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Rotations of Bulk Ions and Impurities in Non-Axisymmetric Toroidal Systems

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Abstract

Neoclassical rotations are investigated theoretically for a plasma consisting of electrons, bulk ions, and impurities in non-axisymmetric toroidal systems. Impurity ions are assumed to lie in the Pfirsch-Schlüter collisionality regime. The results include the case of tokamaks. It is found that the rotations of impurities are genarally different from those of bulk ions. The difference comes from mainly the different diamagnetic flows between bulk ions and impurities. In the $1/\nu$ regime or banana regime the gradient of bulk ion temperature may be another cause for the different rotations between the two species. However, in the region where the configuration is far from the axisymmetry the cause of different rotations is just the different diamagnetic flows. As the poloidal rotations depend on the radial electric field contrary to the axisymmetric system, a relationship, which does not include the radial electric field, can be derived between poloidal and toroidal rotations. By estimating this relationship based on the measured rotations the present neoclassical theory can be validated without knowledge of radial electric field.

Keywords neoclassical theory, non-axisymmetric system, geometric factor, rotation, viscosity, impurity, $1/\nu$ regime, plateau regime, Pfirsch-Schlüter regime

§1. Introduction

In non-axisymmetric toroidal systems, poloidal rotations were measured only near the plasma periphery by using the emission line from impurities[1,2]. Recently, Ida et al. have measured the profiles of poloidal and toroidal rotations in the CHS helical device and JIPPT-IIU tokamak using the charge exchange recombination spectroscopy and compared the results between the two[3]. The radial electric fields derived from the observed poloidal rotation speeds have been compared with neoclassical predictions for the plateau plasma in the CHS heated by the tangential neutral beam injection[4]. Toroidal rotation speeds have been measured in a wide range of densities and magnetic field ripples and it is reported that the observed parallel viscosities agree with the neoclassical predictions within a factor of three[5]. These experiments cast a light on the comprehensive understanding of toroidally confined plasmas. However, it may be worth noting that the experimentally observed rotations are not those of bulk ions but those of impurity ions as the plasma rotations are measured by using an impurity emmission line. Therefore it is important to check the difference of rotations between bulk ions and impurity ions. This was done theoretically in tokamaks by reducing the general neoclassical equations by Hirshman and Sigmar[6] to the case of three component plasma, in which it is shown that the bulk ion rotations can be significantly different from those of impurities in various interesting cases[7]. Thus it is urgent to examine the rotations for bulk ions and impurity ions whether those are the same or different in the non-axisymmetric toroidal devices.

Recently the authors have developed the neoclassical theory for the parallel force balance to investigate the flows, currents, and rotations for a multispecies plasma in general toroidal systems[8]. In this Ref.[8], neoclassical expressions obtained so far in each collisionality regime ($1/\nu$, plateau, and Pfirsch-Schlüter regimes) are unified in terms of the geometric factor. General formula common to all collisionalities are obtained for neoclasical currents (Pfirsch-Schlüter current, bootstrap current, beam driven current, and induction current). Emphasis is placed on the fact that the decisive difference between axisymmetric and non-axisymmetric toroidal systems is attributed to the geometric factor which reflects the breaking of axisymmetry. Explicit expressions for the poloidal and toroidal rotations are given, respectively, in the Boozer coordinate system[9,10)] for a multispecies plasma in any collisionality regime in arbitral magnetic field configuration.

In this paper we examine the rotations for bulk ions and impurities to see the difference between the two species in non-axisymmetric toroidal systems or helical devices. To clear the problem we restrict ourselves to the case of a plasma consisting of electrons, bulk ions, and only one impurity species. Impurities are assumed to be in the Pfirsch-Schlüter collisionality regime as is the usual situation in existing helical devices. The general theory in Ref.[8] is applied to the present investigation. The unified expressions for poloidal and toroidal rotations are reduced to the three species plasma.

In §2 the results for rotations obtained in Ref.[8] are briefly reviewed. General expressions for a multispecies plasma are reduced to the plasma with one impurity species in §3 and a comparison is made between the rotations of bulk ions and impurity ions. In §4 simple forms for rotation are presented for a large aspect ratio and low-β plasma. §5 is devoted to conclusion and discussion. In the Appendix some useful quantities associated with rotations and geometric factors in each collisionality regime are shown.

