ON RELATIVISTIC EFFECTS ON ELECTRON TRANSPORT IN THE BANANA REGIME IN TOKAMAKS

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In the present work, the neoclassical transport theory in tokamaks is re-considered with the relativistic effects for electrons taken into account. Since such effects are important only in high-temperature plasmas, only the low collisional banana regime has been considered. The obtained formulations give a possibility to calculate the electron neoclassical fluxes in very broad range of temperatures.

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INTRODUCTION

A general trend for fusion researches is already oriented to development of the commercial reactor. The focus of researches is set on physics of the processes related to burning plasmas, where the self-heating by fusion reactions is balanced predominantly by transport processes with minimized losses. Quantitative requirements for steady state in burning plasmas are known as Lawson criterion [1], $n\tau_E \geq 5\times 10^{20} s\cdot m^{-3}$, and, as a more modern criterion, the triple product [2], $nT\tau_E \geq 3\times 10^{21} keV\cdot s\cdot m^{-3}$ (here, $T = T_e \cong T_i$ is plasmas temperature, n is the density, $\tau_E = W/P_{inp}$ is the energy confinement time with W as total plasma energy confined and $P_{inp} = P_{fusion} + P_{aux}$ as total input (heating) power balanced by losses from plasmas, $P_{inp} = P_{out}$; numbers are given for the D-T plasma). Actually, these criteria can be reached already in ITER (under construction in Cadarache, France) with plasma density about $10^{20} m^{-3}$ and temperatures about 25…30 keV. At the same time, the fusion reactors based on aneutronic schemes seem more attractive for future due to minimization of the neutron flux which produces a destructive influence on the reactor itself. Being based on other reactions (the best one seems to be a hot D-^{3}He plasmas, where D-D parasitic branch can be minimized by relative concentration), this kind of reactors requires, however, a higher plasma temperatures, in the range of 50…70 keV.

As follows from a simplest estimations, relativistic effects become important for electrons practically in all kinds of fusion plasmas. These effects have been studied in details for radiactive losses from hot plasmas as well for plasma heating by high-frequency waves [3]. And while the neoclassical transport theory is well established [4], the theory of relativistic transport processes for electrons in hot plasmas is still far from completeness. In particular, only several aspects were investigated, such as specific feature of electron-ion energy interchange [5] and influence of relativistic effects on neoclassical radial fluxes in 1/ν-regime for stellarators [6].

In the present work, we continue the same line as in our previous papers [6-8] devoted to stellarators and re-consider now the neoclassical transport theory for electrons in hot tokamak plasmas. Only the low-collisional banana regime is considered.

RELATIVISTIC DRIFT KINETIC EQUATION

As follows from the formulations given in [9, 10], the relativistic drift equations of motion in weakly inhomogeneous magnetic field $\boldsymbol{B}$, where an adiabatic invariance of $\mu = m_e u_\perp^2/2B$ (with $u_\perp = \gamma v_\perp$) is satisfactory, can be represented as

$$\dot{\boldsymbol{X}} = \frac{u_\parallel}{\gamma B}\boldsymbol{B} - \frac{u_\parallel}{\gamma \Omega_e}\boldsymbol{B}\times\nabla\left(\frac{u_\parallel}{B}\right) + \boldsymbol{V}_{\nabla\times\boldsymbol{B}}, \qquad (1)$$

$$\dot{u}_\parallel = \frac{e}{m}E_\parallel - \frac{u_\perp^2}{2\gamma B^2}(\boldsymbol{B}\cdot\nabla B) + (F_{\nabla\times\boldsymbol{B}})_\parallel. \qquad (2)$$

Here all the values correspond to the gyrocenter, $\boldsymbol{X}$ is the drift trajectory, $u_\parallel = \gamma v_\parallel$ is the parallel momentum per unit mass with $\gamma = \sqrt{1+u^2/c^2}$ and $\Omega_{ce} = |eB/m_e c|$ is the electron gyrofrequency. Only the longitudinal electric field, $E_\parallel = \boldsymbol{E}\cdot\boldsymbol{B}/B$, is accounted while the radial electric field leads to a plasma rotation in poloidal direction without any contribution to transport. The terms proportional to $\nabla\times\boldsymbol{B}$, which are related to the poloidal drift perpendicular to the magnetic field line, do not contribute in radial and longitudinal fluxes and are omitted from further consideration. It can be shown also that the second term in $\dot{u}_\parallel$ (so-called the mirror-force in (2)) does not perform any total work (no contribution on current) and can be omitted as well.

