BACKWARD CURRENT AT LOWER HYBRID CURRENT DRIVE BY MEANS OF SHORT PULSES

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In a thermonuclear plasma, lower hybrid waves are absorbed near the plasma boundary. Cohen et al. [1] proposed the use of a train of high power RF pulses, which could penetrate to the plasma core. When the RF pulse is switched on, the fast resonant electrons are accelerated and the RF driven current of density j_d arises. Simultaneously, due to the Faraday's law, an electric field drives the backward Ohmic current of density j_e. The purpose of the present contribution is to give a comprehensive theory of the backward current, viz., to derive its role in the resulting current drive efficiency.

The total current density $j_d + j_e$ parallel to the magnetic field satisfies the skin effect equation

$$\mu_0 \frac{\partial (j_d + j_e)}{\partial t} = \frac{1}{r}\frac{\partial}{\partial r}[r\frac{\partial}{\partial r}(\eta j_e)] \tag{1}$$

where η is plasma resistivity. The characteristic time of current density diffusion is the skin time τ_{sk}

$$\tau_{sk} \approx \mu_0 a^2/\eta, \tag{2}$$

a being the plasma minor radius. Note that $\tau_{sk} \geq 10s$ for present large tokamaks and $\tau_{sk} \approx 10^3 s$ for a reactor. Assume a steady periodic regime of current generation by a train of RF pulses of length τ_p and repetition period τ_r. For the case in question [1], τ_p and τ_r are several orders of magnitude less than τ_{sk}. Consequently, with great accuracy, the total current density is constant and equals the current density j_0 in the time between the RF pulses,

$$j_d + j_e = j_0. \tag{3}$$

Considering the motion of bulk electrons in the induction electric field E, we obtain

$$-E = \frac{1}{\epsilon_0 \omega_{pe}^2}\frac{\partial j_d}{\partial t} + \eta_p\,(j_d - j_0), \tag{4}$$

where η_p is the plasma resistivity with respect to the backward current during τ_p. In the theory of ramp up the poloidal magnetic field, severe restrictions arise due to the runaway electrons accelerated by E [2,3]. Numerical estimates imply that, in our case, $\tau_p \leq 10^{-4}s$ is far too short for the electrons to be accelerated significantly.

The net energy density W_p pumped during τ_p to the poloidal magnetic field is, according to (3) and (4),

$$W_{pol} = -\int_0^{\tau_p} j_0 E dt = j_0 \int_0^{\tau_p} \eta_p\,(j_d - j_0) dt. \tag{5}$$

The energy density W_{pol} is equal to the energy density dissipated during the time without RF. If the integrand in (5) does not change much, we have

$$j_0\eta_p\tau_p(j_d - j_0) = \eta_0 j_0^2(\tau_r - \tau_p), \tag{6}$$

where η_0 is the longitudinal ($\|\mathbf{B}$) Spitzer resistivity. Consequently, for $\tau_r >> \tau_p$,

$$j_d = j_0(1 + \frac{\eta_0\tau_r}{\eta_p\tau_p}). \tag{7}$$

Equation (7) implies that enhanced resistivity η_p improves the current drive efficiency, cf. [2]. The new point here is that, for short powerful RF pulses considered in [1], the value of $j_e/(en_e)$ can approach v_{Te}, the electron thermal velocity. Consequently, the backward current may be unstable, leading to anomalous resistivity η_p. Owing to the slow but non-zero current density diffusion, the actual profile $j_0(r)$ can somewhat differ from that one given by the present theory.

The energy density W_e dissipated by the backward current is calculated similarly to Eqs. (5) and (6). Then, using (7), we obtain

$$W_e = \frac{(j_0\eta_0\tau_r)^2}{\eta_p\tau_p}, \tag{8}$$

a relation needed below.

