Spontaneous spin-up induced by turbulence-driven topological transition of orbits in a tokamak plasma

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Spontaneous spin-up are widely observed in tokamak plasmas, which is crucially important for plasma confinement. A kinetic theory is proposed to show that a toroidal rotation of core plasma is induced by the topological transition of orbits driven by turbulent diffusion in a collisionless tokamak plasma. The theoretical prediction agrees well with the well-known Rice-scaling of intrinsic core plasma flow. This new theory predicts an intrinsic core parallel flow of $\sim 100\text{km/s}$ for the International Thermonuclear Experimental Reactor.

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Spontaneous spin-up is widely observed, from particles to galaxies, for examples, the well-known giant zonal belt on the Jupiter \cite{1}. In a tokamak, a magnetic fusion torus in the shape of a donut, spontaneous toroidal spin-up of core plasma is routinely observed after the Low-confinement to High-confinement (L-H) transition \cite{2-8}. To understand the physical mechanism of this so-called intrinsic flow, various theoretical models have been proposed: the effect of the residual Reynolds stress \cite{9,10} on the momentum redistribution \cite{11}, the effect of turbulence intensity gradient \cite{12} and the effect of thermal ion orbit loss \cite{13-16} on the boundary flow, the effect of Coriolis force on the momentum pinch \cite{17,18}, and the turbulent acceleration \cite{19,20}. However, it is still hard to predict the net core flow \cite{11} in ITER. In this paper, we show that a spontaneous toroidal flow can be induced by the topological transition of ion orbit in a collisionless turbulent tokamak plasma.

The equilibrium magnetic field of a tokamak is written as \( B = I(r) \nabla \zeta + \nabla \zeta \times \nabla \psi(r) \), with \( I = RB_T \), \( B_T \) the toroidal magnetic field, \( R \) the major radius. The poloidal magnetic flux is \( \psi(r) \), with \( r \) essentially the minor radius of the torus. The poloidal magnetic field is given by \( B_P = \psi'/R \), with the prime denoting the derivative with respect to \( r \). \( \zeta \) is the toroidal angle. In this paper, we shall consider a large-aspect-ratio (\( \epsilon \equiv r/R \ll 1 \)) up-down symmetric tokamak.

The ensemble averaged distribution function is \( f(\psi, \theta, \mu, w, t) \), with \( \theta \) the poloidal angle, \( \mu = v_\perp^2/2B \) the magnetic moment, and \( w = v^2/2 + e\phi(\psi)/m \) the energy of the particle; \( v_\perp \) and \( v_\parallel \) are the velocity components perpendicular and parallel to the magnetic field, respectively. \( \phi \) is the electrostatic potential. \( e \) and \( m \) are the charge and mass of the particle, respectively. Since the momentum of plasma is mainly carried by the ions, we consider only the ion dynamics in the following.

We begin with the transport equation,

\[
\partial_t f + \mathcal{L}(f) + \mathcal{T}(f) = S,
\]

with \( S \) the source term; the radial transport term including the effect of turbulence \cite{12,21-23} is given by

\[
\mathcal{T}(f) = -\frac{1}{J} \partial_\psi (J D \partial_\psi f),
\]

where \( D = (\psi')^2 D \), with \( D \) the usual radial diffusivity due to the turbulence. Collisions are ignored since we are concentrating here on the H-mode plasma \cite{12}. 
\[
\mathcal{L}(f) \equiv v_\| \mathbf{b} \cdot \nabla f + \mathbf{V}_d \cdot \nabla f,
\]
(3)

with \( \mathbf{b} = \mathbf{B}/B \), and the guiding-center drift velocity is given by the Alfvén approximation \[24\], \( \mathbf{V}_d = -v_\| \mathbf{b} \times \nabla \rho_\| \), with \( \rho_\| = mv_\|/eB \). The Jacobian of the phase space is given by \( J^{-1} = |v_\| \mathbf{b} \cdot \nabla \theta|/(2\pi) \).