Magnetic field in tokamaks, i.e. in toroidal traps with an axial symmetry, is usually represented [4] as

$$\boldsymbol{B} = I(\psi)\nabla\varphi + \nabla\varphi\times\nabla\psi, \quad I = RB_\varphi. \qquad (3)$$

Here (R,φ,z) are local cylindrical coordinates in torus, ψ is the poloidal magnetic flux that corresponds to the given magnetic surface, $\psi(R,z) = const$, and $I(\psi)$ is the flux-function.

Generally, the linearized relativistic drift kinetic equation (rDKE) for electrons can be written as following:

$$\frac{u_\parallel}{\gamma}\nabla_\parallel f_{e1} + \boldsymbol{V}_{dr}\cdot\nabla F_{MJ} + \dot{u}_\parallel\frac{\partial F_{MJ}}{\partial u_\parallel} = C_e(f_{e1}), \qquad (4)$$

where $f_{e1} = f_e - F_{MJ}$ is the local deviation from thermal equilibrium induced by $E_\parallel$ and neoclassical effects; the thermal equilibrium of relativistic electrons is given by the Maxwell-Jüttner distribution function,

$$F_{MJ}(u) = \frac{n_e}{\pi^{3/2}u_{te}^3}C_{MJ}(\mu_r)e^{-\mu_r(\gamma-1)}, \qquad (5)$$

$$C_{MJ}(\mu_r) = \sqrt{\frac{\pi}{2\mu_r}\frac{e^{-\mu_r}}{K_2(\mu_r)}} \cong 1 - \frac{15}{8\mu_r} + \cdots, \qquad (6)$$

with $u_{te} = \sqrt{2T_e/m_e}$ and $\mu_r = m_e c^2/T_e \gg 1$. The collisions are described by the linearized Coulomb operator, $C_e(f_{e1}) = C_{ee} + C_{ei}$, where C_{ee} is taken as linearized one and C_{ei} handles only the pitch-scattering (Lorentz term). The value $\boldsymbol{V}_{dr}$, given by the second term in (1), describes the radial drift and is responsible for the radial flux in (4).

Using the definitions introduced above, rDKE can be represented in the following form:

$$\frac{u_\parallel}{\gamma}\nabla_\parallel f_{e1} + \boldsymbol{V}_{dr}\cdot\nabla F_{MJ} - \frac{eE_\parallel}{T_e}\frac{u_\parallel}{\gamma}F_{MJ} = C_e(f_{e1}), \quad (7)$$

$$\boldsymbol{V}_{dr}\cdot\nabla F_{MJ} = I\frac{u_\parallel}{\gamma}\nabla_\parallel\left(\frac{u_\parallel}{\Omega_{ce}}\right)\frac{dF_{MJ}}{d\psi}. \quad (8)$$

This equation is quite general and can be applied for any collisional regime in axisymmetrical traps (tokamaks). Below, we consider only the high temperature case, when collisionality is low. For tokamaks, this case is traditionally called the "banana" regime.

RELATIVISTIC ELECTRON TRANSPORT IN BANANA REGIME

In the banana regime, the effective collision frequency for electrons is much lower than the bounce-frequency. The collisions and a presence of gradients (thermodynamic forces) lead to a generation of a diffusive fluxes in both radial and longitudinal directions (the latter is called the bootstrap current). While the radial fluxes in tokamaks are intrinsically ambipolar, the electron component of the bootstrap current can be quite different from the ion component. Apart from this, due to a permanent existence in tokamaks of the inductive electric field, $E_\parallel$, this current is also needs to be estimated correctly, i.e. with account of the relativistic effects.