§2. Poloidal and Toroidal Rotations

In Ref.[8], we have considered the neoclassical particle flows, heat fluxes, currents and rotations based on the moment approach for a multispecies plasma in non-axisymmetric toroidal systems. All the results include the axisymmetric case. Fluctuations, external sources, direct losses of particles (loss cone) and charge transfer losses are not taken into account, but fast ions due to the neutral beam injection and an external inductive electric field are included. We have given the unified expressions, which are common to all collisionalitiy regimes, for the parallel viscosities and parallel heat viscosities in terms of the geometric factor. We show here the results for rotations obtained in Ref.[8] omitting the effects of fast ions and inductive fields. The poloidal and toroidal rotations for species a are given in the Boozer coordinate system[9,10], respectively, by

$$
\begin{aligned}
&< \vec{u}_a \cdot \nabla\theta > \\
&= \frac{J}{J+\iota I}\left(1+\sum_{b=0}^{N} < G_{BS} >_{Nb} L_{2a+1\ 2b+1}\right)\left(\frac{1}{e_a n_a}\frac{dP_a}{d\psi}+\frac{d\phi}{d\psi}\right) \\
&+ \frac{J}{J+\iota I}\sum_{b=0}^{N} < G_{BS} >_{Nb} \left\{L_{2a+1\ 2b+1}\left(\frac{1}{e_b n_b}\frac{dP_b}{d\psi}-\frac{1}{e_a n_a}\frac{dP_a}{d\psi}\right) - L_{2a+1\ 2b+2}\frac{1}{e_b}\frac{dT_b}{d\psi}\right\}, \quad (1)
\end{aligned}
$$

$$
\begin{aligned}
&< \vec{u}_a \cdot \nabla\zeta > \\
&= -\frac{1}{\iota(J+\iota I)}\left(-J\sum_{b=0}^{N} < G_{BS} >_{Nb} L_{2a+1\ 2b+1}+\iota I\right)\left(\frac{1}{e_a n_a}\frac{dP_a}{d\psi}+\frac{d\phi}{d\psi}\right) \\
&+ \frac{J}{\iota(J+\iota I)}\sum_{b=0}^{N} < G_{BS} >_{Nb} \left\{L_{2a+1\ 2b+1}\left(\frac{1}{e_b n_b}\frac{dP_b}{d\psi}-\frac{1}{e_a n_a}\frac{dP_a}{d\psi}\right) - L_{2a+1\ 2b+2}\frac{1}{e_b}\frac{dT_b}{d\psi}\right\} \quad (2)
\end{aligned}
$$

Here, ψ, θ, and ζ are the Boozer coordinates (θ and ζ are poloidal and toroidal angles, respectively) and $< G_{BS} >_{Nb}$ is the normalized geometric factor for species b ($b = 0$ means electrons and ion species are denoted by $b = 1, \cdots, N$). For axisymmetric systems or tokamaks $< G_{BS} >_{Nb}$ is just unity regardless of the magnetic field structure and collisionality, whereas in non-axisymmetric cases $< G_{BS} >_{Nb}$ is less than unity and changes according to the collisionality. In the $1/\nu$ regime the geometric factor changes drastically depending on the field structure[11], in the plateau regime the field dependence is not so strong and $< G_{BS} >_{Nb}$ is small, and in the Pfirsch-Schlüter regime it is very small. However, it should be emphasized that the normalized geometric factor $< G_{BS} >_{Nb}$ approaches unity as the configuration becomes close to the symmetry even if the species b are in the plateau or Pfirsch-Schlüter regime (see Appendix). Other notations in Eqs.(1) and (2) should be referred to Ref.[8].

It is easily seen that the poloidal rotation in the axisymmetric system is independent of the radial electric field [6,12], because $< G_{BS} >_{Nb} = 1$ and $\sum_{b=0}^{N} L_{2a+1\ 2b+1} = -1$ (see Appendix in Ref.[8]), resulting in $1 + \sum_{b=0}^{N} < G_{BS} >_{Nb} L_{2a+1\ 2b+1} = 0$ in Eq.(1). However, in non-axisymmetric systems, poloidal rotations as well as toroidal rotations depend on the radial electric field (see the terms with $d\phi/dr$ in Eqs.(1) and (2)). Accordingly there is a possibility that the impurity rotations both in the poloidal and toroidal directions would be the same as those of bulk ions if $\vec{E} \times \vec{B}$ drifts dominate. Such a situation takes place when $< G_{BS} >_{Nb} \ll 1$ and $(1/e_a n_a) dP_a/d\psi \ll d\phi/dr$. If this is not the case, the difference between diamagnetic flows for different species can be a cause of rotations. It is well known that, in the classical theory, this difference, $(1/e_a n_a)dP_a/d\psi - (1/e_b n_b)dP_b/d\psi$, tends to vanish on the time scale of ion-ion collisions resulting in impurity accumulation at the plasma center[13]. The terms directly proportional to the ion temperature gradient are originated from the parallel heat viscosities.

