Ion Larmour Radius Effect on rf Ponderomotive Forces and Induced Poloidal Flow in Tokamak Plasmas

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Analytical approximations are used to clarify the effect of Larmour radius on rf ponderomotive forces and on poloidal flows induced by them in tokamak plasmas. The electromagnetic force is expressed as a sum of a gradient part and of a wave momentum transfer force, which is proportional to wave dissipation. The first part, called the gradient electromagnetic stress force, is combined with fluid dynamic (Reynolds) stress force, and gyroviscosity is included into viscosity force to model finite ion Larmour radius effects in the momentum response to the rf fields in plasmas. The expressions for the relative magnitude of different forces for kinetic Alfvén waves and fast waves are derived.

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Poloidal flows have been found to play an important role on the creation of internal transport barriers (ITB) which improve the energy confinement in fusion plasmas [1,2]. The improvement is associated with a sudden increase of plasma density and temperature in the central region of a plasma column, with negligible modification outside the region. The temperature gradient, for instance, has a discontinuous variation at the transition between the two regions, indicating a localized change in the heat diffusion coefficient. The transport barriers appear around the middle of the plasma minor radius with a sheared plasma rotation and a negative shear of poloidal magnetic field. The effect of sheared poloidal flows is very well known and usually used in the theory of plasma stability [3]. It may decrease the turbulence level by shearing decorrelation mechanisms [3] and cause ion banana orbit squeezing by the associated radial electric field shear [4,5], modifying therefore the heat diffusion coefficient and the outward radial transport. Several mechanisms have been proposed to generate sheared plasma rotation, including neutral beam injection [1,2], Reynolds stress [3,6], and general ponderomotive forces [7,8] driven by radio-frequency (rf) fields. The idea to use forces driven by Alfvén waves to reduce radial transport has a rather long history [3,9]. Actually, toroidal plasma rotation (and the current drive also) induced by Alfvén waves has already been demonstrated [10,11], showing the possibility to create ITB.

The poloidal flow driven by rf fields has been calculated by many authors in various frequency ranges and with different approximations. The results of calculations of power required to drive substantial flows seem to be rather sensitive to the adopted model and, since many effects act simultaneously, it is not straightforward to pin down the most relevant mechanisms in numerical simulations. Recently, Berry, Jaeger, and Batchelor [12] calculated the poloidal flux induced by fast and ion-Bernstein waves using incompressible and compressible fluid models and a kinetic model, all in slab geometry. They found that, for fast waves, the effect of compressibility is to reduce the contribution of the Reynolds stress to the flow and that kinetic effects give larger ponderomotive force than both fluid models. The authors attribute this somewhat unexpected result to the near cancellation of the electromagnetic and nonlinear pressure forces. Because these results were obtained in slab geometry and by entirely numerical calculations, it is important to verify whether they remain valid in toroidal (or, at least, cylindrical) geometry and to clarify the role played by kinetic effects on the different forces.

In this Letter we investigate analytically the relative magnitudes of the different forces that determine the magnitude of the poloidal flow driven by rf waves, for the case of kinetic Alfvén (KAW), global Alfvén (GAW), and fast magnetosonic (FMSW) waves which dissipate via Landau damping and transit time magnetic pumping. The expressions for the forces are derived in cylindrical geometry using a fluid model. Finite gyroradius effects are taken into account by including gyroviscosity into the viscous force and keeping terms up to second order in the Larmour radius in the kinetic expression for the dielectric tensor. To clarify the influence of kinetic effects on the nonlinear forces, the electromagnetic force is split in a part that depends explicitly on gradients of oscillating currents and fields and a part which is due to wave momentum dissipation. The former can be properly combined with the fluid dynamic stress into what is sometimes loosely mentioned as the rf pressure force.

