Effect of a radial electric field on the transition to improved confinement in a tokamak

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When there is a sufficient number of banana particles in the plasma near the wall in a tokamak, or when a radial potential difference (of arbitrary polarity) is deliberately applied, the radial electric field near the periphery changes abruptly. The rapid rotation associated with this effect and the shear of this rotation can suppress both anomalous and neoclassical transport, causing a transition to the improved confinement mode (the *L–H* transition).

Several pieces of direct experimental evidence found recently point to a relationship between the radial electric field and transport processes. A sharp (severalfold) increase in the poloidal rotation velocity upon the *L–H* transition was recently observed in the *D* III–*D* tokamak.[1] The imposition of a strong negative radial field near the wall by means of a charged electrode in the CCT tokamak[2] has caused the *L–H* transition. The imposition of a positive field in the TEXT tokamak[3] has done the same.

The mechanism proposed below for the onset of a strong radial field involves incorporating in the conservation equation for the longitudinal moment [or in the (equivalent) condition that the radial flux of particles be ambipolar] some additional terms which stem from spatial inertia in anomalous transport in a radially nonuniform electric field and also from an anomalous viscosity.

In the neoclassical theory, the velocity of the poloidal rotation is found by setting the poloidal ion viscosity, averaged over a magnetic surface, $\langle \vec{B}\cdot(\nabla\cdot\overleftrightarrow{\pi})\rangle^{(\mathrm{NEO})}=0$, equal to zero. The neoclassical value of the poloidal revolution velocity of the ions is[2)]

$$v_p^{(\mathrm{NEO})}=v_0^{(\mathrm{NEO})}+u_{pi},$$

$$v_0^{(\mathrm{NEO})}=-\frac{cE_r}{B}=-\frac{cT_i}{eB}\left(\frac{d\ln n}{dr}+k_T\frac{d\ln T_i}{dr}\right), \tag{1}$$

$$u_{pi}=\frac{c}{eBn}\frac{d(nT_i)}{dr},\quad v_p^{(\mathrm{NEO})}=(1-k_T)\frac{c}{eB}\frac{dT_i}{dr},$$

where n is the density, T_i is the ion temperature, and k_T is a numerical factor which depends on the collisionality parameter (in the banana regime, we have $k_T=-0.17$). A deviation of v_p from $v_p^{(\mathrm{NEO})}$ results in a temporal relaxation of the poloidal rotation velocity, with a time scale $\tau\sim\nu_i^{-1}$ (Refs. 4 and 5; ν_i is the ion collision rate) to the value in (1). In the course of this relaxation, the longitudinal viscosity is balanced by the longitudinal inertia, which is proportional to $\partial v_p/\partial t$. When the real anomalous transport is taken into account, on the other hand, additional terms appear in the equation for the longitudinal balance of forces even in the steady-state case:

$$< nm_i\vec{B}\cdot(\vec{u}\cdot\nabla)\vec{u} >= - < \vec{B}\cdot(\nabla\cdot\overleftrightarrow{\pi}) >^{(AN)} - < \vec{B}\cdot(\nabla\cdot\overleftrightarrow{\pi}) >^{(\mathrm{NEO})}, \qquad (2)$$

where $\vec{u}$ is the hydrodynamic velocity of the ions. For simplicity we assume that the magnetic surfaces are nearly circular: $B_t = B_0/(1+\epsilon\cos\vartheta)$, $B_p = \theta(r)B_t$, $\epsilon = r/R \ll 1$. In contrast with the standard neoclassical approach, we incorporate in (2) the anomalous radial transport velocity of the particles, u_r ($u_r \sim D d\ln n/dr$). The term with the anomalous viscosity is

$$- < \vec{B}\cdot(\nabla\cdot\overleftrightarrow{\pi}) >^{(\mathrm{AN})} = < B\,\nabla_\perp\cdot(n\eta\,\nabla_\perp\, u_\parallel(\vartheta)) >, \qquad (3)$$

where η is the anomalous viscosity coefficient. Finding an expression for the hydrodynamic velocity of the ions from the drift kinetic equation, substituting $u_\parallel(\vartheta)$ into (2) and (3), and taking an average over a magnetic surface, we find expressions for the first two terms in (2). When there is an anomalous transport, the neoclassical viscosity in the banana regime contains additional terms of the spatial-inertia type.[6] An expression for $\langle\vec{B}\cdot(\nabla\cdot\overleftrightarrow{\pi})\rangle^{(\mathrm{NEO})}$ can be found by analogy with the case[4,5] of a time variation of v_p, for which a corresponding mechanism was proposed in Ref. 7. Under the condition $|v_p| < \theta c_s$, where $c_s = (2T_i/m_i)^{1/2}$, we have $u_\parallel(\vartheta) = -2\epsilon v_p\cos\vartheta/\theta + \theta v_p$, and we can put Eq. (2) in the form

