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Mechanism of Viscosity Effect on Magnetic Island Rotation

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Abstract. It is shown that plasma viscosity does not influence the magnetic island rotation directly. Nevertheless, it leads to nonstationarity of the plasma velocity. This nonstationarity is the reason of the viscosity effect on island rotation.

Keywords: Magnetic islands, rotation, plasma viscosity, profile function

1. Introduction

The rotating magnetic islands are responsible for the ultimate plasma pressure achievable in long-pulse tokamak discharges [1]. The existing theories and available data (yet incomplete and insufficient for conclusive predictions) give unfavorable scaling for larger devices. Thereby, the physics of magnetic islands is an important area for present-day toroidal experiments with high-beta performance and even more for reactor-size fusion devices similar to ITER (International Thermonuclear Experimental Reactor) [2].

The rotating magnetic islands are the key element of the theory of nonlinear drift-tearing modes in tokamaks and of forced magnetic reconnection in plasmas (see Ref. [3] and the references therein). One of the crucial points there is the problem of determining the island rotation frequency. Basing on semi-intuitive considerations, the authors of Ref. [4] suggested the idea that the perpendicular viscosity is important for this problem. This idea was lively discussed in many subsequent papers (see, e.g., [3, 5-7]). However, till now this problem was studied inconsistently (the surface part of viscosity near the island separatrix and the nonstationarity of the plasma velocity were neglected), so that the mechanism of the viscosity effect on magnetic island rotation and even existence of this effect remained questionable. These subjects are analyzed in the present paper.

2. Magnetic island geometry

We consider magnetic field $\mathbf{B}_0$ with toroidal nested axisymmetric magnetic surfaces, labeled by the coordinate r, and the perturbation producing magnetic island localized near some rational surface $r = r_s$. Near this surface the total magnetic field is $\mathbf{B} = B_0\mathbf{z} - \nabla\psi \times \mathbf{z}$, where $B_0\mathbf{z}$ is the main ("equilibrium") magnetic field at $r = r_s$, and $\mathbf{z}$ is the unit vector determined by $\mathbf{z} = (B_{0\zeta}\boldsymbol{\zeta} + B_{0\theta}\boldsymbol{\theta})/B_0$, $\boldsymbol{\theta}$ and $\boldsymbol{\zeta}$ are the unit vectors along gradients of θ and ζ; θ and ζ are the poloidal and toroidal angles, respectively, $B_{0\zeta}$ and $B_{0\theta}$ are the toroidal and poloidal components of equilibrium magnetic field at the surface $r = r_s$. The island magnetic flux function ψ describes both the perpendicular (to the direction z) component of equilibrium magnetic field at $r \neq r_s$ and the perturbed magnetic field. It is defined by

$$\psi = \widetilde{\psi}\cos\xi - x^2 B_0/2L_s. \tag{1}$$

Here $x = r - r_s$, L_s is the shear length, $\widetilde{\psi}$ is a positive constant characterizing the perturbation amplitude, and $\xi = m(\theta - \zeta/q_s) - \omega t$ is the island cyclic variable, where m is the poloidal mode number, $q_s = q(r_s) = m/n$ is the safety factor at the rational surface, n is the toroidal mode number, and ω is the island rotation frequency.

3. Perpendicular torque balance equation

We use the single-fluid momentum equation

$$\rho_0 dV_y/dt = -j_x B_0/c - \partial\pi_{xy}/\partial x. \tag{2}$$

The direction y is defined by $\mathbf{y} = \mathbf{z} \times \mathbf{x}$, where $\mathbf{y}$ and $\mathbf{x}$ are the unit vectors along gradients of y and x. The explicit form for $\mathbf{y}$ is $\mathbf{y} = (B_{0\zeta}\boldsymbol{\theta} - B_{0\theta}\boldsymbol{\zeta})/B_0$. The function V_y is the y-projection of the plasma velocity $\mathbf{V}$, j_x is the x-projection of the electric current density, π_{xy} is the viscosity tensor component, ρ_0 is the plasma mass density, c is the speed of light, the operator d/dt is defined by $d/dt = \partial/\partial t + \mathbf{V}\cdot\boldsymbol{\nabla}$. The standard approximation $\partial/\partial y \ll \partial/\partial x$ is used. Drift and neoclassical effects are neglected.