The ponderomotive forces in each plasma species α [electrons (e) and ions (i)] are calculated from the collisionless two-fluid momentum equation [14],

$$\frac{\partial}{\partial t}(m_\alpha n_\alpha V^{(\alpha)}_{\theta,\zeta}) + \nabla(m_\alpha n_\alpha \mathbf{V}^{(\alpha)} V^{(\alpha)}_{\theta,\zeta}) = \left(e_\alpha n_\alpha \mathbf{E} + \frac{1}{c}\mathbf{j}^{(\alpha)} \times \mathbf{B}\right)_{\theta,\zeta} - \nabla_s \pi^{(\alpha)}_{s\zeta,\theta}, \qquad (1)$$

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where m_α, e_α are mass and charge of particles, and homogeneous kinetic pressure along magnetic surfaces, $\nabla_{\theta,\zeta} P^{(\alpha)} = 0$, is assumed. Any physical variable $\mathbf{\Phi}$ (density $n_{(\alpha)}$, velocity $\mathbf{V}^{(\alpha)}$, current $\mathbf{j}^{(\alpha)} = e_\alpha n_\alpha \mathbf{V}^{(\alpha)}$, electric $\mathbf{E}$, and magnetic $\mathbf{B}$ fields) is written as a sum of quasi-stationary (represented as $\mathbf{\Phi}$) and oscillating (represented as $\tilde{\mathbf{\Phi}}$) parts. The oscillating part is represented as one wave harmonic, $\tilde{\mathbf{\Phi}} \propto \exp[\mathrm{i}(\int_0^r k_r\, dr + m\theta + n\zeta - \omega t)]$, where ω is the wave frequency, m, n are the poloidal and toroidal wave numbers, and the eikonal approximation is assumed for the radial dependence ($|k_r \mathbf{\Phi}| \gg |\partial \mathbf{\Phi}/\partial r|$ and $|k_r| \gg |m/r|, |n/R_0|$). The cylindrical limit ($r/R_0 \ll 1$) of the pseudotoroidal coordinates (r, θ, ζ) with coaxial magnetic surfaces is used. Following the usual procedure of averaging Eq. (1) over the wave oscillations, we obtain

$$\begin{aligned} F_{FD,\theta,\zeta}^{(\alpha)} &= -\nabla \langle m_\alpha n_\alpha \tilde{\mathbf{V}}^{(\alpha)} \tilde{V}_{\theta,\zeta}^{(\alpha)} \rangle, \\ F_{\mathrm{EM},\theta,\zeta}^{(\alpha)} &= \left\langle e_\alpha \tilde{n}_\alpha \tilde{\mathbf{E}} + \frac{\tilde{\mathbf{j}}^\alpha \times \tilde{\mathbf{B}}}{c} \right\rangle_{\theta,\zeta}, \\ F_{V,\theta,\zeta}^{(\alpha)} &= -\langle \nabla_s \pi_{s\zeta,\theta}^{(\alpha)} \rangle, \end{aligned} \tag{2}$$

where the first force is produced by the fluid dynamic stress (which includes the Reynolds stress), the next ($F_{\mathrm{EM},\theta,\zeta}^{(\alpha)}$) is produced by the electromagnetic stress, and the last one is produced by viscosity. Using continuity and induction equations, the electromagnetic stress force can be represented as the sum [7,8] of a gradient part, $F_{\partial,\theta}^{(\alpha)} = 1/(2r\omega)\,\mathrm{Im}\nabla_r(r\tilde{j}_r E_\theta^*)$ and $F_{\partial,\zeta}^{(\alpha)} = 1/(2\omega)\,\mathrm{Im}\nabla_r(\tilde{j}_r E_\zeta^*)$, and a wave momentum transfer force $W^{(\alpha)}\mathbf{k}/\omega$, which is proportional to wave dissipation, $W^{(\alpha)} = \tilde{\mathbf{j}}^{(\alpha)} \cdot \tilde{\mathbf{E}}$, where $\tilde{j}_s^{(\alpha)} = -\mathrm{i}\omega/(4\pi) \sum_p \varepsilon_{sp}^{(\alpha)} \tilde{E}_p$. The poloidal component of the viscosity force can be represented as the sum of the neoclassical ion viscosity force [3,6], $F_{\mathrm{neo},\theta}^{\pi} = \mu_{\mathrm{neo}} V_{0\theta}$ (μ_{neo} is the ion viscosity coefficient), and the collisionless gyroviscosity force [14],

