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Reduced drift kinetic equation for neoclassical transport of helical plasmas in ultra-low collisionality regime

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Abstract: Based on the original five-dimensional drift kinetic equation, a three-dimensional reduced drift kinetic equation has been obtained to describe the neoclassical transport of helical plasmas in ultra-low collisionality regime where the collision frequency is much lower than the poloidal bounce frequency of the super banana particles. The reduced drift kinetic equation describes the evolution of the distribution function in terms of three independent constants of motion. In contrast to the conventional neoclassical theory, this reduced drift kinetic equation includes the effects of finite super banana width and the effects of finite aspect ratio. Therefore, this reduced drift kinetic equation is suitable for studying the collisional transport of fast ions in helical magnetic confinement systems.

Keywords: reduced drift kinetic equation, neoclassical transport, helical plasmas, effect of finite banana width
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1. INTRODUCTION

The plasma parameters achieved in the experiments of helical confinement systems [1, 2] have demonstrated that the helical confinement systems, in parallel to tokamaks, is an important approach to realize the magnetic confinement fusion reactors. In the reactor relevant helical plasmas, an important component will be the fast ions produced by either neutral beam injection or thermonuclear fusion reactions. To understand the influence of these fast ions on the plasma behavior is one of the major issues in fusion plasma physics research. Recently, confinement of fast ions has attracted more attentions [3, 4, 5] in helical confinement fusion research community. These works have discussed the effect of orbit loss on the confinement of fast ions. In this paper, we will discuss the neoclassical transport behavior of the confined particles; the direct loss which has been discussed [3, 4, 5] is beyond the scope of this paper.

For the fast ions with high energy, both of their slowing-down frequency and pitch angle scattering frequency due to Coulomb collisions are much lower than their poloidal bounce frequency. Therefore, fast ions in helical confinement systems are well in the ultra-low collisionality regime. In addition to the fast ions population, the background plasmas in reactor grade helical systems can also be well within the ultra-low collisionality regime.

In the ultra-low collisionality regime where the collision frequency is much lower than the poloidal bounce frequency of the super banana particles, the effects of the finite super banana width may play an important role in neoclassical transport. However, it is difficult for the conventional neoclassical theory [6] to treat this problem.

Generally, collisional transport is governed by the six-dimensional Fokker-Planck kinetic equation. For neoclassical transport, it is well-known that the problem can be reduced to a five-dimensional one by averaging over the gyro-phase angle; this results in the drift kinetic equation. If the collision frequency is much lower than the toroidal bounce frequency, the problem can be further reduced to a four-dimensional one by averaging over the toroidal angle; this results in the bounce-averaged drift kinetic equation [6]. However, this four-dimensional drift kinetic equation is still difficult to be solved, especially in the ultra-low collisionality regime.

It is well-known that a three-dimensional reduced drift kinetic equation can be obtained to describe the neoclassical transport of tokamak plasmas, with the effect of finite banana width taken into account [7, 8, 9]; and the three-dimensional reduced drift kinetic equation can be numerically solved [10, 11]. Therefore, it is of interest to investigate whether we can obtain a similar three-dimensional reduced drift kinetic equation to describe the neoclassical transport of helical plasmas in the ultra-low collisionality regime.

In this paper, we will develope a reduced drift kinetic equation in terms of three constants of motion; this three-dimensional reduced drift kinetic equation can be applied to investigate the neoclassical transport in ultra-low collisionality regime in helical confinement systems. The remaining part of this paper is organized as follows. In Section 2, we present the derivation of the reduced drift kinetic equation; in Section 3, we discuss the choice of the constants of motion; in Section 4, we make some brief discussions on the results.

