction coefficients have following relations:

$$l_{ij}^{ab} = l_{ji}^{ba}, \quad \sum_b l_{i1}^{ab} = 0, \text{ for } (i=1,2). \tag{17}$$

The 1st order flux-surface-averaged parallel viscosity and heat viscosity are written as follows:[2)]

$$\begin{aligned} \left\langle \vec{B}\cdot\nabla\cdot\overleftrightarrow{\Pi}_{a1} \right\rangle &= \left\langle (P_{\perp a1} - P_{\|a1})\hat{n}\cdot\nabla B \right\rangle, \\ -\left\langle \vec{B}\cdot\nabla\cdot\overleftrightarrow{\Theta}_{a1} \right\rangle &= -\left\langle (\Theta_{\perp a1} - \Theta_{\|a1})\hat{n}\cdot\nabla B \right\rangle. \end{aligned} \tag{18}$$

In Eq.(18), $(P_{\perp a1} - P_{\|a1})$ and $(\Theta_{\perp a1} - \Theta_{\|a1})$ are obtained by solving the 1st order drift kinetic equation for the gyro-averaged distribution function $\bar{f}_{a1}$.

The parallel viscosities for particle species a, which are determined by the 1st order gyro-averaged distribution function $\bar{f}_{a1}$, directly depend on the orbits for the particles species a and the orbits significantly reflect the influence of the magnetic field configuration according to the collisionality of the species. The main physical roles of the parallel viscosities are to transfer the momentum between like particles through the damp of the flow and heat flux in such a direction that the magnetic field strength varies. The parallel viscosities for species a are expressed in terms of the flow and heat flux for species a as follows:[5)]

$$\begin{bmatrix} \left\langle \vec{B} \cdot \nabla \cdot \overleftrightarrow{\Pi}_{a1} \right\rangle \\ -\left\langle \vec{B} \cdot \nabla \cdot \overleftrightarrow{\Theta}_{a1} \right\rangle \end{bmatrix} = \begin{bmatrix} \mu_{a1} & \mu_{a2} \\ \mu_{a2} & \mu_{a3} \end{bmatrix} \begin{bmatrix} \langle \vec{u}_{a1} \cdot \nabla \vartheta^* \rangle \\ -\frac{2}{5P_a} \langle \vec{q}_{a1} \cdot \nabla \vartheta^* \rangle \end{bmatrix}, \tag{19}$$

where

$$\vartheta^* = (I + \langle G_{BS} \rangle_a)\theta + (J - \iota \langle G_{BS} \rangle_a)\zeta, \tag{20}$$

and μ_{aj} $(j = 1 \sim 3)$ are the viscosity coefficients and $\langle G_{BS} \rangle_a$ is the geometric factor for the particle species a. In asymmetric systems, both quantities have different dependences on the magnetic configuration according to the collisionality. Hence, we can understand that the magnitude of the parallel viscosities or the magnitude of the flow and heat flux for species a to be damped is dominated by the viscosity coefficients μ_{aj} according to the collisionality of the species a itself, while the direction of the flow and heat flux for the species a to be damped is determined by the geometric factor $\langle G_{BS} \rangle_a$ depending on the collisionality.

The asymptotic forms of the viscosity coefficients are proportional to f_t/f_c, λ_a/λ_{PL}, and $(\lambda_a/\lambda_{PS})^2$ in the asymptotic limit of the $1/\nu$, plateau, and Pfirsh-Schlüter collisionality regime, respectively, where f_t and f_c are the fraction of trapped and circulating particles, respectively, and λ_a, λ_{PL}, and λ_{PS} are the mean free path for particle species a, characteristic lengths for the plateau and Pfirsh-Schlüter regime, respectively.[5)] Roughly speaking, the magnitude of the viscosity coefficients decrease with increase in the collision frequency. This fact suggests that the differences of flow and heat flux between like particles become small with increase in the collision frequency. The asymptotic form of the geometric factor for species a: $\langle G_{BS}\rangle_a$ depends only on the magnetic field configuration. The dependence on the magnetic field configuration, however, is completely different in different collisionality regimes as well as the viscosity coefficients. In the $1/\nu$ regime, the geometric factor is given by

$$\langle G_{BS}\rangle_{1/\nu} = \frac{1}{f_t}\left\{\langle g_2\rangle - \frac{3\langle B^2\rangle}{4B_{max}^2}\int_0^1 d\lambda \frac{\lambda\langle g_4\rangle}{\langle g_1\rangle}\right\}, \tag{21}$$

where the function g_1 and the fraction of trapped particles f_t are, respectively,

$$g_1 = \sqrt{1-\lambda\frac{B_{max}}{B}}, \quad f_t = 1 - \frac{3}{4}\frac{\langle B^2\rangle}{B_{max}^2}\int_0^1 d\lambda \frac{\lambda}{\langle g_1\rangle}. \tag{22}$$

