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N. Nakajima and M. Okamoto

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Geometrical effects of the magnetic field on the neoclassical flow, current and rotation in general toroidal systems

N.Nakajima and M.Okamoto

National Institute for Fusion Science, Nagoya 464-01, Japan

abstract

In order to clarify geometrical effects of the magnetic field on neoclassical theory, the neoclassical parallel particle flow, heat flux, current and plasma rotation of a multispecies plasma in general toroidal systems are examined in several collisionality regimes. The quantitative and qualitative differences between axisymmetric (tokamaks) and non-axisymmetric toroidal systems (stellarator, heliotron/torsatron) appear mainly through a geometrical factor which prescribes the parallel flow due to the gradients of the density, temperature, and electrostatic potential. In axisymmetric toroidal systems the geometrical factor reduces to the same expression in all collisionality regimes due to axisymmetry. By contrast, in non-axisymmetric toroidal systems, it changes drastically depending on the magnetic field structure and the collisionality regime. Thus, the poloidal flow has the radial electric field dependence. When the geometrical factor is very small, the ion parallel flow almost vanishes and the ion rotation consists of the diamagnetic and $\vec{E} \times \vec{B}$ flows (perpendicular flows).

Keywords: neoclassical theory, non-axisymmetric system, parallel flow, rotation

I. INTRODUCTION

The transport theory of axisymmetric toroidal systems (tokamaks) has been extensively studied by many authors. About 20 years ago, the geometry of the magnetic field was first taken into account in it, where the effects of particles trapped in the weak-field region, i.e., the neoclassical effects were considered.[1] Later, a variational method based on the principle of minimum entropy production was used to obtain a complete set of neoclassical transport coefficients. [2–4] Moreover, as a new approach the moment approach was developed, which is based on the direct solution of the exact moment equations using the closure relations of viscous stress tensors written in terms of the particle flow and heat flux.[5]

Recently the neoclassical theory using the moment approach is extended in many aspects. To study the momentum transport in axisymmetric toroidal systems, the neoclassical theory was extended so that the toroidal flow with $M_t \sim 1$ is allowed where M_t is the toroidal Mach number. [6] The assumption of axisymmetry is excluded to construct the neoclassical theory in non-axisymmetric toroidal systems (stellarator,heliotron/torsatron) in several collisionality regimes. [7–12] Concerning with the anomalous diffusion fluctuations are included in the neoclassical theory. [13] To explain the L-H transition in tokamaks ion orbit losses are taken into acount together with the poloidal viscosity being valid to $M_p \sim 1$ where M_p is the poloidal Mach number. [14,15] The neoclassical current in a plasma with an impurity and a fast ion beam is considered numerically using flux-surface-averaged parallel momentum and heat flux equations to investigate a steady state tokamak. [16] According to the problems to be examined in which some have axisymmetry and others not, various extended neoclassical theories will be constructed. Therefore, it is now meaningful for us to reconsider the neoclassical theory of a multispecies plasma in the general toroidal systems paying attention to the geometrical effects of the magnetic field.

In this paper, in order to clarify geometrical effects of the magnetic field on the neoclassical theory in general toroidal systems, the neoclassical parallel particle flow, heat flux, current and plasma rotation of a multispecies plasma are examined using the moment approach on the basis of the original papers [5,7,9] under the assumptions of no fluctuations, no external sources and losses except for a fast ion beam and an external inductive electric field, steady state, and $| \vec{u}_a | \ll v_{Ta}$ where $\vec{u}_a$ and v_{Ta} are the macro and thermal velocity of species a, respectively. Hence, we might have a point of view of unifying understanding the neoclassical theory in general toroidal systems. Three collisionality regimes, i.e., the $1/\nu$ (in non-axisymmetric toroidal systems) or banana (in axisymmetric toroidal systems), plateau, and Pfirsh-Schlüter collisionality regimes are examined separately.

The flux-surface-averaged parallel momentum and heat flux balance equations are used to obtain the flux-surface-averaged parallel particle flow and heat flux. In these equations, the flux-surface-averaged parallel friction and viscosity and the external inductive electric field balance. The friction coeffecients are classical quantities in the sense that they are independent of the magnetic field $\vec{B}$. The geometrical effects of the magnetic field come from the viscosity terms. [5] In axisymmetric toroidal systems the magnetic field strength B is a function with respect to the flux label, for example, the volume V and the poloidal angle θ, so that there is only one type of trapped particles and the geometrical effects of the magnetic field structure in the flux-surface-averaged parallel viscosity appear through the connection length proportional to $\sqrt{< B^2 > / < (\hat{n} \cdot \nabla B)^2 >}$ where $\hat{n} = \vec{B}/B$ and $<>$ is the flux-surface average, and through the fraction of trapped particles except for some magnetic-surface quantities, for example, rotational transform ι, and $dV/d\Phi_T$ where Φ_T is the toroidal flux. In contrast to the axisymmetric toroidal systems, in non-axisymmetric toroidal systems the strength of the magnetic field B also depends on the toroidal angle ζ, thus there are many types of trapped particles and another type of factor reflecting the magnetic field structure appears in the parallel viscosity, which is called the flux-surface-averaged "geometrical factor" $< G_{BS} >$. [7,17,18]

This factor prescribes the parallel particle flow due to the gradients of the density, temparature, and electrostatic potential. Although this factor reduces to the same expression consisting of well known flux surface quantities in all collisionality regimes in the axisymmetric toroidal systems, in the non-axisymmetric toroidal system it changes drastically depending on the magnetic field structure and the collisionality regime. Therefore, this geometrical factor mainly makes differences to neoclassical results between axisymmetric and non-axisymmetric toroidal systems.

The self adjointness and the momentum conservation of the Coulomb collision operator give common properties to the multispecies plasma in general toroidal systems in spite of the geometrical effects of the magnetic field. The parallel particle flow due to the radial electric field is independent of the particle species and the parallel heat flux is independent of the radial electric field, if there are no other external sources and losses without the momentum conservation, for example, the charge exchange loss. Consequently, the charge neutrality guarantees the neoclassical current (the bootstrap curent) be independent of the radial electric field. Quantitative changes, however, appear in the neoclassical current mainly through the geometrical factor. [7,17–19] The magnitude of the bootstrap current, which is proportional to the flux-surface-averaged geometrical factor, changes considerably in the non-axisymmetric systems according to the magnetic field structure. Moreover, the poloidal and toroidal rotations change both qualitatively and quantitatively. In the axisymmetric toroidal systems, the poloidal rotation in the multispecies plasma is independent of the radial electric field. [20] In non-axisymmetric toroidal systems, however, the poloidal rotation has the radial electric field dependence. The radial electric field dependence becomes larger as the geometrical factor decreases. Consequently, if interactions between thermal ions and fast ions are negligible, then the ion parallel flow almost disappears in non-axisymmetric toroidal systems with a sufficiently small geometrical factor, which contrasts with the fact that the ion parallel flow streams to cancel the ion poloidal flow in axisymmetric toroidal systems.

The organization of this paper is as follows. The magnetic coordinate systems are given in Sec.II. Basically Hamada coordinates [21,22] are used, however, in some calculations Boozer coordinates [23,24] are useful. In Sec.III, flux-surface-averaged parallel momentum and heat flux balance equations of a multispecies plasma in a general toroidal system are shown. The flux-surface-averaged geometrical factor $< G_{BS} >$ and the viscosity coefficients μ are indiceted in the $1/\nu$ (in non-axisymmetric toroidal systems) or banana (in axisymmetric toroidal systems), plateau, and Pfirsh-Schlüter collisionality regimes, separately. The general properties of the flux-surface-averaged parallel particle flow and heat flux, the parallel current, and the poloidal and toroidal rotations are discussed in Sec.IV. The results using the small mass ratio expansion are also shown. In Sec.V, the explicit forms of quantities discussed in Sec.IV are given for a simple electron-ion plasma to obtain clear results from physical views. Discussion and conclusion are given in Sec.VI.

