J. Plasma Physics (1999), *vol.* 62, *part* 2, *pp.* 203–218. Printed in the United Kingdom

Transport of intense LH pulses into a tokamak plasma

L. KRLÍN,[1] R. KLÍMA,[1] P. PAVLO,[1] V. PETRŽÍLKA,[1] V. SVOBODA[2] and J. A. TATARONIS[3]

[1] Institute of Plasma Physics, Academy of Sciences of the Czech Republic, 182 21 Prague 8, PO Box 17, Czech Republic

[2] Czech Technical University, Faculty of Nuclear Sciences and Physical Engineering, Department of Physics, Břehová 7, 115 19 Prague 1, Czech Republic

[3] University of Wisconsin–Madison, WI 53706-1687, USA

(Received 15 October 1998 and in revised form 5 March 1999)

Abstract. Application of lower-hybrid (LH) power in short, intense pulses in the 5–10 GW range should overcome the limiting effects of Landau damping, and thereby permit the penetration of LH power into the interior of large-scale plasmas. We show that, under such very intense LH pulses, wave coupling may deteriorate because of nonlinear density changes due to the ponderomotive force effects in front of the grill. Ponderomotive forces are also likely to induce strong plasma bias and consequent poloidal and toroidal plasma rotation. Although backward electric currents, created in the plasma by intense LH pulses, dissipate a large portion of the radio frequency power absorbed, the current drive efficiency is acceptable. We use a numerical simulation of wave–particle interactions to analyse the applicability of standard quasilinear theory to the case of large energy flux densities. The initial results indicate the existence of important restrictions on the use of the quasilinear approximation. The results of the present paper also indicate that some of the effects considerably alter some ideas of Cohen et al.

1. Introduction

In a thermonuclear tokamak plasma, lower-hybrid (LH) waves are strongly absorbed at the plasma boundary (see e.g. Devoto and Fenstermacher 1990; Pavlo et al. 1991), thereby inhibiting their penetration into the plasma core. To overcome this limitation, Cohen et al. (1990) proposed the use of a train of intense short pulses instead of a continuous launching of LH waves. A pulse power as high as 10 GW for two pulse durations of 10^{-4} s and 10^{-8} s, with an averaged power of 100 MW, was considered by Cohen et al. Aspects of this approach were also discussed by Bertrand et al. (1994). This proposed method to improve the penetration of LH waves raises a number of issues that should be studied thoroughly. Moreover, there are in this problem a variety of novel and important physical issues that, in our opinion, have a broader significance. The purpose of the present paper is to extend and improve the analysis of Cohen et al. (1990) with respect to the main physical phenomena governing the pulsed regime of current drive. In a certain sense, our analysis complements the study and results of Cohen et al. (1990). However, we find results that can considerably modify the intriguing proposal that they reported.

The plan of this paper is as follows. In Sec. 2, we discuss the influence of nonlinear effects on wave coupling. We find that efficient coupling can be achieved either for a sufficiently high temperature of the boundary plasma (30 eV) or at lower temperatures with ultrashort pulses ($\tau_p \approx 30$ ns). In Sec. 3, estimates of the magnitude of the radial electric field generated by the poloidal ponderomotive forces are given. Plasma rotation is suggested by the results we report here. Owing to Faraday's law, the pulsed regime of current drive is inevitably accompanied by the generation of a backward current. This effect is analysed in Sec. 4, with particular emphasis being placed on the resulting current drive efficiency. Section 5 addresses the problem of applicability of the quasilinear approximation for the case of intense wave fluxes. The differences that we find between a direct numerical simulation of wave–particle interactions and the quasilinear approximation calls for a more complete investigation of this problem. In Sec. 6, our results are summarized and discussed.

2. Nonlinear wave coupling

At large LH power levels, toroidal ponderomotive forces in front of an antenna structure may expel plasma from the space near the grill mouth and thus reduce the plasma density. This results in nonlinear changes in the wave coupling. Because of the ponderomotive forces, the boundary plasma density n_b in front of the grill decreases, as suggested by the following expression (Petržílka et al. 1991):

$$n_b = n_0 \exp(-\delta), \tag{1}$$

$$\delta = \frac{W}{T_b}, \tag{2}$$

where W is the ponderomotive potential of the LH wave,

$$W = \frac{e^2 E_0^2}{4 m_e \omega^2}. \tag{3}$$

In (2) and (3), T_b is the sum of boundary temperatures of electrons and ions, e is the charge, m_e is the mass of the electron, ω is the frequency of the LH field, and E_0 is the LH electric field component parallel to the magnetostatic field. The radio frequency (RF) power density flux from the grill mouth into the plasma can be expressed in the form

$$S = \frac{\mathrm{Im}\left(E_0 \dfrac{\partial E_0^*}{\partial x}\right)}{2\mu_0 \omega (1 - N_\parallel^2)}, \tag{4}$$

where $N_\parallel = k_\parallel c/\omega$. Under the assumption that the reflection coefficient R at the plasma boundary is much less than one, we find the following relation between the maximum value of the ponderomotive potential W and S:

$$W_{\max} \approx \frac{\mu_0 c e^2 N_\parallel S}{2 m_e \omega \omega_{pe}}. \tag{5}$$

In eV units, this expression becomes

$$W_{\max}[eV] \approx 8 \times 10^{11} \frac{N_\parallel S}{f_p f}. \tag{6}$$

https://doi.org/10.1017/S0022377899007850 Published online by Cambridge University Press

If we choose the parameters of Cohen et al. (1990), i.e. $N_\parallel = 1.8$, $S = 0.5$ GW m^{-2}, $f = 8$ GHz, with $f_p \geqslant f$, corresponding to $n_b \geqslant 8 \times 10^{17}$ m^{-3}, then (6) yields $W_{\max} \leqslant 10$ eV. This moderate value of $W_{\max}$, which we have also confirmed with numerical computations, is a consequence of the relatively high frequency that we have assumed. As a consequence of this result, the possible deterioration of wave–plasma coupling in 'regime B' of Cohen et al. (1990), with longer pulses ($\tau_p \approx 80$ μs), can be avoided if the boundary plasma temperature T_b exceeds 30 eV. This elevated temperature could arise, for example, from collisional heating and/or parametric instabilities.

