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NONAMBIPOLARITY, ORTHOGONAL CONDUCTIVITY, POLOIDAL FLOW, AND TORQUE

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Abstract

Nonambipolar processes, such as neutral injection onto trapped orbits or ripple-diffusion loss of α-particles, act to charge a plasma. A current j_r across magnetic surfaces must arise in the bulk plasma to maintain charge neutrality. An axisymmetric, neoclassical model of the bulk plasma shows that these currents are carried by the ions and exert a $j_r B_\theta R/c$ torque in the toroidal direction. A driven poloidal flow $V_\theta = E_r'c/B$ must also develop. The average current density $\langle j_r \rangle$ is related to the radial electric field $E_r' = E_r + v_\phi B_\theta/c$ in a frame moving with the plasma via the orthogonal conductivity $\langle j_r \rangle = \sigma_\perp E_r'$, which has the value $\sigma_\perp = (1.65\epsilon^{1/2})(ne^2\nu_{ii}/M\Omega_\theta{}^2)$ in the banana regime. If an ignited plasma loses an appreciable fraction Δ of its thermonuclear α-particles by banana ripple diffusion, then the torque will spin the plasma to sonic rotation in a time $\tau_s = 2\tau_E/\Delta$, τ_E being the energy confinement time.

I. INTRODUCTION

To what extent will nonambipolar processes charge up a tokamak plasma? A real tokamak plasma can be subject to a variety of nonambipolar driving processes: neutral injection onto trapped orbits,[1] ripple-diffusion loss of trapped α-particles,[2] and a direct orbit loss of fast ions. Clearly, a radial current must arise in the bulk plasma to render the total system ambipolar. Our key approximation is to suppose that the bulk plasma is described by an axisymmetric, neoclassical model even though deviations from axisymmetry may be responsible for energetic particle losses. We show that the plasma cannot be in a steady state if it is to generate a radial current; it must be experiencing an angular acceleration resulting from a torque. Equivalently, an empirical toroidal drag force can be introduced, which acts equally on all ions.

Our goal is to develop a simple neoclassical model of a tokamak plasma experiencing an acceleration in the azimuthal direction, and to show that under these circumstances, a radial current develops in the ions as a result of ion-ion collisions. Driven poloidal rotation also arises and we show this velocity to be small and dynamically unimportant. Indeed, the first effect of a nonambipolar process is to create a radial electric field, which implies poloidal rotation. But poloidal rotation is known to decay and it is this decay that exerts a toroidal torque. Goldston et al.[3] have addressed these issues via trapped and circulating particle fluids.

There are two ways to think about the radial current: as a neoclassical polarization current or as a conduction current $j_r = \sigma_\perp E_r'$, where E_r' is the (small) electric field in a frame toroidally rotating with the plasma

$$E_r' = E_r + \omega R\, B_p/C \tag{1}$$

and $\sigma_\perp$ is the orthogonal conductivity.[4] Let dT^* be the torque per toroidal radian exerted on a plasma between magnetic surface ψ and $\psi + d\psi$, where ψ is the poloidal flux. Angular momentum conservation requires

$$dT^* = d\psi \int \frac{Rd\ell}{\nabla\psi} \left(\vec{J}_r \cdot \frac{\vec{\nabla}\psi}{Rc}\right) R = d\psi \int \frac{Rd\ell}{(\nabla\psi)} nMR^2 c \frac{\partial\dot{\Phi}}{\partial\psi} \quad , \tag{2}$$

where

$$\omega = c \frac{\partial\Phi}{\partial\psi} = - cE_r/\nabla\psi = \text{toroidal angular rotation rate} \quad , \tag{3}$$

$$B = \frac{I\hat{\phi}}{R} + + \frac{\hat{\nabla}\psi\times\hat{\phi}}{R} = B_r\hat{\phi} + \vec{B}_p \quad , \tag{4}$$

and ℓ is a poloidal coordinate in the magnetic surface. One can recast (2) to read

$$\frac{dT^*}{d\psi} = - \frac{J_r}{c} = - \int dS \frac{nMc}{B_p^2} \dot{E}_r \quad ,\tag{5}$$

where $dS = rd\ell$ is an area element in the magnetic surface. It is form (5) that interprets the radial current as a neoclassical polarization drift and is a consequence of global angular momentum conservation. This work shows that the plasma accomplishes angular momentum conservation by developing a radial field E_r' in the frame of toroidal rotation and a consequent forced poloidal rotation $v_p = - E_r' c/B_T$.

