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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 the LH power into the interior of large scale plasmas (Cohen *et al.* 1990). We show that, at such very intense LH pulses, the wave coupling may deteriorate because of the 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 plasma by intense LH pulses, dissipate a large portion of the RF power absorbed, the current drive efficiency is acceptable. LH wave scattering on the boundary plasma fluctuations leads to enhanced absorption and wave reflection before the wave reaches the plasma center. We use a numerical simulation of wave - particle interactions to analyze the applicability of standard quasilinear theory to the case of large energy flux densities. The initial results show important restrictions on the use of the quasilinear approximation. The results of the present paper also indicate that some of the effects considerably alter the ideas of Cohen's *et al* (1990).

1. Introduction

In a thermonuclear tokamak plasma, lower hybrid waves (LHW) are strongly absorbed at the plasma boundary (e.g. Devoto *et al.* 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 LHW. 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, has been considered by Cohen *et al.* (1990). Aspects of this approach were also discussed by Bertrand *et al.* (1994). This proposed method to improve the penetration of LH waves opens a number of issues that should be studied thoroughly. Moreover, there are in this problem a variety of novel and important physics issues that have, in our opinion, a broader significance. The purpose of the present paper is to extend and improve the analysis of Cohen *et al.* (1990) relative 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). We find, however, results that can considerably modify the intriguing proposal these researchers reported.

The plan of this paper is as follows. In Sect. 2, we discuss the influence of nonlinear effects on the wave coupling. We find that efficient coupling can be achieved either for sufficiently high temperature of the boundary plasma (30 eV), or at lower temperatures with ultra short pulses ($\tau_p \approx 30ns$). In Sect. 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. Due to Faraday's law, the pulsed regime of current drive is inevitably accompanied by the generation of a backward current. This effect is analyzed in Sect. 4 with particular emphasis placed on the resulting current drive efficiency. Section 5 treats the issue of the negative influence of wave scattering on plasma fluctuations on the transport of powerful wave fluxes into the plasma core. Section 6 addresses the problem of applicability of the quasilinear approximation for the case of intense wave fluxes. The differences that we found between a direct numerical simulation of the wave particle interactions and the quasilinear approximation calls for a more complete investigation of this problem. In Sect. 7, 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 Eqs. (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 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

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

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 $f_p \geq f$, corresponding to $n_b \geq 8 \times 10^{17}$ m^{-3}, Eq. (6) yields $W_{max} \leq 10eV$. This moderate value of $W_{\max}$, which we have also confirmed with numerical computations, is a consequence of the relatively high value of the frequency that we 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 ($\tau_p \approx 80\mu s$) 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 Cohen's et al. "regime A" 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, L_i and L_W are, respectively, characteristic lengths of the plasma displacement and inhomogeneity of $W(z)$. It is natural to set

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

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 Cohen's *et al.* (1990) "regime A" 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 the 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,

$$T_b = T_0(1 + S/S_T). \tag{10}$$

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)),

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

where $\lambda = 2\pi/k_z$ and$k_0 = \omega/c$. We have solved Eq. (12) numerically for $E_1^{(+)}$ and $E_1^{(-)}$ with boundary conditions deep enough inside the plasma where ponderomotive forces are negligible. The RF electromagnetic fields computed 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 this expression, we find, for example, that the boundary temperature with $\bar{S} = S_T$ is twice the temperature with zero LH power, $\bar{S} = 0$. Here, $\bar{S}$ designates the energy flux in front of the grill $S(z)$ averaged over the toroidal coordinate z. If $\bar{S} >> S_T$, the resulting boundary temperature is much higher than the temperature with no LH power. On the other hand, for $\bar{S} << S_T$, the boundary temperature practically does not change as 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 be also 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 effects of S_T on δ. In Fig. 1, we see that the 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 Eq. (1), the plasma density decrease in front of the grill is then stronger, which leads to the deterioration of the coupling and to the 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.* 1996) 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, to 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 (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 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 the subscripts e and i label, respectively, electron and ion components. The induced time-averaged radial electric field can be expressed as a sum of two terms (Petržílka *et al.* 1997),

$$E_{0r} = \frac{1}{en_0}\frac{\partial p_i}{\partial r} + \mathcal{E}^{P}, \qquad (14)$$