Let us now analyse 'regime A' of Cohen et al., with short pulses. The displacement of plasma $z(t)$ along the magnetic field lines owing to the ponderomotive force depends on the ion inertia:

$$m_i \frac{d^2 z}{dt^2} = -\frac{\partial W}{\partial z}. \tag{7}$$

The characteristic time τ_i of the plasma displacement can be approximated as

$$\tau_i \approx \left(\frac{\varepsilon_{0i} L_i L_W}{c^2 W_{\max}}\right)^{1/2}, \tag{8}$$

where $\varepsilon_{0i} = m_i c^2$ is the ion rest energy, and L_i and L_W are respectively the characteristic lengths of the plasma displacement and inhomogeneity of $W(z)$. It is natural to set

$$L_i \approx L_W \approx \frac{1}{k_\parallel} \equiv \frac{c}{\omega N_\parallel}. \tag{9}$$

Then, for a deuterium plasma with $f = 8$ GHz, $N_\parallel = 1.8$ and $W_{\max} \approx 10$ eV, we have $\tau_i \approx 10^{-7}$s. Consequently, the ponderomotive effect is negligible for 'regime A' of Cohen et al. (1990), with short pulses ($\tau_p \approx 30$ ns).

In general, the reduction in wave coupling could be weakened if local plasma heating occurs. We now explore this possibility in more detail. We assume that the boundary plasma temperature T_b increases with increasing launched LH power S, which is consistent with observations on the ASDEX tokamak (Petržílka et al. 1991), in the form given by the following expression:

$$T_b = T_0\left(1 + \frac{S}{S_{\mathrm{T}}}\right). \tag{10}$$

At large values of T_b, and correspondingly higher plasma pressures, the ponderomotive forces are not strong enough to expel plasma with elevated pressure from the space in front of the grill, and therefore to deteriorate the wave coupling.

Consider now a very long grill launching a very narrow spectrum of waves. For this launching configuration, it is sufficient to treat only waves with one k_z, or equivalently a single value of $N_\parallel$. Neglecting higher spatial harmonics, we make the ansatz that the electric field has the form of two oppositely propagating waves:

$$E_z(x,z) = E_1^{(+)}(x)\exp(ik_z z) + E_1^{(-)}(x)\exp(-ik_z z). \tag{11}$$

The governing equation for $E_1^{(\pm)}$ then becomes (Petržílka et al. 1991),

$$\frac{d^2 E_1^{(\pm)}}{dx^2} + (k_0^2 - k_z^2)E_1^{(\pm)} = \frac{n_0(x)}{\lambda n_c}(k_0^2 - k_z^2)\int_0^\lambda \exp[\pm ik_z z - \delta(x,z)]E_z(x,z)\,dz, \tag{12}$$

where $\lambda = 2\pi/k_z$ and $k_0 = \omega/c$. We have solved (12) numerically for $E_1^{(+)}$ and $E_1^{(-)}$

https://doi.org/10.1017/S0022377899007850 Published online by Cambridge University Press

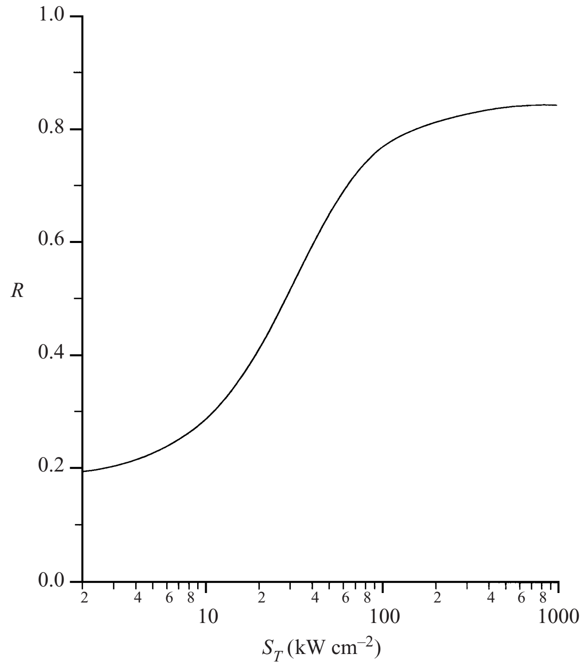

Figure 1. Dependence of the reflection coefficient R on the heating rate S_T at which the boundary temperature T_b doubles. The launched power S = 50 kW cm^{-2}, $N_\parallel$ = 2, the wave frequency f = 8 GHz, the initial boundary temperature T_0 = 10 eV, the initial boundary density normalized to the critical density is $n_0/n_c = 2$, and $n_c = 7.93 \times 10^{17}$ m^{-3}. We note that for low reflection coefficients of about R = 0.2, this launched power of 50 kW cm^{-2} corresponds to a wave electric field amplitude of 5.5 kV cm^{-1} in front of the grill for $N_\parallel$ = 2.

with boundary conditions deep enough inside the plasma where ponderomotive forces are negligible. The RF electromagnetic fields computed in this way can be used to compute the wave reflection coefficient $R_w(z)$. The power reflection coefficient R, averaged over z, is given by (Petržílka et al. 1991)

$$R = \left[\int_0^\lambda \frac{S(z)}{1-|R_w^2(z)|}\,dz\right]^{-1} \int_0^\lambda S(z)\frac{|R_w^2(z)|}{1-|R_w^2(z)|}\,dz, \tag{13}$$

where $S(z)$ is the x component of the Poynting vector of the LH wave transmitted into the plasma.

If we use (10), we find, for example, that the boundary temperature with $\bar{S} = S_\mathrm{T}$ is twice the temperature with zero LH power, $\bar{S} = 0$. Here $\bar{S}$ denotes the energy flux in front of the grill, $S(z)$, averaged over the toroidal coordinate z. If $\bar{S} \gg S_\mathrm{T}$, the resulting boundary temperature is much higher than the temperature with no LH power. On the other hand, for $\bar{S} \ll S_\mathrm{T}$, the boundary temperature practically does not change as the LH power increases.