2. DERIVATION OF THE REDUCED DRIFT KINETIC EQUATION

A. BASIC EQUATIONS

We begin with the five-dimensional drift kinetic equation

$$\partial_t f^{(5)} + \dot{\psi}\, \partial_\psi f^{(5)} + \dot{\theta}\, \partial_\theta f^{(5)} + \dot{\varphi}\, \partial_\varphi f^{(5)}$$

$$= C\left(f^{(5)}\right), \tag{1}$$

where $f^{(5)}$ is the distribution function in the five-dimensional phase space, $(\psi,\ \theta,\ \varphi,\ E,\ \mu)$. $C\left(f^{(5)}\right)$ is the collision term. $\mu = \frac{1}{2}mv_\perp^2/B$; m is the mass of the charged particle; $v_\perp$ is the velocity component perpendicular to the magnetic filed; B is the magnetic field. $E = \mu B + \frac{1}{2}mv_\|^2 + e\Phi$; $v_\|$ is the velocity component parallel to the magnetic filed; e is the electrical charge of the charged particle; Φ is the electrostatic potential. $(\psi,\ \theta,\ \varphi)$ is the magnetic flux coordinates; and the magnetic field is represented as [12]

$$\vec{B} = \nabla\psi \times \nabla\theta + \frac{1}{q(\psi)}\nabla\varphi \times \nabla\psi, \tag{2a}$$

$$\vec{B} = g(\psi)\,\nabla\varphi + I(\psi)\,\nabla\theta + \beta_*(\psi,\theta,\varphi)\,\nabla\psi, \tag{2b}$$

where q is the safety factor. ψ is the toroidal magnetic flux within the flux surface; θ is the poloidal angle; φ is the toroidal angle.

It is well-known that Eq. (1) can be cast into the conservative form in the generalized coordinate system, $(z^1,\ z^2,\ z^3,\ z^4,\ z^5)$,

$$\partial_t f^{(5)} + \frac{1}{J^{(5)}}\frac{\partial}{\partial z^i}\left(J^{(5)}\,\dot{z^i}\,f^{(5)}\right)$$

$$+\frac{1}{J^{(5)}}\frac{\partial}{\partial z^i}\left(J^{(5)}S_z^i\right) = 0, \tag{3}$$

where $J^{(5)}$ is the Jacobian of the five-dimensional phase space, $(z^1, z^2, z^3, z^4, z^5)$; and summation over repeated indices is understood. The contravariant components of the collision induced flux can be written as

$$S_z^i = \left\langle A_z^i\right\rangle_\eta f^{(5)} - \frac{1}{2}\left\langle D_z^{ij}\right\rangle_\eta \frac{\partial}{\partial z^j} f^{(5)}, \tag{4}$$

where the operator, $\langle\cdots\rangle_\eta$, denotes averaging over the gyro-phase angle, η.

The dynamic coefficients used in Eq. (4) are given by

$$A_z^i = \nabla_v z^i \cdot \overrightarrow{A}, \tag{5a}$$

$$D_z^{ij} = \nabla_v z^i \cdot \overleftrightarrow{D} \cdot \nabla_v z^j, \tag{5b}$$

where $\nabla_v = \dfrac{\partial}{\partial \overrightarrow{v}}$. $\overrightarrow{A}$ and $\overleftrightarrow{D}$ are known as the dynamic friction vector and dynamic diffusion tensor in Coulomb collision operator [13].

For the Hamiltonian system, we have

$$\frac{1}{J^{(5)}}\frac{\partial}{\partial z^i}\left(J^{(5)}\,\dot{z^i}\right) = 0. \tag{6}$$

B. FOUR-DIMENSIONAL DRIFT KINETIC EQUATION

Among the five variables, $(z^1, z^2, z^3, z^4, z^5)$, we can choose $z^5 = \varphi$. Using the fact that

$$\dot{\varphi} \gg \frac{\dot{z^i}}{z^i} \sim \frac{C\left(f^{(5)}\right)}{f^{(5)}} \sim \frac{\partial_t f^{(5)}}{f^{(5)}}, (i = 1, 2, 3, 4), \tag{7}$$

and making use of the approximation that $qM \gg 1$ (here M is the number of the toroidal period of the helical winding), we can drop the variable, φ , from the problem under consideration. Using Eq. (7), we obtain the following lowest order form of Eq. (6),

$$\frac{1}{J^{(5)}}\frac{\partial}{\partial \varphi}\left(J^{(5)}\,\dot{\varphi}\right) = 0. \tag{8}$$