§3. Application to a Plasma Containing One Impurity Species

We apply the general expressions for poloidal and toroidal rotations to a plasma consisting of electrons, bulk ions and one species impurity ions. The quantities belonging to electrons, bulk ions, and impurity ions are indicated by the subscripts e, i and I, respectively. Here, we consider heavy impurity ions with a high atomic number to neglect terms of $O(\sqrt{m_e/m_a})$ $(a = i, I)$ and $O(\sqrt{m_i/m_I})$ (m_e, m_i, and m_I are masses of electrons, bulk ions, and impurities, respectively). We define $\alpha \equiv n_I e_I^2/(n_i e_i^2)$ (n_I and n_i are densities of bulk ions and impurities, respectively, and e_I and e_i are electric charges of each species) and assume that α is order of unity or less than unity. Moreover, we consider a situation in which impurity ions lie in the Pfirsh-Schlüter collisionality regime.

In this case, the matrix elements consisting of the friction coefficients $l_{j,k}^{ab}$ $(a, b = e, i, I$ and $j, k = 1, 2)$ and the viscosity coefficients μ_{aj} ($a = e, i, I$ and $j = 1 \sim 3$) are reduced to the following forms:

$$L_{11} \simeq \{(\frac{3}{2}Z_{eff} + \bar{\mu}_{e2})\bar{\mu}_{e2} - (\sqrt{2} + \frac{13}{4}Z_{eff} + \bar{\mu}_{e3})\bar{\mu}_{e1}\}/D_e < 0, \tag{3}$$

$$L_{12} \simeq -\{(\sqrt{2} + \frac{13}{4}Z_{eff})\bar{\mu}_{e2} - \frac{3}{2}Z_{eff}\bar{\mu}_{e3}\}/D_e, \tag{4}$$

$$D_e = (Z_{eff} + \bar{\mu}_{e1})(\sqrt{2} + \frac{13}{4} Z_{eff} + \bar{\mu}_{e3}) - (\frac{3}{2} Z_{eff} + \bar{\mu}_{e2})^2 > 0, \tag{5}$$

$$L_{34} \simeq -(\sqrt{2} + \alpha)\bar{\mu}_{i2}/D_i, \tag{6}$$

$$L_{54} \simeq -\{(\sqrt{2} + \alpha)\bar{\mu}_{i2} + \frac{3}{2}(\bar{\mu}_{i1}\bar{\mu}_{i3} - (\bar{\mu}_{i2})^2)\}/D_i, \tag{7}$$

$$D_i = (\sqrt{2} + \alpha)\bar{\mu}_{i1} + \bar{\mu}_{i1}\bar{\mu}_{i3} - (\bar{\mu}_{i2})^2 > 0, \tag{8}$$

$$L_{13} \simeq -1 - L_{11}, \tag{9}$$

$$L_{14} \simeq \frac{1 + L_{11}}{1 + \alpha}(L_{34} + \alpha L_{54}), \tag{10}$$

$$L_{33} \simeq -1, \tag{11}$$

$$L_{35} \simeq -1, \tag{12}$$

and $L_{j1} = L_{j2} \simeq 0$ for $(j = 3, 5)$, and $L_{j5} = L_{j6} \simeq 0$ for $(j = 1, 3, 5)$. It is worth noting that the coefficients L_{11}, L_{12} and D_e are determined almost by electrons only, whereas the coefficients L_{34}, L_{54} and D_i consist of bulk ion quantities. The friction coefficients l_{jk}^{ab} are expressed explicitly and the dimensionless viscosity coefficients are given by

for $1/\nu$ or banana collisionality regime

$$\bar{\mu}_{a1} = \frac{f_t}{f_c}\{\sqrt{2} - \ln(1 + \sqrt{2}) + c_a\}, \tag{13}$$

$$\bar{\mu}_{a2} = \frac{f_t}{f_c}\{2\sqrt{2} - \frac{5}{2}\ln(1 + \sqrt{2}) + \frac{3}{2}c_a\}, \tag{14}$$

$$\bar{\mu}_{a3} = \frac{f_t}{f_c}\{\frac{39}{8}\sqrt{2} - \frac{25}{4}\ln(1 + \sqrt{2}) + \frac{13}{4}c_a\}, \tag{15}$$

for plateau collisionality regime

$$\bar{\mu}_{a1} = \frac{\lambda_a}{\lambda_{PL}} 2, \tag{16}$$

$$\bar{\mu}_{a2} = -\frac{\lambda_a}{\lambda_{PL}}, \tag{17}$$

$$\bar{\mu}_{a3} = \frac{\lambda_a}{\lambda_{PL}}\frac{13}{2}, \tag{18}$$

for Pfirsh-Schlüter collisionality regime

$$\bar{\mu}_{a1} = (\frac{\lambda_a}{\lambda_{PS}})^2 \frac{205\sqrt{2}/32 + c_a 51/4}{267/40 + c_a 903/80 + c_a^2 36/5}, \quad (19)$$

$$\bar{\mu}_{a2} = -(\frac{\lambda_a}{\lambda_{PS}})^2 \frac{331\sqrt{2}/32 + c_a 57/2}{267/40 + c_a 903/80 + c_a^2 36/5}, \quad (20)$$