For ASDEX, the best fit of the nonlinear reflection curves to experimental data was obtained with S_T = 2 kW cm^{-2} and launched powers up to 4 kW cm^{-2}. Since the typical launched LH power of intense LH pulses would be much higher, about 50 kW cm^{-2}, the corresponding values of S_T would likely also be higher, as assumed in Figs 1 and 2. Figure 1 shows the influence of the value of S_T on the reflection coefficient R, while Fig. 2 shows the effects of S_T on δ. In Fig. 1, we see that the

https://doi.org/10.1017/S0022377899007850 Published online by Cambridge University Press

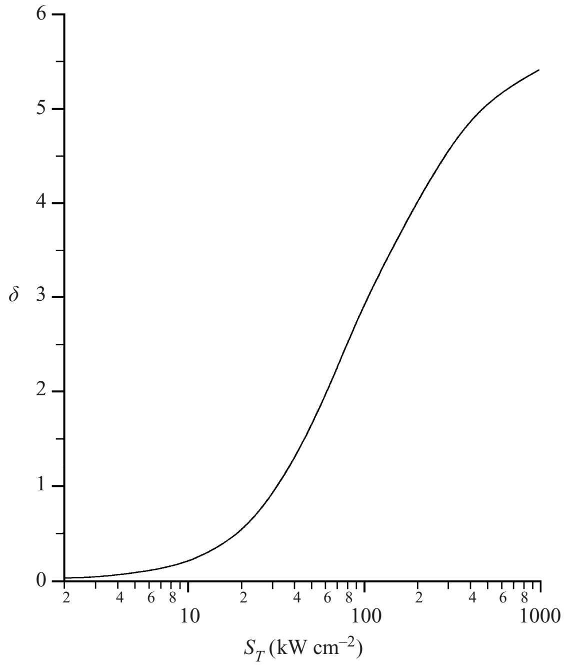

Figure 2. Dependence of the logarithmic boundary density depression $\delta = -\ln(n_b/n_0)$ on the heating rate S_T at which the boundary temperature T_b doubles, for the same parameters as in Fig. 1.

reflection coefficient increases when S_T increases. The reason is that for higher S_T, and therefore for lower boundary temperatures T_b, the quantity δ grows (cf. Fig. 2). According to (1), the plasma density decrease in front of the grill is then stronger, which leads to deterioration of the coupling and to an increase of R.

Resonant electron interactions with strong wave electric fields in front of the grill – either regular (Fuchs et al. 1996) or random (Tataronis et al. 1997) fields – result in strong electron acceleration. This may lead to very high thermal loads on wall components. Nevertheless, this additional resonant acceleration may also further enhance the plasma temperature in front of the grill, and consequently reduce the ponderomotive deterioration of the wave–plasma coupling.

3. Variations in the plasma bias and rotation

As a consequence of wave momentum dissipation, a strong pulsed wave can also exert a strong *poloidal* ponderomotive force (Van Nieuwenhove et al. 1995), in addition to the gradient ponderomotive forces in the axial and toroidal directions. Because of the presence of the strong magnetostatic toroidal field in a tokamak, poloidal forces result in the appearance of strong radial electric fields, which in turn produce plasma rotation. Poloidal ponderomotive forces in front of LH grills would likely arise from wave propagation in the poloidal direction with respect to the toroidal magnetic field. Poloidal wave propagation is a possibility if the mutual phasing of the horizontal waveguide rows of the grill were of a suitable value.

Expressions for the time-averaged radial electric field can be derived from the generalized Ohm's law of the plasma. Assuming a cylindrical plasma model with

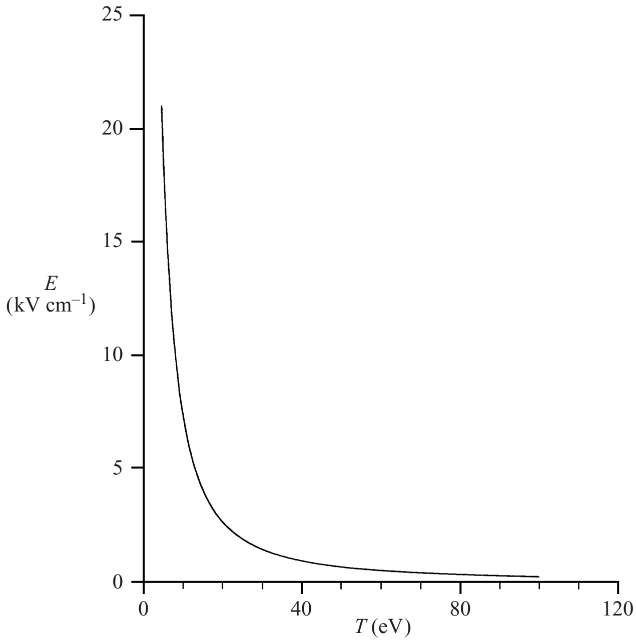

Figure 3. Radial electric field E_r induced by poloidal ponderomotive forces in front of the LH grill as a function of the boundary plasma temperature. The boundary plasma density $n_b = 2 \times 10^{18}$ m^{-3}, the wave frequency $f = 5$ GHz, the toroidal magnetostatic field $B_z = 5$ T, the poloidal magnetostatic field $B_\theta = 0.5$ T, the wave field profile in front of the grill is assumed in the form $[0.1 + (r/a)^s]^t$, $s = 2$, $t = 5$, the plasma minor radius $a = 2$ m, and the wave field amplitude in front of the grill is taken as 5.5 kV cm^{-1}, which corresponds to a coupled wave power of the order of tens of kW cm^{-2}, depending on the wave reflection coefficient.

coordinates (r, θ, z), we let $F^P_{\alpha,\theta}$ and $F^P_{\alpha,z}$ represent the azimuthal and axial components of the LH ponderomotive force $\mathbf{F}^P$. For an electron–ion plasma, $\mathbf{F}^P = \mathbf{F}^P_e + \mathbf{F}^P_i$, where subscripts e and i label electron and ion components respectively. The induced time-averaged radial electric field can be expressed as a sum of two terms (Klíma and Petržílka 1980; Petržílka et al. 1997):

$$E_{0r} = \frac{1}{en_0}\frac{\partial p_i}{\partial r} + \mathscr{E}^P, \tag{14}$$

where p_i designates the scalar partial pressure of the ion fluid, and $\mathscr{E}^P$ represents the component of E_{0r} induced directly by the ponderomotive forces,

$$\mathscr{E}^P = -\frac{B_{0z}}{m_i r^2 n_0 U_{ir}} \int_0^{\mathrm{r}} d\bar{r}\, \bar{r}^2 F^P_\theta + \frac{B_{0\theta}}{m_i r n_0 U_{ir}} \int_0^{\mathrm{r}} d\bar{r}\, \bar{r} F^P_z . \tag{15}$$