Two driving terms are present in (7): the term $\boldsymbol{V}_{dr}\cdot\nabla F_{MJ}$, which describes the radial forces due to the gradients, and the term associated with induced electric field $E_\parallel$. The latter can be accounted by the Spitzer function f_{es}, defined by the equation

$$C_e(f_{es}) = \frac{eE_\parallel}{T_e}\frac{u_\parallel}{\gamma}F_{MJ}. \qquad (9)$$

Regarding now the Spitzer function as a known function, rDKE (7) can be represented now as:

$$\frac{u_\parallel}{\gamma}\nabla_\parallel(f_{e1} - F_e) = C_e(f_{e1} - f_{es}), \qquad (10)$$

$$F_e = -\frac{Iu_\parallel}{\Omega_{ce}}\frac{dF_{MJ}}{d\psi}, \qquad (11)$$

where

$$\frac{dF_{MJ}}{d\psi} = \left[\frac{n_e'}{n_e} - \frac{e\Phi'}{T_e} + \left(\kappa - \frac{3}{2} - \mathcal{R}\right)\frac{T_e'}{T_e}\right]F_{MJ}, \qquad (12)$$

with $A' \equiv dA/d\psi$ and

$$\mathcal{R} = \mu_r\frac{d\ln C_{eMJ}}{d\mu_r} = \frac{15}{8\mu_r} + \mathcal{O}(\mu_r^{-2}). \qquad (13)$$

Note that all definitions obtained so far a very similar to non-relativistic ones, apart from $u_\parallel$ instead of $v_\parallel$ and the additional term $\mathcal{R}$ in (12); see also [6-8].

The electric field $E_\parallel$ no longer appears explicitly, while direct calculation of relativistic Spitzer function can be regarded as the separated problem.

Following the logic used in [4], let's further expand the distribution function f_{e1} in a smallness of the parameter ν_e/ν_b with ν_e as a collisional frequency and $\nu_b \sim v_{te}/2\pi R$ as a bounce frequency: $f_{e1} = f_{e1}^{(0)} + f_{e1}^{(1)} + \cdots$. Then rDKE appears as a pair of equations:

$$\frac{u_\parallel}{\gamma}\nabla_\parallel\left(f_{e1}^{(0)} - F_e\right) = 0, \qquad (14)$$

$$\frac{u_\parallel}{\gamma}\nabla_\parallel f_{e1}^{(1)} = C_e\left(f_{e1}^{(0)} - f_{es}\right). \qquad (15)$$

The first equation in this pair, (14), can be solved by ansatz $f_{e1}^{(0)} = g_e + F_e$, where g_e is the constant at the given magnetic surface labeled by ψ, i.e. $\nabla_\parallel g_e = 0$ and $g_e = g_e(\psi, \mathcal{E}, \mu, \sigma)$ with $\mathcal{E}$, μ, σ as the electron energy, magnetic momentum and the sign of parallel velocity, respectively.

Performing an averaging over the tokamak magnetic surfaces,

$$\langle A\rangle = \int_0^{2\pi}\frac{Ad\vartheta}{\boldsymbol{B}\cdot\nabla\vartheta}\Big/\int_0^{2\pi}\frac{d\vartheta}{\boldsymbol{B}\cdot\nabla\vartheta} \qquad (16)$$

the contributions from trapped electrons are annihilated. As result, the equation for g_e which contains only the contribution from passing electrons is obtained:

$$\langle\frac{\gamma B}{u_\parallel}C_e(g_e + F_e - f_{es})\rangle = 0. \qquad (17)$$