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

$$\mathcal{E}^{P} = -\frac{B_{0z}}{m_i r^2 n_0 U_{ir}}\int_0^r d\bar{r}\,\bar{r}^2 F_\theta^P + \frac{B_{0\theta}}{m_i r n_0 U_{ir}}\int_0^r d\bar{r}\,\bar{r} F_z^P. \qquad (15)$$

In Eq. (15), n_0 is the time-averaged plasma density, and U_{ir} is the radial component of the mean ion fluid velocity. According to Eq. (12), the value that $\mathcal{E}^P$ has at a radial position r depends on the values of two definite integrals carried out from the plasma center, $\bar{r} = 0$, 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 contain results of a numerical evaluation of Eq. (14). The figures show that, for intense LH wave pulses, the poloidal ponderomotive forces may induce large stationary radial electric fields, up to about 10 kV/cm. This value is approximately two orders of magnitude higher than the electric field required for induction of enhanced confinement H modes by plasma biasing.

4. Backward current

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 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

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

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

$$\tau_{sk} \approx \mu_0 a^2/\eta, \qquad (17)$$

a being the plasma minor radius. For present large tokamaks, $\tau_{sk} \geq 10$ s, while for a reactor plasma $\tau_{sk} \approx 10^3$ 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

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,

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

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 due 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_{pol} pumped during τ_p to the poloidal magnetic field is, according to Eqs. (18) and (19),

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

The energy density W_{pol} is equal to the energy density dissipated during the time without RF. If the integrand in Eq. (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 **B**. Consequently, for $\tau_r >> \tau_p$,

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

Equation (22) implies that enhanced resistivity η_p improves the current drive efficiency, cf. Fisch (1987). The new point here is that, for 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 one given by the present theory.

The energy density W_e dissipated by the backward current is calculated similarly to the calculation Eqs. (20) and (21). Thus, with Eq. (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 will 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, $v_a \simeq$ const. The absorbed RF power density P_a can then be expressed as

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

where n_a is the density of the absorbing 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 (Cohen *et al.* 1990). The power density $(-eEn_a v_a)$ obviously equals $(W_e + W_{pol})/\tau_p$. Introducing $\varphi_{EC} = -eE/F_{coll}$, and using Eqs. (20), (23) and (24), we find,

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

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 Eq. (19) can be omitted. This simplification yields

$$\varphi_{EC} = e(j_d - j_0)\eta_p / F_{coll}. \qquad (26)$$

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

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 Eq. (22) in (26) and express the longitudinal resistivity η_0 in terms of the collision time $\tau_s^{e/i}$ used in kinetic theory (e.g. Trubnikov 1965). In the steady state, the plasma $2-D$ model (Karney *et al.* 1979) reveals that $\tau_e(v_a)$ is about a factor 2.5 larger than *in* the $1-D$ model *(*$\tau_e \sim 1/(2+Z_i)$, *e.g.* Klíma *et al.* 1979). It is not clear to what degree the $2-D$ effects, depending on effective ion charge Z_i, will develop under the conditions considered. We now introduce a factor $a_\varphi(Z_i) \simeq 2$ with $1 \leq Z_i \leq 1.5$. This yields

$$\varphi_{EC} = b_\varphi j_0; \quad b_\varphi \simeq \frac{0.4 Z_i}{a_\varphi(Z_i)} \frac{\tau_r}{\tau_p} \frac{1}{en_e v_{Te}} (\frac{v_a}{v_{Te}})^2, \qquad (27)$$

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

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

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

In general, Eqs. (25) and (27) imply the following relation between j_0 and the averaged RF power density absorbed, $< P_a >_\tau = P_a \tau_p / \tau_r$,

$$j_0 \simeq -\frac{1}{2b_\varphi} + [\frac{1}{4b_\varphi^2} + \frac{< P_a >_\tau}{\eta_0}(1 + \frac{\eta_0 \tau_r}{\eta_p \tau_p})^{-1}]^{1/2}. \tag{29}$$

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

$$\eta_{CD} = n_e(10^{20}m^{-3})RJ_0/(\mathcal{V}_0 < P_a >_\tau) \simeq 0.6. \tag{30}$$