Similarly to [3,6], to obtain the equation for the island rotation frequency we use the perpendicular torque balance equation. The last is found by spatial integration of Eq. (2) and can be represented in the form (cf. [3,6]):

$$I_{yNS} = T_{yEM} + T_{yV}. \tag{3}$$

Here T_{yEM} and T_{yV} are the electromagnetic and viscous torques, respectively,

$$T_{yEM} = -(B_0/c)\int ds j_x, \tag{4}$$

$$T_{yV} = -\int ds \partial\pi_{xy}/\partial x, \tag{5}$$

I_{yNS} is the inertia nonstationarity,

$$I_{yNS} = \rho_0 \int ds dV_y/dt, \tag{6}$$

$$\int (...)\, ds \equiv \int_{-\pi}^{\pi} d\xi \int_{-\infty}^{\infty} dx\,(...) = \sum_{\sigma_x} \sigma_x \int_{\widetilde{\psi}}^{-\infty} d\psi \oint \frac{d\xi}{\psi_x}(...), \tag{7}$$

$\psi_x \equiv \partial\psi/\partial x$, σ_x =sgn$x$. The operator $\oint(...)d\xi$ means $\int_{-\pi}^{\pi}(...)d\xi$ outside the separatrix (for $\psi < -\widetilde{\psi}$) and $\int_{-\xi_0}^{\xi_0}(...)d\xi$, where ξ_0=$\cos^{-1}(\psi/\widetilde{\psi})$, inside the island (for $\psi > -\widetilde{\psi}$).

4. The absence of viscosity torque for islands with localized viscosity tensor

We transform Eq. (5) as follows

$$T_{yV} = \sum_{\sigma_x} \sigma_x \int_{\widetilde{\psi}}^{-\infty} d\psi \frac{\partial}{\partial\psi} \oint d\xi \pi_{xy} = -\sum_{\sigma_x} \sigma_x \oint d\xi\, \pi_{xy}|_{\psi=-\infty}. \tag{8}$$

We consider the magnetic islands with localized viscosity tensor, $\pi_{xy} \to 0$ for $\psi \to -\infty$ (the alternative case with $\pi_{xy} \neq 0$ for $\psi \to -\infty$ was considered in [7]). It then follows from Eq. (8) that

$$T_{yV} = 0. \tag{9}$$

In this case Eq. (3) reduces to

$$I_{yNS} = T_{yEM}. \tag{10}$$

Thus, the electromagnetic torque is compensated by the inertia nonstationarity only.

5. Relation between inertia nonstationarity and profile function nonstationarity

The plasma velocity $\mathbf{V}$ is taken in the form $\mathbf{V} = V_\parallel \mathbf{z} + \mathbf{V}_\perp$, where $V_\parallel$ is the parallel velocity, $\mathbf{V}_\perp$ describes the cross-field motion, so that $\mathbf{V}_\perp = c[\mathbf{E}_\perp \times \mathbf{B}]/\mathbf{B}^2$ and $\mathbf{E}_\perp$ is the perpendicular (with respect to $\mathbf{B}$) electric filed. We express the field $\mathbf{E}_\perp$ in terms of the electrostatic potential ϕ by $\mathbf{E}_\perp = -\boldsymbol{\nabla}_\perp\phi$, where $\boldsymbol{\nabla}_\perp = \boldsymbol{\nabla} - \mathbf{b}\nabla_\parallel$ is the perpendicular gradient, $\mathbf{b} = \mathbf{B}/B$, $\nabla_\parallel = \mathbf{b}\cdot\boldsymbol{\nabla}$ is the parallel gradient. The plasma electric conductivity is assumed to be infinite, so that the parallel Ohm's law reduces to $E_\parallel \equiv \mathbf{b}\cdot\mathbf{E} = 0$. On the other hand, $E_\parallel = -\nabla_\parallel \phi - (1/c)\,\partial A_\parallel/\partial t$, where $A_\parallel = -\psi$ is the parallel projection of the vector potential. In addition, we use the expression for the parallel gradient $\nabla_\parallel = k_\parallel \partial/\partial\xi$, where $k_\parallel = -xk_y/L_s$ is the parallel wave number, $k_y = m/r_s$. Introducing k_y, we have allowed for that in the approximation of large aspect ratio $r_s \ll R$, where R is the major torus radius, the rule of differentiating over the variable y is defined by $\partial\ /\partial y = r_s^{-1}\partial\ /\partial\theta$. Then we obtain

$$\phi = (B_0\omega/ck_y)\left[x - h(\psi,t)\right], \qquad (11)$$

where $h(\psi,t)$ is the electrostatic potential profile function. The form of this function will be explained below.