$$-nm_i(1+2q^2)B_0\theta u_r\frac{dv_p}{dr} = -(1+2q^2)\frac{\theta B_0}{r}\frac{d}{dr}\left(rn\eta\frac{dv_p}{dr}\right)$$
$$+\frac{nm_i\sqrt{\epsilon}B_0u_r}{\theta}\frac{dv_p}{dr} + \frac{1,1B_0nm_i\nu_i\sqrt{\epsilon}}{\theta}(v_p - v_p^{(\mathrm{NEO})}), \qquad (4)$$

where $q = \epsilon/\theta$ is the safety factor. From (4) we have

$$\frac{dv_p}{dr} + \alpha(r)(v_p - v_p^{(\mathrm{NEO})}) = \frac{\beta(r)}{m_i\Gamma(r)r}\frac{d}{dr}\left(rn\eta\frac{dv_p}{dr}\right). \qquad (5)$$

Here $\Gamma = nu_r$ is the radial particle flux,

$$\alpha(r) = \frac{1,1\nu_i}{u_r}\left[1+\frac{1+2q^2}{q^2}\epsilon^{3/2}\right]^{-1}, \quad \beta(r) = \left[1+\frac{q^2}{(1+2q^2)\epsilon^{3/2}}\right]^{-1}. \qquad (6)$$

The quantities $\alpha(r)$ and $\beta(r)$ are generally functions of v_p. The quantity $\alpha(a)$ (a is the minor radius) can be expressed in terms of the particular confinement time τ_p: $\alpha(a) = 2n(a)\tau_p\nu_i/\bar{n}a$, where $\bar{n}$ is the average density. Despite the condition $\alpha(a) \gg a^{-1}$, the first and second terms in (5) may be comparable in magnitude upon a sharp change in v_p. (In the *D* III-*D* tokamak, for example, the poloidal rotation velocity near the wall varies with a length scale $L_{v_p} \sim 1$ cm in the *H* mode.) For $\eta/m_i \sim D$, the last term in (5) is on the order of $v_p\epsilon^{-3/2}L_n/L_{v_p}$, where L_n is the length scale of the density variation.

It can be seen from (5) that two types of poloidal rotation profiles are possible near the wall. The first is a smooth v_p profile, with the first and third terms in (5) being unimportant, and with the velocity v_p being equal to $v_p^{(\mathrm{NEO})}$ over the entire distance to the separatrix. Such a profile corresponds to the L confinement mode. In the second case, the neoclassical ion viscosity is balanced by spatial inertia and the anomalous viscosity, and the shear in v_p is large near the separatrix. In this case the shear suppresses the onset of turbulence by the mechanism of Ref. 8 and reduces the radial transport, as is typical of the H mode. The inertial term in (5) causes v_p to increase from the periphery toward the center [if $\alpha(r) = \mathrm{const}$, the increase is exponential] with a length scale $L_{v_p} = \alpha^{-1} \ll a$; the anomalous viscosity causes a spatial relaxation of v_p to the value $v_p^{(\mathrm{NEO})}(r)$. The relation between the two depends on the parameter $L_n \epsilon^{-3/2}/L_{v_p}$ [or L_n/L_{v_p} in the plateau regime, since the third term in (4) is not present]. The specific v_p profile can be constructed from (5).

The transition to the H mode can be caused by the following factors: (1) If there is a sufficient number of banana particles near the separatrix, a negative electric field will be set up and will prevent a penetration of ion bananas beyond the separatrix. The corresponding value of $v_p(a)$ is[9,7] $\sim\sqrt{\epsilon\theta c_s}$. A strong electric field is deliberately set up near the wall, by means of (for example) a charged electrode.[2,3] (3) A high gradient of the density and the temperature (larger than $dv_p^{(\mathrm{NEO})}/dr$) arises near the wall. In all these cases, a rapidly varying v_p arises, according to (5). The effect should be a suppression of turbulence.

At rotation velocities $|v_p| \gtrsim \theta c_s$, which can arise near the wall according to (5), factors $\sim\exp(-v_0^2/\theta^2 c_s^2)$, appear in the last two terms in (4). These factors correspond to an exponential decrease in the number of trapped particles, and the quantity $u_\parallel(\vartheta)$ is related to v_p in a nonlinear way. This case will be analyzed in a separate paper.

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[2)] Since the projection of the average toroidal rotation velocity along $\vec{B}$ onto the poloidal direction is small in all experiments (apparently because of the anomalous skew viscosity), we are ignoring its contribution to v_p.

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