II. DRIFT-KINETIC EQUATION

Neoclassical processes are governed by an axisymmetric drift-kinetic equation. This work departs from previous developments[5] by permitting the plasma to be accelerating in angular rotation speed ω and to have a forced poloidal rotation. The first task will be to transform the standard ion drift kinetic equation to the instantaneous rotating frame. Straightforward algebra then reduces this equation to the mathematical form of a previously solved neoclassical equation. Equating the radial current in the solution to be equal and opposite to the radial current of the nonambipolar driving process closes the system and relates the angular acceleration $\dot{\omega}$ and poloidal flow speed to the driving processes. Our mathematics is brief, full details can be found elsewhere.[6]

It is instructive to compute the driving radial current for the case of neutral coinjection onto a trapped orbit. Figure 1 sketches the geometry. A radial current arises because the banana-averaged magnetic surface for an ion lies inside the corresponding surface for an electron (which is the birth surface for an electron-ion pair). This difference, represented in terms of a flux difference, is, from angular momentum conservation,

$$\delta\psi = - M v_{o\phi} Rc/e \tag{1}$$

and corresponds to an integrated inward radial current J_r per toroidal radian given by

$$J_r = \int \frac{\dot{n} e \ \delta\psi \ Rd\ell}{(\nabla\psi)} = - c \int \frac{Rd\ell}{\nabla\psi} (M\dot{n}Rv_{o\phi}) = - c \frac{dT}{d\psi} \ , \tag{2}$$

where $(M\dot{n}Rv_{o\phi})$ is the mechanical angular momentum input onto trapped orbits and $d\ell$ a poloidal line element along a magnetic surface. Here T denotes input torque per toroidal radian.

Our starting point is the standard drift kinetic equation and a multiple time scale argument. It is well-known that poloidal flow decays rapidly in a tokamak. The time scale of interest is long compared to the poloidal flow decay time [~ ν_i^{-1}] but short compared to the angular acceleration time scale [~ $(2T/M)^{1/2}/R\dot{\omega}$] so that ω, $\dot{\omega}$ and driven poloidal flow can be regarded as constant.

In the lab frame, the drift-kinetic equation is

$$\frac{\partial f}{\partial t} + e \frac{\partial \Phi}{\partial t} \frac{\partial f}{\partial E} + v_\parallel \frac{\vec{\nabla}\psi \times \hat{\phi}}{I} \cdot \vec{\nabla} f + \vec{v}_D \cdot \vec{\nabla} f = C(f) \quad , \tag{3}$$

where the velocity space variables are E,u with

$$E = \frac{1}{2} M v^2 + e\Phi$$

$$\frac{1}{2} M v_\parallel = E - e\Phi - \mu B$$

$$\vec{B} = \frac{I\hat{\phi}}{R} + \frac{\vec{\nabla}\psi \times \hat{\phi}}{R}$$

$$\vec{v}_D = - v_\parallel \left(\hat{\phi} + \frac{\vec{\nabla}\psi \times \hat{\phi}}{I}\right) \times \vec{\nabla} \left(\frac{v_\parallel}{\Omega}\right) \quad . \tag{4}$$

The collision model will be specified below. It is useful to split the potential into two terms

$$\Phi = \Phi_o(\psi) + \Phi_p \quad , \tag{5}$$

where $\Phi_o(\psi)$ is the flux surface average potential and Φ_p the part which varies in the flux surface.