$$\bar{\mu}_{a3} = (\frac{\lambda_a}{\lambda_{PS}})^2 \frac{1045\sqrt{2}/32 + c_a 81}{267/40 + c_a 903/80 + c_a^2 36/5}, \quad (21)$$

where $c_a = Z_{eff}$ if $a = e$ and $c_a = \alpha$ if $a = i$. In the above expressions, f_t is the fraction of trapped particles, $f_c = 1 - f_t$, λ_a is the mean free path of particle species a, λ_{PL} and λ_{PS} are the characteristic lengthes in the plateau and Pfirsh-Schlüter regimes, respectively. These quantities are defined together with the geometric factor in each collisionality regime in Appendix. Note that the dimensionless viscosity coefficients $\bar{\mu}_{e2}$ and $\bar{\mu}_{i2}$ change their signs from negative values to positive values as the collisionality increases from the $1/\nu$ or banana regime to the plateau regime. Then, L_{12}, L_{34} and L_{54} change their signs from negative values to positive values as the collisionality regime shifts from the $1/\nu$ or banana regime to the plateau regime.

We have assumed that impurity ions are in the Pfirsh-Schlüter collisionality regime so that $< G_{BS} >_{NI} \ll 1$. Collisionalities of Electrons and bulk ions are not specified and they are in any collisionality regime of the $1/\nu$, plateau, or Pfirsch-Schlüter regime. Equations (1) and (2) can be simplified under the assumption, $< G_{BS} >_{NI} \ll 1$. Even if the configuration is close to the axisymmetric system ($< G_{BS} >_{NI} \sim 1$) viscosity coefficients are so small that the impurity parallel viscosities are negligible as far as impurities are in the Pfirsch-Schlüter regime. The poloidal and toroidal rotations for bulk ions and impurity ions can be derived from Eqs.(1) and (2) and are given by, respectively,

$$\begin{aligned} < \vec{u}_i \cdot \nabla\theta > &= \frac{J}{J + \iota I}\{(1- < G_{BS} >_{Ni})(\frac{1}{e_i n_i}\frac{dP_i}{d\psi} + \frac{d\phi}{d\psi}) \\ &- < G_{BS} >_{Ni} L_{34}\frac{1}{e_i}\frac{dT_i}{d\psi}\}, \end{aligned} \quad (22)$$

$$\begin{aligned} < \vec{u}_I \cdot \nabla\theta > &= \frac{J}{J + \iota I}\{(1- < G_{BS} >_{Ni})(\frac{1}{e_I n_I}\frac{dP_I}{d\psi} + \frac{d\phi}{d\psi}) \\ &+ < G_{BS} >_{Ni} [-L_{54}\frac{1}{e_i}\frac{dT_i}{d\psi} + \frac{1}{e_I n_I}\frac{dP_I}{d\psi} - \frac{1}{e_i n_i}\frac{dP_i}{d\psi}]\}, \end{aligned} \quad (23)$$

$$
\begin{aligned}
<\vec{u}_i\cdot\nabla\zeta> &= -\frac{1}{\iota}\left(\frac{1}{e_i n_i}\frac{dP_i}{d\psi}+\frac{d\phi}{d\psi}\right)+\frac{J}{\iota(J+\iota I)}\{(1-<G_{BS}>_{N_i})\left(\frac{1}{e_i n_i}\frac{dP_i}{d\psi}+\frac{d\phi}{d\psi}\right)\\
&- <G_{BS}>_{N_i} L_{34}\frac{1}{e_i}\frac{dT_i}{d\psi}\}, \qquad (24)\\
<\vec{u}_I\cdot\nabla\zeta> &= -\frac{1}{\iota}\left(\frac{1}{e_I n_I}\frac{dP_I}{d\psi}+\frac{d\phi}{d\psi}\right)+\frac{J}{\iota(J+\iota I)}\{(1-<G_{BS}>_{N_i})\left(\frac{1}{e_I n_I}\frac{dP_I}{d\psi}+\frac{d\phi}{d\psi}\right)\\
&+ <G_{BS}>_{N_i}[-L_{54}\frac{1}{e_i}\frac{dT_i}{d\psi}+\frac{1}{e_I n_I}\frac{dP_I}{d\psi}-\frac{1}{e_i n_i}\frac{dP_i}{d\psi}]\}. \qquad (25)
\end{aligned}
$$

Equations (22) and (23) are poloidal rotations for bulk ions and impurities, respectively. Toroidal rotations are given by Eqs.(24) and (25) for bulk ions and impurity ions, respectively.