Using Eq. (11), one can find that in terms of (ψ,ξ,t)

$$d/dt = \partial/\partial t - \omega\psi_x h'\partial/\partial\xi, \qquad (12)$$

where the prime denotes the derivative with respect to ψ. On the other hand, in terms of h,

$$V_y = (c/B_0)\partial\phi/\partial x = (\omega/k_y)\left(1 - \psi_x h'\right). \qquad (13)$$

By means of Eqs. (7), (12), and (13), Eq. (6) is transformed to

$$I_{yNS} = -2\pi\rho_0\frac{\omega}{k_y}\sum_{\sigma_x}\sigma_x\int_{\tilde{\psi}}^{-\infty} d\psi\frac{\partial h'}{\partial t}. \qquad (14)$$

Following [7], we assume that $h \equiv 0$ inside the island, i.e. for $\psi > -\tilde{\psi}$. By this reason, we used here $\oint d\xi = 2\pi$. Evidently, the right-hand side of Eq. (14) can be represented in terms of $(\partial h/\partial t)\,|_{x=-\infty}^{x=\infty}$.

6. Equation for electrostatic potential profile function and relation between electromagnetic torque and viscosity

We use the current continuity equation

$$\partial j_x/\partial x + \nabla_\parallel J_\parallel = 0, \qquad (15)$$

where $J_\parallel$ is the parallel electric current density. Multiplying Eq. (15) by ψ_x^{-1} and integrating over ξ, we obtain the ambipolarity condition

$$\frac{\partial}{\partial\psi}\oint j_x d\xi = 0. \qquad (16)$$

Integrating Eq. (16) over ψ and assuming the integration constant to be zero (otherwise the function h is strongly divergent for $\psi \to -\infty$), we obtain

$$\oint j_x d\xi = 0. \qquad (17)$$

Substituting here j_x from Eq. (2) and using Eqs. (12) and (13), we arrive at the equation for h:

$$\rho_0\frac{\omega}{k_y}\left\langle\psi_x^2\right\rangle\frac{\partial h'}{\partial t} = \left\langle\psi_x^2\frac{\partial\pi_{xy}}{\partial\psi}\right\rangle, \qquad (18)$$

where $\langle ...\rangle$ means the averaging over the island magnetic surface defined by

$$\langle ...\rangle = \oint(...)\frac{d\xi}{\psi_x}\bigg/\oint\frac{d\xi}{\psi_x}. \qquad (19)$$

By means of Eqs. (10), (14), and (18), we obtain the relation between the electromagnetic torque and the viscosity:

$$T_{yEM} = -2\pi\sum_{\sigma_x}\sigma_x\int_{\tilde{\psi}}^{-\infty}\frac{d\psi}{\langle\psi_x^2\rangle}\left\langle\psi_x^2\frac{\partial\pi_{xy}}{\partial\psi}\right\rangle. \qquad (20)$$

Thus, while directly the viscosity does not enter the torque balance equation, it contributes into this equation since it causes the inertia nonstationarity.

7. Structure of electrostatic potential profile function and viscosity

In the simplest plasma description the viscosity tensor component π_{xy} is given by

$$\pi_{xy} = -\mu\rho_0\partial V_y/\partial x, \qquad (21)$$

where μ is the viscosity coefficient. By order of magnitude $\mu \simeq \nu_i\rho_i^2$, where ν_i is the ion collision frequency, ρ_i is the ion Larmor radius. In this case Eq. (18) reduces to

$$A_1\partial h'/\partial t = \mu\partial\left(A_3 h''\right)/\partial\psi, \qquad (22)$$

where $A_i = \oint\psi_x^i d\xi$, $i = 1,3$.