With Eq. (8), in parallel to the previous work for tokamak system [9], we can prove two important properties of the Jacobian, which are important when reducing the drift kinetic equation. From Eq. (8), we obtained

$$J^{(5)}\left(z^1, z^2, z^3, z^4, \varphi\right) = J_0^{(5)}\left(z^1, z^2, z^3, z^4\right)\frac{1}{\dot{\varphi}}. \tag{9}$$

Let $J^{(4)}\left(z^1, z^2, z^3, z^4\right)$ denote the Jacobian of the four-dimensional phase space , (z^1, z^2, z^3, z^4). With the aid of Eq. (9) it is not hard to show that

$$J^{(4)}\left(z^1, z^2, z^3, z^4\right) = M J_0^{(5)}\left(z^1, z^2, z^3, z^4\right)\tau_\varphi, \tag{10}$$

τ_φ is the toroidal bounce time in a single helical ripple.

At this point, we define the toroidal-bounce-averaging operator, $\langle\cdots\rangle_\varphi$, as

$$\langle X\rangle_\varphi = \frac{1}{\tau_\varphi}\oint \frac{d\varphi}{\dot{\varphi}} X, \tag{11a}$$

$$\langle X\rangle_\varphi = \frac{1}{\tau_\varphi}\int_0^{2\pi/M} \frac{d\varphi}{\dot{\varphi}} X, \tag{11b}$$

where Eq. (11a) is for particles trapped in helical ripples, while Eq. (11b) is for particles which are not trapped in helical ripples. In Eq. (11), the integrals are performed with (z^1, z^2, z^3, z^4) held constant; this is valid under the condition that $qM \gg 1$.

Clearly, the toroidal bounce time in a single helical ripple, τ_φ , is defined by

$$\langle 1\rangle_\varphi = 1. \tag{12}$$

Now, with the two important properties of the Jacobian [Eq. (9) and Eq. (10)], one can average Eqs. (3-5) over φ by acting $\langle\cdots\rangle_\varphi$ on them. This procedure results in the four-dimensional drift kinetic equation,

$$\partial_t f^{(4)} + \frac{1}{J^{(4)}}\frac{\partial}{\partial z^i}\left(J^{(4)}\left\langle \dot{z^i}\right\rangle_\varphi f^{(4)}\right)$$

$$+\frac{1}{J^{(4)}}\frac{\partial}{\partial z^i}\left[J^{(4)}\left(\left\langle A_z^i\right\rangle_{\eta,\varphi} - \frac{1}{2}\left\langle D_z^{ij}\right\rangle_{\eta,\varphi}\frac{\partial}{\partial z^j}\right) f^{(4)}\right]$$

$$= 0, \tag{13}$$

where $\langle\cdots\rangle_{\eta,\varphi} = \left\langle\langle\cdots\rangle_\eta\right\rangle_\varphi$ was used for abbreviation. $f^{(4)} = f^{(4)}\left(z^1, z^2, z^3, z^4\right)$ is distribution function in the four-dimensional phase space. In order to derive Eq. (13), we have made use of the uniqueness condition for particles trapped in the helical ripple and the periodic condition for particles who are not trapped in the helical ripple, as is similar to what has been discussed in Ref. 6. Of course, it is not hard to show that Eq. (13) is exactly the well-known bounce-averaged drift kinetic equation [6] widely used in the literature, provided that the four variables, (z^1, z^2, z^3, z^4), are chosen in the same way.

C. THREE-DIMENSIONAL REDUCED DRIFT KINETIC EQUATION

According to Ref. 14, this four-dimensional system is also a Hamiltonian system. Therefore, we have

$$\frac{1}{J^{(4)}}\frac{\partial}{\partial z^i}\left(J^{(4)}\left\langle \dot{z^i}\right\rangle_\varphi\right)=0. \tag{14}$$

Now, we choose

$$\left(z^1,z^2,z^3,z^4\right)=\left(c^1,c^2,c^3,\theta\right), \tag{15}$$

where c^1, c^2 and c^3 are three independent constants of motion. At this point, we assume that we can find three constants of motion for charged particles in helical systems. The validity of this assumption will be discussed in the next section.