II. MAGNETIC COORDINATE SYSTEMS

To describe a general non-axisymmetric torus, Hamada coordinates (V, θ, ζ) [21,22] and Boozer coordinates $(\psi, \theta_B, \zeta_B)$ [23,24] are used. In Hamada coordinates the magnetic field $\vec{B}$ and the Jacobian $\sqrt{g}$ are expressed as follows:

$$\vec{B} = \Phi_T' \nabla V \times \nabla\theta + \Phi_P' \nabla\zeta \times \nabla V, \tag{1}$$

$$\sqrt{g} = \frac{1}{\nabla V \times \nabla\theta \cdot \nabla\zeta} = 1, \tag{2}$$

where V is the volume, θ and ζ are the poloidal and toroidal angles, respectively, $'$ means d/dV, and Φ_P and Φ_T are poloidal and toroidal fluxes, respectively. Similary, in Boozer coordinates the magnetic field and the Jacobian are shown by

$$\vec{B} = \nabla\psi \times \nabla\theta_B + \iota \nabla\zeta_B \times \nabla\psi, \tag{3}$$

$$\vec{B} = I\nabla\theta_B + J\nabla\zeta_B + \tilde{\beta}\nabla\psi, \tag{4}$$

$$\sqrt{g_B} = \frac{1}{\nabla\psi \times \nabla\theta_B \cdot \nabla\zeta_B} = \frac{J + \iota I}{B^2}, \tag{5}$$

where $\psi = \Phi_T/2\pi$, θ_B and ζ_B are the poloidal and toroidal angles, respectively, ι is the rotational transform, and $\tilde{\beta}$ is a periodic function with respect to θ_B and ζ_B. $2\pi J$ is the total poloidal current outside the flux surface and $2\pi I$ is the total toroidal current inside the flux surface. In Hamada coordinates the magnetic field is expressed in the contravariant form. In Boozer coordinates, however, it is expressed in both the contravariant and covariant forms. For these characteristics of Boozer coordinates some calculations are performed using Boozer coordinates. The Hamada coordinates are related with Boozer coordinates using the following equations:

$$\begin{aligned} V &= V(\psi), && (6) \\ \theta &= \frac{1}{2\pi}\{\theta_B + \iota G(\psi, \theta_B, \zeta_B)\}, && (7) \\ \zeta &= \frac{1}{2\pi}\{\zeta_B + G(\psi, \theta_B, \zeta_B)\}, && (8) \end{aligned}$$

where the generation function G is a periodic function with respect to θ_B and ζ_B and it is given by

$$\begin{aligned} \vec{B}\cdot\nabla G &= \frac{1}{J+\iota I}\{<B^2> - B^2\}, && (9) \\ <B^2> &= 2\pi\Phi_T'(J+\iota I). && (10) \end{aligned}$$

Above properties are used, for example, as follows:

$$\begin{aligned} & <\vec{B}\times\nabla V\cdot\nabla\theta> \\ =\ & \frac{dV}{d\Phi_T} <\nabla\psi\times\nabla(\theta_B+\iota G)\cdot\tilde{\beta}\nabla\psi + I\nabla\theta_B + J\nabla\zeta_B> \\ =\ & \frac{dV}{d\Phi_T} < \frac{1}{\sqrt{g_B}}(J - \iota I\frac{\partial G}{\partial\zeta_B} + \iota J\frac{\partial G}{\partial\theta_B}) > \\ =\ & \frac{dV}{d\Phi_T}\frac{J<B^2>}{J+\iota I} \\ =\ & 2\pi J. && (11) \end{aligned}$$

In the following expressions, terms with J or I are calculated using Boozer coordinates. For numerical calculations Boozer coordinates are useful [25,26] together with the transformation given by Eqs.(6)-(9). Spectrum broadness of B is different between Hamada and Boozer coordinates. Generally, Boozer coordinates need less number of spectra than Hamada coordinates. Detail discussion on Hamada and Boozer coordinates will be given elsewhere.

III. FLUX-SURFACE-AVERAGED PARALLEL MOMENTUM AND HEAT FLUX BALANCE EQUATIONS

The moment equations to describe the macroscopic plasma behavior are given by Ref.5, thus they are not repeated here. We will consider a multispecies plasma consisting of electrons and N species ions. The particle species are indicated subscript a. The electron is expressed by $a = 0$ and the i-th ion is expressed by $a = i$ ($i = 1 \sim N$). The flow $\vec{u}_a$ and the heat flux $\vec{q}_a$ have the following forms:

$$\vec{u}_a = u_{\|a}\frac{\vec{B}}{B} + \vec{u}_{\perp a}, \tag{12}$$

$$\vec{q}_a = q_{\|a}\frac{\vec{B}}{B} + \vec{q}_{\perp a}. \tag{13}$$

The lowest orders of moment equations are [4,5,7]

$$\nabla \cdot \vec{u}_a = 0, \tag{14}$$

$$\vec{u}_{\perp a} = (\frac{P_a'}{e_a n_a} + \phi')\frac{\vec{B} \times \nabla V}{B^2}, \tag{15}$$

$$\nabla \cdot \vec{q}_a = 0, \tag{16}$$

$$\vec{q}_{\perp a} = \frac{5}{2}\frac{P_a T_a'}{e_a}\frac{\vec{B} \times \nabla V}{B^2}. \tag{17}$$

Using above equations the parallel flow and the heat flux without the flux-surface average are given by

$$Bu_{\|a} = -(\frac{P_a'}{e_a n_a} + \phi')\{g_2 - \frac{B^2}{< B^2 >} < g_2 >\} + \frac{B^2}{< B^2 >} < Bu_{\|a} >, \tag{18}$$

$$Bq_{\|a} = -\frac{5}{2}\frac{P_a T_a'}{e_a}\{g_2 - \frac{B^2}{< B^2 >} < g_2 >\} + \frac{B^2}{< B^2 >} < Bq_{\|a} >, \tag{19}$$

where

$$\vec{B}\cdot\nabla\left(\frac{g_2}{B^2}\right)=\vec{B}\times\nabla V\cdot\nabla\left(\frac{1}{B^2}\right),\quad g_2(B_{max})=0. \tag{20}$$

For axisymmetric toroidal systems g_2 becomes

$$g_2=<G_{BS}>_T(1-\frac{B^2}{B_{max}}), \tag{21}$$

therefore

$$\{g_2-\frac{B^2}{<B^2>}<g_2>\}=<G_{BS}>_T(1-\frac{B^2}{<B^2>}). \tag{22}$$

Here $<G_{BS}>_T$ is defined as follows:

$$<G_{BS}>_T\equiv\frac{J}{\iota}\frac{dV}{d\psi}=\frac{2\pi J}{\Phi'_P}. \tag{23}$$

The meaning of $<G_{BS}>_T$ will be shown later.

To determine the flux-surface-averaged parallel flow $<Bu_{\|a}>$ and heat flux $<Bq_{\|a}>$, the flux-surface-averaged parallel momentum and heat flux balance equations are used: [5,7,9]

$$\begin{bmatrix} <\vec{B}\cdot\nabla\cdot\overleftrightarrow{\Pi}_a> \\ -<\vec{B}\cdot\nabla\cdot\overleftrightarrow{\Theta}_a> \end{bmatrix}=\begin{bmatrix} <\vec{B}\cdot\vec{F}_{a1}> \\ -<\vec{B}\cdot\vec{F}_{a2}> \end{bmatrix}+\begin{bmatrix} <\vec{B}\cdot\vec{F}_{af1}> \\ -<\vec{B}\cdot\vec{F}_{af2}> \end{bmatrix}+\begin{bmatrix} n_a e_a<\vec{B}\cdot\vec{E}^{(A)}> \\ 0 \end{bmatrix}, \tag{24}$$

where $\overleftrightarrow{\Pi}_a$ and $\overleftrightarrow{\Theta}_a$ are the viscosity and heat viscosity tensor, respectively. $\vec{F}_{aj}$ and $\vec{F}_{afj}$ ($j=1$ or 2) are frictions of a species a with thermal species and fast ions due to the neutral beam injection, respectively. Also an external inductive electric field $\vec{E}^{(A)}$ is included.

The flux-surface-averaged parallel friction $<\vec{B}\cdot\vec{F}_{a1}>$ and the heat friction $<\vec{B}\cdot\vec{F}_{a2}>$ with thermal species are given by [5]

$$\begin{bmatrix} <\vec{B}\cdot\vec{F}_{a1}> \\ -<\vec{B}\cdot\vec{F}_{a2}> \end{bmatrix}=\sum_{b=0}^{N}\begin{bmatrix} l_{11}^{ab} & l_{12}^{ab} \\ l_{21}^{ab} & l_{22}^{ab} \end{bmatrix}\begin{bmatrix} <Bu_{\|b}> \\ -\frac{5}{2}\frac{<Bq_{\|b}>}{P_b} \end{bmatrix} \tag{25}$$

where l_{ij}^{ab} $(i,j=1,2)$ are the friction coefficients between species a and b, which are independent of the magnetic field $\vec{B}$. Therefore, the expressions in Eq.(25) for the parallel friction forces in terms of the parallel flow and heat flux are same between axisymmetric and non-axisymmetric toroidal systems, and valid in all collisionality regimes. From the self adjointness of the Coulomb collision operator

$$l_{ij}^{ab} = l_{ji}^{ba}, \tag{26}$$

and the momentum conservation of the Coulomb collision operator gives the following property:

$$\sum_{b=0}^{N} l_{i1}^{ab} = 0 \quad (i = 1, 2). \tag{27}$$

The flux-surface-averaged parallel friction and heat friction with fast ions are expressed by [5,27]

$$\begin{bmatrix} < \vec{B} \cdot \vec{F}_{0f1} > \\ - < \vec{B} \cdot \vec{F}_{0f2} > \end{bmatrix} = \frac{n_e m_e}{\tau_{ee}} \frac{Z_f^2 n_f}{n_e} \begin{bmatrix} 1 \\ \frac{3}{2} \end{bmatrix} < B u_{\parallel f} > \quad \text{for electron,} \tag{28}$$

$$< \vec{B} \cdot \vec{F}_{af1} > = < \vec{B} \cdot \vec{F}_{0f1} > \frac{v_c^3}{\bar{u}_f^3} \frac{n_a e_a^2}{\sum_{b=1}^{N} n_b e_b^2 / m_b} \left(\frac{1}{m_f} + \frac{1}{m_a} \right) \quad \text{for ions,} \tag{29}$$

$$v_c^3 = \frac{3\sqrt{\pi}}{4} \frac{m_e}{m_f} \sum_{b=1}^{N} \frac{n_b e_b^2 m_f}{n_e m_b} v_{Te}^3 \tag{30}$$

$$\frac{1}{\bar{u}_f^3} \equiv \frac{\int \frac{v_\parallel}{v^3} f_f d^3 v}{\int v_\parallel f_f d^3 v} \tag{31}$$

where the quantities with subscript f denote those who belong to fast ions, v_c is the critical velocity at which the drag exerted by the background electrons on the fast ion beam will equal that of the background ions on the fast ion beam. Here the conditions $v_{T0} \gg | \vec{u}_f | \gg v_{Ta}$ for $a = 1 \sim N$ have been assumed where $\vec{u}_f$ is the beam velocity and the heat friction of thermal ions with fast ions is neglected.