To concentrate on the intrinsic rotation, we shall assume that there is no momentum injection, therefore \( S = S(\psi, \mu, w) \) shall be understood. Similarly, we shall also assume that \( \mathcal{D} = \mathcal{D}(\psi, \mu, w) \). In the following, a shorthand notation \( \mathcal{D}(\psi) \) and \( S(\psi) \) shall be used without confusion.

To solve Eq. (1), we transform to the coordinate system, \((P, \mu, w, \theta)\), with

\[
P = \psi - \rho_\| I,
\]
(4)

the toroidal canonical angular momentum \((-eP)\). In this new coordinate system,

\[
\mathcal{L} = \dot{\theta} \partial_\theta,
\]
(5)

with the poloidal angular velocity given by

\[
\dot{\theta} = v_\| \mathbf{b} \cdot \nabla \theta + \mathbf{V}_d \cdot \nabla \theta = v_\| \mathbf{b} \cdot \nabla \theta \partial_\psi P(\psi, \theta, \mu, w).
\]
(6)

Using Eq. (6), one finds

\[
\mathcal{T}(f) = -\frac{1}{\mathcal{J}} \partial_P (\mathcal{J} \mathcal{D} \partial_P f),
\]
(7)

with the Jacobian of the new coordinate system given by \( \mathcal{J}^{-1} = |\dot{\theta}|/(2\pi) \), which is found by using Eq. (6).

Note that \( \partial_t f \sim S \sim \mathcal{T} \sim f/\tau_E \), with \( \tau_E \) the confinement time; \( \mathcal{L} \sim 1/\tau_\theta \), with \( \tau_\theta = \oint d\theta/\dot{\theta} \), the bounce time for trapped particles or the transit time for passing particles. For passing particles, \( \oint d\theta = \int_0^{2\pi} d\theta \); for trapped particles, \( \oint d\theta = \int_{-\theta_b}^{\theta_b} d\theta + \int_{+\theta_b}^{-\theta_b} d\theta \), with \( \pm \theta_b \) the bounce angle. Note that \( \sigma \), the sign of \( v_\| \), is a constant of motion for passing particles, while it is not for trapped particles.

To proceed, we assume that \( \delta \equiv \tau_\theta/\tau_E \ll 1 \). Expanding the distribution function with respect to \( \delta \), we found \( f = F + \delta f \). To \( \mathcal{O}(\delta^0) \), one finds \( \mathcal{L}(F) = 0 \), which demonstrates that the lowest order solution is a constant of motion, \( F = F(P, \mu, w) \).
To the next order, one finds

$$\mathcal{L} (\delta f) + \partial_t F + \mathcal{T} (F) = S.$$  \hspace{1cm} (8)

Orbital-averaging this equation, one finds

$$\partial_t F - \frac{1}{\tau_\theta} \partial_P \left[ \tau_\theta D \left( \overline{\psi} \right) \partial_P F \right] = S \left( \overline{\psi} \right),$$  \hspace{1cm} (9)

which includes the finite-banana-width effects. Note that the phase space element in terms of \((P, \mu, w)\) is \(2\pi dP \times 2\pi d\mu \times \tau_\theta dw; \ \tau_\theta\) is the Jacobian of the constants-of-motion space \((P, \mu, w)\).

The orbital-averaging operator, which is an annihilator of \(\mathcal{L}\), is defined by

$$\overline{A} = \frac{1}{\tau_\theta} \oint d\theta \dot{\theta} A.$$  \hspace{1cm} (10)

Using Eq. (4), one finds

$$\overline{\psi} - \psi = (\rho_\parallel I - \rho_\parallel I) ;$$  \hspace{1cm} (11)

clearly, \(\overline{\psi} = \overline{\psi}_\sigma (P, \mu, w)\). Eq. (9) should be solved for co-current passing \((\sigma = +)\), counter-current passing \((\sigma = -)\), and trapped particles.