Substituting Eq. (15) into Eq. (14), we obtain

$$\frac{1}{J^{(4)}}\frac{\partial}{\partial \theta}\left(J^{(4)}\left\langle \dot{\theta}\right\rangle_\varphi\right)=0. \tag{16}$$

Let J_c denote the Jacobian of the three-dimensional phase space of constants of motion, (c^1, c^2, c^3). Here, we also can prove two properties of the Jacobian.

$$J^{(4)}\left(c^1,c^2,c^3,\theta\right)=J_0^{(4)}\left(c^1,c^2,c^3\right)\frac{1}{\left\langle \dot{\theta}\right\rangle_\varphi}, \tag{17}$$

$$J_c\left(c^1,c^2,c^3\right)=J_0^{(4)}\left(c^1,c^2,c^3\right)\tau_\theta, \tag{18}$$

where τ_θ is the poloidal bounce period.

At this point, we define the poloidal-bounce-averaging operator, $\langle\cdots\rangle_\theta$, as

$$\langle X\rangle_\theta=\frac{1}{\tau_\theta}\oint\frac{d\theta}{\left\langle \dot{\theta}\right\rangle_\varphi}X, \tag{19}$$

where the loop integral is evaluated along the real toroidal-bounce-averaged guiding center orbit over the whole poloidal bounce period. The poloidal bounce period is defined by

$$\langle 1\rangle_\theta=1. \tag{20}$$

In this four-dimensional phase space, the toroidal-bounce-averaged drift kinetic equation is written as

$$\partial_t f^{(4)}+\frac{1}{J^{(4)}}\frac{\partial}{\partial\theta}\left(J^{(4)}\left\langle\dot{\theta}\right\rangle_\varphi f^{(4)}\right)$$

$$+\frac{1}{J^{(4)}}\frac{\partial}{\partial c^i}\left[J^{(4)}\left(\left\langle A_c^i\right\rangle_{\eta,\varphi}-\frac{1}{2}\left\langle D_c^{ij}\right\rangle_{\eta,\varphi}\frac{\partial}{\partial c^j}\right)f^{(4)}\right]$$

$$=0, \tag{21}$$

where

$$A_c^i=\nabla_v c^i\cdot\vec{A}, \tag{22a}$$

$$D_c^{ij}=\nabla_v c^i\cdot\overleftrightarrow{D}\cdot\nabla_v c^j, \tag{22b}$$

Consider the ultra-low collisionality regime where

$$\frac{\partial_t f^{(4)}}{f^{(4)}}\sim\frac{C\left(f^{(4)}\right)}{f^{(4)}}\ll\left\langle\dot{\theta}\right\rangle. \tag{23}$$

Under the condition, Eq. (23), we can expand the distribution function as

$$f^{(4)}\left(c^1,c^2,c^3,\theta,t\right)=f_c\left(c^1,c^2,c^3,t\right)$$

$$+f_1\left(c^1,c^2,c^3,\theta,t\right), \tag{24}$$

with f_1 as the perturbation term.

Substitution of Eq. (24) into Eq. (21) gives

$$\partial_t f_c+\frac{1}{J^{(4)}}\frac{\partial}{\partial c^i}\left[J^{(4)}\left(\left\langle A_c^i\right\rangle_{\eta,\varphi}-\frac{1}{2}\left\langle D_c^{ij}\right\rangle_{\eta,\varphi}\frac{\partial}{\partial c^j}\right)f_c\right]$$

$$=-\frac{1}{J^{(4)}}\frac{\partial}{\partial\theta}\left(J^{(4)}\left\langle\dot{\theta}\right\rangle f_1\right). \tag{25}$$