The function g_2 is the solution of Eq.(13), and the function g_4 is given by

$$\vec{B}\cdot\nabla\left(\frac{g_4}{g_1}\right) = \vec{B}\times\nabla\psi\cdot\nabla\left(\frac{1}{g_1}\right), \quad g_4(B=B_{max}) = 0. \tag{23}$$

The functions g_2 and g_4 appearing Eq.(21) are related with a function g_3 as follows:

$$g_2 = \frac{3}{2B_{max}}B^2\int_0^{B_{max}/B} d\lambda\frac{g_1}{B}g_3, \quad g_4 = -2B_{max}g_1\frac{\partial}{\partial\lambda}\left(\frac{g_1}{B}g_3\right), \tag{24}$$

where the function g_3 is a solution of the following equation:

$$\vec{B}\cdot\nabla\left(\frac{g_1}{B}g_3\right) = \vec{B}\times\nabla\psi\cdot\nabla\left(\frac{g_1}{B}\right), \quad g_3(B=B_{max}) = 0. \tag{25}$$

The function g_3 appears in solving the lowest order drift kinetic equation in the $1/\nu$ collisionality regime:[1-3)]

$$\frac{v_\parallel}{B}\vec{B}\cdot\nabla\bar{f}_{a1c}^{(0)} + \vec{v}_{Da}\cdot\nabla\psi\frac{\partial f_{aM}}{\partial\psi} = 0, \tag{26}$$

where the radial drift of the guiding center $\vec{v}_{Da}\cdot\nabla\psi$ is given by

$$\vec{v}_{Da}\cdot\nabla\psi = \frac{v_\parallel}{\Omega_a}\vec{B}\times\nabla\psi\cdot\nabla\left(\frac{v_\parallel}{B}\right). \tag{27}$$

Taking account of the functional forms of Eqs.(26) and (27), and $v_\parallel = \sigma v g_1$, $(\sigma = \pm 1)$, and considering the fact that both circulating and trapped particles exist in the $1/\nu$ regime, we see that the function g_3 appears to reflect the radial drift of the guiding center of both trapped and circulating particles. The trapped particles not drifting off a flux surface have a residual momentum on the flux surface, which comes from differences between the parallel and anti-parallel flows on the flux surface. The differences are brought about by the density gradient, the temperature gradient changing the banana radii, and the radial electric field making the banana radii and the turning points change. Since the viscosities can transfer the momentum between like particles, this residual momentum of trapped particles is transferred to circulating particles.[7] As the parallel viscosities essentially damp the flow and heat flux in such a direction that the magnetic field strength varies, the geometric factor acts so as to damp the flow and heat flux in such a direction that the variation of the magnetic field strength is the largest. The geometric factor in the Pfirsh-Schlüter regime is given by[8]

$$\langle G_{BS}\rangle_{PS} = \frac{\left\langle \dfrac{(\hat{n}\cdot\nabla B)}{2B}\hat{n}\cdot\nabla\left[B^2\langle g_2\rangle - \left\langle B^2\right\rangle g_2\right]\right\rangle}{\langle(\hat{n}\cdot\nabla B)^2\rangle} \tag{28}$$

In the Pfirsh-Schlüter collisionality regime no trapped particles exist essentially due to high collisionality. Thus, the geometric factor is expressed in terms of g_2, which reflects the radial drift integrated by the pitch angle variable λ as shown in Eq.(24). The geometric factor in this limit is so small that we can neglect this factor in Eq.(19), which means that the parallel viscosities damp the parallel flow and heat flux in the Pfirsh-Schlüter collisionality regime (this is more clear in Eq.(32)). This fact is consistent with the particle motions in this high collisional regime, i.e., the fact that the parallel flow is small due to high collisionality. The asymptotic form of the geometric factor in the plateau regime is given in Ref.5. As similar as the viscosity coefficients, the geometric factor roughly decreases with increase in the collision frequency. Thus, the direction of flow and heat flux of species a to be damped by the parallel viscosities significantly changes according to its own collisionality in asymmetric systems. If we construct an adequate formulae for the viscosity coefficients and the geometric factor connecting all the collisionality regimes, Eq.(19) may hold for an arbitrary collisionality in general toroidal systems.

