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Tokamak electrical conductivity modified by electrostatic trapping in the applied electric field

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Passing electrons traveling in the same direction as the applied Ohmic electric field, near the trapped-passing boundary, can get electrostatically trapped by the applied Ohmic potential, reverse their parallel flow direction, and become passing particles traveling in the opposite direction. The number of these electrons is proportional to the applied voltage and is small. But, the amount of change in their time-averaged parallel velocity, by reversing their flow direction, is about twice the average parallel velocity of the barely passing electrons and is large. Hence, these electrostatically trapped particles can significantly contribute to the Ohmic electrical current. The new electrical conductivity is calculated in the banana regime using a pitch angle scattering operator, and the values are found to be significantly enhanced from previous neoclassical theories.

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Establishing an accurate theoretical expression for electrical conductivity is an important task in understanding tokamak plasma behavior. Without an accurate electrical conductivity calculation, one cannot have confident analyses of current profile and, thus, of basic plasma equilibrium and stability. Despite the turbulent behavior of tokamak plasmas, it has been evidenced experimentally that the electrical conductivity in tokamaks may not be affected by plasma turbulence, and thus neoclassical analysis is a reliable tool for the subject.

In previous analyses,¹⁻³ it was shown that electrical conductivity values in a tokamak were significantly reduced from the result of Spitzer⁴ due to the electron trapping effect in the inhomogeneous magnetic field and also due to the collisional drag effect by the trapped electrons. However, in order to simplify the analyses, the previous theories used an approximation based on the conservation of particle kinetic energy, instead of conservation of the total energy, which includes the applied potential energy. Since the applied Ohmic voltage is usually small compared to the particle's kinetic energy, the approximation is not expected to yield much error for the particles deeply passing or trapped. But, for the particles near the trapped-passing boundary, their parallel kinetic energy becomes small near the maximum- B point and the modification of the particle orbit by the applied voltage becomes significant. As a result, some of the passing electrons traveling in the same direction as the applied Ohmic electric field ($\vec{E}_{Ohm}$) near the trapped-passing boundary are forced by the applied electric field to turn around and become part of the oppositely traveling

passing electrons (see Fig. 1). We call these particles "Ohmically reversing" electrons in the present work. The motion of these electrons is characterized by an applied-electric-field driven electron flux in velocity space across the trapped-passing boundary, which crosses the boundaries from the co-streaming (relative to $\vec{E}_{Ohm}$) passing region to the trapped region and then to the counter-streaming passing region. This extra particle flux (eventually balanced by the collisional counter flux) across the trapped-passing boundary is found to generate an additional Ohmic electrical current and enhance the neoclassical conductivity. Previous analyses did not allow for such a particle motion due to the approximation that the particle kinetic energy is a constant of motion.

In Spitzer's analyses of electrical conductivity⁴ in a homogeneous magnetic field geometry, this Ohmically reversing particle effect was included naturally. However, the average parallel (to $\vec{E}_{Ohm}$) velocity of the Ohmically reversing electrons was vanishingly small in a homogeneous magnetic field, and thus they did not contribute to the electrical conductivity. The contribution from the Ohmically reversing electrons becomes significant only in an inhomogeneous magnetic field geometry, where the time-averaged parallel velocities of those particles do not vanish, but are rather large. It is, therefore, important to include these particles in the analyses to obtain correct values for electrical conductivity in an inhomogeneous magnetic field. In previous neoclassical theories, which used the Spitzer equation as the basis of analyses, the effect of these Ohmically reversing electrons was implicitly included in the Spitzer function, but only to the extent of

homogeneous magnetic field problems where such an effect gave negligible contribution.