$$F_{\mathrm{GV},\theta}^{(i)} \approx \frac{m_i}{2} \rho_i^2 \omega_{ci} \frac{\partial}{\partial r} \left\langle \tilde{n}_i \left(\frac{\partial \tilde{V}_r^{(i)}}{\partial r} - \frac{\partial \tilde{V}_\theta^{(i)}}{r \partial \theta} \right) \right\rangle, \tag{3}$$

in which the second-order finite ion Larmour radius ($\rho_i = v_{Ti}/\omega_{ci}$) effect is taken into account. The poloidal rotation velocity $V_{0\theta}$ is found by balancing the sum of ponderomotive forces on electrons and ions with the poloidal component of the total ion viscosity force F_θ^π.

The relations between the oscillating current and velocity and the components of the rf electric fields are determined by the dielectric tensor $\varepsilon_{sp}^{(\alpha)}$ (see, for example, Refs. [8,13]):

$$\begin{gathered} \varepsilon_{11} \approx \varepsilon_{11}^{(i)} = \varepsilon_1 \left(1 - \frac{3}{4} k_r^2 \rho_i^2\right); \qquad \varepsilon_1 = \varepsilon_1^{(i)} = \frac{\omega_{pi}^2}{\omega_{ci}^2 - \omega^2}, \qquad \varepsilon_{12} = -\varepsilon_{21} = \mathrm{i}\varepsilon_1 \frac{\omega}{\omega_{ci}}; \qquad \varepsilon_{12}^{(i)} = \mathrm{i}\varepsilon_1 \frac{\omega_{ci}}{\omega}; \\ \varepsilon_{22} = \varepsilon_{22}^{(i)} + \varepsilon_{22}^{(e)} = \varepsilon_1 \left(1 - \frac{11}{4} k_r^2 \rho_i^2\right) - 2 \frac{\omega_{pe}^2}{\omega^2} \eta_e k_r^2 \rho_e^2; \qquad \varepsilon_{32}^{(e)} = -\varepsilon_{23}^{(e)} = \mathrm{i} \frac{\omega_{pi}^2 k_r}{\omega_{ci} \omega k_\parallel} (1 - \eta_e); \\ \varepsilon_{33}^{(e)} = \frac{\omega_{pe}^2}{k_\parallel^2 v_{Te}^2} (1 - \eta_e) \qquad \eta_e = \frac{S_{\mathrm{ph}}}{\sqrt{\pi}} \int \frac{\exp(-S^2)}{S_{\mathrm{ph}} - S}\, dS; \qquad S_{\mathrm{ph}} = \frac{\omega}{\sqrt{2}\,|k_\parallel| v_{Te}}. \end{gathered} \tag{4}$$

The relative magnitudes of the different forces depend on the polarization relationships of the components of the rf electric field; solving the Maxwell equations, we have

$$\frac{E_r}{E_b} \approx \frac{N_r N_b + \varepsilon_{12}}{N_b^2 + N_\parallel^2 - \varepsilon_{11}}; \qquad \frac{E_\parallel}{E_b} \approx \frac{(\varepsilon_{11} - N_\parallel^2)(\varepsilon_{22} - N^2) - \varepsilon_{12}\varepsilon_{21} - N_b^2 \varepsilon_{22}^{(e)}}{N_r^2 N_b N_\parallel + (\varepsilon_{23} + N_b N_\parallel)(N_b^2 + N_\parallel^2 - \varepsilon_{11}) + \varepsilon_{21} N_r N_\parallel}, \tag{5}$$

where $\mathbf{N} = \mathbf{k}c/\omega$ is the refractive index. Note that, to estimate the value of the poloidal component of the ponderomotive forces, we will neglect the difference between poloidal and binormal components of the oscillating current and electric field (for example, $\tilde{E}_b = \tilde{E}_\theta B_\zeta / B_0 - \tilde{E}_\zeta B_\theta / B_0 \approx \tilde{E}_\theta$) because of $\tilde{E}_r, \tilde{E}_b \gg \tilde{E}_\parallel$ and $B_\theta \ll B_0$ (B_θ is the poloidal component of the magnetic field $\mathbf{B}_0$).