It is noted that if $\vec{E}\times\vec{B}$ drifts are dominant the rotation speeds are the same for bulk ions and impurity ions both in the poloidal and toroidal directions. However when diamagnetic flows are comparable with the $\vec{E}\times\vec{B}$ drifts there are diffenrences of rotations between the two species ;

$$
\begin{aligned}
&<\vec{u}_i\cdot\nabla\theta> - <\vec{u}_I\cdot\nabla\theta>\\
&= \frac{J}{J+\iota I}\{\frac{1}{e_i n_i}\frac{dP_i}{d\psi}-\frac{1}{e_I n_I}\frac{dP_I}{d\psi} - <G_{BS}>_{N_i}(L_{34}-L_{54})\frac{1}{e_i}\frac{dT_i}{d\psi}\}, \qquad (26)\\
&<\vec{u}_i\cdot\nabla\zeta> - <\vec{u}_I\cdot\nabla\zeta>\\
&= -\frac{I}{J+\iota I}\{\frac{1}{e_i n_i}\frac{dP_i}{d\psi}-\frac{1}{e_I n_I}\frac{dP_I}{d\psi} + <G_{BS}>_{N_i}\frac{J}{\iota I}(L_{34}-L_{54})\frac{1}{e_i}\frac{dT_i}{d\psi}\}. \qquad (27)
\end{aligned}
$$

It is seen that the two effects contribute to the differences of rotations ; one is the different diamagnetic flows for the two species and the another is the thermodynamic force due to the bulk ion temperature gradient originated from the parallel heat viscosities. The impurity temperature T_I relaxes to become the bulk ion temperature T_i ($T_I \simeq T_i$) on the time scale of ion-ion collisions, but impurity density profile is often different from that of bulk ions to produce the different diamagnetic flows. If the gradient of bulk ion temperature is steep, the second effect becomes significant as far as bulk ions lie in the $1/\nu$ or banana regime and $< G_{BS} >_{N_i}$ is not so small, because the difference between two coefficients, $L_{34}-L_{54}(=-\frac{3}{2}K_2$, where K_2 is the notation of Ref.[7]), remains finite in the $1/\nu$ or banana regime but very small in the plateau and Pfirsch-Schlüter regimes.

It is interesting that the poloidal rotations depend on the radial electric field. We can obtain a relationship between poloidal and toroidal rotations, which does not include the radial electric field. For impurity ions this relationship is obtained by eliminating $d\phi/d\psi$ from Eqs.(23) and (25) ;

$$\begin{aligned}
& \iota\frac{J}{J+\iota I}(1-<B_{BS}>_{Ni})<\vec{u}_I\cdot\nabla\zeta> \\
= & \{\frac{J}{J+\iota I}(1-<G_{BS}>_{Ni})-1\}<\vec{u}_I\cdot\nabla\theta> \\
- & \frac{J}{J+\iota I}<G_{BS}>_{Ni}(L_{54}\frac{1}{e_i}\frac{dT_i}{d\psi}+\frac{1}{e_in_i}\frac{dP_i}{d\psi}-\frac{1}{e_In_I}\frac{dP_I}{d\psi}).
\end{aligned} \tag{28}$$

If this relationship holds between poloidal and toroidal rotations observed in the experiment, we can have a confirmation for the validity of the neoclassical theory in this paper, which gives the relations of rotations with radial electric fields, without knowledge of the radial electric field.

§4. Simple Case for a Large Aspect Ratio and Low-β Plasma

In order to obtain simpler forms and compare the results with those in tokamaks, we consider a large aspect ratio, low-β toroidal plasma using a cylindrical coordinates (r,θ,ζ). The averaged minor radius $\bar{r}$ and the averaged poloidal field $\bar{B}_\theta$ are defined as $\bar{r}^2=V/(2\pi^2R)$ and $\bar{B}_\theta=\bar{r}B_\zeta\iota/R$, respectively, where R is the major radius and $B_\zeta\simeq$ const is the toroidal field. The toroidal flux is $\Phi_T=2\pi\psi\simeq\pi\bar{r}^2B_\zeta$. From the definition of J, $J\simeq RB_\zeta$. Using the relation $<B^2>=2\pi(J+\iota I)\mathrm{d}\Phi_T/\mathrm{d}V$, we obtain

$$\begin{aligned}
\bar{u}_{\theta i} = & \frac{B_\zeta^2}{<B^2>}\{(1-<G_{BS}>_{Ni})(\frac{1}{2}\rho_{iT}v_{Ti}L_{Pi}^{-1}-\frac{E_r}{B_\zeta}) \\
- & <G_{BS}>_{Ni}\frac{1}{2}\rho_{iT}v_{Ti}L_{34}L_{Ti}^{-1}\},
\end{aligned} \tag{29}$$

$$\begin{aligned}
\bar{u}_{\theta I} = & \frac{B_\zeta^2}{<B^2>}\{(1-<G_{BS}>_{Ni})(\frac{1}{2}\rho_{iT}v_{Ti}\frac{e_iT_I}{e_IT_i}L_{PI}^{-1}-\frac{E_r}{B_\zeta}) \\
+ & <G_{BS}>_{Ni}\frac{1}{2}\rho_{iT}v_{Ti}[-L_{54}L_{Ti}^{-1}+\frac{e_iT_I}{e_IT_i}L_{PI}^{-1}-L_{Pi}^{-1}]\},
\end{aligned} \tag{30}$$