The flux-surface-averaged parallel viscosity and the parallel heat viscosity have the following general form

$$\begin{bmatrix} < \vec{B} \cdot \nabla \cdot \overleftrightarrow{\Pi}_a > \\ - < \vec{B} \cdot \nabla \cdot \overleftrightarrow{\Theta}_a > \end{bmatrix} = \frac{n_a m_a}{\tau_{aa}} \begin{bmatrix} \mu_{a1} & \mu_{a2} \\ \mu_{a2} & \mu_{a3} \end{bmatrix} \begin{bmatrix} < B u_{\parallel a} > \\ -\frac{2}{5} \frac{< B q_{\parallel a} >}{p_a} \end{bmatrix}$$

$$+ \quad \frac{n_a m_a}{\tau_{aa}} < G_{BS} > \begin{bmatrix} \mu_{a1} & \mu_{a2} \\ \mu_{a2} & \mu_{a3} \end{bmatrix} \begin{bmatrix} \frac{T_a}{e_a}(\frac{P_a'}{P_a} + \frac{e_a \phi'}{T_a}) \\ -\frac{T_a'}{e_a} \end{bmatrix}. \qquad (32)$$

where μ_{aj} $(j = 1 \sim 3)$ are the non-dimensional viscosity coefficients which reflect the self adjointness of the Coulomb collision operator and $< G_{BS} >$ is the flux-surface-averaged geometrical factor. The above general form is obtained by modifying the original results. [5,7,9] The viscosity coefficients and the flux-surface-averaged geometrical factor suffer the geometrical dependences of the magnetic field. Hence, the difference of the parallel flows between axisymmetric and non-axisymmetric toroidal systems, and collisionality dependence of the parallel flows appear through the parallel viscocity. Here, the $1/\nu$ collisionality regime (in non-axisymmetric toroidal systems) or the banana collisionality regime (in axisymmetric toroidal systems), the plateau collisionality regime, and the Pfirsh-Schlüter collisionality regime are treated separately.

In the $1/\nu$ or banana collisionality regime, the viscosity coefficients are [5,7]

$$\mu_{aj} = \frac{f_t}{f_c} \bar{\mu}_{aj} \quad (j = 1 \sim 3) \qquad (33)$$

where f_t are the fraction of trapped particles and $f_c = 1 - f_t$. The coefficients $\bar{\mu}_{aj}$ $(j = 1 \sim 3)$ are not dependent on the magnetic field but dependent on the particle constitution. The flux-surface-averaged geometrical factor $< G_{BS} >$ is given by the following equations: [7]

$$< G_{BS} > \;=\; \frac{1}{f_t} \left\{ \langle g_2 \rangle - \frac{3}{4} \frac{\langle B^2 \rangle}{B_{max}^2} \int_0^1 \frac{\langle g_4 \rangle}{\langle g_1 \rangle} \lambda \, d\lambda \right\}, \qquad (34)$$

$$f_t \;=\; 1 - \frac{3}{4} \frac{\langle B^2 \rangle}{B_{max}^2} \int_0^1 \frac{1}{\langle g_1 \rangle} \lambda \, d\lambda, \qquad (35)$$

$$g_1 \;=\; \sqrt{1 - \lambda \frac{B}{B_{max}}}, \qquad (36)$$

$$\vec{B} \cdot \nabla \left(\frac{g_4}{g_1} \right) \;=\; \vec{B} \times \nabla V \cdot \nabla \left(\frac{1}{g_1} \right), \quad g_4(B_{max}) = 0, \qquad (37)$$

where g_2 is given by Eq.(20). For axisymmetric toroidal systems

$$g_4 =< G_{BS} >_T (1 - \frac{g_1}{\sqrt{1 - \lambda}}). \qquad (38)$$

From this result and Eq.(21) we have

$$< G_{BS} >=< G_{BS} >_T \tag{39}$$

for axisymmetric toroidal systems, so that we use the definition given by Eq.(23).

In the plateau collisionality regime, the viscosity coefficients μ_{aj} $(j = 1 \sim 3)$ are expressed as follows: [5,9,11]

$$\mu_{aj} = \frac{\lambda_a}{\lambda_{PL}} \bar{\mu}_{aj} \quad (j = 1 \sim 3), \tag{40}$$

$$\frac{1}{\lambda_{PL}} \equiv \frac{\Phi_T'(\iota\mu_p + \mu_t)}{< B^2 >}, \tag{41}$$

where $\lambda_a = \tau_{aa} v_{ta}$ is the mean free path of species a, $\bar{\mu}_{a1} = 2$, $\bar{\mu}_{a2} = -1$, $\bar{\mu}_{a3} = 13/2$, and λ_{PL} is a characteristic length of the magnetic field in the plateau collisionality regime. Note that $\bar{\mu}_{aj}$ $(j = 1 \sim 3)$ is independent of the particle constitution. μ_p and μ_t are given by

$$\mu_p = \frac{\sqrt{\pi}}{2} \langle \frac{\hat{n} \cdot \nabla B}{\Phi_T'} \sum_{m \neq 0, n \neq 0} \frac{\partial B_{mn}(\theta, \zeta)/\partial \theta}{2\pi \mid \iota m + n \mid} \rangle, \tag{42}$$

$$\mu_t = \frac{\sqrt{\pi}}{2} \langle \frac{\hat{n} \cdot \nabla B}{\Phi_T'} \sum_{m \neq 0, n \neq 0} \frac{\partial B_{mn}(\theta, \zeta)/\partial \zeta}{2\pi \mid \iota m + n \mid} \rangle, \tag{43}$$

where $\hat{n} = \vec{B}/B$ and $B = \sum_{m,n} B_{mn}(\theta, \zeta) = \sum_{m,n} B_{mn} cos(2\pi m\theta + 2\pi n\zeta + \varphi_{mn})$ is assumed. The flux-surface-averaged geometrical factor $< G_{BS} >$ is given by

$$< G_{BS} >= \frac{dV}{d\psi} \frac{J\mu_p - I\mu_t}{\iota\mu_p + \mu_t} \tag{44}$$

For axisymmetric toroidal systems, $\mu_t = 0$, thus the flux-surface-averaged geometrical facor $< G_{BS} >$ reduces to $< G_{BS} >_T$ as well as in the $1/\nu$ or banana collisionality regime. It is noted that for non-axisymmetric toroidal systems $\mu_p = 0$ and $\mu_t \neq 0$ on the magnetic axis so that $< G_{BS} >$ reduces to $-IdV/d\psi$ on it.

In the Pfirsh-Schlüter collisionality regime, the viscosity coefficients are [5,7]

$$\mu_{aj} = (\frac{\lambda_a}{\lambda_{PS}})^2 \bar{\mu}_{aj} \cdot (j = 1 \sim 3), \tag{45}$$

$$\frac{1}{{\lambda_{PS}}^2} \equiv \frac{3}{2} \frac{< (\hat{n} \cdot \nabla B)^2 >}{< B^2 >^{\cdot}}, \tag{46}$$

where $\bar{\mu}_{aj}$ $(j = 1 \sim 3)$ are not dependent on the magnetic field but dependent on the particle constitution as well as in the $1/\nu$ collisionality regime. λ_{PS} is the characteristic length in the Pfirsh-Schlüter collisionality regime. The flux-surface-averaged geometrical factor is given by

$$< G_{BS} >= 2\pi \frac{< \frac{1}{B}(J\frac{\partial B}{\partial \theta} - I\frac{\partial B}{\partial \zeta})(\hat{n} \cdot \nabla B) >}{< (\hat{n} \cdot \nabla B)^2 >} \tag{47}$$

As well as in the $1/\nu$ or banana collisionality regime and in the plateau collisionality regime, for axisymmetric toroidal systems $< G_{BS} >$ reduces to $< G_{BS} >_T$.