Both the inhomogeneity of the RF power absorption found in Cohen *et al.* (1990) and our results in Sec. 6 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 $< P_a >_\tau = 6 \times 10^5\mathrm{Wm}^{-3}$. Retaining all the above parameters, we find $j_0 \simeq 2.8 \times 10^6\mathrm{Am}^{-2}$ 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.(29), we have

$$j_0 \simeq 1.2 \times 10^7 Am^{-2}, \quad \varphi_{EC} \simeq 2. \tag{31}$$

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. (25) implies

$$j_0 \simeq [\frac{< P_a >_\tau}{\eta_0}(1 + \frac{\eta_0 \tau_r}{\eta_p \tau_p})^{-1}]^{1/2}. \tag{32}$$

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 $n_e(10^{20}) j_0/(2\pi < P_a >_\tau)$ 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 Eq. (18) in Klíma *et al.* (1979), the distribution function of the resonant electrons is stable with respect to the Parail - Pogutse instability for the specific parameters considered here.

5. LH wave scattering at the boundary

Scattering of lower hybrid waves on low frequency density fluctuations at the plasma boundary (Andrews *et al.* 1988) can also influence the pulse scheme. We explore this effect here.

As the density fluctuations have a very small wave vector component $k_{\|}$ along the magnetostatic field **B**, they do not change the lower hybrid $P(N_{\|})$ spectra directly, but only rotate the initially radial perpendicular component of the LH wave vector $\mathbf{k}_{\perp}$ during the scattering process. This leads to the appearance of a nonzero poloidal component k_{θ} of the wave vector k. The resulting poloidal component k_{θ} can significantly contribute to the $k_{\|}$ during the radial propagation of the LH wave in the sheared tokamak magnetostatic field after the wave scattering on the density fluctuations at the plasma boundary.

Our analysis of this scattering effect proceeds analogously to that presented by Andrews *et al.* (1988). However, the parameters are now related to the launching of LH pulses into a large tokamak or a reactor relevant plasma. We assume here that a narrow turbulent scattering layer extends near the plasma boundary from $r = r_1$ to $r = r_0$, $r_0 > r_1$. Because of scattering, the wave spectrum radiated by the grill $P_p(N_{\|}, \sin\psi, r_0) = P_b(N_{\|})\delta(\sin\psi)$ is changed to

$$P_p(N_{\|}, \sin\psi, r_1) = P_b(N_{\|})F(N_{\|}, \sin\psi), \tag{33}$$

In this expression, ψ is defined by the equation $N_{\perp}\cos\psi = N_r$, where $N_{\perp} = k_{\perp}c/\omega$, $N_r = k_r c/\omega$, k_r is the radial component of k, $P_b(N_{\|})$ is the power spectrum radiated by the grill and $F(N_{\|}, \sin\psi)$ is determined by the characteristics of the scattering layer. The function $F(N_{\|}, \sin\psi)$ is given by the equation (Andrews *et al.* 1988)

$$\cos\psi\frac{\partial f}{\partial x} = \frac{1}{l_s}\frac{\partial^2 F}{\partial \psi^2}, \tag{34}$$

where l_s is defined by

$$l_s^{-1} = \xi_0\frac{\sqrt{\pi}}{8}\left[\left(\frac{\delta n_e}{n_e}\right)^2\right]\left(1 + \frac{3\omega_{pe}^2\xi_0^2c^2}{2\omega^2\Omega_e^2N_{\|}^2}\right). \tag{35}$$

Here, ξ_0 is the typical wave vector magnitude of the drift wave turbulence, n_e is the plasma density and Ω_e is the electron cyclotron frequency. Outside of the scattering layer and in the interior of the tokamak plasma, the spectrum can be determined by the geometrical transformation described in Andrews *et al.* (1988) or Petržílka (1988). For further details, we refer to Andrews *et al.* (1988).