In tokamaks, where the toroidal rotation of both electrons and ions due to the collisional drag of plasma components can be a significant factor and cannot be ignored in transport physics, an accuracy of the models related to the parallel momentum conservation is one of the most important points. Since the linearized collision operator for electrons contains both the "field" part and the integral part, there is no way to solve (17) analytically without some sort of simplification. Fortunately, it can be simplified without significant loss of the physical content if one describes the α/β collisions using a "drifting Maxwellian" model which conserves both the number of particles and the longitudinal momentum:

$$C_{\alpha\beta}(f_{\alpha1}) = \nu_D^{\alpha\beta}(u)\left[\mathcal{L}(f_{\alpha1}) + \frac{m_\alpha u_\parallel}{T_\beta}V_{\beta\parallel}F_{\beta MJ}\right]. \quad (18)$$

Here $V_{\alpha\parallel}$ is the parallel drift velocity for the test particles α and $\nu_D^{\alpha\beta} = (2/u^2)D_{\theta\theta}^{\alpha\beta}(u)$ is the deflection frequency due to pitch scattering on the field particles β; for definition of $D_{\theta\theta}^{\alpha\beta}(u)$ see [11].

Applying (15) for e/e and e/i collisions, the model operator for electrons in this approach can be written s following:

$$C_e(f_{e1}-f_{es}) = \left(\nu_D^{ee}+\nu_D^{ei}\right)\mathcal{L}(f_{e1}-f_{es}) + \frac{m_e u_\parallel}{T_e}\left(\nu_D^{ee}V_{e\parallel}+\nu_D^{ei}V_{i\parallel}\right)F_{MJ}, \quad (19)$$

with

$$\nu_D^{ei}(u) = \nu_{e0}Z_{eff}\frac{\gamma u_{te}^3}{u^3} \quad (20)$$

for e/i deflections; $V_{e\parallel}$ and $V_{i\parallel}$ are the parallel drift velocities of electrons and ions respectively. Here

$$\mathcal{L} = \frac{1}{2}\frac{\partial}{\partial\xi}(1-\xi^2)\frac{\partial}{\partial\xi} \quad (21)$$

is the Lorentz operator with $\xi = u_\parallel/u$ as a pitch. In the following, however, it is more convenient to use instead of ξ the normalized magnetic moment $\lambda = (1-\xi^2)/b$ with $b = B/B_{max}$, where B_{max} is the maximum of B at the given magnetic surface.

Since $\mathcal{L}u_\parallel = -u_\parallel$ and hence $\mathcal{L}F_e = -F_e$, our kinetic problem (16) can now be rewritten as following:

$$\frac{\partial}{\partial\lambda}\lambda\langle u_\parallel\rangle\frac{\partial g_e}{\partial\lambda} = -\frac{u^2}{2}S_e(u,\psi)F_{MJ}, \quad (22)$$

$$S_e(u,\psi) = \frac{bI(\psi)}{\Omega_{ce}}\frac{d\ln F_{MJ}}{d\psi} + \frac{m_e}{T_e}\langle b\,\frac{\nu_D^{ee}V_{e\parallel}+\nu_D^{ei}V_{i\parallel}}{\nu_D^{ee}+\nu_D^{ei}}\rangle + \langle b\hat{f}_s\rangle, \quad (23)$$

with $f_s = u_\parallel \hat{f}_s F_{MJ}$. Note that in steady state tokamak plasmas the toroidal rotation velocities of electrons and ions are usually approximately equal to each other, $V_{e\parallel} \simeq V_{i\parallel}$.

The kinetic problem is now simplified to an ordinary differential equation due to small Larmor radius and the transport ordering. Since g_e vanishes for the trapped particles, we need to consider only the passing particles with $0 \le \lambda \le 1$ (this condition follows directly from the definition of λ). Hence, $g_e = H(1-\lambda)U_\parallel S_e F_{MJ}$, where $H(x)$ is the Heaviside step function. In the passing region $H(x) = 1$, while in the trapped region, $1 < \lambda \le \lambda_{max} \equiv B_{max}/B_{min}$, $H(x) = 0$. For convenience, the value $U_\parallel(\lambda,u,\psi)$ is introduced,

$$U_\parallel(\lambda,u,\psi) = \frac{u^2}{2}\int_\lambda^1\frac{d\lambda'}{\langle u_\parallel(\lambda')\rangle}. \quad (24)$$

In a large-aspect-ratio torus, $R/a \gg 1$, the magnetic field strenght is almost constant, and the trapped-passing boundary is located at $\lambda_{max} \simeq 1+\epsilon$ with $\epsilon = r/R \ll 1$. In this case the quantity $U_\parallel$ is approximately equal to the parallel momentum of electron, $U_\parallel = u_e$, in most of velocity space, i.e. if $\epsilon \to 0$, then $U_\parallel \to u_e$.