It follows from above discussions that the geometrical effects of the magnetic field on the flux-surface-averaged parallel particle flow and heat flux come from the flux-surface-averaged geometrical factor $< G_{BS} >$ and the viscosity coefficients $\mu_{aj}(j = 1 \sim 3)$. In axisymmetric toroidal systems, the geometrical factor reduces to the same expression given by Eq.(23) in all collisionality regimes. In contrast with it, in non-axisymmetric toroidal systems, the expression of the geometrical factor changes according to the collisionality regime as well as the viscosity coefficients, i.e., it has the dependence of the fraction of trapped particles in the $1/\nu$ collisionality regime, and the connection length dependence in the plateau and Pfirsh-Schlüter collisionality regimes. The normalized geometrical factor $< G_{BS} >_N$, which is defined as $< G_{BS} > / < G_{BS} >_T$ and generally smaller than unity, may indicate the deviation of the general toroidal system from axisymmetric toroidal system. The normalized geometrical factor changes drastically according to the magnetic field structure in the $1/\nu$ collisionality regime and it becomes very small and even negative. [7,17–19] In the plateau and Pfirsh-Schüter collisionality regimes, it is considerably small. Then, in the general toroidal systems where the normalized geometrical factor is small enough, considerably different properties of the parallel flow, heat flux, current, and plasma rotation from those of axisymmetric toroidal systems are expected. At this point we will discuss in the next section.

IV. PARALLEL FLOW, CURRENT, AND PLASMA ROTATION

To solve Eq.(24), we define the following quantities ($a, b = 0 \sim N$):

$$\vec{F}_t \equiv \{F_{t\ i}\}: \quad F_{t\ 2a+1} \equiv <\vec{B}\cdot\vec{F}_{a1}>, \quad F_{t\ 2a+2} \equiv -<\vec{B}\cdot\vec{F}_{a2}>, \tag{48}$$

$$\vec{F}_f \equiv \{F_{f\ i}\}: \quad F_{f\ 2a+1} \equiv <\vec{B}\cdot\vec{F}_{af1}>, \quad F_{f\ 2a+2} \equiv -<\vec{B}\cdot\vec{F}_{af2}>\delta_{a0}, \tag{49}$$

$$\vec{E}_A \equiv \{E_{A\ i}\}: \quad E_{A\ 2a+1} \equiv e_a n_a <\vec{B}\cdot\vec{E}^{(A)}>, \quad E_{A\ 2a+2} \equiv 0, \tag{50}$$

$$\vec{V} \equiv \{V_i\}: \quad V_{2a+1} \equiv <\vec{B}\cdot\nabla\cdot\overleftrightarrow{\Pi}_a>, \quad V_{2a+2} \equiv -<\vec{B}\cdot\nabla\cdot\overleftrightarrow{\Theta}_a>, \tag{51}$$

$$\vec{G}_1 \equiv \{G_{1\ i}\}: \quad G_{1\ 2a+1} \equiv \frac{P_a'}{e_a n_a}, \quad G_{1\ 2a+2} \equiv -\frac{T_a'}{e_a}, \tag{52}$$

$$\vec{G}_2 \equiv \{G_{2\ i}\}: \quad G_{2\ 2a+1} \equiv \phi', \quad G_{2\ 2a+2} \equiv 0, \tag{53}$$

$$\vec{X} \equiv \{X_i\}: \quad X_{2a+1} \equiv <Bu_{\|a}>, \quad X_{2a+2} \equiv -\frac{2}{5}\frac{<Bq_{\|a}>}{P_a}, \tag{54}$$

$$\overleftrightarrow{L}_t \equiv \{l_{ij}\}: \quad l_{k+2a\ k'+2b} \equiv l_{kk'}^{ab} \quad (k, k' = 1, 2), \tag{55}$$

$$\overleftrightarrow{\mu}_a \equiv \frac{n_a m_a}{\tau_{aa}}\begin{bmatrix} \mu_{a1} & \mu_{a2} \\ \mu_{a2} & \mu_{a3} \end{bmatrix}, \quad \overleftrightarrow{\mu} \equiv \begin{bmatrix} \overleftrightarrow{\mu}_0 & & & \\ & \overleftrightarrow{\mu}_1 & & \\ & & \cdot & \\ & & & \overleftrightarrow{\mu}_N \end{bmatrix}. \tag{56}$$

Using above definitions, Eq.(24) becomes

$$\overleftrightarrow{V} = \vec{F}_t + \vec{F}_f + \vec{E}_A, \tag{57}$$

where

$$\vec{F}_t = \overleftrightarrow{L}_t \vec{X}, \tag{58}$$

$$\overleftrightarrow{V} = \overleftrightarrow{\mu}\{\vec{X} + <G_{BS}>(\vec{G}_1 + \vec{G}_2)\}. \tag{59}$$

Solving Eq.(57), we have the following flux-surface-averaged parallel flow and heat flux:

$$\vec{X} = <G_{BS}>\overleftrightarrow{L}(\vec{G}_1 + \vec{G}_2) - \overleftrightarrow{N}\vec{F}_f - \overleftrightarrow{N}\vec{E}_A, \tag{60}$$

where

$$\overleftrightarrow{L} \equiv \overleftrightarrow{N}\overleftrightarrow{\mu} = \{L_{ij}\}, \tag{61}$$

$$\overleftrightarrow{N} \equiv \overleftrightarrow{M}^{-1} = \{N_{ij}\} = \{N_{ji}\}, \tag{62}$$

$$\overleftrightarrow{M} \equiv \overleftrightarrow{L}_t - \overleftrightarrow{\mu} = \{M_{ij}\} = \{M_{ji}\}. \tag{63}$$

Properties of the friction coefficients l_{ij}^{ab} given by Eq.(27) are succeeded in matrix $\overleftrightarrow{M}$ as follows:

$$\sum_{b=0}^{N} M_{2a+1\ 2b+1} = -\mu_{a1} \quad (a = 0 \sim N), \tag{64}$$

$$\sum_{b=0}^{N} M_{2a+2\ 2b+1} = -\mu_{a2} \quad (a = 0 \sim N), \tag{65}$$

Using Eqs.(64) and (65), we see

$$\begin{aligned}
\sum_{b=0}^{N} L_{i\ 2b+1} &= \sum_{b=0}^{N} \{N_{i\ 2b+1}\mu_{b1} + N_{i\ 2b+2}\mu_{b2}\} \\
&= \frac{1}{|\overleftrightarrow{M}|}\sum_{b=0}^{N}\{\tilde{M}_{2b+1\ i}\mu_{b1} + \tilde{M}_{2b+2\ i}\mu_{b2}\} \\
&= -\frac{1}{|\overleftrightarrow{M}|}\sum_{b=0}^{N}\sum_{j=1}^{2N+2} \tilde{M}_{j\ i} M_{j\ 2b+1} \\
&= -\sum_{b=0}^{N} \delta_{i\ 2b+1} \\
&= \begin{cases} -1 & \text{for } i = 1, 3, \cdots, 2N+1 \\ 0 & \text{for } i = 2, 4, \cdots, 2N+2 \end{cases}
\end{aligned} \tag{66}$$

where the linear algebraic relations $N_{ij} = \tilde{M}_{ji}/\,|\overleftrightarrow{M}|$, and $\sum_{k=1}^{2N+2} \tilde{M}_{ki} M_{kj} = \delta_{ij}\,|\overleftrightarrow{M}|$ ($\tilde{M}_{ij}$ is the cofactor of the matrix $\overleftrightarrow{M}$) are used. Then,

$$\overleftrightarrow{L}\,\vec{G}_2 = \begin{cases} -\phi' & \text{for odd } i \\ 0 & \text{for even } i \end{cases} \tag{67}$$

Consequently, the flux-surface-averaged parallel particle flow and heat flux have the following forms:

$$< Bu_{\|a} > \ = \ - < G_{BS} > \phi' + < G_{BS} > \sum_{b=0}^{N}\{L_{2a+1\ 2b+1}\frac{P_b'}{e_b n_b} - L_{2a+1\ 2b+2}\frac{T_b'}{e_b}\}$$

$$
\begin{aligned}
& - \sum_{b=0}^{N} N_{2a+1\ 2b+1}\{< \vec{B}\cdot\vec{F}_{bf1} > + e_b n_b < \vec{B}\cdot\vec{E}^{(A)} >\}, \\
& + N_{2a+1\ 2} < \vec{B}\cdot\vec{F}_{0f2} >, \qquad (68)
\end{aligned}
$$

$$
\begin{aligned}
-\frac{2}{5}\frac{< Bq_{\|a} >}{P_a} &= < G_{BS} > \sum_{b=0}^{N}\{L_{2a+2\ 2b+1}\frac{P_b'}{e_b n_b} - L_{2a+2\ 2b+2}\frac{T_b'}{e_b}\} \\
&- \sum_{b=0}^{N} N_{2a+2\ 2b+1} < \vec{B}\cdot\vec{F}_{bf1} > + N_{2a+2\ 2} < \vec{B}\cdot\vec{F}_{0f2} > . \qquad (69)
\end{aligned}
$$