To analyze the ion Larmour radius effect on the forces given by Eq. (2) for the different waves, we consider the approximate dispersion relations,

$$\begin{aligned} N_{r\mathrm{KA}}^2 &\approx (\varepsilon_{11} - N_\parallel^2)\varepsilon_{33}/\varepsilon_{11}, \\ N_{r\mathrm{GA}}^2 &\approx \varepsilon_{12}^2/(\varepsilon_{11} - N_\parallel^2), \\ N_{r\mathrm{FMS}}^2 &\approx (\varepsilon_{11} - N_\parallel^2 - N_b^2), \end{aligned} \tag{6}$$

and the polarization coefficients for KAW, GAW, and FMSW, respectively,

$$\begin{aligned} \left(\frac{E_r}{E_b}\right)_{\mathrm{KA,GA}} &\approx \frac{N_r N_b + \varepsilon_{12}}{N_b^2}, \\ \left(\frac{E_r}{E_b}\right)_{\mathrm{FMS}} &\approx -\frac{N_r N_b + \varepsilon_{12}}{\varepsilon_1}. \end{aligned} \tag{7}$$

We begin the analysis of the poloidal driving forces by assuming dissipation only on electrons and defining the wave momentum transfer force via electron Landau damping [7,8],

$$F_{\mathrm{MT},\theta}^{(e)} = \frac{m}{r\omega} W^{(e)} \approx \frac{m}{8\pi r} \mathrm{Im}\varepsilon_{33}^{(e)} |E_\parallel|^2, \tag{8}$$

which is valid for the kinetic and global Alfvén waves. For the FMSW, this force can be calculated using the relation $E_\parallel/E_b \approx -\varepsilon_{32}/\varepsilon_{33}$. The other driving forces will be compared with this momentum transfer force.

Combining the gradient electromagnetic stress force with the fluid dynamic stress force ($F^{(\alpha)}_{\partial,\theta} + F^{(\alpha)}_{FD,\theta}$), we obtain the total gradient force,

$$F^{(\alpha)}_{\nabla,\theta} \approx \frac{1}{r^2}\frac{\partial}{\partial r}\left\{\frac{r^2}{8\pi}\mathrm{Re}\left[\varepsilon^{(\alpha)}_{12}\left(1 + \frac{\omega^2}{\omega^2_{p\alpha}}\varepsilon^{(\alpha)*}_{22}\right)|E_b^2| + \frac{\omega^2}{\omega^2_{p\alpha}}\varepsilon^{(\alpha)}_{11}\varepsilon^{(\alpha)}_{12}|E_r^2| + \varepsilon^{(\alpha)}_{11}\left(1 + \frac{\omega^2}{\omega^2_{p\alpha}}\varepsilon^{(\alpha)*}_{22} + \frac{\omega^2_{c\alpha}}{\omega^2_{p\alpha}}\frac{\varepsilon_1^2}{\varepsilon^{(\alpha)}_{11}}\right)E_r E_b^*\right]\right\}. \tag{9}$$

In the zero Larmour radius limit ($\rho_{e,i} \to 0$), this force and gyroviscosity force $F^{(i)}_{\mathrm{GV},\theta}$ equal zero independently of the type of waves, and the poloidal rotation is defined only by the momentum transfer force. If only the Reynolds stress term is taken into account in Eq. (1) ($\omega^2 \gg \omega^2_{ci}$), the force given by Eq. (9) will be an overestimate. This seems somewhat in disagreement with the conclusions in Ref. [12].