Poloidal-bounce-averaging Eq. (25) [acting Eq. (25) with the poloidal-bounce-averaging operator defined in Eq. (19)], with the aid of Eq. (17) and Eq. (18), we obtain the following reduced drift kinetic equation that determines the distribution function in the three-dimensional phase space of constants of motion,

$$\partial_t f_c+\frac{1}{J_c}\frac{\partial}{\partial c^i}\left[J_c\left(\left\langle A_c^i\right\rangle_{\eta,\varphi,\theta}-\frac{1}{2}\left\langle D_c^{ij}\right\rangle_{\eta,\varphi,\theta}\frac{\partial}{\partial c^j}\right)f_c\right]$$

$$=0, \tag{26}$$

where we have used the uniqueness condition and the periodic condition when annihilating the right-hand-side of Eq. (25) by poloidal-bounce-averaging. Also $\langle\cdots\rangle_{\eta,\varphi,\theta}=\left\langle\langle\cdots\rangle_{\eta,\varphi}\right\rangle_\theta$ has been used for abbreviation. Note that $S_c^i=\left(\left\langle A_c^i\right\rangle_{\eta,\varphi,\theta}-\frac{1}{2}\left\langle D_c^{ij}\right\rangle_{\eta,\varphi,\theta}\frac{\partial}{\partial c^j}\right)f_c$ are contravariant components of collision-induced fluxes in the $\left(c^1,c^2,c^3\right)$ space.

Note that Eq. (26) is a continuity equation in the three-dimensional phase space of constants of motion, and the form of the reduced drift kinetic equation is independent of the specific choice of the three constants of motion. The toroidal-bounce-averaging operator defined here only performs time-average along the guiding center orbits, and the poloidal-bounce-averaging operator defined in this paper merely performs the time-average along the toroidal-bounce-averaged guiding center orbits.

3. CHOICE OF THE CONSTANTS OF MOTION

Having established the reduced drift kinetic equation, which does not depend on the specific choice of the constants of motion, we are in a position to discuss the choice of the constants of motion in order to make it convenient for us to apply the reduced drift kinetic equation.

In addition to E and μ, we have to find the third constant of motion that is necessary for our formalism. Generally, this can not be accomplished in helical systems lacking of the azimuthal symmetry. However, it is well-known [14, 15] that there exists an adiabatic invariant, I_σ, in helical systems under the approximations that $1/qM \ll 1$ and $\varepsilon_t/\varepsilon_h qM \ll 1$ (here, ε_t is inverse of the aspect ratio and ε_h is the amplitude of the helical ripple); and fortunately, these approximations can be applied to a large class of machines.

$$I_\sigma = \oint d\varphi \frac{g}{B} mv_\parallel + \sigma \frac{2\pi}{M} e\psi_P, \tag{27}$$

where $\sigma = 0$ for particles trapped in helical ripples; $\sigma = \pm$ for particles who are not trapped in helical ripples. ψ_P is the poloidal magnetic flux, which is defined by $\dfrac{d\psi}{d\psi_P} = q$.

Of course, there are cyclic changes in this adiabatic invariant; however, these cyclic changes in the adiabatic invariant are not important for the transport behavior of the plasmas. Therefore, we ignore these cyclic changes in the adiabatic invariant and take it as an approximate constant of motion.

At this point, we clarify the classification of particles according to their guiding center orbits projected to the minor cross section. Particles can be classified into three types. The first one is trapped particle that is always trapped in the helical ripple during the entire poloidal bounce period; the second one is passing particle that is never trapped in the helical ripple during the entire poloidal bounce period; the last one is transition particle that is trapped in the helical ripple during one part of the entire poloidal bounce period but escapes from the helical ripple during another part of the entire poloidal bounce period. Transition particles can be subclassified in two groups: transition particles in trapping state (trapped in the helical ripple) and transition particles in passing state (not trapped in the helical ripple). For the details of the theory of guiding center orbits in helical systems, we refer the readers to Refs. 12-15.

Clearly, there is a jump in I_σ when transition particles cross the de-trapping (retrapping) points. Therefore, it is not convenient here to use I_σ as the third constant of motion in addition to E and μ.

One possible choice of the third constant of motion is that $c^3 = \langle\psi\rangle_\theta$; this is similar to the previous discussions on the reduced drift kinetic equation for tokamak plasmas [7, 9]. In helical systems, however, it will be a very involving task to calculate explicitly $\nabla_v c^3$ which is necessary when evaluating the dynamic coefficients used in the reduced drift kinetic equation, if c^3 is defined in this way.