Here n_0 is the time-averaged plasma density and U_{ir} is the radial component of the mean ion fluid velocity. According to (12), the value that $\mathscr{E}^P$ has at a radial position r depends on the values of two definite integrals from the plasma centre ($\bar{r} = 0$), to $\bar{r} = r$. However, because of the nature of the dissipation processes in the plasma and the geometry of the LH cones, the LH-wave electric field and the associated ponderomotive force attain their largest values near the grill region. Figures 3–5 show the results of a numerical evaluation of (14). They show that, for intense LH-

https://doi.org/10.1017/S0022377899007850 Published online by Cambridge University Press

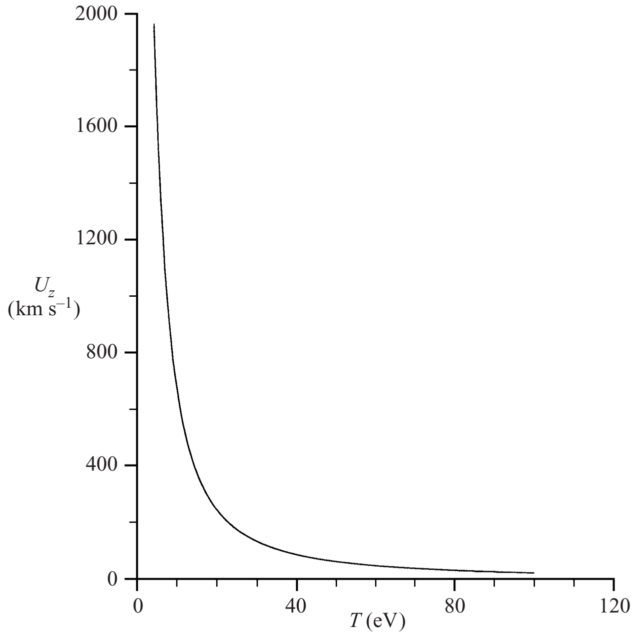

Figure 4. The toroidal velocity U_z of plasma rotation induced by poloidal ponderomotive forces in front of the LH grill as a function of the boundary plasma temperature for the same parameters as in Fig. 3.

wave pulses, the poloidal ponderomotive forces may induce large stationary radial electric fields, up to about 10 kV cm^{-1}.

An electric field of this magnitude has obvious implications regarding plasma biasing and confinement. This effect can be used for changing the plasma potential, and in this way for inducing higher confinement modes (H modes). A great advantage of this method is that there is no need for electrodes in the plasma environment, which are currently used for changing the plasma bias and inducing H modes in tokamaks. Further, it is possible to easily change also the sign of the plasma bias, since the sign of the induced radial electric field depends on the sign of the poloidal wavenumber, which is determined by the value of the mutual poloidal phasing of the horizontal waveguide rows. Of course, the problem of the pulsed character of the biasing electric field remains open.

The associated induced radial plasma flow is rather small, and therefore has little impact on the plasma density in front of the LH grill and on the wave coupling. However, the plasma bias may influence the fluctuations, and thus indirectly affect the radial transport and the plasma density in front of the LH grill, and therefore also the scattering of LH wave on the density fluctuations at the plasma boundary (Andrews et al. 1985; Petržílka 1988).

A more detailed study of these mutually coupled nonlinear effects is beyond the scope of the present paper.

In this connection, we only note that the relative amount of wave energy scattered by density fluctuations just at the plasma boundary is acceptably low, in accordance with the results of Cohen et al. (1990). However, the poloidal wave vector is changed by the scattering process, and this significantly influences the wave

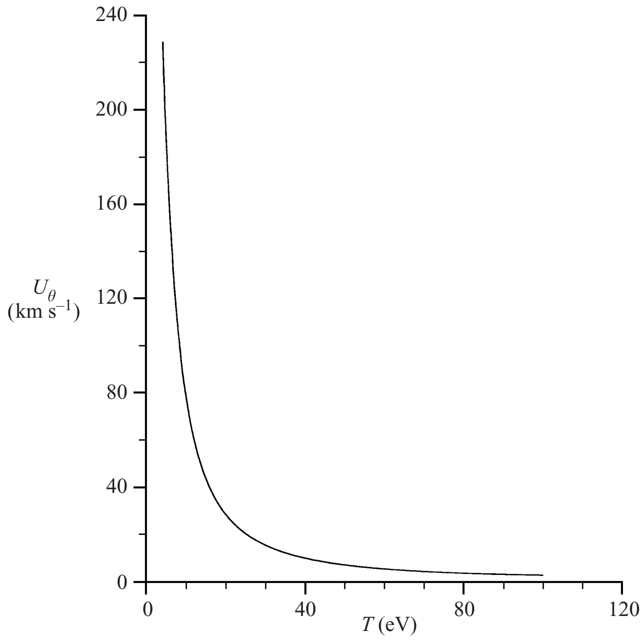

Figure 5. The poloidal velocity U_θ of plasma rotation induced by poloidal ponderomotive forces in front of the LH grill as a function of the boundary plasma temperature for the same parameters as in Fig. 3.

propagation characteristics in the subsequent wave propagation in the sheared magnetostatic field (Krlín et al. 1998).

4. Backward current

When the RF pulse is switched on, fast resonant electrons are accelerated and an RF-driven current of density j_d arises. Simultaneously, owing to Faraday's law, an electric field drives a backward Ohmic current of density j_e. The total current density $j_d + j_e$ parallel to the magnetic field satisfies the skin-effect equation

$$\mu_0 \frac{\partial(j_d + j_e)}{\partial t} = \frac{1}{r}\frac{\partial}{\partial r}\left[r\frac{\partial}{\partial r}(\eta j_e)\right], \tag{16}$$

where η is the plasma resistivity. The characteristic time of the resulting current density diffusion is the skin time $\tau_{\rm sk}$, which can be expressed as

$$\tau_{\rm sk} \approx \mu_0 a^2/\eta, \tag{17}$$

a being the plasma minor radius. For present-day large tokamaks, $\tau_{\rm sk} \geqslant 10$ s, while for a reactor plasma, $\tau_{\rm sk} \approx 10^3$ s. 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 (Cohen et al. 1990), τ_p and τ_r are several orders of magnitude less than $\tau_{\rm sk}$. Consequently, with great accuracy, the total current density is constant and equals the current density j_0 in the time between the RF pulses:

$$j_d + j_e = j_0. \tag{18}$$

https://doi.org/10.1017/S0022377899007850 Published online by Cambridge University Press

Considering the motion of bulk electrons under the influence of the induction electric field E, we find

$$-E = \frac{1}{\epsilon_0 \omega_{pe}^2} \frac{\partial j_d}{\partial t} + \eta_p \, (j_d - j_0), \tag{19}$$

where η_p is the plasma resistivity with respect to the backward current during the time interval τ_p. In the theory of 'ramp-up' of the poloidal magnetic field, severe restrictions arise owing to the runaway electrons accelerated by the electric field E (Fisch 1987; Kolesnichenko et al. 1989). Numerical estimates imply that in our case, τ_p, which is less than 10^{-4} s, is far too short for the electrons to be accelerated significantly.

