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EFFECTS OF FINITE DRIFT ORBIT WIDTH AND RF-INDUCED SPATIAL TRANSPORT ON PLASMA ROTATION BY ICRH

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ICRH heating have been seen to affect the plasma rotation in several tokamaks [1-3]. Detailed measurements of the toroidal velocity profiles during ICRH heating were recently made in the last JET campaigns [4, 5]. Owing to charge separation through resonant ions moving across magnetic flux surfaces and momentum transfer through friction between the resonant ions and the background plasma, toroidal and poloidal acceleration of the bulk plasma can be induced by ICRH [6-9]. The RF-induced transport, finite drift orbit width and orbit topology of resonant ions are all found to play an important role for the torque imparted by the fast ions to the bulk plasma during high power ICRH c. f. [10]. For example heating with waves propagating parallel to the current leads to an inward drift of the turning points of trapped resonant ions towards the mid-plane. As these ions reach the mid-plane the orbits will de-trap, preferentially into co-current passing orbits which may ultimately be displaced to the low field side of the magnetic axis. They will then, through friction, transfer a net momentum to the plasma whereas the trapped ions give a bipolar torque.

1. RF-heating in a toroidal plasma

When discussing RF-interactions in a nearly collisionless plasma it is convenient to express the distribution functions of the resonant ion species in terms of the guiding centre invariants of the unperturbed orbits. We use the invariants: the energy, $E \equiv mv^2/2$; the canonical angular momentum, $P_\phi = mRv_\phi + eZ\Psi$, where Ψ is the poloidal magnetic flux function; and $\Lambda=\mu/E$, where μ is the magnetic moment. In our coordinate system (r, ϕ, θ) the toroidal magnetic field is in the positive ϕ-direction and the plasma current in the negative ϕ-direction, producing a positive poloidal magnetic flux $2\pi\Psi$. Decorrelated RF-interactions with a single toroidal mode describes a one dimensional diffusion process along the characteristics

$$\left(E,\Lambda,P_\phi\right)=\left(E(t),\ \frac{n\omega_{c0}}{\omega}+\frac{A_2}{E(t)},\ \frac{n_\phi}{\omega}E(t)+A_1\right) \tag{1}$$

where A_1 and A_2 are integration constants changing with Coulomb collisions. Since RF-heating essentially increases the perpendicular energy of the ions, passing orbits are turned into trapped orbits. As the energy further increases the trapped orbits will increase their width and the turning points will be displaced both poloidally towards the cyclotron resonance and across the flux surfaces. Across the flux surfaces the displacement is given by [6]

$$\Delta\psi_T = \frac{1}{eZ}\frac{n_\phi}{\omega}\Delta E \tag{2}$$

For $n_\phi>0$ the turning points are drifting away from the mid-plane as the energy increases. For $n_\phi<0$ the turning points of the trapped ions are drifting towards the mid-plane and towards the cyclotron resonance as their energy increases. When the turning points meet in the mid-plane they will detrap into counter passing ions. If the energy is further increased the orbits will be shifted at the low field side of the cyclotron resonance $\omega = n\omega_c$.

2. Toroidal rotation

The electric field arising from the charge separation produces a polarisation drift. The time integrated polarisation current caused by the acceleration of the background plasma approximately equals the time integrated current across the flux surfaces of the resonant ions but has opposite sign. The toroidal torque on the bulk plasma, $T_{j\times B}$, from them becomes

$$T_{j\times B} = -\int_{t_1}^{t_2} dt \int_{V(\psi)}^{V(\psi+\Delta\Psi)} R\, j_{F\psi} B_\theta dV \tag{3}$$

where $j_{F\psi}$ is the radial current associated with the drift of the ions. Assuming RB_θ = constant, the change in the torque on the resonant ions inside the volume element with the width $d\Psi = B_\theta R dl_{\nabla\Psi}$ becomes $\Delta T_{j\times B} = -d\Psi[Q(\Psi,t_2) - Q(\Psi,t_1)]$, where $Q(\Psi,t)$ is the charge inside the flux surface Ψ. The total change in toroidal angular momentum of the plasma becomes