For the present, we propose to choose the third constant of motion as

$$c^3 = \psi_0. \tag{28}$$

Determination of ψ_0 will be described as follows.

Following Ref. 3 and Ref. 14, we specify the fields by

$$B = \bar{B}_0 \left[\varepsilon_0(\psi,\theta) + \varepsilon_1(\psi,\theta)\cos(2\theta - M\varphi)\right], \tag{29a}$$

$$\Phi = \Phi_e(\psi), \tag{29b}$$

where $\bar{B}_0$ is a constant. And we define

$$k^2(\psi,\theta,E,\mu) = \frac{E - \mu B_0 + \mu B_1 - e\Phi_e}{2\mu B_1}, \tag{30}$$

where $B_0 = \bar{B}_0\ \varepsilon_0(\psi,\theta)$, and $B_1 = \bar{B}_0\ \varepsilon_1(\psi,\theta)$. With the model fields given above, analytical functions can be found for

$$I_\sigma = I_\sigma(\psi,\theta,E,\mu). \tag{31}$$

For the details of the functional form of $I_\sigma(\psi,\theta,E,\mu)$, we refer the readers to Ref. 3 and Ref. 12.

With the definition of k^2 given in Eq. (30), it is clear that when $k^2 < 1$, particles are trapped in the helical ripple; when $k^2 \geq 1$, particles are not trapped in the helical ripple. And the de-trapping (retrapping) points of the transition particles satisfy the condition,

$$k^2 = 1. \tag{32}$$

Note that in addition to the model fields we adopted here, there are alternative models for the fields [15, 18]. But this difference does not influence the method we present in this paper.

For passing particles, trapped particles and transition particles in their trapping state, ψ_0 is determined by

$$I_\sigma(\psi,\theta,E,\mu)=I_\sigma^*(\psi_0,E,\mu), \quad (33a)$$

$$I_\sigma^*(\psi_0,E,\mu)\equiv I_\sigma(\psi_0,\theta,E,\mu)_{\theta=\pi}, \quad (33b)$$

where $\theta=0$ corresponds to the outside midplane.

For transition particles in their passing state, ψ_0 is determined by

$$k^2(\psi_T,\theta_T,E,\mu)=1, \quad (34a)$$

$$I_\pm^T(\psi_T,\theta_T,E,\mu)=I_\pm(\psi,\theta,E,\mu), \quad (34b)$$

$$I_0^T(\psi_T,\theta_T,E,\mu)=I_0^*(\psi_0,E,\mu), \quad (34c)$$

$$I_0^*(\psi_0,E,\mu)\equiv I_0(\psi_0,\theta,E,\mu)_{\theta=\pi}. \quad (34d)$$

Essentially, our choice of ψ_0 is the generalization of Ref. 6 and Ref. 8. Clearly, ψ_0 is the value of ψ when the particle crosses the inner-side midplane ($\theta=\pi$). This choice is valid for the model fields we have adopted here; for different model fields, ψ_0 can be chosen in the similar way.

Now, we have chosen the three constants of motion

$$\left(c^1,c^2,c^3\right)=(E,\mu,\psi_0). \quad (35)$$

In order to calculate the dynamic coefficients used in the reduced drift kinetic equation, we have to calculate $\nabla_v c^3$. Note that

$$\nabla_v c^3=\nabla_v\psi_0(\psi,\theta,E,\mu)$$

$$=\nabla_v E\partial_E\psi_0+\nabla_v\mu\partial_\mu\psi_0, \quad (36)$$

we only have to calculate $\partial_E\psi_0$ and $\partial_\mu\psi_0$. The mathematical manipulation is lengthy but straightforward. Here, we only present the resulting formulae.