The net energy density $W_{\rm pol}$ pumped during τ_p to the poloidal magnetic field is, according to (18) and (19),

$$W_{\rm pol} = -\int_0^{\tau_p} j_0 E \, dt = j_0 \int_0^{\tau_p} \eta_p \, (j_d - j_0) \, dt. \tag{20}$$

The energy density $W_{\rm pol}$ is equal to the energy density dissipated during the time without RF. If the integrand in (20) does not change significantly, we have

$$j_0 \eta_p \tau_p (j_d - j_0) = \eta_0 j_0^2 (\tau_r - \tau_p), \tag{21}$$

where η_0 is the Spitzer resistivity parallel to $\mathbf{B}$. Consequently, for $\tau_r \gg \tau_p$,

$$j_d = j_0 \left(1 + \frac{\eta_0 \tau_r}{\eta_p \tau_p}\right). \tag{22}$$

Equation (22) implies that enhanced resistivity η_p improves the current drive efficiency (cf. Fisch 1987). The new point here is that, for the short powerful RF pulses considered by Cohen et al. (1990), the value of j_e/en_e can approach v_{Te}, the electron thermal velocity. Consequently, the backward current may be unstable, leading to an anomalous resistivity η_p. Owing to the slow but non-zero current density diffusion, the actual profile $j_0(r)$ can differ somewhat from that given by the present theory.

The energy density W_e dissipated by the backward current is calculated similarly to the calculation with (20) and (21). Thus, with (22), we find

$$W_e = \frac{(j_0 \eta_0 \tau_r)^2}{\eta_p \tau_p}, \tag{23}$$

which is a relation that we shall need below.

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

$$P_a = n_a v_a (F_{\rm coll} - eE), \tag{24}$$

where n_a is the density of the absorbing electrons in question and $F_{\rm coll}$ 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 (Cohen et al. 1990). The power density $(-eEn_a v_a)$ obviously equals $(W_e + W_{\rm pol})/\tau_p$. Introducing $\varphi_{EC} = -eE/F_{\rm coll}$, and using (20), (23)

https://doi.org/10.1017/S0022377899007850 Published online by Cambridge University Press

and (24), we find

$$\frac{P_a\varphi_{EC}}{1+\varphi_{EC}} = j_0^2\eta_0\frac{\tau_r}{\tau_p}\left(1+\frac{\eta_0\tau_r}{\eta_p\tau_p}\right). \tag{25}$$

Note that φ_{EC} is essentially the ratio of the energies lost 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 Fisch 1987). Therefore the term proportional to the time derivative of j_d in (19) can be omitted. This simplification yields

$$\varphi_{EC} = \frac{e(j_d-j_0)\eta_p}{F_{\text{coll}}}. \tag{26}$$

The collision time $\tau_e(v_a)$ in the relation $F_{\text{coll}} = -mv_a/\tau_e(v_a)$ is, with the conditions considered by Cohen et al. (1990), not much less than τ_r, i.e. $\tau_p \ll \tau_e(v_a) < \tau_r$.

It is unknown how the electron distribution function will be affected by a train of powerful RF pulses and what the actual value of $\tau_e(v_a)$ will be. Therefore the following analysis should be viewed as an approximation. We substitute (22) into (26) and express the longitudinal resistivity η_0 in terms of the collision time $\tau_s^{e/i}$ used in kinetic theory (see e.g. Trubnikov 1965). In the steady state, the plasma 2D model (Karney and Fisch 1979) reveals that $\tau_e(v_a)$ is about a factor 2.5 larger than in the 1D model ($\tau_e \propto 1/(2+Z_i)$, see e.g. Klíma and Longinov 1979). It is not clear to what degree the 2D effects, depending on effective ion charge Z_i, will develop under the conditions considered. We now introduce a factor $a_\varphi(Z_i) \approx 2$, with $1 \leqslant Z_i \leqslant 1.5$. This yields

$$\varphi_{EC} = b_\varphi j_0, \tag{27a}$$

$$b_\varphi \simeq \frac{0.4Z_i}{a_\varphi(Z_i)}\frac{\tau_r}{\tau_p}\frac{1}{en_e v_{Te}}\left(\frac{v_a}{v_{Te}}\right)^2, \tag{27b}$$

where n_e is the density of electrons. According to (22) and (27a,b), the explicit dependence of φ_{EC} on η_p is

$$\varphi_{EC} \approx \left(1+\frac{\eta_0\tau_r}{\eta_p\tau_p}\right)^{-1}\frac{0.4Z_i}{a_\varphi(Z_i)}\frac{\tau_r}{\tau_p}\frac{v_a^2}{v_{Te}^2}\frac{j_d}{en_e v_{Te}}. \tag{28}$$

Assume for the moment that $\varphi_{EC} \ll 1$ and $\eta_p = \eta_0$. Equations (25) and (28) then lead to known results (cf. Fisch 1987, equation (3.7) or (2.31) and the text below them).