Now we will consider the effect of symmetry. When the helical symmetry, i.e., $Q = Q(\psi,\vartheta)$ holds, where $\vartheta = L\theta - M\zeta$, we see

$$\vec{B}\times\nabla\psi\cdot\nabla Q = \langle G_{BS}\rangle_{sym}\vec{B}\cdot\nabla Q, \tag{29}$$

$$\langle G_{BS}\rangle_{sym} = \frac{LJ+MI}{L\iota - M}. \tag{30}$$

Thus, the symmetry implies that the variation perpendicular to both a flux surface and a magnetic field line on it is proportional to the variation along the magnetic field line with the proportional coefficient of a surface quantity $\langle G_{BS}\rangle_{sym}$ given by Eq.(30). This is a natural result because a symmetry makes an ignorable coordinate on the two dimensional flux surface. Alternatively, the symmetry makes all the magnetic field line be equivalent with each other, which leads to the fact that all the single particle orbits do not drift off the flux surface. Thus, the geometric factor in each asymptotic limit have the same form given by Eq.(30), so that Eq.(19) becomes

$$\begin{bmatrix} \left\langle \vec{B}\cdot\nabla\cdot\overleftrightarrow{\Pi}_{a1}\right\rangle \\ -\left\langle \vec{B}\cdot\nabla\cdot\overleftrightarrow{\Theta}_{a1}\right\rangle \end{bmatrix} = \frac{J+\iota I}{L\iota - M}\begin{bmatrix} \mu_{a1} & \mu_{a2} \\ \mu_{a2} & \mu_{a3} \end{bmatrix}\begin{bmatrix} \langle \vec{u}_{a1}\cdot\nabla\vartheta\rangle \\ -\frac{2}{5P_a}\langle \vec{q}_{a1}\cdot\nabla\vartheta\rangle \end{bmatrix}. \quad (31)$$

The toroidal symmetry (axisymmetry) correspond to $M=0$ case, and the poloidal symmetry to $L=0$ case. Hence, we see that in symmetric systems, the flow and heat flux in the symmetric direction do not enter in the expression for the flux-surface-averaged parallel viscosities, i.e., they are not damped by the parallel viscosities regardless of the collisionality regime of particle species, which directly reflects the fact that the parallel viscosities damp the flow and heat flux in such a direction that the magnetic field strength varies. Thus, in symmetric systems, the direction of the flow and heat flux to be damped by the parallel viscosities is independent of the particle species and their collisionality regimes, while the magnitude of the flow and heat flux to be damped is dependent of the collisionality regime of particle species through the viscosity coefficients.

Using the relations of parallel and perpendicular components with poloidal and toroidal ones for the flow and heat flux:

$$\begin{aligned} Bu_{\|a1} &= I\vec{u}_{a1}\cdot\nabla\theta + J\vec{u}_{a1}\cdot\nabla\zeta, \quad A_{a1}(\psi) = \vec{u}_{a1}\cdot\nabla\theta - \iota\vec{u}_{a1}\cdot\nabla\zeta, \\ Bq_{\|a1} &= I\vec{q}_{a1}\cdot\nabla\theta + J\vec{q}_{a1}\cdot\nabla\zeta, \quad A_{a2}(\psi) = \frac{2}{5P_a}\left[\vec{q}_{a1}\cdot\nabla\theta - \iota\vec{q}_{a1}\cdot\nabla\zeta\right], \end{aligned} \quad (32)$$

we obtain another form of Eq.(19):

$$\begin{bmatrix} \left\langle \vec{B}\cdot\nabla\cdot\overleftrightarrow{\Pi}_{a}\right\rangle \\ -\left\langle \vec{B}\cdot\nabla\cdot\overleftrightarrow{\Theta}_{a}\right\rangle \end{bmatrix} = \begin{bmatrix} \mu_{a1} & \mu_{a2} \\ \mu_{a2} & \mu_{a3} \end{bmatrix}\begin{bmatrix} \left\langle Bu_{\|a}\right\rangle + \langle G_{BS}\rangle_a A_{a1}(\psi) \\ -\frac{2}{5P_a}\left\langle Bq_{\|a}\right\rangle - \langle G_{BS}\rangle_a A_{a2}(\psi) \end{bmatrix}. \quad (33)$$

This form is used to determine the flux-surface-averaged parallel flow and heat flux in the next section.