For passing particles, trapped particles and transition particles in their trapping state, we have found

$$\partial_E\psi_0=\left(\partial_E I_\sigma-\partial_E I_\sigma^*\right)/\partial_{\psi_0}I_\sigma^*, \quad (37a)$$

$$\partial_\mu\psi_0=\left(\partial_\mu I_\sigma-\partial_\mu I_\sigma^*\right)/\partial_{\psi_0}I_\sigma^*. \quad (37b)$$

For transition particles in their passing state, we have found

$$\partial_E\psi_0=\frac{\partial_E I_0^T+\partial_E\psi_T\partial_{\psi_T}I_0^T+\partial_E\theta_T\partial_{\theta_T}I_0^T-\partial_E I_0^*}{\partial_{\psi_0}I_0^*}, \quad (38a)$$

$$\partial_\mu\psi_0=\frac{\partial_\mu I_0^T+\partial_\mu\psi_T\partial_{\psi_T}I_0^T+\partial_\mu\theta_T\partial_{\theta_T}I_0^T-\partial_\mu I_0^*}{\partial_{\psi_0}I_0^*}, \quad (38b)$$

where

$$\partial_E\psi_T=\triangle_{E,\theta_T}/\triangle, \quad (39a)$$

$$\partial_\mu\psi_T=\triangle_{\mu,\theta_T}/\triangle, \quad (39b)$$

$$\partial_E\theta_T=\triangle_{E,\psi_T}/\triangle, \quad (39c)$$

$$\partial_\mu\theta_T=\triangle_{\mu,\psi_T}/\triangle, \quad (39d)$$

with

$$\triangle=\left|\begin{array}{cc}\partial_{\psi_T}I_\pm^T & \partial_{\theta_T}I_\pm^T\\ \partial_{\psi_T}k^2 & \partial_{\theta_T}k^2\end{array}\right|, \quad (40a)$$

$$\triangle_{\alpha,\beta}=\left|\begin{array}{cc}\partial_\beta I_\pm^T & -\partial_\alpha I_\pm^T+\partial_\alpha I_\pm\\ \partial_\beta k^2 & -\partial_\alpha k^2\end{array}\right|, \quad (40b)$$

where $\alpha=E,\mu$; $\beta=\psi_T,\theta_T$.

4. DISCUSSIONS

Based on the original five-dimensional drift kinetic equation, we have established the three-dimensional reduced drift kinetic equation for neoclassical transport of helical plasmas in ultra-low collisionality regime; the derivation of the reduced drift kinetic equation is based on a plausible assumption that we can find an approximate constant of motion, in addition to energy and magnetic moment. The reduced drift kinetic equation is a continuity equation governing the distribution of confined particles in the three-dimensional phase space of constants of motion. It is clear that the effects of finite super banana width and finite aspect ratio have been included in this equation.

Although the reduced drift kinetic equation derived in this paper is independent of the specific choice of the three constants of motion, we have discussed a convenient choice of the three constants of motion for the potential application of the reduced drift kinetic equation. And we have presented the necessary formulae for calculating the dynamic coefficients used in the reduced drift kinetic equation.

One of the potential applications of the reduced drift kinetic equation is that it may provide an alternative way to numerical simulations of neoclassical transport of helical plasmas in ultra-low collisionality regime. Since the reduced drift kinetic equation is a three-dimensional continuity equation, it will not be difficult to numerically solve it. It is being under consideration to develope a computer code to numerically solve this reduced drift kinetic equation. Once the equation is solved for the distribution in the three-dimensional phase space of constants of motion, the macroscopic physical quantities can be found immediately.

Generally, there are two ways to obtain the transport fluxes in real space from the reduced drift kinetic equation [Eq. (26)]. In the first way, using f_c, we can obtain the distribution function $f(\psi, \theta, E, \mu)$; then we can obtain the macroscopic transport fluxes by taking moments of $f(\psi, \theta, E, \mu)$ in the conventional way. In the second way, the transport fluxes can be obtained directly by transformation from the (c^1, c^2, c^3) space to the conventional (ψ, θ, E, μ) space; the velocity-dependent collision-induced transport flux in the direction of $\nabla\psi$ in the conventional (ψ, θ, E, μ) space, Υ^ψ, can be obtained by using the general theory of coordinates transformation, since we know S_c^i, the contravariant components of the collision-induced flux in the (c^1, c^2, c^3) space; then we can obtain the macroscopic particle (energy) flux by integrating Υ^ψ $(\Upsilon^\psi mv^2/2)$ over the velocity space and averaging the results over the magnetic flux surface.

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