In general, (25) and (27) imply the following relation between j_0 and the averaged RF power density absorbed, $\langle P_a\rangle_r = P_a\tau_p/\tau_r$,

$$j_0 \approx -\frac{1}{2b_\varphi} + \left[\frac{1}{4b_\varphi^2} + \frac{\langle P_a\rangle_r}{\eta_0}\left(1+\frac{\eta_0\tau_r}{\eta_p\tau_p}\right)^{-1}\right]^{1/2}. \tag{29}$$

Considering the case of ITER-like parameters studied by Cohen et al. (1990), we assume that $n_e = 7\times10^{19}$ m^{-3}, $T_e = 30$ keV, $N_\parallel = 1.8$, $Z_i = 1.5$, $\eta_p = \eta_0 = 2.74\times10^{-10}$ Ω m, $R = 8$ m, and a cross-section $\mathscr{S}_0 = 10$ m^2 of the toroidal current $J_0 = j_0\mathscr{S}_0$. The corresponding volume $\mathscr{V}_0 = 2\pi R\mathscr{S}_0 \approx 500$ m^3. Following Cohen et al. (1990), we suppose that $\mathscr{V}_0 P_a = 9$ GW and $\mathscr{V}_0\langle P_a\rangle_r = 100$ MW. Consequently, we have $\tau_r/\tau_p = 90$, $v_{Te} = 7.26\times10^7$ m s^{-1}, $v_a/v_{Te} = 2.3$, and, from (27) with $a(Z_i) = 2$, $b_\varphi \approx 1.7\times10^{-7}$ (in m^2 A^{-1} units). Using (29), we obtain $j_0 \approx 1.1\times10^6$ A m^{-2}, $J_0 \approx 11$ MA and $\varphi_{EC} = b_\varphi j_0 \approx 0.2$. The conventional efficiency is

$$\eta_{CD} = \frac{n_e(10^{20}\ \mathrm{m}^{-3})RJ_0}{\mathscr{V}_0\langle P_a\rangle_r} \approx 0.6. \tag{30}$$

Both the inhomogeneity of the RF power absorption found in Cohen et al. (1990) and our results in Sec. 5 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 assume here that it is three times larger: $\langle P_a\rangle_r = 6\times10^5$ W m^{-3}. Retaining all the above parameters, we find $j_0 \approx 2.8\times10^6$ A m^{-2} and $\varphi_{EC} \approx 0.5$. Note that the mean drift velocity of bulk electrons creating the backward current j_e is about $\frac{1}{3}v_{Te}$, approaching the threshold of Buneman instability. Assume for a moment that in the case considered, the resistivity for the backward current is anomalous, namely $\eta_p \approx 10\,\eta_0$. Using (29) again, we have

$$j_0 \approx 1.2\times10^7\ \mathrm{A\ m}^{-2}, \qquad \varphi_{EC} \approx 2. \tag{31}$$

The unnecessarily high j_0 can be reduced, for example, by diminishing the RF pulse length τ_p. Note that for $\varphi_{EC} \gg 1$, (25) implies

$$j_0 \approx \left[\frac{\langle P_a\rangle_r}{\eta_0}\left(1+\frac{\eta_0\tau_r}{\eta_p\tau_p}\right)^{-1}\right]^{1/2}. \tag{32}$$

In this case, almost all the RF power is dissipated by the backward current, $P_a \approx j_e^2\eta_p$. Nevertheless, the corresponding value of $n_e(10^{20})j_0/(2\pi\langle P_a\rangle_r)$ can be quite high, because the Ohmic current drive acting in the time intervals between the RF pulses is very efficient. The question is whether such a high value of φ_{EC} can be reached in a fusion-relevant experiment.

According to the inequality specified by (18) in Klíma and Longinov (1979), the distribution function of the resonant electrons is stable with respect to the Parail–Pogutse instability for the specific parameters considered here.

5. Applicability of the quasilinear approximation for the case of intense wave pulses

The quasilinear approach is considered to be an excellent tool for the description of LH-wave–plasma interaction. Nevertheless, the quasilinear approximation itself has been developed on the basis of a perturbation analysis, i.e. on the assumption that the changes in particle velocities during the wave–particle interaction are small.

In case of large wave power fluxes, a possible change of the energy of particles during their single transit through the LH wave cone can easily constitute a significant fraction of their original energy. This makes the reliability of the quasilinear approach questionable. Since the model of Cohen et al. (1990) of the interaction of intensive pulses with the thermonuclear plasma depends in some degree on the quasilinear description (QLD), it is advisable to test the validity of the QLD by means of direct numerical simulation of this interaction.

The simulation that we carried out is based on the equations of motion of particles in the tokamak geometry for a prescribed form of the launched LH wave spectrum. Our earlier numerical code, successfully used already for other plasma waves (Krlín et al. 1997), has been employed. This code is based on the Hamiltonian formalism, which enables us to take into account all features of the particle dynamics.

The Hamiltonian describing the motion of a particle in the tokamak magnetic field (with circular cross-sections of magnetic surfaces) and in the fields of an LH wave, under the electrostatic approximation, is

$$H = \omega_{c0} P_1 \left[1 - \frac{r(P_2)}{R_0}\cos\tilde{\beta}\right] + \frac{P_3^2}{2m}\left[1 - 2\frac{r(P_2)}{R_0}\cos\tilde{\beta}\right] + e\Psi_0 \cos\left[k_r\left(\frac{2P_2}{eB_0}\right)^{1/2} + m_p\left(Q_2 + \frac{Q_3}{qR_0}\right) + \frac{k_\parallel Q_3}{\left(1 - 2\dfrac{r}{R_0}\cos\tilde{\beta}\right)^{1/2}} - \omega t\right]. \quad (33)$$

The definitions of the canonical coordinates and of the other symbols in (33) are as follows:

$$P_1 = \tfrac{1}{2} eB_0\rho_c^2, \qquad P_2 = \tfrac{1}{2} eB_0 r^2,$$
$$P_3^2\left(1 - 2\frac{r}{R_0}\cos\tilde{\beta}\right) = m_e^2 v_\parallel^2,$$
$$Q_1 = \omega_{c0} t, \qquad \tilde{\beta} = \theta = Q_2 + \frac{Q_3}{qR_0},$$
$$Q_3 = R_0\phi, \qquad r = \left(\frac{2P_2}{eB_0}\right)^{1/2}.$$

R_0 and a are respectively the major and minor radii of the tokamak, ρ_c and r are respectively the Larmor and guiding-centre radii, θ and ϕ are respectively the poloidal and toroidal angles, e is the particle charge, and m_e is the particle mass. Ψ_0, k_r, m_p and n_t are the wave amplitude, the radial wave vector component, and the poloidal and toroidal wavenumbers respectively, and ω is the angular frequency of the wave.