One can see that Eq. (22) has no stationary solutions with $h' \neq 0$. The standard approach to analyzing Eq. (22) (cf., e.g., [7]) is based on the idea that this equation has approximate solutions which are stationary everywhere with the exception of a narrow "nonstationarity layer" near the separatrix. Let Δ_ψ be the characteristic scale of the nonstationarity layer in the magnetic flux space ψ. Then, for $\psi < -\tilde{\psi} - \Delta_\psi$ it follows from Eq. (22) that

$$h'' = C/A_3, \qquad (23)$$

where C is an integration constant. Requiring that $h' \to 0$ for $\psi \to -\infty$, we obtain from Eq. (23)

$$h' = C\widetilde{g}(\psi), \tag{24}$$

where $\widetilde{g}(\psi) = \int_{-\infty}^{\psi} d\psi/A_3$. Using Eq. (13), from the condition $V_y \to V_0$ for $\psi \to -\infty$, where V_0 is the equilibrium value of V_y, we find $C = 2\pi B_0(1 - k_y V_0/\omega)/L_s$. If the nonstationarity layer scale Δ_ψ is smaller than the ion Larmor radius scale $\Delta_\rho \simeq \rho_i \left(\widetilde{\psi} B_0/L_s\right)^{1/2}$, $\Delta_\psi < \Delta_\rho$, the expression for π_{xy} of form (21) should be modified by allowing for the terms with higher derivatives of V_y corresponding to the hyperviscosity (see in detail [8]).

Equations (23) and (24) are compatible with the above condition $h \equiv 0$ inside the island if one takes

$$\oint d\xi \psi_x \frac{\partial \pi_{xy}}{\partial \psi} = -2\pi\mu\rho_0 \frac{B_0(\omega - k_y V_0)}{L_s k_y} [\hat{\delta}(\hat{\psi}) + A_{3s}\widetilde{g}_s \hat{\delta}'(\hat{\psi})], \tag{25}$$

where $A_{3s} = A_3(-\widetilde{\psi})$, $\widetilde{g}_s = \widetilde{g}(-\widetilde{\psi})$, $\hat{\psi} = \psi + \widetilde{\psi}$, $\hat{\delta}$ is a function like the δ-function which is nonzero only for $\hat{\psi} < 0$, localized in the nonstationarity layer and normalized by the condition $\int_{-\infty}^{0} \hat{\delta}(\hat{\psi}) d\hat{\psi} = 1$. As an example of the function $\hat{\delta}(\hat{\psi})$ one can take the "half-Gaussian",

$$\hat{\delta}(\hat{\psi}) = 2\pi^{-1/2} \Delta_\psi^{-1} \exp\left(-\hat{\psi}^2/\Delta_\psi^2\right). \tag{26}$$

Exact knowledge of $\hat{\delta}(\hat{\psi})$ is not too important for our problem since, as it will be shown below, our final result depends on Δ_ψ logarithmically.

8. Dispersion relation

Using Eqs. (25) and (26), Eq. (20) reduces to

$$T_{yEM} = -\frac{\pi^2 g(1)}{3} \mu\rho_0 \frac{\omega - k_y V_0}{k_y} \left(\frac{B_0}{L_s \widetilde{\psi}}\right)^{1/2} \ln\left(\frac{\widetilde{\psi}}{\Delta_\psi}\right), \tag{27}$$

where $g(1) \equiv -8\sigma_x \widetilde{g}_s \widetilde{\psi}^{1/2} (B_0/L_s)^{3/2} = 0.869$ (see details in Ref. [8]). The value T_{yEM} is calculated by the following standard manner. Integrating by parts and using the current continuity equation (15), we transform Eq. (4) to

$$T_{yEM} = \frac{k_y \widetilde{\psi}}{c} \int J_\parallel \sin\xi ds. \tag{28}$$

By means of Eq. (1) and the parallel Ampere's law $J_\parallel = (c/4\pi)\, \mathbf{b} \cdot \mathbf{\nabla} \times \mathbf{B}$ we express the parallel current $J_\parallel$ in terms of $\partial^2\psi/\partial x^2$ and use the known matching condition at infinity

$$(\partial\psi/\partial x)|_{-\infty}^{\infty} = \widetilde{\psi}\left(\Delta_c' \cos\xi + \Delta_s' \sin\xi\right), \tag{29}$$

where Δ_c' and Δ_s' are the cos-part and sin-part of the matching parameter, respectively. Then Eq. (28) reduces to

$$T_{yEM} = k_y \Delta_s' \widetilde{\psi}^2/4. \tag{30}$$

From Eqs. (27) and (30) we obtain the equation

$$\omega - k_y V_0 = -\frac{3k_y^2 v_A^2 w^5}{32\pi g(1) L_s^2 \mu \ln\left(\widetilde{\psi}/\Delta_\psi\right)} \Delta_s', \tag{31}$$

where $v_A^2 = B_0^2/4\pi\rho_0$ is the Alfvén velocity squared and $w = 2\left(\widetilde{\psi} L_s/B_0\right)^{1/2}$ is the island halfwidth. This equation can be called a dispersion relation for the magnetic island.