Figures 6 - 8 show numerical results of our analysis of the scattering effects. Figure 6 shows the radial dependence of the LH spectral power density $P(N_{\|})$, which is obtained by integration of the function $P_p(N_{\|}, \sin\psi, r)$, defined by Eq. (33), over ψ at various radii. It can be seen that in the presence of the wave scattering on plasma density fluctuations at the plasma boundary, the spectrum of LH waves broadens. This spectral broadening increases wave absorption. In addition, spectral broadening can lead to reflection of LH waves before they penetrate deeper into the plasma, as shown by Figs. 7 and 8. Figure 7 shows the radial dependency of the wave power $P(r)$, obtained by integration of the function $P_p(N_{\|}, \sin\psi, r_1)$, defined by Eq. (33), over $N_{\|}$ and ψ. Figure 8 shows the dependency of the wave power $P_1 = P(r_2)$ on the plasma boundary temperature, where r_2 is chosen as the radius inside the plasma just after the scattering layer. The relative amount of the wave energy scattered by density fluctuations just at the plasma boundary is acceptably low, as Fig. 8 shows. This is in accordance with the results of Cohen *et al.* (1990). However, the poloidal wave vector is changed by the scattering process, which significantly influences the wave propagation characteristics in the subsequent wave propagation in the sheared magnetostatic field, cf. Figs. 6 and 7.

6. Applicability of the quasilinear approximation for the case of intensive wave pulses

The quasilinear approach is considered as an excellent tool for the description of LHW-plasma interaction. Nevertheless, the quasilinear approximation itself has been developed on the basis of the perturbation analysis, i.e., on assumption that the change of the particle velocities during the wave-particle interaction is small.

In case of large wave power fluxes, a possible change of the energy of particles during their single transit through the LHW 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 QLD by means of direct numerical simulation (DNS) 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 LHW 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 and in the fields of a lower hybrid wave, under the electrostatic approximation, is

$$H = \omega_{c0} P_1 \left[1 - \frac{r(P_2)}{R_0} \cos\tilde{\beta}\right] + \frac{P_3^2}{2m}\left[1 - 2\frac{r(P_2)}{R_0}\cos\tilde{\beta}\right] +$$

$$+e\Psi_0 \cos\left[k_r\sqrt{\frac{2P_2}{eB_0}} + m_p\left(Q_2 + \frac{Q_3}{qR_0}\right) + \frac{k_\parallel Q_3}{\left(1 - 2\frac{r}{R_0}\cos\tilde{\beta}\right)^{1/2}} - \omega t\right].\quad(36)$$

The definitions of the canonical coordinates and of the other symbols in Eq. (21) are the following:

$$P_1 = \frac{1}{2}eB_0\rho_c^2,\quad P_2 = \frac{1}{2}eB_0r^2,\quad P_3^2(1 - 2\frac{r}{R_0}\cos\tilde{\beta}) = m_e^2 v_\parallel^2,$$

$$Q_1 = \omega_{c0}t,\quad \tilde{\beta} = \theta = Q_2 + \frac{Q_3}{qR_0},\quad Q_3 = R_0\phi,$$

$$r = \sqrt{\frac{2P_2}{eB_0}}.$$

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

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

$$|E_\parallel| = \sqrt{\sum_{i=1}^{M} E_i^2} = \left(2\mu_0 c N_\parallel \frac{\omega}{\omega_{pe}} S\right)^{1/2}. \quad (37)$$

For example, in case of a single wave, with the parameters that have been assumed by Cohen *et al.* (1990), i.e., with the frequency $f = 8\text{GHz}$, $N_\parallel = 1.8$ and the energy flux density 0.5GW/m^2, and with the density $n_e = 10^{20}\text{m}^{-3}$, Eq. (37) yields $E_\parallel \approx 3\times 10^5 Vm^{-1}$, which corresponds to $\Psi_0 = 10^3$ V. For M waves, $E_i \approx 3\times10^5/\sqrt{M}\text{Vm}^{-1}$.

The spatial distribution of LHW can be obtained by ray tracing. However, at the present stage, our primary interest is more the qualitative than an accurate quantitative analysis of the LHW-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.