Let us consider the fast electrons absorbing the RF energy. Suppose that the interval of their velocity components v_a parallel to the magnetic field is very narrow, $v_a \simeq const.$. The absorbed RF power density P_a can then be expressed as

$$P_a = n_a v_a(F_{coll} - eE), \tag{9}$$

where n_a is the concentration of the electrons in question and F_{coll} is the corresponding friction force due to collisions with other particles. We neglect the transient dissipation needed for establishing the nonlinear and collisional deformation of the electron distribution function [1]. The power density $(-eEn_a v_a)$ obviously equals $(W_e + W_{pol})/\tau_p$. Introducing $\varphi_{EC} = -eE/F_{coll}$, using (5), (8) and (9), we obtain

$$\frac{P_a\varphi_{EC}}{1+\varphi_{EC}} = j_0^2\eta_0\frac{\tau_r}{\tau_p}(1 + \frac{\eta_0\tau_r}{\eta_p\tau_p}). \tag{10}$$

Note that φ_{EC} is essentially the ratio of the energies lossed by the resonant electrons due to the induction electric field and due to the collisional drag, respectively (cf. $P_{el}/(P_{in} - P_{el})$ in [2]). Therefore, the term with $\frac{\partial}{\partial t}$ in Eq. (4) can be omitted and

$$\varphi_{EC} = e(j_d - j_0)\eta_p/F_{coll}. \tag{11}$$

The collision time $\tau_e(v_a)$ in the relation $F_{coll} = -mv_a/\tau_e(v_a)$ is, at conditions considered in [1], not much less than τ_r, viz., $\tau_p << \tau_e(v_a) < \tau_r$. It is not known, what happens with the electron distribution function under the action of the train of powerful RF pulses and what will be the actual value of $\tau_e(v_a)$. Therefore, the following considerations are rather estimates.

We use (7) in (11) and express the longitudinal resistivity η_0 in terms of the collision time $\tau_s^{e/i}$ used in the kinetic theory (e.g. [4]). In the steady state, $2-D$ model [5] gives $\tau_e(v_a)$ about a factor 2.5 larger than the $1-D$ model ($\tau_e \sim 1/(2+Z_i)$, e.g. [6]). It is not clear in what degree the $2-D$ effects depending on effective ion charge Z_i will develop under conditions considered. We introduce a factor $a_\varphi(Z_i) \simeq 2$ for $1 \leq Z_i \leq 1.5$ and obtain

$$\varphi_{EC} = b_\varphi j_0; \quad b_\varphi \simeq \frac{0.4 Z_i}{a_\varphi(Z_i)} \frac{\tau_r}{\tau_p} \frac{1}{e n_e v_{Te}} (\frac{v_a}{v_{Te}})^2, \tag{12}$$

where n_e is the concentration of electrons. According to (7) and (12), the explicit dependence of φ_{EC} on η_p is

$$\varphi_{EC} \simeq (1 + \frac{\eta_0 \tau_r}{\eta_p \tau_p})^{-1} \frac{0.4 Z_i}{a_\varphi(Z_i)} \frac{\tau_r}{\tau_p} \frac{v_a^2}{v_{Te}^2} \frac{j_d}{e n_e v_{Te}}. \tag{13}$$

Assume for a moment that $\varphi_{EC} << 1$ and $\eta_p = \eta_0$; Eqs. (10) and (13) then lead to known results, cf. Ref. [2], Eq. (3.7) or Eqs. (2.31) and the text below them.

In general, Eqs. (10) and (12) imply the following relation between j_0 and the averaged RF power density absorbed, $< P_a >_r = P_a \tau_p / \tau_r$:

$$j_0 \simeq -\frac{1}{2 b_\varphi} + [\frac{1}{4 b_\varphi^2} + \frac{< P_a >_r}{\eta_0} (1 + \frac{\eta_0 \tau_r}{\eta_p \tau_p})^{-1}]^{1/2}. \tag{14}$$

Considering the case of ITER-like parameters in [1], we assume that $n_e = 7 \times 10^{19} m^{-3}$, $T_e = 30 keV$, $N_\| = 1.8$, $Z_i = 1.5$, $\eta_p = \eta_0 = 2.74 \times 10^{-10} \Omega m$, $R = 8m$ and the cross-section $\mathcal{S}_0$ of the toroidal current $J_0 = j_0 \mathcal{S}_0$, $\mathcal{S}_0 = 10 m^2$. The corresponding volume $\mathcal{V}_0 = 2\pi R \mathcal{S}_0 \simeq 500 m^3$. Following [1], we suppose that $\mathcal{V}_0 P_a = 9 GW$ and $\mathcal{V}_0 < P_a >_r = 100 MW$. Consequently, we have $\tau_r/\tau_p = 90$, $v_{Te} = 7.26 \times 10^7 m s^{-1}$, $v_a/v_{Te} = 2.3$ and, from (12) with $a(Z_i) = 2$, $b_\varphi \simeq 1.7 \times 10^{-7}$ (in $m^2 A^{-1}$ units). Using (14), we obtain $j_0 \simeq 1.1 \times 10^6 A m^{-2}$, $J_0 \simeq 11 MA$, and $\varphi_{EC} = b_\varphi j_0 \simeq 0.2$. The conventional efficiency is

$$\eta_{CD} = n_e (10^{20} m^{-3}) R J_0 / (\mathcal{V}_0 < P_a >_r) \simeq 0.6. \tag{15}$$