A rotating plasma distribution function is a Maxwellian in the rotating frame and must be constructed from particle motion invariants E, μ, and the angular momentum ψ'

$$\psi' = \frac{c}{e} L = \psi + MRv_\phi c/e \quad . \tag{6}$$

Let us construct ε and $v_p = v_\parallel - \omega R$

$$\varepsilon = E - e\Phi_o(\psi') + \frac{1}{2} M\omega^2 R_o^2 = \frac{1}{2} M(\vec{v}-\omega R\hat{\phi})^2 - \frac{1}{2} M\omega^2(R^2-R_o^2) + e\Phi_p$$

$$\frac{1}{2} Mv_p^2 = \varepsilon + \frac{1}{2} M\omega^2(R^2 - R_o^2) - e\Phi_p - \mu B$$

where

$$\omega = c \frac{\partial\Phi}{\partial\psi} \tag{8}$$

is an angular rotation frequency which increases slowly with time ($\dot{\omega} > 0$). Our zero order distribution function is

$$f_o = n_o (M/2\pi T)^{3/2} e^{-\varepsilon/T} \quad . \tag{9}$$

and by Fig. 1

$$R = R_o + r\cos\theta \quad . \tag{10}$$

Spatial gradients in n,T, and ω are ignored; we focus solely on the consequences of nonzero $\dot{\omega}$. Quasineutral neutrality determines the poloidal potential variation

$$e\Phi_p = \frac{1}{2} M\omega^2(R^2 - R_o^2)T_e/(T_e + T_i) + e\Delta\Phi_p \quad , \tag{11}$$

where $\Delta\Phi_p$ arises because of nonzero $\dot{\omega}$.

Transformation of the kinetic equation from E,μ to ϵ,μ variables yields

$$\frac{v_p \partial g}{qR_o \partial\theta} + C(g) = \frac{r\sin\theta}{qR_o^2 T} (\epsilon - \frac{3}{2} Mv_p^2) u_o f_o$$

$$- \frac{M\dot{\omega}Rv_p}{T} f_o + \frac{v_p}{qR_o} \frac{\partial}{\partial\theta} (\frac{e\Delta\Phi_p}{T}) f_o \quad , \tag{12}$$

where the distribution has been separated into three parts

$$f = f_o - \frac{M u_o R_o v_p}{TR} f_o + g \quad . \tag{13}$$

The term in u_o represents a shift in the velocity of the Maxwellian, corresponding to incompressible parallel (i.e., poloidal) flow. As a result, g must have no velocity moment

$$\int d^3\vec{v}\ v_p g = 0 \quad . \tag{14}$$

Equation (11) has been simplified by taking the limit $\omega \to 0$ after terms in $\dot{\omega}$ have been evaluated. The kinetic equation has been linearized in $\dot{\omega}$, u_o, and g. A momentum-conserving pitch-angle scattering operator represents ion-ion collision

$$C(g) = \nu \frac{Mv_p}{B} \frac{\partial}{\partial\mu} \mu v_p \frac{\partial}{\partial\mu} g + \nu \frac{Mwv_p}{T} f_o \quad , \tag{15}$$

$$\bar{\nu}w = \int d\vec{v}\ v_p \nu g$$

$$\bar{\nu} = \int d\vec{v}\ (Mv_p^{\ 2}/T)\ \nu f_o$$

$$w = w_o\ R_o/R \quad , \tag{16}$$

where ν depends only on velocity magnitude in a manner appropriate to ion-ion collisions.