In view of the complicated trajectories of the particles in 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 have been excluded. If the field is continuous 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 RF field has a rectangular form, a rapid loss of correlation occurs for the potential $\Psi_0 \geq 1\mathrm{V}$. 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 LHW field on the particle velocity distribution, we have discretized the velocity space and for each $v_{\|}$, we calculated a collection of trajectories for randomly generated phases of waves. The results presented here are for a magnetic field line with $q = 2$, and for $v_{\perp} = 0$. Figure 9 gives the diffusion coefficient in $v_{\|}$ space (full lines). The parameters Cohen *et al.* (1990) are assumed here [cf. the paragraph following Eq. (37)], but with four values of wave potential amplitudes as indicated. The largest value of the amplitude corresponds to the energy flux $0.5\mathrm{GW/m}^2$. The values shown are weighted by the ratio of the RF segment of the magnetic surface to the whole magnetic surface. They can be directly compared with the quasilinear values (dashed lines). For each $v_{\|}$, an ensemble of 500 random phase samples was used. The number of modes is $M = 10$. 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 of its magnitude. We remark that, in 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 $D_{QL} \propto \Psi_0^2$. Broadening of the diffusion coefficient given by the curve $\Psi_0 = 1000V$ in Fig. 9 can be compared with the estimate of Cohen*et al.* (1990), viz., their Eq. (2). For the same set of parameters, this estimate is about one half of the velocities interval ($1.38 \leq v_{\|} \leq 1.94$) implied by Fig. 9. Obviously, non-resonant electrons become accelerated. From this point of view, there exists some similarity with the recent paper of Fuchs *et al.* (1996). Both broadening of the diffusion coefficient and the decrease of its magnitude will result in a more intensive damping of LHW during their penetration into the plasma core and might therefore represent a serious obstacle for Cohen's proposal.

A limitation of our model is that it 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).

7. Conclusion

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 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 the 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 the 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 the 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 leads to extremely efficient Ohmic current drive during the time between the RF pulses. For a small tokamak with high anomalous resistance, 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.

We have examined the effect of LH scattering. The LH wave scattering on the boundary plasma fluctuations can significantly broaden $P(N_{\|})$ spectra, which leads to enhanced wave absorption and reflection of LH waves before they reach the plasma center. This effect may be namely strong at low plasma boundary temperatures and high plasma densities.

We have analyzed the applicability of quasilinear theory. The diffusion coefficient appears to significantly differ from that predicted by the quasilinear theory. This may represent a serious obstacle for the applicability of the proposal of Cohen's *et al.* (1990). Nevertheless, our analysis must be considered as only first approximation.

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 itself.

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REFERENCES

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

Bertrand P et al. *21st Conf. on Contr. Fus. Plasma Phys.*, 1994, Part III, 1082.

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

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

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

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 Šov. 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 in the right-hand side of inequality (18) or, approximately, ω_{He}/ω_{pe} must be instead of Ω there.

Kolesnichenko Ya I, Parail V V and Pereverzev G V in *Problems of Plasma Theory*, Vol. 17, Energoatomizdat, Moscow 1989 (in Russian).

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

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

Pavlo P, Krlín L and Petržílka V *Second Europhysics Topical Conf. on Radiofrequency Heating and Current Drive of Fusion Devices*, Brussels 1998 EPS Conf. Abstracts, Vol 22A, 201.

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

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

Petržílka V, Leuterer F, Söldner GF 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: Fast electrons generated by lower hybrid waves at the tokamak plasma edge, in *Proceedings of the International Sherwood'97 Fusion Theory Conference, Madison*, Wisconsin, USA, April 28 - 30, 1997, paper 1C28.

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

FIGURE CAPTIONS

Fig. 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, $N_{\|} = 2$, wave frequency $f = 8$ GHz, the initial boundary temperature $T_0 = 10$ eV, the initial boundary density normalized to the critical density is $n_0/n_c = 2$, $n_c = 7.93 \times 10^{17}$ m^{-3}. We note that for low reflection coefficients of about $R = 0.2$, this launched power of 50 kW/cm^2 corresponds to the wave electric field amplitude 5.5 kV/cm in front of the grill for $N_{\|} = 2$.

Fig. 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 parameters of Fig. 1.

Fig. 3. Radial electric field E_r, induced by poloidal ponderomotive forces in front of LH grill, in dependence on the boundary plasma temperature. The boundary plasma density is $n_b = 2 \times 10^{18}$ 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 assumed as 5.5 kV/cm, which corresponds to the coupled wave power of the order of tens of kilowatts per cm^2, depending on the wave reflection coefficient.

Fig. 4. The toroidal velocity U_z of the plasma rotation, induced by poloidal ponderomotive forces in front of LH grill, in dependence on the boundary plasma temperature, for the parameters of Fig. 3.

Fig. 5. The poloidal velocity U_θ of the plasma rotation, induced by poloidal ponderomotive forces in front of LH grill, in dependence on the boundary plasma temperature, for the parameters of Fig. 3.