Combining all pieces together, we thus obtain the electron distribution function: Combining all pieces together, we thus obtain the electron distribution function,

$$f_{e1} = -\frac{bI(\psi)}{\Omega_{ce}}\left(\frac{u_\parallel}{b} - H(1-\lambda)U_\parallel\right)\frac{dF_{MJ}}{d\psi} + \langle\frac{m_e b}{T_e}\frac{\nu_D^{ee}V_{e\parallel}+\nu_D^{ei}V_{i\parallel}}{\nu_D^{ee}+\nu_D^{ei}} + b\hat{f}_s\rangle H(1-\lambda)U_\parallel F_{MJ}, \quad (25)$$

which is required for calculation of the neoclassical fluxes in tokamaks with relativistic effects accounted.

RADIAL NEOCLASSICAL FLUXES OF RELATIVISTIC ELECTRON IN BANANA REGIME

For calculation of the neoclassical fluxes, we apply the standard definition for both radial and longitudinal fluxes. And most important from physics are the radial heat flux, $\langle\mathbf{q}_e\cdot\nabla\psi\rangle$ and longitudinal electron flux, i.e. the electron component of the bootstrap current, $j_{be} = -e\langle\mathbf{\Gamma}_e\cdot\mathbf{b}\rangle$. Important is that, contrary to stellarators (e.g. see [6,7]), the radial and longitudinal fluxes in tokamaks are coupled.

Generally, the neoclassical relativistic fluxes have to be defined from kinetics with linear plasma response f_{e1} induced by thermodynamic forces, i.e. the radial gradients and the inductive longitudinal electric field. All necessary information has been already obtained and the particle flux can be calculated as:

$$\langle\mathbf{\Gamma}_e\cdot\nabla\psi\rangle = \langle\frac{I(\psi)}{eB}\int m_e u_\parallel C_e(f_{e1}-f_s)d^3u\rangle. \quad (26)$$

Similarly, the radial energy flux is:

$$\langle\mathbf{Q}_e\cdot\nabla\psi\rangle = \langle\frac{I(\psi)T_e}{eB}\int m_e u_\parallel\,\kappa\,C_e(f_{e1}-f_s)d^3u\rangle, \quad (27)$$

where $\kappa = \mu_r(\gamma-1)$ is the normalized kinetic energy of electron (equal to v^2/v_{te}^2 in a non-relativistic limit). Now, extracting the mechanical and advective contributions from the energy flux, the heat flux can be found as $q_e^\psi = Q_e^\psi - T_e\Gamma_e^\psi - W_e V_e^\psi$, where

$$W_e = \int d^3u\, m_e c^2(\gamma-1)F_{MJ} = \left(\frac{3}{2}+\mathcal{R}\right)n_e T_e \quad (28)$$

is the energy density related to the Maxwell-Jüttner thermal equilibrium, and $V_e^\psi = \Gamma_e^\psi/n_e$ is the radial flow velocity. Finally, the relativistic heat flux is [6,7]:

$$\langle\mathbf{q}_e\cdot\nabla\psi\rangle = \langle\mathbf{Q}_e\cdot\nabla\psi\rangle - \left(\frac{5}{2}+\mathcal{R}\right)T_e\langle\mathbf{\Gamma}_e\cdot\nabla\psi\rangle. \quad (29)$$