$$
\begin{aligned}
\bar{u}_{\zeta i} &= \{1-\frac{B_\zeta^2}{<B^2>}(1-<G_{BS}>_{Ni})\}(\frac{E_r}{\bar{B}_\theta}-\frac{1}{2}\rho_{iP}v_{Ti}L_{Pi}^{-1}) \\
&- \frac{B_\zeta^2}{<B^2>}<G_{BS}>_{Ni}\frac{1}{2}\rho_{iP}v_{Ti}L_{34}L_{Ti}^{-1}, \qquad (31)\\
\bar{u}_{\zeta I} &= \{1-\frac{B_\zeta^2}{<B^2>}(1-<G_{BS}>_{Ni})\}(\frac{E_r}{\bar{B}_\theta}-\frac{1}{2}\rho_{iP}v_{Ti}\frac{e_iT_I}{e_IT_i}L_{PI}^{-1}) \\
&+ \frac{B_\zeta^2}{<B^2>}<G_{BS}>_{Ni}\frac{1}{2}\rho_{iP}v_{Ti}[-L_{54}L_{Ti}^{-1}+\frac{e_iT_I}{e_IT_i}L_{PI}^{-1}-L_{Pi}^{-1}], \qquad (32)
\end{aligned}
$$

where $L_{Pi}^{-1}=\frac{d}{d\bar{r}}\ln P_i$, $L_{PI}^{-1}=\frac{d}{d\bar{r}}\ln P_I$, $L_{Ti}^{-1}=\frac{d}{d\bar{r}}\ln T_i$, $E_r=-\frac{d}{d\bar{r}}\phi$, $v_{Ti}=\sqrt{2T_i/m_i}$, $\omega_{iTc}=e_iB_\zeta/m_i$, $\rho_{iT}=v_{Ti}/\omega_{iTc}$, $\omega_{iPc}=e_i\bar{B}_\theta/m_i$, and $\rho_{iP}=v_{Ti}/\omega_{iPc}$.

As in axisymmetric toroidal systems, $<G_{BS}>_{Na}=1$ in all collisionality regimes for all particle species, Eqs.(29) to (32) are reduced to:

$$
\begin{aligned}
\bar{u}_{\theta i} &= -\frac{B_\zeta^2}{<B^2>}\frac{1}{2}\rho_{iT}v_{Ti}L_{34}L_{Ti}^{-1}, \qquad (33)\\
\bar{u}_{\theta I} &= \frac{B_\zeta^2}{<B^2>}\frac{1}{2}\rho_{iT}v_{Ti}[-L_{54}L_{Ti}^{-1}+\frac{e_iT_I}{e_IT_i}L_{PI}^{-1}-L_{Pi}^{-1}], \qquad (34)\\
\bar{u}_{\zeta i} &= \frac{E_r}{\bar{B}_\theta}-\frac{1}{2}\rho_{iP}v_{Ti}L_{Pi}^{-1}-\frac{B_\zeta^2}{<B^2>}\frac{1}{2}\rho_{iP}v_{Ti}L_{34}L_{Ti}^{-1}, \qquad (35)\\
\bar{u}_{\zeta I} &= \frac{E_r}{\bar{B}_\theta}-\frac{1}{2}\rho_{iP}v_{Ti}\frac{e_iT_I}{e_IT_i}L_{PI}^{-1} \\
&+ \frac{B_\zeta^2}{<B^2>}\frac{1}{2}\rho_{iP}v_{Ti}[-L_{54}L_{Ti}^{-1}+\frac{e_iT_I}{e_IT_i}L_{PI}^{-1}-L_{Pi}^{-1}]. \qquad (36)
\end{aligned}
$$

Equations (33) to (36) reproduce the results in Ref.[7].

As mentioned previously, in non-axisymmetric toroidal systems the normalized geometric factor for bulk ions $<G_{BS}>_{Ni}$ becomes small in the plateau regime and very small in the Pfirsh-Schlüter regime. In the limit of vanishing $<G_{BS}>_{Ni}$, the poloidal rotation velocities of bulk ions and impurities are expressed as follows:

$$
\begin{aligned}
\bar{u}_{\theta i} &= \frac{B_\zeta^2}{<B^2>}[\frac{1}{2}\rho_{iT}v_{Ti}L_{Pi}^{-1}-\frac{E_r}{B_\zeta}], \qquad (37)\\
\bar{u}_{\theta I} &= \frac{B_\zeta^2}{<B^2>}[\frac{1}{2}\rho_{iT}v_{Ti}\frac{e_iT_I}{e_IT_i}L_{PI}^{-1}-\frac{E_r}{B_\zeta}]. \qquad (38)
\end{aligned}
$$

Thus, the ploidal rotations of bulk ions and impurity ions are the usaul poloidal components of the diamagnetic and $\vec{E}\times\vec{B}$ drifts and the difference of their poloidal rotations is only the difference of the poloidal component of the diamagnetic drift velocity.