Qualitative estimation of the amount of additional current generation by these Ohmically reversing electrons can be easily obtained. The fraction of the Ohmically reversing electrons is roughly the thickness $\Delta v_{\parallel}/v$ of the pitch angle layer, near the trapped-passing boundary, occupied by these particles. If the electrons are within this layer, the amount of electrostatic deceleration in a bounce time is large enough to force them to reverse their flow direction near the maximum Mod- B point, satisfying $m\Delta v_{\parallel} \sim eE_{\parallel}\tau_b$. Thus the number density Δn is roughly given as $\Delta n \sim (eE_{\parallel}\tau_b/mv)$. Since the magnitude of the parallel velocity change for these particles near the trapped-passing boundary is $(r/R)^{1/2}v$, the rate of current generation is $dj/dt \sim (e\Delta n/\tau_b)(r/R)^{1/2}v \sim (ne^2/m)(r/R)^{1/2}E_{\parallel}$. This current generation will eventually have to be balanced by Coulomb collisions, and we get the relation $(r/R)^{1/2}(ne^2/m)E_{\parallel} \sim \nu_c j$, where ν_c is the Coulomb 90°-collision frequency. Hence the amount of current generated by the Ohmically reversing electrons is $j \sim (r/R)^{1/2}(ne^2/m\nu_c)E_{\parallel}$, which is of the same order of magnitude as the usual neoclassical reduction current.

We define ϵ and λ as the total energy and the magnetic moment normalized to the kinetic energy W per unit mass, respectively: $\epsilon = W + q\phi/m$, $\lambda = (\mu/W) = v_{\perp}^2 / 2WB$. Then ϵ is a constant of motion, but λ is not. If we define θ as a poloidal angle and $\sigma \equiv v_{\parallel}/|v_{\parallel}|$, the electron motion in $(\lambda, \sigma, \theta)$ -space is as described in Fig. 2. We can see from Fig. 2 that the passing electrons crossing the usual, purely magnetic, trapped-passing

boundary (dotted line) from the $\sigma = 1$ side bounce back and cross the other trapped-passing boundary at the $\sigma = -1$ side into the counter-streaming passing regime. The angle at which the Ohmically reversing electrons meet the usual trapped-passing boundary is a function of particle kinetic energy and thus the combined electro-magnetic trapped-passing boundary cannot be defined uniquely. In other words, the higher energy electron motions are less affected and the lower energy ones are more affected. For this reason, we will still keep the usual purely magnetic definition of trapped-passing boundary in the present work.

Since the Ohmically reversing motions can be characterized by the rapid change of particle pitch angles into and out of the trapped region (in the usual sense), in order to study their effect on the electrical current, we need to calculate accurately the fast exchange rate of the particles between the trapped and passing regions. For this purpose, it is more convenient to treat the trapped and passing particles as separate species, but collisionally shaped together into a Maxwellian, to the lowest order. Notice here that as far as other dissipational processes (other than the Coulomb collisions) are concerned, trapped and passing particles respond differently and they are two different groups of particles.

We note here that the present problem may also be handled without separating the trapped and passing electrons into two different species, as long as the fast particle flow across the trapped-passing boundary is evaluated accurately outside of the drift kinetic equation and used, as a boundary condition, in the determination of the perturbed distribution

function. But, as will be discussed later, such an approach may be more complicated than the one used in the present analyses.

The usual drift-kinetic equation is written as

$$\vec{v}_{||} \cdot \nabla f_{\alpha} = C(f_{\alpha}) + \nabla \cdot \vec{\Gamma}_{\alpha} , \quad (1)$$

where $\alpha = t$ for trapped electrons and $\alpha = p$ for passing electrons, and Γ_{α} is the particle flux in the velocity space driven by the applied electric field. In terms of v and $\xi = v_{||}/v$, $\nabla \cdot \vec{\Gamma}$ can be expressed as

$$\nabla \cdot \vec{\Gamma} = \frac{eE_{||}}{m} \left(\frac{v_{||}}{v} \frac{\partial f_{\alpha}}{\partial v} + \frac{1-\xi^2}{v} \frac{\partial f_{\alpha}}{\partial \xi} \right) , \quad (2)$$

where $E_{||}$ is the parallel component of the applied Ohmic electric field. In order to avoid the runaway problem, the applied Ohmic electric field has to be small. This leads us to the small- $E_{||}$ ordering $\hat{E}_{||} \equiv (eE_{||}/mv\omega_b) \ll 1$, and we expand f in $\hat{E}_{||}$: $f = f_{\alpha 0} + f_{\alpha 1} + f_{\alpha 2} + \dots$. Then the H-theorem assures us a Maxwellian solution for the combined f_o : $f_o = f_{t0} + f_{p0} = f_M = \text{Maxwellian}$, and