The particle and heat neoclassical radial fluxes can be calculated from (26) and (28) using relativistic collisional operator. However, obtained results are somewhat cumbersome, so it is more illustrative to look at the limit of large Z, which corresponds to pitch-angle scattering of electrons (Lorentz term in collisional operator) and $\nu_D^{ee} \ll \nu_D^{ei}$. In this particular limit the fluxes can be calculated and then expanded by the order of μ_r^{-1}. The first term in such expansion is the non-relativistic limit (NR), and the second term is the first-order relativistic correction (rel).
For the radial particle flux we obtain:

$$\langle\mathbf{\Gamma}_e\cdot\nabla\psi\rangle \approx \langle\mathbf{\Gamma}_e{}^{NR}\cdot\nabla\psi\rangle + \frac{1}{\mu_r}\langle\delta\mathbf{\Gamma}_e{}^{rel}\cdot\nabla\psi\rangle, \quad (30)$$

$$\langle \delta \mathbf{\Gamma}_e{}^{rel} \cdot \nabla\psi \rangle = -f_t \frac{n_e T_e I^2}{m_e \Omega_{ce}^2 \tau_{ei}} \times \left[\left(1 + \frac{T_i}{Z T_e}\right) \frac{n_e'}{n_e} - \frac{5}{4} \frac{T_e'}{T_e} - \frac{0.173}{Z T_e} \frac{T_i'}{T_i} \right], \quad (31)$$

and for the radial heat flux we obtain:

$$\langle \mathbf{q}_e \cdot \nabla\psi \rangle = \langle \mathbf{q}_e{}^{NR} \cdot \nabla\psi \rangle + \frac{1}{\mu_r} \langle \delta \mathbf{q}_e{}^{rel} \cdot \nabla\psi \rangle, \quad (32)$$

$$\langle \delta \mathbf{q}_e{}^{rel} \cdot \nabla\psi \rangle = -f_t \frac{n_e T_e^2 I^2}{m_e \Omega_{ce}^2 \tau_{ei}} \times \left[-\frac{9}{4} \left(1 + \frac{T_i}{Z T_e}\right) \frac{n_e'}{n_e} + \frac{29}{4} \frac{T_e'}{T_e} + \frac{39}{Z T_e} \frac{T_i'}{T_i} \right]. \quad (33)$$

Here f_t is the "effective fraction of trapped particles", $f_t \simeq 1.46\sqrt{\epsilon}$, and respective non-relativistic radial fluxes can be found elsewhere (for example, see [4]).

SUMMARY

In the present work, relativistic effects in the electron neoclassical transport in tokamaks were considered. Only the low collisional "banana" regime was investigated. Starting from the general form of guiding center equations and the drift kinetic equations with the relativistic corrections taken into account, the standard method for calculation of particle and heat radial fluxes was revisited. As the main result, the relativistic corrections for the radial neoclassical electron fluxes were obtained. The Lorentz invariance is lost, however the advantage of the approach used is a possibility to include the obtained formulations in any standard transport code based on non-relativistic transport theory.

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О РЕЛЯТИВИСТСКИХ ВЛИЯНИЯХ НА ЭЛЕКТРОННЫЙ ТРАНСПОРТ В БАНАНОВОМ РЕЖИМЕ В ТОКАМАКАХ

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Неоклассическая теория переноса в токамаке пересматривается с учетом релятивистских эффектов для электронов. Поскольку такие эффекты важны только в высокотемпературной плазме, был рассмотрен только слабостолкновительный банановый режим. Полученные формулировки дают возможность рассчитать электронные неоклассические потоки в очень широком диапазоне температур.

ПРО РЕЛЯТИВІСТСЬКИЙ ВПЛИВ НА ЕЛЕКТРОННИЙ ТРАНСПОРТ У БАНАНОВОМУ РЕЖИМІ В ТОКАМАКАХ

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Неокласична теорія транспорту в токамаках розглядається з урахуванням релятивістських ефектів для електронів. Оскільки такі ефекти важливі лише у високотемпературній плазмі, розглянуто лише слабозіткневий банановий режим. Отримані формулювання дають можливість обчислювати електронні неокласичні потоки в дуже широкому діапазоні температур.