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The use of unidirectional lower hybrid (LH) waves to drive currents in plasmas has been well-established in numerous experiments. Two new features of LH current drive have emerged recently. The first, observed on JT-60 [1], was that the LH current drive efficiency increased as the volume averaged electron temperature was increased. The second, observed on JET [2], was that the current drive efficiency increased as the LH waves were used in conjunction with fast Alfvén waves (FAW). The increase in the efficiency occurred at higher volume averaged electron temperatures while at low temperatures the FAW did not seem to affect the LH current drive efficiency [2]. In these experiments the FAW encountered a resonance/mode-conversion layer and part of its energy flow could convert to ion-Bernstein waves (IBW). These must also be considered in understanding the observed increase in current drive efficiency [3,4]. The launched FAW spectrum was that of a monopole antenna so that the FAW/IBW by themselves would not lead to any plasma current. In this paper we consider the possible interactions of the FAW/IBW with the electron tail produced by the unidirectional LH waves. This is done with a view towards understanding JET results and exploring other synergistic current drive possibilities.

Relativistic quasilinear diffusion coefficient

For waves with frequencies, ω, below the electron cyclotron frequency, the relativistic resonant diffusion equation is [5]:

$$\frac{\partial f_0}{\partial t} = \frac{\partial}{\partial p_{||}} D \frac{\partial}{\partial p_{||}} f_0 \tag{1}$$

where f_0 is the space-averaged and gyroangle-averaged distribution function, $p_{||}$ is the momentum along the ambient magnetic field $\vec{B}_0$, and D is the quasilinear diffusion coefficient.

$$D = \frac{\pi}{V} e^2 \int \frac{d^3k}{(2\pi)^3} \delta\left(k_{||} - \frac{\omega}{v_{||}}\right) D_0 \tag{2}$$

where V is the appropriate volume of space over which the averaging takes place and

$$D_0 = \frac{1}{|v_{||}|}\left[\frac{v_\perp^2}{v_{||}^2} J_1^2 |E_{ky}|^2 + J_0^2 |E_{kz}|^2 + \frac{2v_\perp}{v_{||}} J_0 J_1 \mathrm{Im}(E_{ky}^* E_{kz})\right] \tag{3}$$

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$v_\perp = p_\perp/m\gamma$, $v_\parallel = p_\parallel/m\gamma$, $\gamma^2 = 1 + p^2/m^2c^2$, and $\vec{E}_k$ is the Fourier transformed electric field (z is along $\vec{B}_0$). The argument of the Bessel functions is $k_\perp v_\perp/\Omega_e$ where Ω_e is the relativistic electron cyclotron frequency.

D_0 can be expressed in one of two convenient forms: $D_0 = D_{01}|E_{kz}|^2$, and $D_0 = D_{02}|E_{ky}|^2$ where

$$D_{01} = \frac{1}{|v_\parallel|}\left[\left\{J_0 - \frac{v_\perp}{v_\parallel}J_1 \mathrm{Im}\left(\frac{E_{ky}}{E_{kz}}\right)\right\}^2 + \left\{\frac{v_\perp}{v_\parallel}J_1 \mathrm{Re}\left(\frac{E_{ky}}{E_{kz}}\right)\right\}^2\right] \tag{4a}$$

$$D_{02} = \frac{1}{|v_\parallel|}\left[\left\{\frac{v_\perp}{v_\parallel}J_1 + J_0 \mathrm{Im}\left(\frac{E_{kz}}{E_{ky}}\right)\right\}^2 + \left\{J_0 \mathrm{Re}\left(\frac{E_{kz}}{E_{ky}}\right)\right\}^2\right] \tag{4b}$$

D_{01} and D_{02} are completely determined by the local dispersion relation for the appropriate rf wave. The amplitudes $|E_{kz}|^2$ and $|E_{ky}|^2$ are determined from the external rf excitation structure and energy flow spectra into the plasma. For LH waves, (4a) is appropriate, while for FAW (4b) is a suitable form. The term proportional to J_0 corresponds to electron Landau damping (ELD) while the term proportional to J_1 corresponds to damping by transit time magnetic pumping (TTMP). For the FAW these two effects oppose each other; for low $v_\perp$'s ELD dominates over TTMP while at higher $v_\perp$'s TTMP becomes bigger than ELD.

Considering $k_\perp v_\perp/\Omega_e \ll 1$, and assuming that $|E_{ky}^{LH}/E_{kz}^{LH}| \ll 1$ for LH waves and $|E_{kz}^{FW}/E_{ky}^{FW}| \ll 1$ for FAW, it is easy to see that:

$$\frac{D_{02}^{FW}}{D_{01}^{LH}} \approx \left(\frac{v_\perp}{v_\parallel}\right)^2 \left(\frac{k_\perp^{FW} v_\perp}{2\Omega_e}\right)^2 \tag{5}$$

Thus, for comparable energy flow densities, giving $|E_{kz}^{LH}|^2 \approx |E_{ky}^{FW}|^2$, (5) indicates that the FAW diffusion coefficient is small compared to the LH diffusion coefficient. For mode converted ion-Bernstein waves (IBW) the appropriate form would be given also by (4b). But as IBW's propagate away from the mode conversion layer their $|k_\parallel|$ (and $|k_\perp|$) enhances significantly leading to electron Landau damping [6]. The dominant contribution to the diffusion coefficient comes from the J_0 term in (4b). Under JET-like conditions, for example, D_{02} for IBW's is about an order of magnitude smaller than D_{01} for LHW's. The FAW D_{02} is five to six orders of magnitude smaller. We thus must conclude that in JET the FAW by itself can only affect the LHCD current drive by heating the plasma—similar to what was observed in JT60 [1] (but where the plasma heating was by neutral beam injection).