It is shown in Eqs.(68) and (69) that the flux-surface-averaged geometrical factor $< G_{BS} >$ prescribes the parallel particle flow and the parallel heat flux due to the thermodynamical forces consisting of the gradients of the density, temperature, and electrostatic potential. Other driving terms, i.e., the friction with fast ions and the external inductive electric field are independent of the geometrical factor. If there are no other external momentum sources and losses without the momentum conservation, for example, the charge-exchange loss, then the self adjointness and the momentum conservation of the Coulomb collision operator guarantee that the effects of the radial electric filed on the parallel particle flow are independent of the particles species and that the parallel heat flux is independent of the radial electric field. Therefore, under the condition of the charge neutrality $\sum_{a=0}^{N} e_a n_a = 0$ the neoclassical current is independent of the radial electric field:

$$
\begin{aligned}
BJ_{\|} &= \sum_{a=0} e_a n_a B u_{\|a} + < BJ_{\|f} > \\
&= BJ_{\|PS} + < BJ_{\|SP} > + < BJ_{\|BS} > + < BJ_{\|OH} >, \qquad (70)
\end{aligned}
$$

where

$$
\begin{aligned}
BJ_{\|PS} &= -\sum_{a=0}^{N} P_a'\{g_2 - \frac{B^2}{< B^2 >} < g_2 >\} \\
&\quad : \text{ Pfirsh-Schlüter current,} \qquad (71) \\
< BJ_{\|SP} > &= -\sum_{a,b=0}^{N} e_a n_a N_{2a+1\ 2b+1} e_b n_b < \vec{B}\cdot\vec{E}^{(A)} > \\
&\quad : \text{ Spitzer current,} \qquad (72)
\end{aligned}
$$

$$< BJ_{\|BS} > \;=\; < G_{BS} > \sum_{a,b=0}^{N} \{ L_{2a+1\ 2b+1} \frac{e_a n_a}{e_b n_b} P_b' - L_{2a+1\ 2b+2} \frac{e_a}{e_b} n_a T_b' \}$$

$$: \text{ Bootstrap current,} \tag{73}$$

$$< BJ_{\|OH} > \;=\; < BJ_{\|} >_f - \sum_{a=0}^{N} e_a n_a \{ \sum_{b=0}^{N} N_{2a+1\ 2b+1} < \vec{B} \cdot \vec{F}_{bf1} > - N_{2a+1\ 2} < \vec{B} \cdot \vec{F}_{0f2} > \}$$

$$: \text{ Ohkawa current.} \tag{74}$$

Here $< BJ_{\|f} >= Z_f e n_f < Bu_{\|f} >$ is used and the Pfirsh-Schlüter current is defined as $< BJ_{\|PS} >= 0$. In axisymmetric systems, the geometrical factor of the Pfirsh-Schlüter current is given by Eq.(22).

Using Eqs.(12) and (15), we have the expressions of the plasma rotation in the Boozer coordinates $(\psi, \theta_B, \zeta_B)$:

$$< \vec{u}_a \cdot \nabla \theta_B > \;=\; \frac{\iota}{J + \iota I} < Bu_{\|a} > + \frac{J}{J + \iota I} (\frac{1}{e_a n_a} \frac{dP_a}{d\psi} + \frac{d\phi}{d\psi}), \tag{75}$$

$$< \vec{u}_a \cdot \nabla \zeta_B > \;=\; \frac{1}{J + \iota I} < Bu_{\|a} > - \frac{I}{J + \iota I} (\frac{1}{e_a n_a} \frac{dP_a}{d\psi} + \frac{d\phi}{d\psi}). \tag{76}$$

Then, the poloidal and the toroidal rotations are connected by the following equations:

$$< \vec{u}_a \cdot \nabla \theta_B > - \iota < \vec{u}_a \cdot \nabla \zeta_B > \;=\; (\frac{1}{e_a n_a} \frac{dP_a}{d\psi} + \frac{d\phi}{d\psi}), \tag{77}$$

$$I < \vec{u}_a \cdot \nabla \theta_B > + J < \vec{u}_a \cdot \nabla \zeta_B > \;=\; < Bu_{\|a} > . \tag{78}$$

Substitution of Eq.(68) into Eqs.(75) and (76) gives the poloidal and toroidal rotations of the multispecies plasma in the general toroidal system:

$$\begin{aligned} &< \vec{u}_a \cdot \nabla \theta_B > \\ = \;& \frac{J}{J + \iota I} (1 - \frac{< G_{BS} >}{< G_{BS} >_T}) (\frac{1}{e_a n_a} \frac{dP_a}{d\psi} + \frac{d\phi}{d\psi}) \\ + \;& \frac{J}{J + \iota I} \frac{< G_{BS} >}{< G_{BS} >_T} \sum_{b=0}^{N} \{ L_{2a+1\ 2b+1} (\frac{1}{e_b n_b} \frac{dP_b}{d\psi} - \frac{1}{e_a n_a} \frac{dP_a}{d\psi}) - L_{2a+1\ 2b+2} \frac{1}{e_b} \frac{dT_b}{d\psi} \} \\ - \;& \frac{\iota}{J + \iota I} \Big(\sum_{b=0}^{N} N_{2a+1\ 2b+1} \{ < \vec{B} \cdot \vec{F}_{bf1} > + e_a n_a e_b n_b < \vec{B} \cdot \vec{E}^{(A)} > \} \\ & - N_{2a+1\ 2} < \vec{B} \cdot \vec{F}_{0f2} > \Big), \end{aligned} \tag{79}$$

$$
\begin{aligned}
&< \vec{u}_a \cdot \nabla \zeta_B > \\
= & -\frac{1}{\iota(J+\iota I)} (J \frac{< G_{BS} >}{< G_{BS} >_T} + \iota I) \, (\frac{1}{e_a n_a} \frac{dP_a}{d\psi} + \frac{d\phi}{d\psi}) \\
+ & \frac{J}{\iota(J+\iota I)} \frac{< G_{BS} >}{< G_{BS} >_T} \sum_{b=0}^{N} \{ L_{2a+1 \ 2b+1} (\frac{1}{e_b n_b} \frac{dP_b}{d\psi} - \frac{1}{e_a n_a} \frac{dP_a}{d\psi}) - L_{2a+1 \ 2b+2} \frac{1}{e_b} \frac{dT_b}{d\psi} \} \\
- & \frac{1}{J+\iota I} \left(\sum_{b=0}^{N} N_{2a+1 \ 2b+1} \{ < \vec{B} \cdot \vec{F}_{bf1} > + e_a n_a e_b n_b < \vec{B} \cdot \vec{E}^{(A)} > \} \right. \\
& \left. - N_{2a+1 \ 2} < \vec{B} \cdot \vec{F}_{0f2} > \right), \qquad (80)
\end{aligned}
$$

where Eq.(66) is used. It follows from Eq.(79) that in the axisymmetric toroidal systems without external sources and losses which do not guarantee the momentum conservation, the flux-surface-averaged poloidal flow is independent of the radial electric field because $< G_{BS} > = < G_{BS} >_T$.

To obtain more physical results, we use the small mass ratio expansion, i.e., $m_0/m_a \ll 1$ for $a = 1 \sim N$. Neglecting terms with the order of $\sqrt{m_0/m_a}$ for $a = 1 \sim N$, we obtain the following parallel flow and heat flux:

$$
\begin{aligned}
< B u_{\|0} > &= - < G_{BS} > \phi' + < G_{BS} > \sum_{b=0}^{N} \{ L_{1 \ 2b+1} \frac{P_b'}{e_b n_b} - L_{1 \ 2b+2} \frac{T_b'}{e_b} \} \\
& - N_{11} \{ < \vec{B} \cdot \vec{F}_{0f1} > + e_0 n_0 < \vec{B} \cdot \vec{E}^{(A)} > \}, \\
& - \sum_{b=1}^{N} N_{1 \ 2b+1} < \vec{B} \cdot \vec{F}_{bf1} > + N_{12} < \vec{B} \cdot \vec{F}_{0f2} >, \qquad (81) \\
-\frac{2}{5} \frac{< B q_{\|0} >}{P_0} &= < G_{BS} > \sum_{b=0}^{N} \{ L_{2 \ 2b+1} \frac{P_b'}{e_b n_b} - L_{2 \ 2b+2} \frac{T_b'}{e_b} \} \\
& - \sum_{b=0}^{N} N_{2 \ 2b+1} < \vec{B} \cdot \vec{F}_{bf1} > + N_{22} < \vec{B} \cdot \vec{F}_{0f2} >, \quad \text{for electron} \qquad (82) \\
< B u_{\|a} > &= - < G_{BS} > \phi' + < G_{BS} > \sum_{b=1}^{N} \{ L_{2a+1 \ 2b+1} \frac{P_b'}{e_b n_b} - L_{2a+1 \ 2b+2} \frac{T_b'}{e_b} \} \\
& - \sum_{b=1}^{N} N_{2a+1 \ 2b+1} < \vec{B} \cdot \vec{F}_{bf1} >, \qquad (83) \\
-\frac{2}{5} \frac{< B q_{\|a} >}{P_a} &= < G_{BS} > \sum_{b=1}^{N} \{ L_{2a+2 \ 2b+1} \frac{P_b'}{e_b n_b} - L_{2a+2 \ 2b+2} \frac{T_b'}{e_b} \} \\
& - \sum_{b=1}^{N} N_{2a+2 \ 2b+1} < \vec{B} \cdot \vec{F}_{bf1} >, \qquad \text{for ions } (a = 1 \sim N) \qquad (84)
\end{aligned}
$$

where the friction term of the thermal ions with the fast ions $< \vec{B} \cdot \vec{F}_{bf1} >$ is retained, because the term $(v_c^3/\bar{u}_f^3)(n_a e_a^2/\sum_{b=1}^N n_b e_b^2/m_b)(1/m_f + 1/m_a)$ in Eq.(29) may becomes large ($\sim \sqrt{m_a/m_0}$) for high charge number impurities if $v_c^3/\bar{u}_f^3 \sim 1$. From Eqs.(81)-(84) we see that the external inductive electric field mainly contributes to the parallel flow of light electrons. Although fast ions due to the neutral beam injection contributes to both electrons and thermal ions, if $v_c \ll | \vec{u}_f |$ or $v_c^3/\bar{u}_f^3 \ll 1$, then the contribution of fast ions to thermal ions disappears. Thus, the Spitzer current and Ohkawa current in the multispecies plasma are determined only by the electron responce. Moreover, in the non-axisymmetric toroidal systems with the very small geometrical factor, the parallel flow and the parallel heat flux almost disappear. Hence, the ion rotations reduce to the $\vec{E} \times \vec{B}$ and diamagnetic drifts.