§4. Physical mechanism of E_ψ-driven current

Solving Eq.(15) by using Eqs.(16) and (33), we have the following solutions:

$$\left\langle B u_{\|a1} \right\rangle = \sum_b \langle G_{BS}\rangle_b \{L_{2a+1\ 2b+1} A_{b1}(\psi) - L_{2a+1\ 2b+2} A_{b2}(\psi)\} - \sum_b N_{2a+1\ 2b+1} e_b n_b \left\langle \vec{B}\cdot\vec{E}^{(A)} \right\rangle, \tag{34}$$

$$-\frac{2}{5P_a}\left\langle B q_{\|a1}\right\rangle = \sum_b \langle G_{BS}\rangle_b \{L_{2a+2\ 2b+1} A_{b1}(\psi) - L_{2a+2\ 2b+2} A_{b2}(\psi)\}. \tag{35}$$

The neoclassical current or bootstrap current is given by

$$\left\langle B J_{\|BS} \right\rangle = \sum_{a,b} \langle G_{BS}\rangle_b e_a n_a \{L_{2a+1\ 2b+1} A_{a1}(\psi) - L_{2a+1\ 2b+2} A_{a2}(\psi)\}, \tag{36}$$

where $L_{2a+k\ 2b+k}$ $(k = 1,2)$ and $N_{2a+1\ 2b+1}$ are transport matrices defined in Ref.5. The transport matrix $L_{2a+k\ 2b+1}$ $(k = 1,2)$ has an important property:

$$\sum_b L_{2a+k\ 2b+1} = \begin{cases} -1 & \text{for} \quad k=1 \\ 0 & \text{for} \quad k=2 \end{cases} \tag{37}$$

This property stems from the total momentum conservation of the Coulomb collision operator given by Eq.(17).[5)]

The thermodynamic force $A_{a1}(\psi)$ given by Eq.(6) includes the radial electric field $E_\psi = -d\phi/d\psi$ and so we may expect that the neoclassical current or bootstrap current given by Eq.(36) include the component which is directly driven by the radial electric field E_ψ. However, as is well known the E_ψ-driven neoclassical current does not exist in axisymmetric systems[1-2)] and in non-axisymmetric systems with all particle species having the same collisionality.[3-4)] In order to clarify why the E_ψ-driven neoclassical current vanishes in such systems, let us consider helical symmetric systems. As is shown in the previous section, in symmetric systems, the direction of the flow and heat flux to be damped by the parallel viscosities is the same for all particle species regardless of their collisionalities, which is reflected by the fact that the geometric factor of species a: $\langle G_{BS}\rangle_a$ has the common expression given by Eq.(30). Therefore, using Eq.(30) and Eq.(37), we obtain the flux-averaged-parallel flow and heat flux in the helical symmetric systems:

$$\left\langle B u_{\|a1} \right\rangle = \langle G_{BS}\rangle_{sym} E_\psi$$

$$+\ \langle G_{BS}\rangle_{sym}\sum_b\left\{L_{2a+1\ 2b+1}\frac{1}{e_bn_b}\frac{dP_b}{d\psi}-L_{2a+1\ 2b+2}\frac{1}{e_b}\frac{dT_b}{d\psi}\right\}$$

$$-\ \sum_b N_{2a+1\ 2b+1}e_bn_b\left\langle\vec{B}\cdot\vec{E}^{(A)}\right\rangle, \tag{38}$$

$$-\frac{2}{5P_a}\left\langle Bq_{\|a1}\right\rangle = \langle G_{BS}\rangle_{sym}\sum_b\left\{L_{2a+2\ 2b+1}\frac{1}{e_bn_b}\frac{dP_b}{d\psi}-L_{2a+2\ 2b+2}\frac{1}{e_b}\frac{dT_b}{d\psi}\right\}. \tag{39}$$

The role of E_ψ is similar to the perpendicular flow and heat flux, i.e., the E_ψ-driven parallel flow is independent of particle species and the parallel heat flux is independent of E_ψ. Thus, E_ψ-driven flow in an arbitrary direction does not depend on particle species, and owing to the charge neutrality $\sum_a e_an_a = 0$ the neoclassical current given by Eq.(36) is independent of the radial electric field:

$$\left\langle BJ_{\|BS}\right\rangle = \langle G_{BS}\rangle_{sym}\sum_{a,b}e_an_a\left\{L_{2a+1\ 2b+1}\frac{1}{e_bn_b}\frac{dP_b}{d\psi}-L_{2a+1\ 2b+2}\frac{1}{e_b}\frac{dT_b}{d\psi}\right\}. \tag{40}$$