For applications, the case of interest is when the equilibrium poloidal plasma velocity vanishes. In this case $k_y V_0 = -nU_\zeta/R$, where $U_\zeta \approx V_\parallel$ is the equilibrium toroidal plasma velocity.

9. Qualitative analysis of dispersion relation and energy conservation law

After calculation of the value Δ_s', the dispersion relation (31) together with the generalized Rutherford equation for the island width can be used for analyzing the problem of electromagnetic interaction between magnetic islands and the resistive wall (cf. Ref. [7]). Preliminary qualitative conclusions can be obtained by using expression for Δ_s' given in Ref. [9]. According to [9], in the case of resistive wall $\text{sgn}\Delta_s' = \text{sgn}\omega$. Therefore, if plasma is nonrotating for $r = r_s$, $V_0 = 0$, dispersion relation (31) is not satisfied. Thus, rotating magnetic islands can not be driven in nonrotating plasma because of their interaction with a resistive wall. On the other hand, in the presence of plasma rotation, $V_0 \neq 0$, it is possible to drive the magnetic islands for $k_y V_0/\omega > 1$. Thereby, the reason for driving the magnetic islands interacting with a resistive wall is the plasma rotation. One can suggest that such a driving should lead to excluding this reason, i.e. to braking the plasma rotation. This suggestion is in correspondence with the fact that the magnetic islands considered are related to the velocity profile nonstationarity. Such a braking of plasma rotation can be interpreted as an effective friction of rotating plasma with the nonrotating wall.

The effect of plasma rotation breaking can be seen also using the energy conservation law following from Eq. (2):

$$\frac{\partial}{\partial t}\frac{1}{2}\int \rho_0 V_y^2 r dr d\theta =$$

$$-\mu\rho_0 r_s \int \left(\frac{\partial V_y}{\partial x}\right)^2 ds + \frac{\omega r_s}{k_y} T_{yEM}, \tag{32}$$

where integration in the left-hand side is performed over all plasma volume (per length unit). For $\Delta_\psi \simeq \Delta_\rho$, from Eq. (32) one can find the estimate

$$\frac{\partial \ln \overline{V}_y}{\partial t} \simeq -\nu_i \frac{\rho_i}{a}\left(\frac{k_y V_0}{\omega} - 1\right)^2, \tag{33}$$

where $\overline{V}_y$ is a characteristic quasistationary plasma rotation velocity, a is the wall radius.

10. Discussion

In terms of the torque balance equation, the ideology of preceding studies of the viscosity effect on island rotation can be explained as follows (cf., e.g., Ref. [3]). The velocity profile was assumed to be exactly stationary, so that the left-hand side of Eq. (3) was neglected. The viscous torque T_{yV} was represented in the form of first equality (8), but the full region of integration over ψ was replaced by the region outside the nonstationarity layer. In other words, only the volume viscous torque was taken into account, while the surface viscous torque in the vicinity of the island separatrix was neglected. Then, instead of the identical zero given by Eq. (9), a finite value for the viscous torque was obtained. As a result of these two mistakes, a relation between ω and Δ'_s was derived. This relation coincides qualitatively with our dispersion relation (31) (cf. our Eq. (31) with Eq. (31) of Ref. [3]). This fact is not surprising since the velocity nonstationarity is also related to the surface part of the viscosity [see Eq. (18)].

Turning to [10], one can think that the islands studied can be suppressed by the feedback technique based on the finite wall resistivity effect. Mathematically, the role of feedback should consist in modification of the above condition $\mathrm{sgn}\Delta'_s = \mathrm{sgn}\omega$ by $\mathrm{sgn}\Delta'_s = \mathrm{sgn}(\omega - k_y V_0)$.

The approach presented can be generalized by including the finite plasma resistivity and drift as well as neoclassical effects.

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