V. A SIMPLE ELECTRON-ION PLASMA CASE

In this section, we consider a simple plasma consisting of electrons and ions with effective ionic charge number Z using the small mass ratio expansion. In order to indicate quantities belonging to the electron and the ion, subscripts e and i are used. Neglecting terms of $O(\sqrt{m_e/m_i})$, we introduce the non-dimensional friction coefficients $\bar{l}_{ij}^{ab}$, i.e., $\bar{l}_{11}^{ee} = -Z$, $\bar{l}_{12}^{ee} = -\frac{3}{2}Z$, $\bar{l}_{22}^{ee} = -(\sqrt{2} + \frac{14}{3}Z)$, and $\bar{l}_{22}^{ii} = -\sqrt{2}$. The viscosity coefficients in the plateau collisionality regime are given in Eq.(40), which are independent of the particle constitution. The viscosity coefficients in the $1/\nu$ and Pfirsh-Schlüter collisionality regimes are given by [5,7]

$$\bar{\mu}_{a1} = \sqrt{2} - ln(1+\sqrt{2}) + Z\delta_{ae}, \tag{85}$$

$$\bar{\mu}_{a2} = 2\sqrt{2} - \frac{5}{2}ln(1+\sqrt{2}) + \frac{3}{2}Z\delta_{ae}, \tag{86}$$

$$\bar{\mu}_{a3} = \frac{39}{8}\sqrt{2} - \frac{25}{4}ln(1+\sqrt{2}) + \frac{13}{4}Z\delta_{ae}, \tag{87}$$

and

$$\bar{\mu}_{a1} = \frac{3.020 + 4.250Z\delta_{ae}}{2.225 + 5.321Z\delta_{ae} + 2.400Z^2\delta_{ae}}, \tag{88}$$

$$\bar{\mu}_{a2} = -\frac{4.876 + 9.500Z\delta_{ae}}{2.225 + 5.321Z\delta_{ae} + 2.400Z^2\delta_{ae}}, \tag{89}$$

$$\bar{\mu}_{a3} = \frac{15.394 + 27.000Z\delta_{ae}}{2.225 + 5.321Z\delta_{ae} + 2.400Z^2\delta_{ae}}, \tag{90}$$

respectively, where $\delta_{ae} = 1$ for electron and 0 for ion. It is found out from Eq.(29) and straightforward calculations that for the simple electron and ion plasma using the small mass ratio expansion the friction of thermal ions with fast ions are negligible even if $v_c^3/\bar{u}_f^3 \sim 1$. Thus, we have the following flux-surface-averaged parallel particle flows from Eq.(68):

$$\begin{aligned} < Bu_{\|e} > &= < G_{BS} > \{-\phi' - L_{11}\frac{P_e'}{en_e} - (L_{11}+1)\frac{P_i'}{en_e} + L_{12}\frac{T_e'}{e} - L_{34}(L_{11}+1)\frac{T_i'}{Ze}\} \\ &+ [1 - F_{NC}]\frac{1}{en_e}Z_f en_f < Bu_{\|f} > - \frac{\sigma_{NC}}{en_e} < BE_\|^{(A)} >, \end{aligned} \tag{91}$$

$$< Bu_{\|i} > = < G_{BS} > \{-\phi' - \frac{P_i'}{en_e} - L_{34}\frac{T_i'}{Ze}\}, \tag{92}$$

where

$$\sigma_{NC} = \frac{e^2 n_e \tau_{ei}}{m_e}\Lambda_{NC}^0, \tag{93}$$

$$\Lambda_{NC}^0 = -\frac{Z(l_{22}^{ee} - \mu_{e3})}{D}, \tag{94}$$

$$L_{11} = \frac{\mu_{e1}(l_{22}^{ee} - \mu_{e3}) - \mu_{e2}(l_{12}^{ee} - \mu_{e2})}{D}, \tag{95}$$

$$L_{12} = \frac{\mu_{e2}l_{22}^{ee} - \mu_{e3}l_{12}^{ee}}{D}, \tag{96}$$

$$L_{34} = -\frac{\mu_{i2}l_{22}^{ii}}{\mu_{i1}(l_{22}^{ii} - \mu_{i3}) + \mu_{i2}^2}, \tag{97}$$

$$F_{NC} = 1 - \frac{Z_f}{Z}\{\Lambda_{NC}^0 + \frac{3}{2}\Lambda_{NC}^1\}, \tag{98}$$

$$\Lambda_{NC}^1 = -\frac{Z(l_{12}^{ee} - \mu_{e2})}{D}, \tag{99}$$

$$D = (l_{11}^{ee} - \mu_{e1})(l_{22}^{ee} - \mu_{e3}) - (l_{12}^{ee} - \mu_{e2})^2. \tag{100}$$

For the viscocity coefficients μ_{aj} and the flux-surface-averaged geometrical factor $< G_{BS} >$ Eqs.(20), (33)-(37), and (85)-(87) (in the $1/\nu$ or banana collisionality regime), Eqs.(40)-(44) (in the plateau collisionality regime), Eqs.(45)-(47) and (88)-(90) (in the Pfirsh-Schlüter collisionality regime) are used.

As well as in Sec.IV, the parallel current is obtained:

$$\begin{aligned} BJ_{\|} &= en_e B(u_{\|i} - u_{\|e}) + < BJ_{\|f} > \\ &= BJ_{\|PS} + < BJ_{\|SP} > + < BJ_{\|BS} > + < BJ_{\|OH} >, \end{aligned} \tag{101}$$

where

$$BJ_{\|PS} = -(P_e' + P_i')\{g_2 - \frac{B^2}{< B^2 >} < g_2 >\} \quad : \text{ Pfirsh-Schlüter current,} \tag{102}$$

$$< BJ_{\|SP} > = \sigma_{NC} < BE_{\|}^{(A)} > \quad : \text{ Spitzer current,} \tag{103}$$

$$< BJ_{\|BS} > = < G_{BS} > \{L_{11}(P_e' + P_i') - L_{12} n_e T_e' + L_{11} L_{34} n_i T i'\} \quad : \text{ Bootstrap current,} \tag{104}$$

$$< BJ_{\|OH} > = F_{NC} < BJ_{\|} >_f \quad : \text{ Ohkawa current.} \tag{105}$$

Neoclassical electric conductivity σ_{NC} and the ratio of Ohkawa current to fast ion current F_{NC} in the $1/\nu$ collisionality regime have the same forms as ones in tokamak. [27,28] The reason of which comes from the fact that the interaction between thermal particles and the parallel external inductive electric field $E_{\|}^{(A)}$ or fast ions occures in the velocity space, so that in these quantities the difference of the magnetic field structure appears through the fraction of trapped particles f_t. Contrastively, the bootstrap current has signficantly different magnetic field dependence through the geometrical factor $< G_{BS} >$, which of axisymmetric cases is given by Eq.(23). [7,17-19]