Equations (12-17) can be reduced to a previously solved neoclassical equation by the substitution

$$g = \{H(\mu,\theta)[w_o + u_o - \frac{\omega R_o}{\nu}] \frac{Mv^2}{2T} + \frac{Mv_p}{T} (w - \frac{\dot{\omega} R}{\nu}) \tag{17}$$

$$- \frac{e\Delta\Phi_p}{T} + \frac{r}{R_o} [\frac{\omega R_o}{\nu} + 3(u_o + w_o)] \frac{M}{2T} (v_p\cos\theta - \nu q R_o \sin\theta)\} f_o \quad .$$

This equation is

$$\frac{\partial H}{\partial\theta} - \nu \frac{MqR_o}{B} \frac{\partial}{\partial\mu} (\mu v_p \frac{\partial H}{\partial\mu}) = (\frac{r}{R_o}) \frac{\sin\theta}{v_p} \quad . \tag{18}$$

Equation (18) has the same mathematical form as Eq. (19) of Hirshman and Sigmar.[7] One can follow their method, recast (18) in terms of plateau velocity variables, and directly solve for H in the plateau regime

$$H = (\frac{r}{R_o}) (\frac{1}{\nu q R_o v^2})^{1/3} \int_o^\infty \sin[\theta - \frac{v_p \tau}{(v^2 q R_o \nu)^{1/3}}] e^{-\tau^3/6} d\tau \qquad (19)$$

In the banana regime, Rutherford's[8] interative method is used. Once H is known, the cross-surface ion current can be computed via

$$J_r = n_o Mc \int_o^{2\pi} R d\theta \; \{ \sum_\sigma \sigma \int_o^\infty \frac{2\pi d\varepsilon}{M} \int_o^{u_{max}(\sigma)} d\mu \; v_\parallel \frac{\partial g}{\partial \theta} \} \quad . \qquad (20)$$

Equations (14) and (16) can be used to eliminate $\dot{\omega}$ and w_o. The results are expressed in terms of the orthogonal conductivity

$$j_r = \frac{J_r}{2\pi r R_o} = \sigma_\perp E_r' \quad , \qquad (21)$$

where

$$\sigma_\perp = \begin{cases} \frac{3}{\sqrt{2}} K' \frac{n_o e^2 \nu_{i,eff}}{M \; \Omega_p^2} (\frac{r}{R})^2 & \text{Banana} \\ \sqrt{\pi/2} \; (\frac{n_o e^2}{M \Omega_p^2}) (\frac{T}{M})^{1/2} \frac{1}{Rq} (\frac{r}{R})^2 & \text{Plateau} \end{cases} \qquad (22)$$

and

$$K' \doteq 1 - \int_o^1 \frac{dk}{k^2} [\frac{\pi}{2E(k)} - 1] \doteq 0.702 \quad .$$

We used

$$u_o = \frac{cE_r' R_o}{\nabla \psi} \qquad (23)$$

to express the radial electric field in the frame rotating with the plasma.

In steady-state neoclassical theory, an ion current across magnetic surfaces does not arise from momentum-conserving ion-ion collisions. Our treatment, in which the rotation rate increases with time, does generate such a current. Equations (21-23) show that a poloidal rotation necessarily occurs. Conversely, any nonambipolar process will first generate a poloidal rotation whose decay by collisional processes exerts an azimuthal torque.

III. APPLICATIONS AND DISCUSSION

How large will the poloidal flow velocity be? Let us estimate the poloidal Mach number $M_p = (cE_r'/B_\theta)(M/2T)^{1/2}$. Straightforward algebra yields

$$M_p = \sin\alpha \left(\frac{T}{E_b}\right)^{1/2} \frac{(dP_b/dV)}{nT\bar{\nu}} \left(\frac{R}{r}\right)^{1/2} \qquad (24)$$

for the banana regime. Here E_b denotes the beam energy and dP_b/dV the beam injection power per unit volume onto trapped orbits, and $\sin\alpha$ is the beam injection angle with respect to the perpendicular. For midplane injection, the criterion for trapped orbits is $\sin^2\alpha < 2r/R_o$. Using the estimate $(dP_b/dV)/3nT \sim 1/\tau_E$, one finds

$$M_p \lesssim 3\,(T/E_b)^{1/2}(\bar{\nu}\tau_E)^{-1} << 1 \quad , \qquad (25)$$

since $\bar{\nu}\tau_E >> 1$ for all real tokamaks. Because the dynamical effect of poloidal rotation enters[9] as M_p^2, the dynamical consequences of forced poloidal rotation are negligible. In particular, driven poloidal rotation falls far short of the values needed to account for density asymmettries.[10]