Now, let us consider why the parallel flow due to E_ψ is independent of the particle species in symmetric toroidal systems. To do this, we substitute Eqs.(38) and (39) into Eq.(33), which reveals that the flow and heat flux appearing in the parallel viscosities do not include E_ψ in symmetric systems. This fact physically means that the parallel viscosities do not damp the flow and heat flux due to the radial electric field E_ψ. This result is well known in axisymmetric systems in a different expression, i.e., the poloidal rotation, which appears in the parallel viscosities, is independent of E_ψ. The reason why the radial electric field E_ψ must not be damped by the parallel viscosities is explained as follows: In symmetric systems, an arbitrary total angular momentum in the symmetric direction must be conserved in the steady state up to the order of ε^2, and the change in the angular momentum occurs by small off-diagonal components of the parallel viscosities on the ε^3 order and non-Coulombic frictions (charge exchanges etc.), as is well known in axisymmetric systems.[1-2] Corresponding to this fact, all the radial particle fluxes (classical, banana-plateau, and Pfirsh-Schlüter fluxes) of the order of ε^2 are intrinsically ambipolar by the total momentum conservation due to the Coulomb collision, which means that the electrostatic potential ϕ can not be determined in the steady state. Hence, arbitrariness of the angular momentum in the symmetric direction may be attributed to arbitrariness of the flow due to the radial electric field $E_\psi = -d\phi/d\psi$. Thus, the flow and heat flux due to

E_ψ must not appear in the parallel viscosities of the order of ε, otherwise they suffer from damping, yielding the contradiction to the total angular momentum conservation.

As is discussed above, the parallel viscosities in symmetric systems are described only by the flow and heat flux in such a direction that the symmetry does not exist, and such the flow and heat flux consist of both the parallel and perpendicular components. On the other hand, the E_ψ-driven perpendicular flow ($\vec{E} \times \vec{B}$ drift) is independent of particle species and the perpendicular heat flux is independent of the E_ψ as shown in Eqs.(8) and (9). Therefore, in order to satisfy the requirement by symmetry that the flow and heat flux due to E_ψ do not appear in the parallel viscosities, the E_ψ-driven parallel flow must be independent of particle species and E_ψ-driven parallel heat flux must vanish. Consequently, E_ψ-driven flows in an arbitrary direction are independent of particle species and the E_ψ-driven current vanishes completely due to the charge neutrality.

On the other hand, in asymmetric systems, there is no conservation of the angular momentum in any direction. Corresponding to it, there are trapped particles drifting off a flux surface, which contribute to the $1/\nu$ ripple diffusion in the low collisionality regime ($1/\nu$ and plateau regime[4]). The $1/\nu$ ripple diffusion comes from the radial particle flux in the 2nd order of ε. The electrostatic potential ϕ or the radial electric field E_ψ can be determined by the ambipolar condition of the radial current by the $1/\nu$ ripple diffusion in the steady-state (on the order of ε^2). Therefore, the radial electric field E_ψ is not arbitrary in contrast with the symmetric case. This fact suggests that the E_ψ-driven flow and heat flux may be damped by the parallel viscosities. Indeed, the substitution Eqs.(34) and (35) into Eq.(33) shows that the flow and heat flux due to E_ψ are damped by the parallel viscosities except for the situation that all particle species have the same collisionality. As is shown in the previous section, in asymmetric systems the direction of the flow and heat flux of a species to be damped by the parallel viscosities are dependent on the collisionality of the species itself through the geometric factor. Consequently, when dominant species such as electrons and primary ions have different collisionalities, the E_ψ-driven parallel flows are different between the species, resulting in the fact that the E_ψ-driven neoclassical net current can flow in asymmetric systems in spite of the charge neutrality.

Let us consider a 3-fluids plasma consisting of electrons, primary ions, and impurity ions. If impurity ions with the charge number of z_I exist in the Pfirsh-Schlüter collisionality regime, Eq.(36) reduces to