To proceed, we make general comments on Eq. (9). Since the orbital time is much longer than the diffusion time, it is essentially the guiding-center orbit rather than the particle itself that is diffused by the turbulence. When the passing particles, which carry the parallel momentum in the core, diffuse outward, both the co-passing orbit and the counter-passing orbit may undergo a topological transition into trapped orbit; this topological transition of orbits, which are specified by the three constants of motion \(25\), \((P, \mu, w)\), is schematically shown in Fig. 1.

The spontaneous co-current spin-up of core plasma after L-H transition of a tokamak can be explained as follows. After L-H transition, the plasma enters the collisionless state, and particles in the boundary region can complete their trapping (banana) orbit, which dominates the orbit loss; when the free-moving co/counter-passing particles diffuse from the core to the boundary region, a co-passing orbit or a counter-passing orbit may change to a trapped orbit through the topological transition; therefore, the co-moving particles have larger confinement region \([P_3, P_0]\) in Fig. 1] than the counter-moving particles \([P_2, P_0]\) in Fig. 1]; this \(\sigma\) asymmetry induced by the topological transition clearly generates a co-current flow in the core region. We note here that this spontaneous spin-up of core plasma
FIG. 1. Guiding-center orbits with given $(\mu, w)$ in $R - \psi$ space are horizontal hyperbolas, with their tips given by $[R_b(\mu, w), P_i]$; $\psi_E$ is the curve $\psi = \psi(R, \theta = 0)$ defined by equilibrium field; a trapped/passing particle has 2/1 (a topological number) cross points with the right branch of the curve $\psi_E$. Diffusion of co-passing orbits from $(P_3, +)$ to $(P_1, +)$ and from $(P_2, -)$ to $(P_1, -)$ induces a topological transition at the trapping-passing boundary, from $(P_1, +)$ to $(P_0, T)$.

in a turbulent tokamak induced by the topological transition of orbits may be taken as a classical counterpart of topological phase transition in condensed matter physics \[26, 27\].

However, when directly solving Eq. (9) for trapped particles, co-counter-passing particles, this topological transition incurs an mathematical inconvenience, and a connection formula must be introduced in the trapping-passing boundary in the phase space; this inconvenience can be avoided by the following modification.

Define $F_\sigma = F$. The Jacobian of the constant-of-motion space is changed from $\tau_\theta$ to $\tau_{\theta, \sigma}$ accordingly. For passing particles, $\tau_{\theta, \sigma} = \tau_\theta$. For trapped particles, $\tau_{\theta, -}$ ($\tau_{\theta, +}$), the inner (outer) half period of trapped particles is given by

$$\tau_{\theta, \pm} = \int_{\pm \theta_b}^{\pm \theta_b} d\theta / \dot{\theta}.$$  \hfill (12)

Thus Eq. (9) can be re-written as

$$\partial_t F_\sigma (\psi, \mu, w, t) - \frac{1}{\tau_{\theta, \sigma}} \partial_P \left[ \tau_{\theta, \sigma} D (\psi) \partial_P F_\sigma \right] = S (\psi),$$  \hfill (13)

which shall be solved for $\sigma = \pm$, with the orbital-average prescribed as before.

Let $\psi_O (P, \mu, w)$ be the outermost radial position that a particle labeled by $(\mu, w, \sigma)$ launched from $(\psi, \theta)$ can reach. The conservation of the canonical toroidal angular momentum and the energy demands,

$$\psi - \rho_\parallel (\psi, \theta, \mu, w, \sigma) I = \psi_O - \rho_\parallel O I = P,$$  \hfill (14)
\[
\frac{1}{2}v^2 + \mu B (\psi, \theta) + \frac{e}{m} \phi (\psi) = \frac{1}{2}v^2_\parallel + \mu B (\psi, 0) + \frac{e}{m} \phi (\psi_\circ) = w. \quad (15)
\]

Clearly, the above two equations define a function \( \psi_L \) through \( \psi_L (\psi_\circ, \theta, \mu, w) = \psi \); therefore, the orbit loss condition is given by

\[
\psi \geq \psi_L (\psi_\circ, \theta, \mu, w, \sigma)_{\psi_\circ=\psi_a}, \quad (16)
\]

with \( \psi_a \) the poloidal magnetic flux at the boundary of the torus. Clearly, a boundary condition

\[
[f]_{\psi \geq \psi_L} = 0, \quad (17)
\]

should be associated with Eq. (1).