$$\left[\Delta T_\phi(\Psi,t_2) - \Delta T_\phi(\Psi,t_1)\right]_{plasma} = \int_{t_1}^{t_2} dt \int_{V(\psi)}^{V(\psi+\Delta\Psi)} m \frac{\partial\left(n_F \langle R v_\phi \rangle\right)_{coll}}{\partial t} dV - d\Psi[Q(\Psi,t_2) - Q(\Psi,t_1)] \tag{4}$$

where index coll for Coulomb collisions and plasma for the background plasma.

As the high-energy tails on the distribution functions develop there will be a transient change in the toroidal momentum due to the change in angular velocity of the heated ions and the change in their density profiles. In the absence of a momentum input by the wave the conservation of angular momentum causes the resonant ions and the background plasma to move in opposite directions with a differential rotation at the inner and outer part of the plasma due to orbit broadening. For waves having no toroidal angular momentum there will be no net torque on the plasma as the distribution function reaches steady state, except for that arising from lost ions. However, for a bipolar torque the toroidal velocity may still become unidirectional since the momentum input near the edge is lost faster than near the centre [9]. Interactions with waves having a finite toroidal mode number produce continuous momentum input and will thereby produce a larger net velocity. For typical JET parameters the initial torque due to RF-interactions is typically an order of magnitude lower than that for NBI with the same power. The initial change in torque by RF is a few times larger than the continuous torque arising from waves with a finite toroidal mode number.

3. Numerical simulation with the SELFO code

The RF-induced torque during minority ^{3}He heating in a ^{4}He plasma is calculated with the SELFO code [11] for the following parameters: B_0=3.44T, n_e=2.42×10^{19}×(1-0.9$(r/a)^2$)m^{-3}, n_{4He}/n_e=0.207, n_C/n_e=0.093, n_{3He}/n_e=0.015, T_e=T_i=5keV, I_p=2MA, R_0=2.97m, a=1.2m, P=5.0MW and f=35MHz, resulting in a central slowing down time of 0.4s.

For waves with n_ϕ = -15 most of the high-energy ions have trapped or counter-passing orbits. There is a strong asymmetry between co- and counter-passing ions, which increases with energy. The ions with the highest energies have counter-passing orbits localised on the low field side of the magnetic axis. For n_ϕ = 15 most of the high-energy ions are trapped. There is still an asymmetric high-energy tail of counter passing ions. In spite of the outward RF-induced flux there is an asymmetric high-energy tail of de-trapped passing ions caused by Coulomb collisions and RF-diffusion.

The RF-heating produces in general a differential torque. During the first few milliseconds the torque is relatively strong but then rapidly relaxes. This torque originates both from transport of ions across flux surfaces and from friction between resonant ions and background ions. The friction between resonant ions and electrons does not become important until after the high-energy tail has been built up. The time histories of the total torques, T, integrated over the plasma volume are shown in Fig. 1 for on-axis heating with $n_\phi = 15$ and $n_\phi = -15$, where a positive torque is in our coordinate system in the counter-current direction. A positive net torque is obtained for $n_\phi = 15$ and a negative for $n_\phi = -15$ corresponding to the momentum absorbed by the wave. In Fig. 2 the integrated torque density on the background plasma for the various components of $T(\Psi) = \int d\Psi m\, d(Rv_\phi)/dt$ are shown at $t = 0.5$s when the distribution function is close to that of steady state. For $n_\phi = 15$ the torque density from the ion-ion collisions is in the same direction as that from the outward drift, whereas the friction by high-energy resonant ions on electrons produces a torque in the opposite direction. The momentum absorbed by the passing resonant ions is partially transferred to the background plasma by ion-ion collisions. For heating with co-current propagating waves having $n_\phi = -15$ the torque is dominated by a co-current torque density caused by the friction between resonant ions and electrons. The net momentum required to sustain the frictional torque by the resonant high-energy ions comes from the momentum absorbed from the wave.