$$f_{t0} = f_M \Theta(\xi_c - |\xi|) , \quad (3)$$

$$f_{p0} = f_M \Theta(|\xi| - \xi_c) , \quad (4)$$

where ξ_c is the magnitude of the pitch angle corresponding to the trapped-passing boundary, and Θ is the heavy-side step function. The first order version of Eq. (1) becomes, using Eq. (2):

$$\vec{v}_{||} \cdot \nabla f_{\alpha 1} - C(f_{\alpha 1}) = - \frac{eE_{||}}{T} v_{||} f_{\alpha 0} + \frac{eE_{||}}{mv} (1-\xi^2) \frac{\partial f_{\alpha 0}}{\partial \xi} . \quad (5)$$

The first term in the right-hand side of Eq. (5) represents the usual value of $\nabla \cdot \Gamma$ obtained by considering the trapped and passing particles as one Maxwellian species. Any divergence-free flow, no matter how large it is, cannot contribute to the first term and cannot be calculated from the drift kinetic equation if the trapped and passing particles are treated as one species. In determining the perturbed distribution function from the differential equation (drift-kinetic equation), the flux of particles across the trapped-passing boundary is often an important factor, due to the different properties between the trapped and passing particles. By considering the trapped and passing particles as two different species, such a particle exchange between the trapped and passing regions can be easily evaluated, as can be seen from the second term in the right-hand side of Eq. (5).

If the trapped and passing particles are treated as one species, such a driven flow will have to be calculated outside of the drift-kinetic equation and used as an additional boundary condition in determining f_1 at the trapped-passing boundary. This additional flow cannot be easily handled within the framework of the usual variational approach^{1,2} if the trapped and passing particles are treated as one species, because the electrical force on the present Ohmically reversing particles are balanced by both collisions and bounce motions. This is unlike the usual cases, where there are clear-cut relationships between the forces and the class of particles: the electric force on the passing particles is mainly balanced by collisions and on the trapped particles by bounce motions. Determination of the Ohmically reversing particle flow requires determination of the perturbed distribution

function for the Ohmically reversing particles. However, in this situation, the attempt to describe the perturbed distribution function for the Ohmically reversing particles, within the one species model, can become highly tedious and erroneous. By treating the trapped particles and the passing particles as two separate species, we can easily determine the divergence-free, Ohmically reversing particle flow across the trapped-passing boundary, and incorporate it into the drift-kinetic equation as the second term in the right-hand side of Eq. (5). We can treat the second term as if it is an additional scattering term localized at the trapped-passing boundary. We note here that including the particle flow across the trapped-passing boundary in this manner in the drift-kinetic equation is not usual, but mainly a matter of mathematical convenience.

Using Eqs. (3) and (4), we can write the first order equations as

$$\begin{aligned}\vec{v}_{\parallel} \cdot \nabla f_{p1} - C(f_{p1}) &= -\frac{eE_{\parallel}}{T} v_{\parallel} f_M - \frac{eE_{\parallel}}{m} f_M \frac{2\xi}{Bv} (1 - \xi^2) \delta(\lambda - \lambda_c), \\ \vec{v}_{\parallel} \cdot \nabla f_{i1} - C(f_{i1}) &= -\frac{eE_{\parallel}}{T} v_{\parallel} f_M + \frac{eE_{\parallel}}{m} f_M \frac{2\xi}{Bv} (1 - \xi^2) \delta(\lambda - \lambda_c),\end{aligned}$$

where $\lambda_c = \lambda$ at $\xi = \xi_c$. It is more convenient to transform the above equations, using a scheme similar to one used in Ref. 1,

$$\begin{aligned}f_{\alpha 1} = & -\frac{e}{T} f_M \int_0^{\ell_p} \frac{d\ell_p}{B_p} \left[B E_{\parallel} - \frac{B^2 \langle E_{\parallel} B \rangle}{\langle B^2 \rangle} \right] \\ & + v_{\parallel} f_{\alpha} \frac{B \langle E_{\parallel} B \rangle}{\langle B^2 \rangle} + H_{\alpha}, \quad \alpha = p \text{ or } t,\end{aligned}$$

where the ion flow has been assumed to be zero for simplicity, ℓ_p is the length of the field line in the poloidal direction, $\langle \dots \rangle$ represents the flux-surface

average, and f_s is the Spitzer function⁴ defined as $C(v_{||}f_s) = (e/T) v_{||}f_M$.