Substitution of Eqs.(91) and (92) into Eqs.(75) and (76) gives the poloidal and toroidal rotations for electrons and ions:

$$< \vec{u_e} \cdot \nabla\theta_B >$$
$$= -\frac{J}{J + \iota I} \frac{1}{en_e < G_{BS} >_T} \{< BJ_{\|SP} > + (1 - \frac{1}{F_{NC}}) < BJ_{\|OH} > + (1 + \frac{1}{L_{11}}) < BJ_{\|BS} >\}$$

$$-\frac{J}{J+\iota I}\frac{<G_{BS}>}{<G_{BS}>_T}\frac{L_{12}}{L_{11}}\frac{1}{e}\frac{dT_e}{d\psi}-\frac{J}{J+\iota I}(1-\frac{<G_{BS}>}{<G_{BS}>_T})(\frac{1}{en_e}\frac{dP_e}{d\psi}-\frac{d\phi}{d\psi}), \tag{106}$$

$$\begin{aligned}
&<\vec{u_e}\cdot\nabla\zeta_B> \\
=\ & -\frac{J}{\iota(J+\iota I)}\frac{1}{en_e<G_{BS}>_T}\{<BJ_{\|SP}>+(1-\frac{1}{F_{NC}})<BJ_{\|OH}>+(1+\frac{1}{L_{11}})<BJ_{\|BS}>\} \\
&-\frac{J}{\iota(J+\iota I)}\frac{<G_{BS}>}{<G_{BS}>_T}\frac{L_{12}}{L_{11}}\frac{1}{e}\frac{dT_e}{d\psi}+\frac{1}{\iota(J+\iota I)}(J\frac{<G_{BS}>}{<G_{BS}>_T}+\iota I)(\frac{1}{en_e}\frac{dP_e}{d\psi}-\frac{d\phi}{d\psi}),
\end{aligned} \tag{107}$$

$$\begin{aligned}
&<\vec{u_i}\cdot\nabla\theta_B> \\
=\ & -\frac{J}{J+\iota I}\frac{<G_{BS}>}{<G_{BS}>_T}L_{34}\frac{1}{eZ}\frac{dT_i}{d\psi}+\frac{J}{J+\iota I}(1-\frac{<G_{BS}>}{<G_{BS}>_T})(\frac{1}{en_e}\frac{dP_i}{d\psi}+\frac{d\phi}{d\psi}),
\end{aligned} \tag{108}$$

$$\begin{aligned}
&<\vec{u_i}\cdot\nabla\zeta_B> \\
=\ & -\frac{J}{\iota(J+\iota I)}\frac{<G_{BS}>}{<G_{BS}>_T}L_{34}\frac{1}{eZ}\frac{dT_i}{d\psi}-\frac{1}{\iota(J+\iota I)}(J\frac{<G_{BS}>}{<G_{BS}>_T}+\iota I)(\frac{1}{en_e}\frac{dP_i}{d\psi}+\frac{d\phi}{d\psi}).
\end{aligned} \tag{109}$$

For axisymmetric case, $<G_{BS}>$ reduces to $<G_{BS}>_T$, therefore, Eqs.(106)-(109) change as follows:

$$\begin{aligned}
&<\vec{u_e}\cdot\nabla\theta_B> \\
=\ & -\frac{J}{J+\iota I}\frac{1}{en_e<G_{BS}>_T}\{<BJ_{\|SP}>+(1-\frac{1}{F_{NC}})<BJ_{\|OH}>+(1+\frac{1}{L_{11}})<BJ_{\|BS}>\} \\
&-\frac{J}{J+\iota I}\frac{L_{12}}{L_{11}}\frac{1}{e}\frac{dT_e}{d\psi},
\end{aligned} \tag{110}$$

$$\begin{aligned}
&<\vec{u_e}\cdot\nabla\zeta_B> \\
=\ & -\frac{J}{\iota(J+\iota I)}\frac{1}{en_e<G_{BS}>_T}\{<BJ_{\|SP}>+(1-\frac{1}{F_{NC}})<BJ_{\|OH}>+(1+\frac{1}{L_{11}})<BJ_{\|BS}>\} \\
&-\frac{J}{\iota(J+\iota I)}\frac{L_{12}}{L_{11}}\frac{1}{e}\frac{dT_e}{d\psi}+\frac{1}{\iota}(\frac{1}{en_e}\frac{dP_e}{d\psi}-\frac{d\phi}{d\psi}),
\end{aligned} \tag{111}$$

$$\begin{aligned}
&<\vec{u_i}\cdot\nabla\theta_B> \\
=\ & -\frac{J}{J+\iota I}L_{34}\frac{1}{eZ}\frac{dT_i}{d\psi},
\end{aligned} \tag{112}$$

$$\begin{aligned}
&<\vec{u_i}\cdot\nabla\zeta_B> \\
=\ & -\frac{J}{\iota(J+\iota I)}L_{34}\frac{1}{eZ}\frac{dT_i}{d\psi}-\frac{1}{\iota}(\frac{1}{en_e}\frac{dP_i}{d\psi}+\frac{d\phi}{d\psi}).
\end{aligned} \tag{113}$$

From the comparison between Eqs.(106)-(109) (in non-axisymmetric toroidal systems) and Eqs.(110)-(113) (in axisymmetric toroidal systems), the effects of axisymmetry on the plasma rotation are found out. If the system has axisymmetry like a tokamak, then the geometrical factor $< G_{BS} >$ reduces to $< G_{BS} >_T$, resulting in the contribution of the radial electric field $d\phi/d\psi$ to the poloidal rotations of both ions and electrons disappears in all collisionality regimes as shown in the multispecies plasma. [20] Contrastively, non-axisymmetric plasmas are driven poloidally by the radial electric field.

The flux-surface-averaged geometrical factor in non-axisymmetric toroidal systems $< G_{BS} >$ is generally smaller than that of axisymmetric toroidal systems $< G_{BS} >_T$. The normalized flux-surface-averaged geometrical factor $< G_{BS} >_N = < G_{BS} > / < G_{BS} >_T$ is advantage to investigate the plasma rotation. In the $1/\nu$ collisionality regime $< G_{BS} >_N$ changes drastically according to the magnetic axis shift and the shaping of the toroidally-averaged magnetic surface, i.e., the ellipticity. [19] $< G_{BS} >_N$ in the plateau collisionality regime also has the dependence of the magnetic field structure, however, its magnitude is considerably small. In the Pfirsh-Schlüter collisionality regime $< G_{BS} >_N$ is very small. The detail discussion on the geometrical factor in each collisionality regime will be given elsewhere. Depending upon the plasma parameters, i.e., the density n_a and the temperature T_a, and the structure of the magnetic field, the collisionality regime of the non-axisymmetric plasma changes variously. Generally speaking, in the periphery and near the magnetic axis the plasma may belong to the plateau collisionality regimes, and in interspace the plasma may belong to the $1/\nu$ or plateau collisionality regime. Moreover, near the outermost surface the collisionality regime may be Pfirsh-Schlüter. Therefore the poloidal and toroidal rotations of ions in the periphery and near the magnetic axis reduce to

$$< \vec{u}_i \cdot \nabla\theta_B > \;\simeq\; \frac{J}{J+\iota I}\left(\frac{1}{en_e}\frac{dP_i}{d\psi} + \frac{d\phi}{d\psi}\right), \tag{114}$$

$$< \vec{u}_i \cdot \nabla\zeta_B > \;\simeq\; -\frac{I}{J+\iota I}\left(\frac{1}{en_e}\frac{dP_i}{d\psi} + \frac{d\phi}{d\psi}\right). \tag{115}$$

Note that above equations are valid in all plasma region as long as $< G_{BS} >_N$ is small whether the plasma belongs to the plateau collisionality regime or the $1/\nu$ collisionality regime. From Eq.(92) we see that the flux-surface-averaged geometrical factor $< G_{BS} >$

prescribes the magnitude of the parallel ion flow for given density n_i, temperature T_i, and electrostatic potential ϕ profile. Hence, in non-axisymmetric systems with small $< G_{BS} >_N$ the parallel ion flow almost disappears (Compare between Eqs.(75)-(76) and Eqs.(114)-(115)). Consequently, in non-axisymmetric systems the poloidal and toroidal rotation of ions is usual diamagnetic and $\vec{E} \times \vec{B}$ drifts. This fact is true for electons if the neutral beam injection and the external inductive electric field do not exist. Using the cylindrical coordinates (r, θ, ζ), we can get the following simple expressions:

$$< \vec{u}_i \cdot \hat{e}_\theta > \sim \frac{B_\zeta}{B^2}(\frac{1}{en_e}\frac{dP_i}{dr} + \frac{d\phi}{dr}), \tag{116}$$

$$< \vec{u}_i \cdot \hat{e}_\zeta > \sim -\frac{B_\theta}{B^2}(\frac{1}{en_e}\frac{dP_i}{dr} + \frac{d\phi}{dr}), \tag{117}$$

where $\hat{e}_\theta$ and $\hat{e}_\zeta$ are the unit vectors in the poloidal and toroidal direction, respectively. From above equations,

$$\frac{|< \vec{u}_i \cdot \hat{e}_\zeta >|}{|< \vec{u}_i \cdot \hat{e}_\theta >|} = \frac{B_\theta}{B_\zeta} \ll 1. \tag{118}$$

In non-axisymmetric systems the poloidal rotation is dominant.