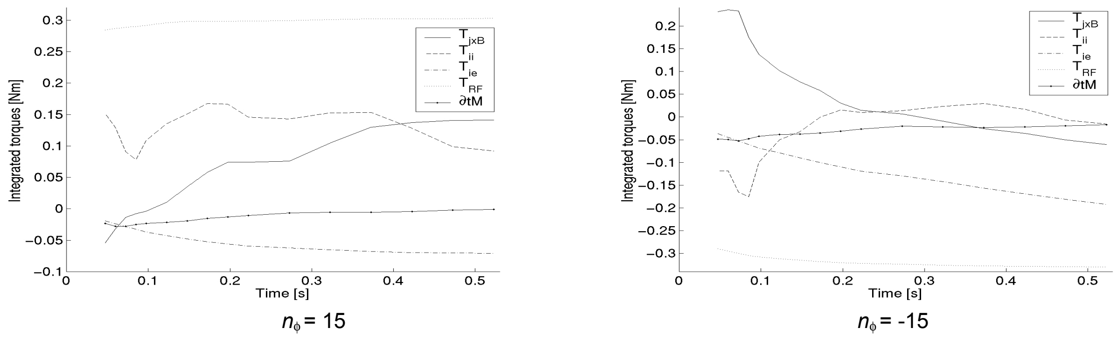


Fig. 1 Time evolutions of the integrated torques T_{jxB}(full line), T_{ii} (dashed line), T_{ie} (dot line), T_{RF} (dot-dashed line) and the change in momentum of the resonant ions (full line with points).

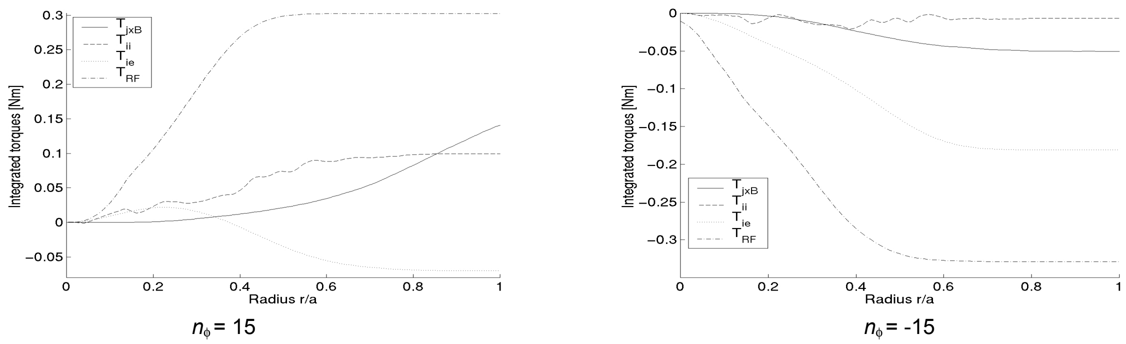
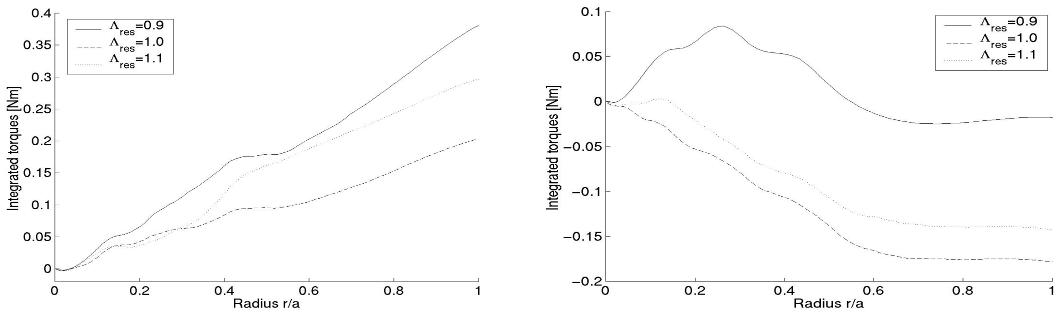


Fig. 2. The radial variation of the integrated torques T_{jxB}(full line), T_{ii} (dashed line), T_{ie} (dot line), T_{RF} (dot-dashed line) and the change in momentum of the resonant ions.