Note that \( \psi_\circ \) is also a constant of motion; the above discussions also define the orbit loss condition in an alternative form,

\[
P \geq P_L (\mu, w, \sigma), \quad (18)
\]

which specifies the outer boundary condition of Eq. (13),

\[
[F_\sigma]_{P \geq P_L} = 0. \quad (19)
\]

We observed that \( \bar{\psi} (P_{\min}, \mu, w, \sigma) \), with \( P_{\min} \) the minimum value of \( P \), is weakly dependent on \( \sigma \) (see, Fig. 1). The inner boundary condition for Eq. (13) is given by

\[
[\tau_{\theta, \sigma} D (\bar{\psi}) \partial_P F_\sigma]_{P=P_{\min}} = 0, \quad (20)
\]

which is similar to the natural boundary condition used previously in neoclassical transport theory [28].

Eq. (13) is readily integrated with the above boundary conditions. The steady-state solution is given by

\[
h = -\frac{1}{\tau_{\theta, \sigma} D (\bar{\psi})} \int_{\bar{\psi}(P_{\min})}^{\bar{\psi}(P)} \partial_P \bar{\psi} \bar{\psi} \tau_{\theta, \sigma} S (\bar{\psi}), \quad (21)
\]

\[
F_\sigma [\bar{\psi} (P, \mu, w)] = \int_{\bar{\psi}(P_{\min})}^{\bar{\psi}(P)} \partial_P \bar{\psi} \bar{\psi} \partial_P h \left[ \bar{\psi} (P, \mu, w), \mu, w \right], \quad (22)
\]

Clearly, this solution is weakly anisotropic in the core region where \( P < P_L \), while it is strongly anisotropic in the boundary region where \( P \geq P_L \). The width of the boundary region is given by

\[
\Delta \psi_B = \psi_a - \psi_B \sim |\rho| I. \quad (23)
\]
Since the orbit loss is dominated by the initially counter-moving trapped particles, it can be estimated that
\[- \partial_\psi T_B \sqrt{\epsilon v_{th,B}} I/\Omega = T_B,\]  
with \(\partial_\psi T_B\) the radial gradient of temperature in the boundary region, and \(T = \frac{1}{2}mv_{th}^2\).

The temperature \(T\) is given by \(p = nT = \int d^3v \frac{1}{2}mv^2F\), with \(n = \int d^3v F\) the density, and \(p\) the pressure. Note that \(\Delta \psi_B \sim I \sqrt{\epsilon v_{th,B}}/\Omega\) is the half banana-width of the typical trapped particles launched from the typical orbit-loss boundary labeled by \(\psi_B\), where the ion temperature is \(T_B = \frac{1}{2}mv_{th}^2\). Note that \(T_B\) is typically the pedestal temperature of the H-mode plasma.

Therefore, one finds the typical parallel velocity of the trapped particle at the orbit-loss boundary
\[\sqrt{\epsilon v_{th,B}} = -2a_e \partial_\psi T_B.\]  

The ion parallel flow, \(u\), is defined by
\[nu = \left\langle \int d^3v v_\parallel F_\sigma \right\rangle.\]  

Note that
\[d^3v = \sum_\sigma \frac{2\pi B}{|v_\parallel|} d\mu dw.\]  

The magnetic-flux-surface average operator is given by
\[\langle A \rangle = \frac{1}{\mathcal{F}} \oint \frac{d\theta}{B \cdot \nabla_\theta} A.\]  

The parallel momentum due to the upper/lower end of integral in Eq. \(22\) can be found in the following way. Define
\[F_{\text{ped}} = -\int_{\psi(P_{\text{max}})}^{\psi(P_L)} \partial_\psi P d\psi h \left[ \psi(P,\mu,w) , \mu,w \right],\]  

with \(P_{\text{max}} = \psi_a - \rho_\parallel (\psi,\theta,\mu,w) I\). Clearly \(F = F_\sigma - F_{\text{ped}}\) represents the solution with the boundary condition \(f|_{\psi=\psi_a} = 0\), which ignores the effect of ion orbit-loss and therefore contains a the weak anisotropy; the strong anisotropy due to the orbit-loss boundary condition is contained in \(F_{\text{ped}}\), which can be taken as a "pedestal" of anisotropy. This pedestal of distribution anisotropy is schematically shown in Fig. 2.
FIG. 2. Pedestal of the distribution anisotropy.