$$\begin{aligned}\left\langle BJ_{\|BS}\right\rangle &= L_{11}\left[\langle G_{BS}\rangle_e - \langle G_{BS}\rangle_i\right] en_e E_\psi \\ &+ L_{11}\left[\langle G_{BS}\rangle_e \frac{dP_e}{d\psi} + \langle G_{BS}\rangle_i \frac{n_e}{z_i n_i}\frac{dP_i}{d\psi}\right] - L_{12}\langle G_{BS}\rangle_e n_e \frac{dT_e}{d\psi} \\ &+ \langle G_{BS}\rangle_i \left[\left(\frac{L_{11}-\alpha}{1+\alpha} + \frac{z_I n_I}{n_e}\right) L_{34} + \left(\frac{\alpha(1+L_{11})}{1+\alpha} - \frac{z_I n_I}{n_e}\right) L_{54}\right] \frac{n_e}{z_i}\frac{dT_i}{d\psi}\end{aligned} \quad (41)$$

where the subscripts e, i, and I indicate electrons, primary ions, and impurity ions, respectively, and $\alpha = z_I^2 n_I/n_i \ll 1$ or < 1. In deriving Eq.(41), the relation resulting from the total momentum conservation given by Eq.(37) and the small mass ratio expansion $\sqrt{m_e/m_i}, \sqrt{m_e/m_I} \ll 1$ have been used. The viscosity coefficients of the impurity ion μ_{Ik} are very small, because impurity ions exist in the Pfirsh-Schlüter regime. Consequently, thermodynamic forces from impurity ions vanish. The geometric factor of particle species a: $\langle G_{BS}\rangle_a$ depends only on its own collisionality and have the asymptotic form depending only on the magnetic field configuration in each asymptotic limit. Therefore, when electrons and primary ions exist in the same collisionality regime E_ψ-driven neoclassical current disappears. The transport coefficients L_{11} and L_{12} are mainly dominated by electrons and L_{34} and L_{54} are dominated by primary ions. When electrons are in the plateau regime and primary ions exist in $1/\nu$ regime, $L_{11}, L_{12} \ll L_{34}, L_{54}$ and the effect of E_ψ-driven current is negligible as well as pressure-driven neoclassical current. The contribution of E_ψ-driven current is most effective when electrons and primary ions are in the $1/\nu$ and plateau regime, respectively. Since $e\phi/T \sim 1$ in usual plasmas, the E_ψ-driven current is comparable with the pressure-driven current. Note that if the radial electric field E_ψ is determined by the $1/\nu$ ripple diffusion, the direction of the E_ψ-driven current is always opposite to that of the pressure-driven neoclassical current.

§5. Conclusion and Discussion

Physical mechanism of the newly found E_ψ-driven neoclassical parallel current in asymmetric systems[5)] has been clarified. This current can flow when dominant particle species such as electrons and primary ions have different collisionalities. In symmetric systems, this type of current does not exist under any condition because symmetry-breaking plays an essential role for this current to exist.

It has been emphasized that the crucial difference of the property of parallel viscosities between symmetric and asymmetric systems comes from the direction of the flow and heat flux to be damped. In asymmetric systems this direction of a particle species changes according to its own collisionality through the geometric factor, while the direction is independent of the collisionality in symmetric systems. When the collisionalities of dominant species such as electrons and primary ions are different in asymmetric toroidal systems, the directions of the flow and heat flux to be damped by the parallel viscosities become different between both species. As a result of it, the flow and heat flux due to E_ψ are damped by the parallel viscosities with different rates according to the collisionality of each species, which leads to the fact that the E_ψ-driven parallel flow depends on the particle species. This is the most important point indicated in this paper. Consequently, the E_ψ-driven neoclassical parallel current can flow in spite of the charge neutrality. Simple analysis for a three fluids plasma has shown that this current is most effective when electrons and primary ions exist in the $1/\nu$ and plateau collisionality regimes, respectively, under the situation that impurity ions exist in Pfirsh-Schlüter regime. It has also indicated that the radial electric field E_ψ determined by the $1/\nu$ ripple diffusion always reduces the pressure-driven bootstrap current.

It should be noted that all the present results stem from influences of symmetry and symmetry-breaking in the system on the particle orbits. The existence of a symmetry vanishes the E_ψ-driven parallel current on a flux surface because of the resultant conservation of an arbitrary angular momentum. The symmetry makes the radial particle fluxes across a flux surface be intrinsic ambipolar, which leads to arbitrariness of E_ψ. In contrast to it, symmetry-breaking does not make E_ψ-driven parallel current on a flux surface vanish generally and does not make the radial particle fluxes be intrinsic ambipolar, which leads to determination of E_ψ.

Acknowledgments

The authors are grateful to Professor T.Sato, Professor M.Wakatani, and Doctor K.Watanabe for their fruitful discussions and interest in this work. One of the authors

(N.N) wishes to thank Doctor K.C.Shaing for his fruitful discussions and comments.

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