In axisymmetric systems, the ion poloidal rotation of a simple electron-ion plasma is caused only by the temperature gradient. If $dT_i/d\psi$ is negligible, the toroidal rotation is dominant and Eqs.(112) and (113) reduce to

$$< \vec{u}_i \cdot \nabla\theta_B > \sim 0, \tag{119}$$

$$< \vec{u}_i \cdot \nabla\zeta_B > \sim -\frac{1}{\iota}(\frac{1}{en_e}\frac{dP_i}{d\psi} + \frac{d\phi}{d\psi}). \tag{120}$$

In this case the parallel ion flow given by Eq.(92) becomes

$$< Bu_{\parallel i} > \sim -\frac{J}{\iota}(\frac{1}{en_e}\frac{dP_i}{d\psi} + \frac{d\phi}{d\psi}). \tag{121}$$

Substitution of Eq.(121) into Eq.(75) gives $< \vec{u}_i \cdot \nabla\theta_B > \sim 0$. Hence the parallel ion flow is made so that it cancels the poloidal rotation. In the cylindrical coordinates Eq.(120) becomes

$$< \vec{u}_i \cdot \hat{e}_\zeta > \sim -\frac{1}{B_\theta}(\frac{1}{en_e}\frac{dP_i}{dr} + \frac{d\phi}{dr}), \tag{122}$$

thus for the similar profile of n_e, T_i, and ϕ

$$\frac{|<\vec{u}_i \cdot \hat{e}_\zeta>_{helical}|}{|<\vec{u}_i \cdot \hat{e}_\zeta>_{tokamak}|} \sim (\frac{B_\theta}{B_\zeta})^2 \ll 1, \tag{123}$$

$$\frac{|<\vec{u}_i \cdot \hat{e}_\theta>_{helical}|}{|<\vec{u}_i \cdot \hat{e}_\zeta>_{tokamak}|} \sim \frac{B_\theta}{B_\zeta} \ll 1. \tag{124}$$

Therefore, the toroidal flow in axisymmetric systems is larger than the poloidal and toroidal flows in the non-axisymmetric systems.

VI. CONCLUSION AND DISCUSSION

The neoclassical parallel particle flow, parallel heat flux, current and rotation of a multispecies plasma in general toroidal systems are studied in several collisionality regimes, in order to clarify geometrical effects of the magnetic field structure on the neoclassical theory. Fast ions due to the neutral beam injection and an external inductive electric field are included as external momentum sources, where fast ions contribute to the thermal species as friction forces. The flux-surface-averaged parallel particle flow and parallel heat flux are obtained by using the flux-surface-averaged parallel momentum and heat flux balance equations, where the parallel viscosity, friction, and inductive electric field balance. The firction coefficients are classical in the sense that they are independent of the magnetic field. Thus, the geometrical effects of the magnetic field come from the viscosity terms. [5] In axisymmetric toroidal systems the magnetic field strength is a function with respect to the flux label and the poloidal angle. Hence, there is only one type of trapped particles and the geometrical effects of the magnetic field structure in the flux-surface-averaged parallel viscosity appear through the connection length and the fraction of trapped particles except for some magnetic-surface quantities. In contrast to it, in non-axisymmetric toroidal systems the strength of the magnetic field also depends on the toroidal angle, thus there are many types of trapped particles and the geometrical factor reflecting the magnetic field structure $< G_{BS} >$ appears as another type of factor. [7,18,19] This factor prescribes the parallel particle flow due to the gradients of density, temparature, and electrostatic potential. In axisymmetric toroidal systems, this factor reduces to the same expression in all collisionality regimes. In non-axisymmetric toroidal systems, however, its expression and magnitude changes drastically depending on the magnetic field structure and the collisionality regime as well as the viscosity coefficients. The normalized geometrical factor $< G_{BS} >_N \equiv < G_{BS} > / < G_{BS} >_T$ ($< G_{BS} >_T$ is the expression of geometrical factor in axisymmetric toroidal systems) has very small magnitude in Pfirsh-Schlüter regime which may be realized near the outermost magnetic surface. In the plateau regime which may be realized both near the magnetic axis and in the periphery of the plasma, or whole region except for near the outermost surface, $< G_{BS} >_N$ has a considerable small magnitude. $< G_{BS} >_N$ in the $1/\nu$ collisonality regime widely changes according to the magnetic field structure by the magnetic axis shift and the toroidally-averaged magnetic surface shaping, i.e., the ellipticity. [19] Therefore, this geometrical factor mainly makes differences between axisymmetric and non-axisymmetric toroidal systems.

In the situation considered where there are no other external sources and losses without the momentum conservation, for example, the charge exchange loss, the self adjointness and the momentum conservation of the Coulomb collision operator give common properties to the multispecies plasma in spite of the geometrical effects of the magnetic field. The parallel particle flow due to the radial electric field is independent of the particle species and the parallel heat flux is independent of the radial electric field. Consequently, the charge neutrality guarantees the neoclassical current (the bootstrap curent) be independent of the radial electric field. Quantitative changes, however, appear in the neoclassical current mainly through the geometrical factor. [7,17–19] The magnitude of the bootstrap current, which is proportional to the flux-surface-averaged geometrical factor, changes considerably in the non-axisymmetric toroidal systems. Moreover, the poloidal and toroidal rotations change both qualitatively and quantitatively. In axisymmetric toroidal systems, the poloidal rotation in the multispecies plasma is independent of the radial electric field. [20] In non-axisymmetric systems, however, the poloidal rotation has the radial electric field dependence. The smaller the geometrical factor becomes, the larger the effects of the radial electric field does. Hence, for the non-axisymmetric magnetic field configuration where the geometrical factor is small, the ion parallel flow almost disappears and ion rotation is described in terms of the perpendicular flow, i.e., the usual diamagnetic and $\vec{E} \times \vec{B}$ drifts, if the interaction between thermal ions and fast ions is negligible ($v_c \ll | \vec{u}_f |$ where v_c and $\vec{u}_f$ are the critical and fast ions velocity, respectively). This is true for electrons if there are no fast ions and the external inductive electric field. These results contrast with the fact that the ion parallel flow streams to cancel the ion poloidal flow in axisymmetric toroidal systems. This is the manifestation whether the toroidal momentum conservation exists or not. It is found out from the analysis of a simple electron-ion plasma that in non-axisymmetric toroidal systems with sufficiently small geometric factor the poloidal rotation is larger than the toroidal rotation and that for similar plasma parameters the toroidal rotation in the axisymmetric toroidal system is larger than the poloidal rotation in the non-axisymmetric toroidal systems.

The numerical calculation is now in progress taking the high charge number impurities into consideration under the condition $v_c \sim | \vec{u}_f |$ where v_c and $\vec{u}_f$ are the critical and the fast ion velocities, respectively. This treatment is very important because the plasma rotation is measured using impurity emmision line.

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REFERENCES

1 A.A.Galeev and R.Z.Sagdeev, Sov.Phys.-JETP **26**, 233 (1968)

2 M.N.Rosenbluth, R.D.Hazeltine, and F.L.Hinton, Phys.Fluids **15**, 116 (1972)

3 R.D.Hazeltine, F.L.Hinton, and M.N.Rosenbluth, Phys.Fluids **16**, 1645 (1973)

4 F.L.Hinton and R.D.Hazeltine, Rev.Mod.Phys. **48**, 239 (1976)

5 S.P.Hirshman and D.J.Sigmar, Nucl.Fusion **21**, 1079 (1981)

6 F.L.Hinton and S.K.Wong, Phys.Fluid **28**, 3082 (1985)

7 K.C.Shaing and J.D.Callen, Phys.Fluid **26**, 3315 (1983)

8 K.C.Shaing, Phys.Fluid **26**, 3164 (1983)

9 K.C.Shaing, S.P.Hirshman, and J.D.Callen, Phys.Fluid **29**, 521 (1986)

10 K.C.Shaing, Phys.Fluid **29**, 2231 (1986)

11 M.Coronado and H.Wobig, Phys.Fluid **29**, 527 (1986)

12 M.Coronado and H.Wobig, Phys.Fluid **30**, 3171 (1987)

13 K.C.Shaing, Phys.Fluids **31**, 2249 (1988)

14 K.C.Shaing and E.C.Crume Jr., Phys.Rev.Lett **63**, 2369 (1989)

15 K.C.Shaing, E.C.Crume Jr., and W.A.Houlberg, Phys.Fluid B **2**, 1492 (1990)

16 M.Kikuchi, M.Azumi, S.Tsuji, K.Tani, and H.Kubo, Nucl.Fusion **30**, 343 (1990)

17 K.C.Shaing, E.C.Crume Jr., J.S.Tolliver, S.P.Hirshman, and W.I.van Rij, Phys.Fluids B **1**, 148 (1989)

18 K.C.Shaing, B.A.Carreras, N.Dominguez, V.E.Lynch, and J.S.Tolliver, Phys.Fluids B **1**, 1663 (1989)

19 N.Nakajima, M.Okamoto, J.Todoroki, Y.Nakamura, and M.Wakatani, Nucl.Fusion **29**, 605 (1989)

20 R.D.Hazeltine and A.A.Ware, Nucl.Fusion **16**, 538 (1976)

21 S.Hamada, Nucl.Fusion **2**, 23 (1962)

22 G.Bateman, *MHD instabilities* (MIT, Cambridge, 1980), p.125

23 A.H.Boozer, Phys.Fluids **23**, 904 (1980)

24 A.H.Boozer, Phys.Fluids **25**, 520 (1982)

25 G.Kuo-Petravic, A.H.Boozer, J.A.Rome, and R.H.Fowler, J.Comput.Phys. **51**, 261 (1983)

26 J.A.Rome, J.Comput.Phys. **82**, 348 (1989)

27 Y.B.Kim, UWPR 88-2 (1988)

28 N.Nakajima and M.Okamoto, J.Phys.Soc.Jpn **59**, 3595 (1990)

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