By changing the position of the resonance we have found that high-field side off-axis heating in general produces higher losses of high-energy ions intercepted by the wall due to wider banana orbits located further out than on-axis heating, provided the power density is sufficient to produce high-energy tails. The de-trapping is also less asymmetric leading to more co-passing orbits. For off-axis heating at the low field side more trapped ions are formed and the losses are lower. High-energy ions with counter-passing orbits located on the low field side of the magnetic axis appear for $n_\phi = -15$. The total torques for heating at the high field side, $R_C = 2.7$m, and at the low field side, $R_C = 3.3$m, are compared with on axis heating in Fig. 3 at $t = 0.5$s for $n_\phi = 15$ and $n_\phi = -15$. For $n_\phi = 15$ the counter-current torque dominates for all cases independent of the position of the cyclotron resonance. The torque density extends out to the plasma boundary caused by ions transported outwards by the RF-induced spatial drift. For n_ϕ = -15 a counter current torque appears near the centre for off-axis at the low field side and further out the torque reverses. For on-axis and high-field side heating co-current passing ions are formed transferring a major part of the torque through friction by ion-electron collisions.



n_ϕ = 15 n_ϕ = -15

Fig. 3 The radial variation of the integrated total torques of the resonant ions for Λ_{res} = 0.9, R_c = 2.7*m*, (full line), Λ_{res} = 1 R_c = 2.97*m* (dashed line) and Λ_{res}= 1.1, R_c = 3.3*m*, (dotted line).

Symmetric spectra produces a bipolar torque with the strength increasing with $|n_\phi|$. A co-current torque arises in the centre from friction between the electrons and co-current passing high-energy ions whereas the counter-current torque further out arises from the RF-induced outward drift.

4. Conclusions

The drift of resonant ions across magnetic flux surfaces gives rise to a torque on the background plasma. Due to the finite orbit width and the finite Doppler shift of the cyclotron resonance, the inward and outward drifts, and hence, the torques do not cancel even for toroidally symmetric wave spectra. In particular for a symmetric spectrum the drift is directed inwards near the centre and outwards further out. Although the torque is much less than the corresponding one for NBI (100keV), the RF-induced torque is more localised in the centre of the plasma and has often a bipolar structure and can therefore still give rise to a non-negligible shear flow.

References

[1] Eriksson L.-G., et al, *Plasma Physics and Controlled Fusion*, **34**, 863 (1992).
[2] Eriksson L.-G., et al, *Plasma Physics and Controlled Fusion*, **39**, 27 (1997).
[3] Rice, J. E., et al, *Nuclear Fusion,* **39**, 1175 (1999).
[4] Notredaeme, J.-M., et al, *28th EPS Conf on Controlled Fusion and Plasma Physics*, Madeira 2001, P2, 080.
[5] Eriksson L.-G., private communications.
[6] Hellsten, T., *Plasma Physics and Controlled fusion,* **31**, 1391-1406 (1989).
[7] Chang, C. S., et al, *Physics of Plasma*, **6**, 1969-77 (1999).
[8] Chan, V.S., Chiu, S. C., Lin-Liu, Y.R. and Omelchenko, Y. A., in Radio Frequency Power in Plasmas ed by Bernabei, S and Paoletti, F. APS 1-56396- 861-4/99 pp 45.
[9] Perkins, F. W,.White, R. B., Bonoli, P. T., Chan V. S., Plasma Physics, **8**, 2181, (2001)
[10] Hedin, J., Hellsten, T., and Eriksson, L.-G., to appear in *Nuclear Fusion* **42**, (2002).
[11] Hedin, J., Hellsten, T., and Carlsson, J., *Proc. of Joint Varennna-Lausanne Workshop* "*Theory of Fusion Plasmas*" pp. 467, Varenna (1998) ISBN 88-7794-167-7.