The core momentum due to the weak anisotropy contained in $F$ can be evaluated in a straightforward way by using the method previously developed [29, 30],

$$\langle \int d^3 v v_{\|} F \rangle = -1.6 \frac{e^{3/2}}{\Omega_p} (-n e \partial_r \phi - \partial_r p),$$ (30)

which may be ignored; by using the ion radial force balance equation, one finds that this equation gives a small correction term to the toroidal rotation [30].

The parallel momentum “pedestal” due to the lower end of integral in Eq. (22), which represents the effect of boundary trapped ion orbit loss on core passing ion through the topological transition, can be evaluated as

$$[\text{nmu}]_{\text{ped}} = \langle \int d^3 v v_{\|} F_{\text{ped}} \rangle \sim -\epsilon^{3/2} \frac{1}{\Omega_p} n \partial_r T_B,$$ (31)

where $\Omega_p = \Omega B_p / B$, and we have used Eq. (25).

Clearly the core plasma momentum is approximately

$$[\text{nmu}]_{\psi \leq \psi_B} = [\text{nmu}]_{\text{ped}} + [\text{nmu}]_c \sim -\epsilon^{3/2} \frac{1}{\Omega_p} n \partial_r T_B.$$ (32)

The pedestal structure found in this paper is schematically shown in Fig. 3.

For the DIII-D experimental observation [5] of the anomalous co-current momentum source at the edge of the Deuteron H-mode plasma. The main parameters of the DIII-D experiments are $R/a = 1.6m / 0.6m$; $n \sim 5 \times 10^{19} / m^3$; the poloidal magnetic field at the edge is $\sim 0.15 T$esla. With these parameters, following the above analysis, $T_B \sim 1 keV$ is estimated. The intrinsic parallel flow estimated by using the above theory and parameters is $70 km/s$, which agrees well with the experimental observations [5].

Eq. (32) predicts a scaling of the intrinsic parallel flow of the core plasma

$$u \sim \epsilon^{3/2} T / I_p,$$ (33)
with $I_p$ the plasma current. This is consistent with the well-known Rice-scaling \cite{3} of intrinsic toroidal rotation of core plasma, $u \sim p/I_p$.

For a typical ITER plasma, $R/a = 6.2m/2m$, $B_T = 5.3Tesla$, $B_P = 0.3Tesla$ at the edge. $T_B \sim 3keV$ may be estimated. Eq. (32) predicts an intrinsic core parallel flow of $\sim 100km/s$.

In conclusion, by solving the transport equation with the ion orbit loss boundary condition, we have proposed a kinetic theory of spontaneous core parallel flow induced by turbulence-driven topological transition of orbits in a tokamak H-mode plasma. The proposed theory is consistent with the well-known Rice-scaling \cite{3} of intrinsic core plasma flow; it predicts a $\sim 100km/s$ intrinsic parallel flow for a typical ITER core plasma. The key point of the proposed mechanism of spontaneous core plasma spin-up in an H-mode tokamak is as follows. Although the orbit loss is dominated by the trapped particles in the boundary region, it affects the distribution of passing particles, which carry the parallel momentum, in the core region; when a co-passing orbit or a counter-passing orbit diffuse to the boundary region, it may change to a trapped orbit through the topological transition. In this way, an asymmetric confinement of passing ions is induced, therefore a spontaneous toroidal flow is maintained, since the confinement region of the co-passing ions is larger than the counter-passing ions.
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