Using $< B^2 > \simeq B_\zeta^2 + \bar{B}_\theta^2$, two limiting case of the toroidal rotations of bulk ions and impurity ions are considered. When $\bar{B}_\theta^2 / < B^2 > \gg < G_{BS} >_{Ni}$, they become

$$\bar{u}_{\zeta i} = \frac{\bar{B}_\theta^2}{< B^2 >} (\frac{E_r}{\bar{B}_\theta} - \frac{1}{2} \rho_{iP} v_{Ti} L_{Pi}^{-1}), \tag{39}$$

$$\bar{u}_{\zeta I} = \frac{\bar{B}_\theta^2}{< B^2 >} (\frac{E_r}{\bar{B}_\theta} - \frac{1}{2} \rho_{iP} v_{Ti} \frac{e_i T_I}{e_I T_i} L_{PI}^{-1}). \tag{40}$$

In this limit, the toroidal rotations are reduced to the toroidal component of the usual diamagnetic and $\vec{E} \times \vec{B}$ drifts. In another limit, i.e., $\bar{B}_\theta^2 / < B^2 > \ll < G_{BS} >_{Ni}$, the toroidal rotations of bulk ions and impurity ions are given by

$$\bar{u}_{\zeta i} = < G_{BS} >_{Ni} \{\frac{E_r}{\bar{B}_\theta} - \frac{1}{2} \rho_i P v_{Ti} [L_{Pi}^{-1} - L_{34} L_{Ti}^{-1}]\}, \tag{41}$$

$$\bar{u}_{\zeta I} = < G_{BS} >_{Ni} \{\frac{E_r}{\bar{B}_\theta} - \frac{1}{2} \rho_i P v_{Ti} [L_{PI}^{-1} - L_{54} L_{Ti}^{-1}]\}. \tag{42}$$

Therefore, for the same parameters of E_r, B_ζ, $\bar{B}_\theta$, L_{Pi}, etc., the toroidal rotation speeds of bulk ions and impurity ions in non-axisymmetric systems are smaller than ones in axisymmetric systems by the amount of $< G_{BS} >_{Ni}$.

§5. Conclusion and Discussion

Neoclassical rotations for bulk ions and impurity ions have been obtained under the assumption that the impurities are in the Pfirsch-Schlüter regime in non-axisymmetric toroidal systems. The results include the ones in tokamaks. The main difference between the rotations of bulk ions and impurity ions is originated from the difference of diamagnetic flows between the two species. As $T_I \simeq T_i$ the difference between diamagnetic flows of bulk ions and impurities comes from the different profiles between the two species. If bulk ions lie in the $1/\nu$ regime the steep gradient of bulk ion temperature may cause a difference of rotations between the two species as indicated in Ref.[7].

Rotations of bulk ions and impurities are given by Eqs.(22) to (25) in general toroidal devices. Simple forms of Eqs.(37) to (42) are applicable if and only if $< B_{BS} >_{Ni}$

is small. It should be noted that the geometric factor $< G_{BS} >_{N_t}$ is a surface quantity and changes with radius. Even in non-axisymmetric systems there is a situation in which the configuration becomes almost axisymmetric near the magnetic axis. Such a situation can be seen in the CHS configuration, where the field ripple strength on the axis decreases rapidly as the axis is shifted inward by the control of applied dipole fields[5)]. The configuration with no ripple on the magnetic axis is almost axisymmetric and $< G_{BS} >_{N_t}$ approaches unity towards the axis regardless of collisionality.

What we have obtained in the present paper are the ralations connecting the poloidal or toroidal rotation with the radial electric field. The present theory does not give any information on what determines the radial electric fields or rotations. However, it is interesting that the neoclassical relations between rotations and radial electric fields presented in this paper can be tested without knowledge of the radial electric field by comparing the radial electric field estimated by the observed poloidal rotations with that estimated by the observed toroidal rotations via Eqs.(23) and (25). In other words we can have a comfirmation for the validity of the present neoclassical theoty if the relationship by Eq.(28) holds for observed poloidal and toroidal rotations. This is a beneficial point for non-axisymmetric systems or helical devices. Such the verification is impossible in the case of tokamaks, because the poliodal rotation is independent of the radial electric field in axisymmetric systems.