Using Eq. (7) and $k_b \ll k_r, \omega^2 \ll \omega^2_{ci}$ in Eq. (9), we have, for both the KAW and GAW,

$$F^{(i)}_{\nabla,\theta} \approx \frac{3}{32\pi r^2}\frac{\partial}{\partial r} \times\left[r^2\rho_i^2\varepsilon_1\left(\frac{\omega}{\omega_{ci}}\frac{k_r^2}{k_b^2}\mathrm{Im}k_r^2 - \frac{k_r}{k_b}\mathrm{Re}k_r^2\right)|E_b|^2\right], \tag{10}$$

where the first term is important for KAW and the second term is important for GAW. Estimating the relation $E_\parallel/E_b$ from Eq. (5) as $(N_\parallel^2 - \varepsilon_{11})/(N_\parallel N_b)$ and $\partial/\partial r$ as $-2|\mathrm{Im}k_r|$, we get the ratio of the total gradient force to the wave momentum transfer force for KAW,

$$(F^{(i)}_{\nabla,\theta}/F^{(i)}_{\mathrm{MT},\theta})_{\mathrm{KA}} \approx -\frac{3}{8}\frac{T_i}{T_e}\frac{k_\parallel^2}{k_b^2}\left(\frac{2\omega}{\omega_{ci}}\frac{|\mathrm{Im}k_r|}{k_b} - 1\right). \tag{11}$$

This ratio can be large because the wave dissipation length $|\mathrm{Im}k_r|^{-1}$ is the order of few Larmour radii. Note that this gradient force changes direction at the mode conversion point [8] (where $\varepsilon_1(r_A) = N_\parallel^2$) because we have a wave dissipation maximum at this point.

For the fast waves the dissipation length is larger than plasma radius $\mathrm{Im}k_r \ll 1/a$ and the parameter of the radial plasma inhomogeneity l_r ($\partial/\partial r \approx -1/l_r$) may be important. For GAW the small ion Larmour radius effect becomes relevant in the gradient force because $\omega^2 \ll \omega^2_{ci}$,

$$(F^{(i)}_{\nabla,\theta}/F^{(i)}_{\mathrm{MT},\theta})_{\mathrm{GA}} \approx \frac{3}{8}\frac{T_i}{T_e}\frac{k_r}{k_b^2 l_r}\frac{\omega^2}{\omega^2_{ci}}, \tag{12}$$

Using relations in Eq. (7) and $|\mathrm{Im}k_r| \ll |k_b| < |\mathrm{Re}k_r|$, we find for FMSW,

$$\begin{aligned} F^{(i)}_{\nabla,\theta} &\approx \frac{3}{32\pi r^2}\frac{\partial}{\partial r}\left[r^2\varepsilon_1\rho_i^2 k_r\left(k_b + \frac{11}{3}\frac{\omega^2(k_b + 2\,\mathrm{Im}k_r\,\omega_{ci}/\omega)}{\omega^2_{ci} - \omega^2}\right)|E_b|^2\right] \\ &\approx -\frac{1}{4}\frac{T_i}{T_e}\left(\frac{3\omega^2}{\omega^2_{ci} - \omega^2} + \frac{11\omega^4}{(\omega^2_{ci} - \omega^2)^2}\right)\frac{1}{k_r l_r}\frac{F^{(i)}_{\mathrm{MT},\theta}}{\mathrm{Im}\eta}. \end{aligned} \tag{13}$$

This force is large near ion-cyclotron resonance and it is also larger than the momentum transfer force for large phase velocity, $\omega/k_\parallel \gg v_{Te}$, because electron dissipation is exponentially small, $-\mathrm{Im}\eta \approx \sqrt{\pi/2}\,\omega/(k_\parallel v_{Te})\exp[-\omega^2/(k_\parallel v_{Te})^2] \ll 1$. The result is in qualitative agreement with those of Ref. [12].

In the last step, we analyze the gyroviscosity force (3) and compare it with the wave momentum transfer force (8). Estimating the density fluctuations from the continuity equation, $\tilde{n}_i \approx (k_r\tilde{j}_r^{(i)} + k_b\tilde{j}_b^{(i)})/(e_i\omega)$, we obtain the value of the gyroviscosity force,

$$F^{(i)}_{\mathrm{GV},\theta} \approx \frac{-\omega\,\omega_{ci}\rho_i^2}{8\pi}\frac{\partial}{\partial r}\mathrm{Im}\left[\frac{k_r k_b}{\omega^2_{pi}}\varepsilon^{(i)}_{12}\left(\varepsilon^{(i)}_{11}|E_r|^2 + \varepsilon^{(i)*}_{22}|E_b|^2 + \frac{\varepsilon^{(i)}_{11}}{\varepsilon^{(i)}_{12}}\varepsilon^{(i)*}_{22}E_r E_b^* + \varepsilon^{(i)}_{12}E_r^* E_b\right)\right]. \tag{14}$$