Loss of α-particles from an ignited tokamak exerts a torque. Equation (15) states that each α-particle lost imparts an angular momentum impulse ΔT^* $= Ze\psi_o/c$, where ψ_o is the poloidal flux difference between the α birth-point and the plasma boundary. A simple balance between angular acceleration and torque

$$\left(\frac{Ze\psi_o}{c}\right) \dot{N}_\alpha \Delta = NM \left(\frac{2T}{M}\right)^{1/2} R\tau_s^{-1} \qquad (26)$$

defines the spin-up time τ_s. Let us define

$$\tau_E = 3NT/\dot{N}_\alpha E_\alpha$$

$$\psi_o = B_T a^2 \ln q_L / 2(q_L - 1) \qquad (27)$$

Δ = fraction of α's lost.

Then, straightforward algebra yields

$$\tau_s = \frac{\tau_E}{\Delta}\left(\frac{2(q_L-1)}{\ln q_L}\right)\left(\frac{R}{3a}\right)\left(\frac{\rho_\alpha}{a}\right)\left(\frac{E_\alpha}{T}\right)^{1/2}\left(\frac{M}{M_\alpha}\right)^{1/2} \approx \frac{2\tau_E}{\Delta} \quad , \qquad (28)$$

where the final form rests on estimates of reactor parameters. A reactor losing α-particles should spontaneously start to rotate.

Nonambipolar processes do not create large plasma potentials. Orthogonal conductivity (22), which reflects the damping of poloidal flows, keeps the radial electric field E_r' and poloidal rotation low. But nonambipolar processes do torque a tokamak, and the limiting rotation rate (and concomitant radial electric field in the lab frame) depends on a viscosity whose physics is observed to be not neoclassical.

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Figure Caption

Fig. 1 a) Geometry for calculations. The poloidal field lies in the $\hat{\theta}$ direction while the toroidal field is in the $\hat{\phi}$. The toroidal current is in the $-\hat{\phi}$ direction. 1b) Banana orbit for a coinjected ion born at A and traveling parallel to the toroidal current but anti-parallel to the magnetic field.

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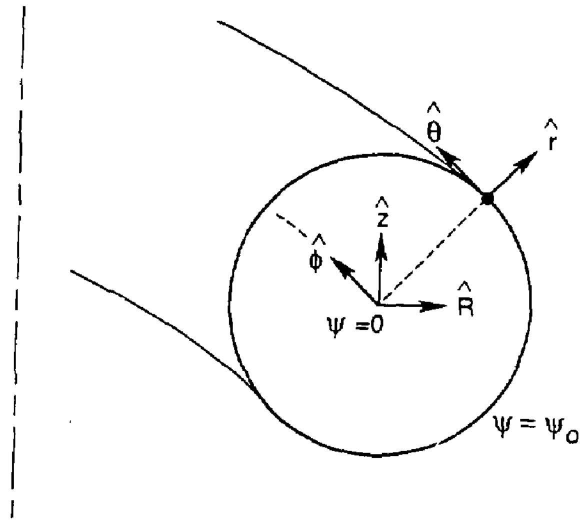
(a)



(b)

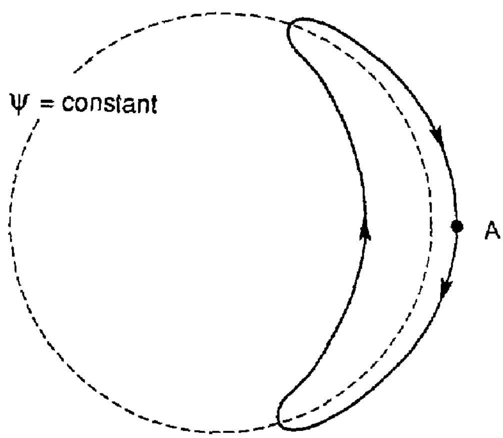


Fig. 1

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