For simplicity, we have assumed a rectangular LH wave spectrum centred around $k_\parallel$, with full width $\Delta k_\parallel$. The continuous spectrum is replaced by an equidistant discrete spectrum with M modes of equal potentials $\Psi_m = \Psi_0 M^{-1/2}$. The amplitudes of the spectrum were determined by the total wave power flux S, which, for a narrow wave spectrum, can be expressed as

$$|E_\parallel| = \left(\sum_{i=1}^{M} E_i^2\right)^{1/2} = \left(2\mu_0 c N_\parallel \frac{\omega}{\omega_{pe}} S\right)^{1/2}. \quad (34)$$

For example, in the case of a single wave, with the parameters that have been assumed by Cohen et al. (1990), i.e. with frequency $f = 8$ GHz, $N_{\parallel,0} = 1.8$, energy flux density 0.5 GW m^{-2} and density $n_e = 10^{20}$ m^{-3}, (34) yields $E_\parallel \approx 3\times10^5$ V m^{-1}, which corresponds to $\Psi_0 = 10^3$ V. For M waves, $E_i \approx 3\times10^5 M^{-1/2}$ V m^{-1}.

The spatial distribution of LH waves can be obtained by ray tracing. However, at the present stage, our primary interest is more a qualitative than an accurate quantitative analysis of the LH-wave–plasma interaction. Therefore, instead of LH cones, we simply assume that some portions of the plasma volume are filled with the RF field. These are defined as N toroidal segments of length $l = R_0\Delta\phi$ and height $h = a\Delta\theta$. Here Nlh corresponds to the considered grill area, which in our case is 18 m^2.

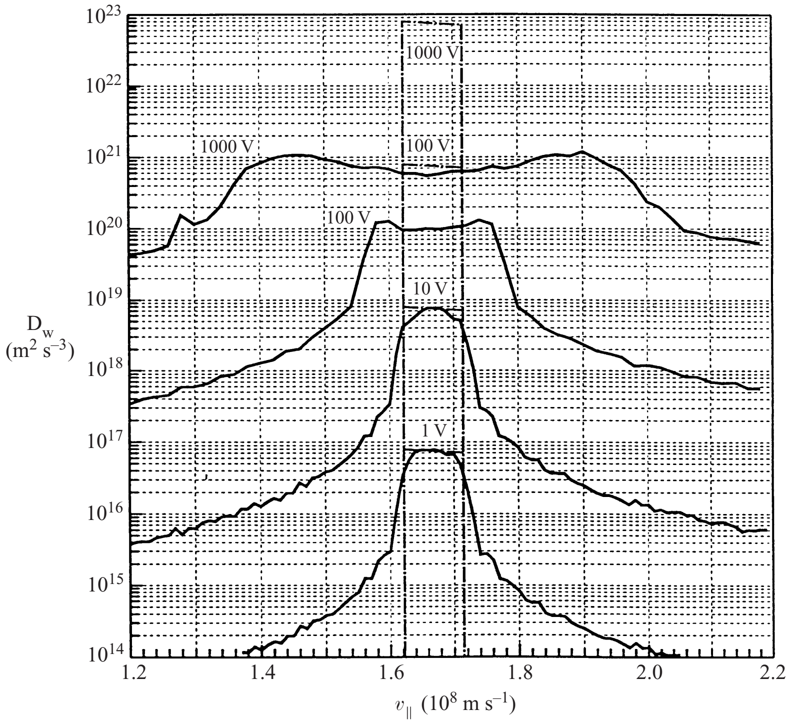

Figure 6. The full lines show the results of direct numerical simulation of the diffusion coefficient for various values of the potential Ψ_0: 1 V, 10 V, 100 V and 1000 V(Ψ_0 = 1000 V corresponds to an energy flux 0.5 GW m^{-2}). $\Delta N_\parallel/N_{\parallel,0} = 0.06$, 500 random-phase samples, and the number of modes $M = 10$. The dot-dashed lines show the values from the quasilinear approximation.

In view of the complicated trajectories of the particles in a tokamak geometry, we can assume that any correlation between subsequent transits through the same RF field segment will be lost. In fact, we have verified this assumption on a model case with just one segment, a single wave and a circular toroidal orbit (with no rotational transform). This model is analytically tractable. Thus the role of numerical errors has been excluded. If the field is homogeneous along the trajectory, the motion is completely regular. However, if there exists a field-free region of just a few wavelengths and if, as an approximation, the spatial envelope of the RF field has a rectangular form, a rapid loss of correlation occurs for the potential $\Psi_0 \geqslant 1$ V. More details will be presented elsewhere. Therefore, to obtain a statistically correct picture, it is sufficient to follow a large enough number of particles for one transit through the RF segment and randomly chosen phases of the waves. This should be done for any magnetic field line passing through the segment, a representative number of perpendicular velocities, and any parallel velocity.

To obtain an estimate of the effect of the strong LH wave field on the particle velocity distribution, we have discretized the velocity space, and for each $v_\parallel$, we have calculated a collection of trajectories for randomly generated phases of waves. The results presented here are for ITER-like parameters (B_0 = 5.7 T, R_0 = 8.1 m, a = 2.8 m), for a magnetic field line with $q = 2$ at r_0 = 2.4 m (following the chosen form of $q(r)$) and for $v_\perp = 0$. Figure 6 gives the diffusion coefficient D_w in $v_\parallel$ space

https://doi.org/10.1017/S0022377899007850 Published online by Cambridge University Press

(full lines). The parameters of Cohen et al. (1990) are assumed here (cf. the text following (34)), but with four values of wave potential amplitudes as indicated. The quasilinear spectral width was chosen as $\Delta N_{\parallel}/N_{\parallel,0} = 0.06$. The largest value of the amplitude corresponds to an energy flux of 0.5 GW m^{-2}. The values shown are weighted by the ratio of the RF segment of the magnetic surface to the whole magnetic surface. They can be compared directly with the quasilinear values (dot-dashed lines). For each $v_{\parallel}$, an ensemble of 500 random-phase samples was used. The number of modes $M = 10$. Increasing further the number of samples and modes has only a negligible influence on the overall results.

The most striking difference with the quasilinear approximation consists in a dramatic broadening of the diffusion coefficient, combined with a decrease in its magnitude. We remark that, in a cylindrical geometry, this effect has been studied in more detail by Pavlo et al. (1998). This broadening of the diffusion coefficient occurs for the potential Ψ_0 greater than about 100 V, in comparison with the quasilinear values where $D_{QL} \propto \Psi_0^2$. The broadening of the diffusion coefficient given by the curve $\Psi_0 = 1000$ V in Fig. 6 can be compared with the estimate of Cohen et al. (1990), namely their equation (2). For the same set of parameters, this estimate is about one-half of the velocity interval ($1.38 \leqslant v_{\parallel} \leqslant 1.94$) implied by Fig. 6. Obviously, non-resonant electrons become accelerated. From this point of view, there exists some similarity with the paper of Fuchs et al. (1996). Both the broadening of the diffusion coefficient and the decrease in its magnitude will result in a stronger damping of LH waves during their penetration into the plasma core, and might therefore represent a serious obstacle for the proposal of Cohen et al.