Using the above transformation of $f_{\alpha 1}$ into H_α , we obtain

$$\begin{aligned}\bar{v}_{||} \cdot \nabla H_p - C(H_p) &= -\bar{v}_{||} \cdot \nabla v_{||}f_s \frac{B \langle E_{||}B \rangle}{\langle B^2 \rangle} - \frac{eE_{||}}{m} f_M \frac{2\xi}{Bv} (1 - \xi^2) \delta(\lambda - \lambda_c), \\ \bar{v}_{||} \cdot \nabla H_i - C(H_i) &= -\bar{v}_{||} \cdot \nabla v_{||}f_s \frac{B \langle E_{||}B \rangle}{\langle B^2 \rangle} - \frac{eE_{||}}{m} f_M \frac{2\xi}{Bv} (1 - \xi^2) \delta(\lambda - \lambda_c).\end{aligned}$$

We can further simplify the above equations by defining g and g_i as follows:

$$C(g) = \frac{eE_{||}}{m} f_M \frac{2\xi}{Bv} (1 - \xi^2) \delta(\lambda - \lambda_c),$$

$$\bar{v}_{||} \cdot \nabla g_i = -\frac{eE_{||}}{m} f_M \frac{4\xi}{Bv} (1 - \xi^2) \delta(\lambda - \lambda_c),$$

and transforming H_α into F_α , $H_\alpha = F_\alpha + g$. The first order equations are then finally transformed into

$$\bar{v} \cdot \nabla F_p - C(F_p) = -\bar{v}_{||} \cdot \nabla \left[v_{||}f_s \frac{B \langle E_{||}B \rangle}{\langle B^2 \rangle} + g \right], \quad (6)$$

$$\bar{v} \cdot \nabla F_i - C(F_i) = -\bar{v}_{||} \cdot \nabla \left[v_{||}f_s \frac{B \langle E_{||}B \rangle}{\langle B^2 \rangle} + g - g_i \right]. \quad (7)$$

The function g represents the effect of additional electron flux due to the Ohmically reversing motions, and thus $v_{||}f_s B \langle E_{||}B \rangle / \langle B^2 \rangle + g$ constitutes a modified Spitzer function.

In the banana regime of collisionality, we adopt the small collision frequency ordering, $\nu_* \equiv (\nu_e R q / v \epsilon^{3/2}) \ll 1$, where q is the safety factor, $\epsilon \equiv r/R$ is the inverse aspect ratio, and ν_e is the electron 90°-scattering frequency. And we expand F_α in ν_* : $F_\alpha = F_\alpha^{(0)} + F_\alpha^{(1)} + F_\alpha^{(2)} + \dots$, where the superscripts represent the corresponding orders in ν_* . Then, to the

lowest order in ν_* , Eqs. (6) and (7) become

$$\begin{aligned}\vec{v} \cdot \nabla F_p^{(0)} &= -\vec{v}_{\parallel} \cdot \nabla \left(v_{\parallel} f_s \frac{\langle B E_{\parallel} B \rangle}{\langle B^2 \rangle} + g \right), \\ \vec{v} \cdot \nabla F_i^{(0)} &= -\vec{v}_{\parallel} \cdot \nabla \left(v_{\parallel} f_s \frac{\langle B E_{\parallel} B \rangle}{\langle B^2 \rangle} + g - g_i \right),\end{aligned}$$

and the solution $F_{\alpha}^{(0)}$ is trivially obtained as follows:

$$\begin{aligned}F_p^{(0)} &= -v_{\parallel} f_s \frac{B \langle E_{\parallel} B \rangle}{\langle B^2 \rangle} - g + G_p, \\ F_i^{(0)} &= -v_{\parallel} f_s \frac{B \langle E_{\parallel} B \rangle}{\langle B^2 \rangle} - g + g_i + G_i,\end{aligned}$$

where the function G_{α} is independent of the poloidal angle θ and is determined from the first order(in ν_*) version of Eqs. (6) and (7) in the form of a solubility condition

$$\{C(F_{\alpha}^{(0)})\}_b = 0, \quad (8)$$

where $\{\dots\}_b$ represents the bounce average operator.

The function $H_{\alpha}^{(0)}$ is then given by the relation $H_{\alpha}^{(0)} = F_{\alpha}^{(0)} + g$, and thus

$$\begin{aligned}H_p^{(0)} &= -v_{\parallel} f_s \frac{B \langle E_{\parallel} B \rangle}{\langle B^2 \rangle} + G_p, \\ H_i^{(0)} &= -v_{\parallel} f_s \frac{B \langle E_{\parallel} B \rangle}{\langle B^2 \rangle} + g_i + G_i.\end{aligned}$$

And the current density is calculated as:

$$\begin{aligned}\left\langle \frac{j - j_s}{h} \right\rangle &= -e \left\langle \frac{1}{h} \int d^3v \frac{H}{f_M} v_{\parallel} f_M \right\rangle = -e \left\langle \frac{1}{h} \int d^3v \frac{H}{f_M} \frac{T}{e} C(v_{\parallel} f_s) \right\rangle \\ &\cong -T \left\langle \frac{1}{h} \int d^3v v_{\parallel} \frac{f_s}{f_M} C(H^{(0)}) \right\rangle,\end{aligned}$$

where the self-adjointness property of the Coulomb collision operator has been used and j_s is the Spitzer current density defined as

$$\langle \frac{j_s}{h} \rangle = -e \langle \frac{1}{h^2} \int d^3v v_{\parallel}^2 f_s \rangle \frac{B_o \langle E_{\parallel} B \rangle}{\langle B^2 \rangle}.$$

The expressions for $H_p^{(0)}$ and $H_t^{(0)}$ are then to be substituted into the above equation. However, from the definition of g_t and G_t , it can be easily seen that g_t and G_t are even functions of $v_{\parallel}$ and, therefore, cannot contribute to $\langle (j - j_s)/h \rangle$. Thus our neoclassical current density becomes

$$\langle \frac{j - j_s}{h} \rangle \cong -T \langle \frac{1}{h} \int d^3v v_{\parallel} \frac{f_s}{f_M} C \left[-v_{\parallel} f_s \frac{B \langle E_{\parallel} B \rangle}{\langle B^2 \rangle} + G_p \Theta(\lambda_c - \lambda) \right] \rangle. \quad (9)$$

It remains to determine the function G_p from the constraint equation (8):

$$\left\{ C(-v_{\parallel} f_s \frac{B \langle E_{\parallel} B \rangle}{\langle B^2 \rangle} + G_p) + \frac{e E_{\parallel}}{m} f_M \frac{2\xi}{Bv} (1 - \xi^2) \delta(\lambda - \lambda_c) \right\}_b = 0.$$

By subtracting out G_o , the known part of G_p from Ref. 1, which satisfies

$$\left\{ C(-v_{\parallel} f_s \frac{B \langle E_{\parallel} B \rangle}{\langle B^2 \rangle} + G_o) \right\}_b = 0,$$

we obtain

$$\left\{ C(G_p^*) + \frac{e E_{\parallel}}{m} f_M \frac{2\xi}{Bv} (1 - \xi^2) \delta(\lambda - \lambda_c) \right\}_b = 0, \quad (10)$$

where $G_p^* = G_p - G_o$. Since a discontinuous structure of the second derivative of G_p^* in λ is responsible for the δ -function, we approximate the collision operator as a simplified pitch angle scattering operator in this letter:

$$C(G_p^*) \equiv \nu_e \frac{2\xi}{B} \frac{\partial}{\partial \lambda} \xi \lambda \frac{\partial}{\partial \lambda} G_p^* \simeq \nu_e \frac{2\xi^2 \lambda}{B} \frac{\partial^2}{\partial \lambda^2} G_p^*,$$

where $\nu_e = \nu_{ei} + \nu_{ee}$ contains both electron-ion collision frequency ν_{ei} and the electron-electron collision frequency ν_{ee} . Then Eq. (10) trivially yields

$$C(G_p^*) \simeq - \frac{2e \langle E_{\parallel} B \rangle}{mv B} \lambda \frac{\xi^2}{\langle \xi \rangle} f_M \delta(\lambda - \lambda_c).$$

Equation (9) then becomes

$$\begin{aligned} \left\langle \frac{j - j_s}{h} \right\rangle &\simeq -T \left\langle \frac{1}{h} \int_{\text{pass}} d^3v v_{\parallel} \frac{f_s}{f_M} C \left[-\frac{v_{\parallel} f_s B_o \langle E_{\parallel} B \rangle}{h \langle B^2 \rangle} + G_o + G_p^* \right] \right\rangle \\ &= \sigma_o \frac{B_o \langle E_{\parallel} B \rangle}{\langle B^2 \rangle} - T \left\langle \frac{1}{h} \int d^3v v_{\parallel} \frac{f_s}{f_M} C(G_p^*) \right\rangle \\ &\simeq \sigma_o \frac{B_o \langle E_{\parallel} B \rangle}{\langle B^2 \rangle} \\ &\quad + \frac{2Te \langle E_{\parallel} B \rangle}{m \langle \xi_c \rangle} \left\langle \frac{1}{hB} \int \frac{d^3v}{v} \lambda \xi^2 v_{\parallel} f_s \delta(\lambda - \lambda_c) \right\rangle. \end{aligned}$$

Here σ_o is the usual neoclassical reduction amount from the Spitzer conductivity σ_s as given in Ref. 1, $\sigma_o = -1.46(1 + 0.34/Z_i) \sqrt{r/R} \sigma_s$, where Z_i is the ionic charge number. By evaluating the velocity integral and the flux-surface-average in a straightforward manner, and using the numerical tabulation of Ref. 4 for f_s , we find that the second term in $\langle (j - j_s)/h \rangle$ becomes

$$0.387 \left(1 + \frac{0.413}{Z_i} \right) \sqrt{\frac{r}{R}} \sigma_s \frac{B_o \langle E_{\parallel} B \rangle}{\langle B^2 \rangle},$$

where in the parentheses, the first term is from ν_{ei} and the Z_i -dependent term is from ν_{ee} (the $1/Z_i$ dependence is an approximate fit).

Combining the above two terms, we finally obtain

$$\left\langle \frac{j}{h} \right\rangle = \left[1 - 1.07 \left(1 + \frac{0.32}{Z_i} \right) \sqrt{\frac{r}{R}} + \left(0.07 + \frac{0.34}{Z_i} \right) \frac{r}{R} \right] \sigma_s \frac{B_o \langle E_{\parallel} B \rangle}{\langle B^2 \rangle},$$

where the term proportional to r/R is added here to make $\langle j/h \rangle \rightarrow 0$ as $r/R \rightarrow 1$. As a conclusion, the electrical conductivity in the banana regime calculated here, modified by the electrostatic trapping effect in the applied Ohmic field, is significantly enhanced from the previous result. For $Z_i = 1$, the conductivity values are about 28% closer to the value of Spitzer and Härm.⁴

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Figures

FIG. 1. Trapped-passing boundaries in $(v_{||}/v)$ and poloidal angle space. (a) Boundaries without the modification by the applied electric field. (b) Boundaries with the modification. Electrons in the shaded region are Ohmically reversing. And the electrons in the dotted area are truly trapped.

FIG. 2. Exaggerated illustration of the Ohmically reversing electron orbits in $(\lambda, \sigma, \theta)$ -space. The dotted areas are the usual magnetically trapped regions.

Fig. 1

Fig. 2

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