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Appendix. Quantities appearing in the viscosities and the geometric factors

We show here some useful quantities associated with the viscosities and geometric factors obtained in Ref.[8]. The fraction of the trapped particles is given by

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

where

$$g_1 = \sqrt{1 - \lambda\frac{B}{B_{max}}}. \tag{A.2}$$

The mean free path of particle species a is

$$\lambda_a = \tau_{aa} v_{Ta}, \tag{A.3}$$

where v_{Ta} is the thermal velocity and τ_{aa} is the Braginskii collision time: $\tau_{aa}^{-1} = 4/(3\sqrt{\pi})\cdot 4\pi n_a e_a^4 \ln\Lambda/(m_a^2 v_{Ta}^3)$. The characteristic length of the magnetic field in the plateau regime λ_{PL} is defined as follows:

$$\frac{1}{\lambda_{PL}} \equiv \frac{\Phi_T'(\iota\mu_p + \mu_t)}{< B^2 >}, \tag{A.4}$$

where

$$\mu_p = \frac{\sqrt{\pi}}{2}\langle\frac{\hat{n}\cdot\nabla B}{\Phi_T'}\sum_{m\neq 0, n\neq 0}\frac{\partial B_{mn}(\theta_H,\zeta_H)/\partial\theta_H}{2\pi \mid \iota m + n \mid}\rangle, \tag{A.5}$$

$$\mu_t = \frac{\sqrt{\pi}}{2}\langle\frac{\hat{n}\cdot\nabla B}{\Phi_T'}\sum_{m\neq 0, n\neq 0}\frac{\partial B_{mn}(\theta_H,\zeta_H)/\partial\zeta_H}{2\pi \mid \iota m + n \mid}\rangle. \tag{A.6}$$

$\hat{n} = \vec{B}/B$ and $B = \sum_{m,n} B_{mn}(\theta_H,\zeta_H) = \sum_{m,n} B_{mn}^H \cos(2\pi m\theta_H + 2\pi n\zeta_H + \varphi_{mn})$. In these expressions, spectra of B in terms of the Hamada coordinates[14)] (V,θ_H,ζ_H) are used. These spectra can be calculated from ones in the Boozer coordinates[9,10)] (ψ,θ,ζ) by the transformation.[11)]

The characteristic length of the magnetic field in the Pfirsh-Schlüter regime λ_{PS} is defined as follows:

$$\frac{1}{\lambda_{PS}^2} \equiv \frac{3}{2}\frac{< (\hat{n}\cdot\nabla B)^2 >}{< B^2 >}. \tag{A.7}$$

The geometric factor $< G_{BS} >_a$ in the $1/\nu$ or banana collisionality regime is given by the following equations:[15)]

$$< G_{BS} >_a = \frac{1}{f_t}\left\{\langle g_2\rangle - \frac{3}{4}\frac{\langle B^2\rangle}{B_{max}^2}\int_0^1 \frac{\langle g_4\rangle}{\langle g_1\rangle}\lambda\, d\lambda\right\}, \tag{A.8}$$

$$\vec{B}\cdot\nabla\left(\frac{g_2}{B^2}\right) = \vec{B}\times\nabla V\cdot\nabla\left(\frac{1}{B^2}\right), \quad g_2(B_{max})=0, \tag{A.9}$$
$$\vec{B}\cdot\nabla\left(\frac{g_4}{g_1}\right) = \vec{B}\times\nabla V\cdot\nabla\left(\frac{1}{g_1}\right), \quad g_4(B_{max})=0. \tag{A.10}$$

For axisymmetric toroidal systems, $< G_{BS} >_a$ becomes $< G_{BS} >_T$ where $< G_{BS} >_T$ $= J/\iota \cdot dV/d\psi$. In other collisionality regimes (plateau and Pfirsch-Schüter regimes) the geometric factors are constructed so as to result in $< G_{BS} >_T$ in the limit of axisymmetry.

In the plateau collisionality regime,[16] the geometric factor is defined as

$$\begin{aligned} < G_{BS} >_a &= \frac{< \vec{B}\times\nabla V\cdot\nabla\theta_H > \mu_p + < \vec{B}\times\nabla V\cdot\nabla\zeta_H > \mu_t}{\Phi'_T(\iota\mu_p+\mu_t)} \\ &= \frac{dV}{d\psi}\frac{J\mu_p - I\mu_t}{\iota\mu_p+\mu_t}. \end{aligned} \tag{A.11}$$

As well as $1/\nu$ collisionality regime, $< G_{BS} >_a$ becomes $< G_{BS} >_T$ in axisymmetric torus.

In the Pfirsh-Schlüter collisionality regime,[15] we define the the geometric factor as follows:

$$< G_{BS} >_a = 2\pi\frac{< \frac{1}{B}(J\frac{\partial B}{\partial\theta_H} - I\frac{\partial B}{\partial\zeta_H})(\hat{n}\cdot\nabla B) >}{< (\hat{n}\cdot\nabla B)^2 >} \tag{A.12}$$

As well as in the $1/\nu$ and plateau regimes, $< G_{BS} >_a$ of the Pfirsch-Schlüter regime becomes $< G_{BS} >_T$ for axisymmetric toroidal systems.

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