Then, using the approximations for KAW as in Eq. (11), we estimate this force,

$$(F^{(i)}_{\mathrm{GV},\theta}/F^{(i)}_{\mathrm{MT},\theta})_{\mathrm{KA}} \approx (T_i/T_e)k_\parallel^2/k_b^2, \tag{15}$$

which is therefore the same order of the momentum transfer force in Eq. (8). Note that the gyroviscosity force driven by GAW is very small.

Using $-\mathrm{Im}(E_r/E_b) \approx \omega/\omega_{ci} < 1$ and $|\mathrm{Im}k_r| \ll |k_b| < |\mathrm{Re}k_r|$ for FMSW, we find

$$(F^{(i)}_{\mathrm{GV},\theta})_{\mathrm{FMS}} \approx -\frac{T_i}{T_e}\frac{\omega^2\omega^2_{ci}}{(\omega^2_{ci} - \omega^2)^2}\frac{k_b^2}{k_r^3 l_r}\frac{F^{(i)}_{\mathrm{MT},\theta}}{\mathrm{Im}\eta}, \tag{16}$$

which is much smaller than the gradient force (13) driven by the FMSW.

Finally, we estimate the finite ion Larmor radius effect on power requirements to drive sufficient flow to create the internal transport barriers by the kinetic Alfvén waves in TCABR [the Tokamak Chauffage Alfvén wave experiment in Brazil (for parameters see Ref. [8])]. The KAW ponderomotive force can drive poloidal flow in TCABR near the mode conversion surface $r_A \approx 0.7a$ in a narrow radial layer [8] $\Delta \approx (ac^2/\omega_{pe}^2)^{1/3}(v_{Te}/v_A)^{4/3} \approx 2$ cm, which accounts for the structure of the Airy function for KAW at r_A. Note that the flow velocity changes sign at the mode conversion point. Balancing the poloidal force (11) with the neoclassical viscous force, we find for the parameters of TCABR ($a = 18$ cm, $R_0 = 61$ cm, $B_t = 1$ T, $q_0 = 1.01$, $T_{i/e} = 400/500$ eV, $n_0 = 2 \times 10^{13}$ cm^{-3}, and all profiles are assumed parabolic) the poloidal plasma flow velocity for wave dissipation of 400 kW, $U_{\theta,\max} \approx 6$ km/s, and radial electric field, $B_0 U_\theta / c \approx 60$ V/cm, can also be estimated. This strongly sheared profile of the poloidal velocity is similar to the one in the conditions of the internal transport barriers in the Test Fusion Tokamak Reactor experiment [2]. In this case, a sheared flow $2U_{\theta,\max}/\Delta$ will be enough to overcome the turbulence threshold [3], $\Delta\omega_{k'}/(\Delta x_{k'} k'_\theta)$, where $\Delta\omega_{k'}$ is the order of drift frequency ≈100 kHz and the decorrelation length $\Delta x_{k'}$ is the order of the inverted turbulence wave number $k_\theta'^{-1}$. Note that without the Larmor radius effect the sheared flow (less by one-third) is not enough to overcome the turbulence threshold.

In conclusion, we can say the following: A general form of time-averaged poloidal ponderomotive forces induced by fast and kinetic Alfvén waves in geometric optics approximation is analyzed on the basis of the collisionless two fluid (ions and electrons) magnetohydrodynamics equation and it is shown that accounting only the Reynolds stress term in Eq. (1) can overestimate the plasma flow. It is found that the finite ion Larmor radius effect plays a fundamental role in ponderomotive forces because of a radial plasma inhomogeneity that can drive a poloidal flow, which is larger than a flow driven by a wave momentum transfer force. This effect can be used to drive the poloidal flow in tokamaks by the kinetic Alfvén waves with low phase velocity and by fast waves with phase velocity larger than thermal velocity.

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