Both the inhomogeneity of the RF power absorption found in [1] and the results in [7] show that the absorbed RF power density in some region of the plasma torus can be considerably higher than its average over the plasma cross-section. Therefore, we choose three times larger $< P_a >_r = 6 \times 10^5 W m^{-3}$. Retaining all the above parameters, we find $j_0 \simeq 2.8 \times 10^6 A m^{-2}$ and $\varphi_{EC} \simeq 0.5$. Note that the mean drift velocity of bulk electrons creating the backward current j_e is about $v_{Te}/3$, approaching the threshold of Buneman's instability. Assume for a moment that, in the case considered, the resistivity for the backward current is anomalous, viz., $\eta_p \simeq 10\ \eta_0$. Using again Eq.(14), we have

$$j_0 \simeq 1.2 \times 10^7 A m^{-2}, \quad \varphi_{EC} \simeq 2. \tag{16}$$

The unnecessarily high j_0 can be reduced, e.g., by diminishing the RF pulse length τ_p. Note that, for $\varphi_{EC} >> 1$, Eq. (10) implies

$$j_0 \simeq [\frac{< P_a >_r}{\eta_0}(1 + \frac{\eta_0 \tau_r}{\eta_p \tau_p})^{-1}]^{1/2}. \quad (17)$$

In this case, almost all the RF power is dissipated by the backward current, $P_a \simeq j_e^2 \eta_p$. Nevertheless, the corresponding value of $n_e(10^{20}) j_0 / (2\pi < P_a >_r)$ can be quite high. The question is, whether such high value of φ_{EC} can be reached in a fusion relevant experiment.

According to the inequality (18) in Ref. [6], the distribution function of resonant electrons is stable with respect to the Parail - Pogutse instability in the case of parameters cosidered here.

Conclusions

The lower hybrid current drive by means of short pulses of high RF power is interesting not only for the wave penetration, but also for the efficiency. Although a considerable part of the RF power is dissipated via the backward current Joule heating, the resulting current drive efficiency is quite good. The reasons are:

- The higher the pulsed RF power, the stronger is the induction electric field and the less is the role of the collisional dissipation (of resonant electrons) in the absorption of RF power;
- the energy pumped during the RF pulse to the magnetic field is used for extremly efficient Ohmic current drive in the time between the RF pulses.

At high anomalous resistance and a small tokamak, τ_{sk} may be less than τ_p. If the well-known $\mathcal{L}/\mathcal{R}$ time of the tokamak is much larger than τ_r, the above considerations can be repeated *mutatis mutandis* for the plasma torus as a whole.

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References

[1] R.H. Cohen et al.: 13th Conf. Plasma Phys. and Contr. Fusion (1990), Vol. 1, 805.

[2] N.J. Fisch: Rev. Mod. Phys. **59** (1987) 175.

[3] Ya.I. Kolesnichenko, V.V. Parail and G.V. Pereverzev: in Problems of Plasma Theory, Vol. 17, ENERGOATOMIZDAT, Moscow 1989 (in Russian).

[4] B.A. Trubnikov: in Reviews of Plasma Phys., Vol. 1, Consultants Bureau, New York 1965.

[5] C.F.F. Karney and N.J. Fisch: Phys. Fluids **22** (1979) 1817.

[6] R. Klíma and A.V. Longinov: Sov. J. Plasma Phys. **5** (1979) 277; in the notation used there, the factor $(\omega_{pe}^2 + \omega_{He}^2)/[\omega_{He}(\omega_{pe}^2 - \Omega^2)]$ must be supplemented in the right-hand side of inequality (18) or, approximately, ω_{He}/ω_{pe} must be instead of Ω there.

[7] L. Krlín, P. Pavlo and V. Petržílka: *This Conference*.