A limitation of our model is that it is not self-consistent. Moreover, the effects of collisions must be included in order that the model be a complete analogy to the original quasilinear theory and to the original proposal of Cohen et al. (1990).

6. Conclusions

A thorough analysis of the interesting proposal of Cohen et al. (1990) brings out several new phenomena that accompany the interaction of powerful wave fluxes with plasma. Among them, the following have been discussed and evaluated in this paper.

We have explored ponderomotive force effects at antennas. We have found that all our results concerning the nonlinear reflection coefficient of the LH wave, the plasma bias and plasma rotation induced by the LH wave are critically dependent on the value of the boundary plasma temperature in front of the grill. For plasma temperatures of about 10 eV in front of the grill, the reflection coefficient of the LH wave would be unacceptably high. On the other hand, for a boundary plasma temperature of about 30 eV or higher, the value of the nonlinear reflection coefficient will approach the values according to the linear theory. Similarly, the plasma bias and the corresponding plasma rotation decrease with growing plasma temperature. The possibility of growth of the plasma temperature in front of the grill is supported by experiments (see e.g. Petržílka et al. 1991).

We have explored the induced backward current and its effect on current drive efficiency. Although a large portion of the RF power is lost via backward current Joule heating, the current drive efficiency is still acceptable. The reason is that, with powerful RF pulses, the energy pumped into the poloidal magnetic field increases.

This energy is spent on the extremely efficient Ohmic current drive during the time between the RF pulses. If the high-density backward current leads to anomalous resistivity, the current drive efficiency increases considerably. For a small tokamak with high anomalous resistance, $\tau_{\rm sk}$ may be less than τ_p. If the well-known $\mathscr{L}/\mathscr{R}$ 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.

We have analysed the applicability of quasilinear theory. The diffusion coefficient appears to significantly differ from that predicted by the quasilinear theory.

Some of the effects mentioned here may represent serious obstacles for the proposal of Cohen et al. (1990). Nevertheless, we consider our study preliminary rather than a definitive answer. Significantly more work is necessary. Moreover, some phenomena that have been inspired by the work of Cohen et al. appear to be very interesting in themselves.

Note added in proof

The problem of the quasilinear approximation in the regime of strong fields has also been discussed by Fuchs et al. (1985).

Acknowledgements

The authors thank the referee for valuable suggestions. This work is supported by the Grant Agency of the Czech Republic, Grants 202/96/1350, 1355 and 97/0778, by the US–Czech Science and Technology Program, Grant 93068, and by the US Department of Energy, Grant DE-FG02-88ER53264.

References

Andrews, P. L., Chan, V. S. and Liu, C. S. 1985 *Phys. Fluids* **28**, 1148.

Bertrand, P., Ghizzo, A., Karttunen, S. J., Pättikangas, T. J. H., Salomaa, R. R. E., Shoucri, M. and Shkarofski, I. 1994 In: *Proceedings of 21st Conference on Controlled Fusion Plasma Physics*, Part III, p. 1082.

Cohen, R. H., Rognlien, T. D., Bonoli, P. T. and Porkolab, M. 1990 In: *Proceedings of 13th Conference on Plasma Physics and Controlled Fusion*, Vol. 1, p. 805.

Devoto, R. S. and Fenstermacher, M. E. 1990 *Nucl. Fusion* **30**, 362.

Fisch, N. J. 1987 *Rev. Mod. Phys.* **59**, 175.

Fuchs, V., Krapchev, V., Ram, A. and Bers, A. 1985 *Physica* **14D**, 141.

Fuchs, V. M., Goniche, M., Demers, Y. P., Jacquet, P. and Mailloux, J. 1996 *Phys. Plasmas* **3**, 4023.

Karney, C. N. F. and Fisch, N. J. 1987 *Phys. Fluids* **22**, 1817.

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

Klíma, R. and Petržílka, V. A. 1980 *Czech. J. Phys.* **B30**, 1002.

Kolesnichenko, Ya. I., Parail, V. V. and Pereverzev, G. V. 1989 In: *Problems of Plasma Theory*, Vol. 17 (ed. B. B. Kadomtsev), p. 3. Energoatomizdat, Moscow (in Russian).

Krlín, L., Pavlo, P. and Malijevský, I. 1997 *Phys. Lett.* **226A**, 80.

Krlín, L., Klíma, R., Pavlo, P., Petržílka, V., Svoboda, V. and Tataronis, J. A. 1998 Institute of Plasma Physics, Academy of Sciences of the Czech Republic, Report IPPCZ 360.

Pavlo, P., Krlín, L. and Tlučhoř, Z. 1991 *Nucl. Fusion* **31**, 711.

Pavlo, P., Krlín, L. and Petržílka, V. 1998 In: *Proceedings of 2nd Europhysics Topical Conference on Radiofrequency Heating and Current Drive of Fusion Devices, Brussels*. EPS Conference Abstracts, Vol. 22A, p. 201.

Petržílka, V. 1988 *Czech. J. Phys.* **38**, 937.

Petržílka, V., Leuterer, F., Söldner, G. F. X., Giannone, L. and Schubert, R. 1991 *Nucl. Fusion* **31**, 1758.

Petržílka, V., Tataronis, J. A., Söldner, F. X. and Krlín, L. 1997 *Phys. Plasmas* **4**, 3222.

Tataronis, J. A., Petržílka, V., Fuchs, V. and Krlín, L. 1997 In: *Proceedings of the International Sherwood '97 Fusion Theory Conference, Madison, WI, April 28–30*, Paper 1C28.

Trubnikov, B. A. 1965 In: *Reviews of Plasma Physics*, Vol. 1 (ed. M. A. Leontovich), p. 105. Consultants Bureau, New York.

Van Nieuwenhove, R., Petržílka, V. and Tataronis, J. A. 1995 *J. Plasma Phys.* **54**, 245.

https://doi.org/10.1017/S0022377899007850 Published online by Cambridge University Press