LHCD in the presence of FAW/IBW

The LHCD is relatively peaked off-axis [2],[7]. The small $k_\parallel$ mode-converted spectrum of IBW's, if of sufficient amplitude, can act to extend the current drive spectrum in the tail of the electron distribution function beyond the LHW spectrum. The IBW spectrum that propagates further, upshifts to near thermal velocities and can also heat the plasma. Similarly, the transmitted FAW power that returns to the resonance layer is also mode-converted to IBW's and damps on electrons. Here we are only interested in the interaction of the IBW's with electrons in the LHCD layer and assume that there the IBW's are sufficiently intense. An analytic, nonrelativistic model developed just for the LH waves [8] could then be used to roughly estimate the changes in current drive efficiency. From this model we obtained:

$$\frac{j}{p_d} = \frac{2(v_2^2 - v_1^2)}{(3Z_i + 5)\ell n(a_3) - (Z_i + 1)\ell n(a_1)} \tag{6}$$

where

$$a_m = \frac{v_2^2 + mT_p}{v_1^2 + mT_p} \tag{7a}$$

$$T_p = \frac{T_\perp}{T_e} = \left(\frac{\alpha - 1}{\alpha + 1}\right)\left\{\frac{v_1^{2-\alpha}(v_2^{\alpha+1} - v_1^{\alpha+1}) - v_1^2(v_2 - v_1)}{2(\alpha - 1)(v_2 - v_1) - v_1^{2-\alpha}(v_2^{\alpha-1} - v_1^{\alpha-1})}\right\} \tag{7b}$$

$$\alpha = \frac{2(1 + Z_i)}{2 + Z_i} \tag{7c}$$

Z_i is the effective ion charge of the plasma, T_e is the bulk electron temperature, v_1 and v_2 are the lower and upper bounds, respectively, in parallel velocity (normalized to thermal velocity $(T_e/m_e)^{1/2}$) of the LH tail; j is the normalized current density and p_d is the normalized power dissipated [4,8]. The current drive efficiency is then given by [4]:

$$\eta = n_{20} R_m \left(\frac{I}{P_d}\right)_{A/W} \approx \frac{31.2}{\ell n\Lambda}\left(\frac{j}{p_d}\right)\left(\frac{v_{Te}}{c}\right)^2 \tag{8}$$

where $\ell n\Lambda$ is the Coulomb logarithm, and $v_{Te} = \sqrt{T_e/m_e}$ is the electron thermal velocity.

Assuming that the initital LH tail extends from $v_1 = 4$ to $v_2 = 8$ with $Z_i = 1$, then from (6) and (7): $\frac{j}{p_d} = 26$ and $T_p = 9.6$; for $T_e = 1keV$ this gives $\eta = 0.11$ and $T_\perp = 9.6keV$. Now suppose that the FAW/IBM effectively extend this tail to $v_2 = 12$. Then from (6) and (7) $\frac{j}{p_d} = 48$ and $T_p = 20$; for $T_e = 1keV$ this gives: $\eta = 0.2$ and $T_\perp = 20keV$, while for $T_e = 2keV$ we obtain: $\eta = 0.4$ and $T_\perp = 40keV$. Comparing these we see that an extension of the original LH tail would lead to an increase in the perpendicular temperature of the tail and an enhancement in the current drive efficiency. If, in addition, the bulk electron temperature is increased, then the current drive efficiency is enhanced by a larger factor compared to the case when the bulk electron temperature remains unchanged. There is also a corresponding increase in the tail perpendicular temperature.

A next step beyond this simple model is to go to a 2-D (momentum space) Fokker-Planck code including a quasilinear operator with the coefficient (4), assuming $|E_{ky}^{IB}|^2 \approx |E_{kz}^{LH}|^2$. Using the code CQL3d [7] on one flux surface, we illustrate in Table I the effect (at zero loop voltage) of the IBW under two different conditions. In row five only tail electrons resonate with the IBW, in row six the IBW is allowed to penetrate to $3v_{Te}$ which is somewhat below the LH spectrum lower bound at $v_1 = 3.5v_{Te}$. (In principle, the IBW could go down to about $1.8v_{Te}$ given enough $k_\parallel$-upshift.) The LH "base" case is shown in row one. The main result, seen on comparison of rows one and six, is that the current and current drive efficiency have gone up. We note that this synergism occurs even though the IBW diffusion coefficient is about an order of magnitude below the LH one for the same reference fields. Rows three and four indicate that we have used unidirectional IBW spectra. For obvious reasons, this should be the preferred configuration. For a symmetric IBW spectrum the IBW current from rows three and four should be subtracted form the currents of, respectively, rows five and six. In addition the IBW power dissipated is doubled. The net result for cases five and six is that (I/P) is reduced, respectively, to 0.76 and 0.45.

Table I

	T_{e0} [keV]	v_1^{LH}/v_{Te}	v_1^{IBW}/v_{Te}	$J[A/cm^2]$	$I/P[A/W]$	$T_\perp/T_{e0}$
LH	4	3.5	-	280	0.73	8
LH	7	3.5	-	81	0.87	3.6
IBW	7	-	3.5	10.5	0.57	1
IBW	7	-	3	79	0.5	2.5
LH+IBW	7	3.5	3.5	91	1.04	3.8
LH+IBW	7	3.5	3	304	0.88	4

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