Fig. 6. Broadening of LH spectral power density $P(N_{\|})$ [a.u.] at various minor radii r for the case of LH wave, which was partially scattered by plasma fluctuations at the plasma boundary. The minor radius $a = 100$ cm, plasma temperature at the boundary $T_b = 30$ eV, magnetostatic field $B = 5$ T, LH wave frequency $f = 5$ GHz, and central plasma density $n(0) = 10^{20}$ m^{-3}, $N_{\|}$ is the parallel index of refraction - denoted as N_{par} in the Figure.

Fig. 7. Radial dependency of the wave power $P(r)$ [a.u.] transmitted to various minor radii r, for the same parameters as in Fig. 6. The incident power $P_i = 8$ [a.u.].

Fig. 8. Dependency of the wave power P_1 [a.u.], transmitted through the narrow scattering fluctuation slab at the boundary, on the plasma boundary temperature. The incident power $P_i = 8$ [a.u.], the power P_1 is given just after passing the scattering layer. The parameters are the same as in Fig. 6.

Fig. 9. The diffusion coefficient for the potential Ψ_0 indicated (i.e. for $\Psi_0 = 1\ V, 10\ V, 100\ V, 1000\ V$. The last value corresponds to the energy flux $0.5 GW/m^2$.)

Fig. 1

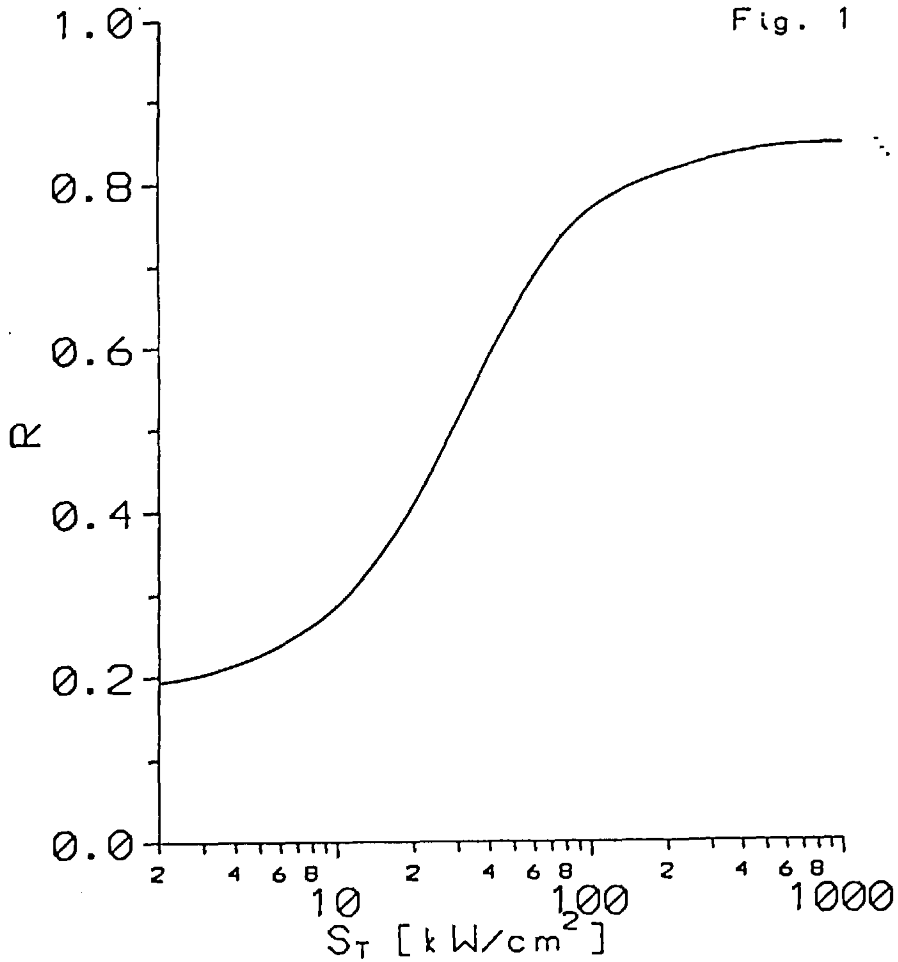


Fig. 2

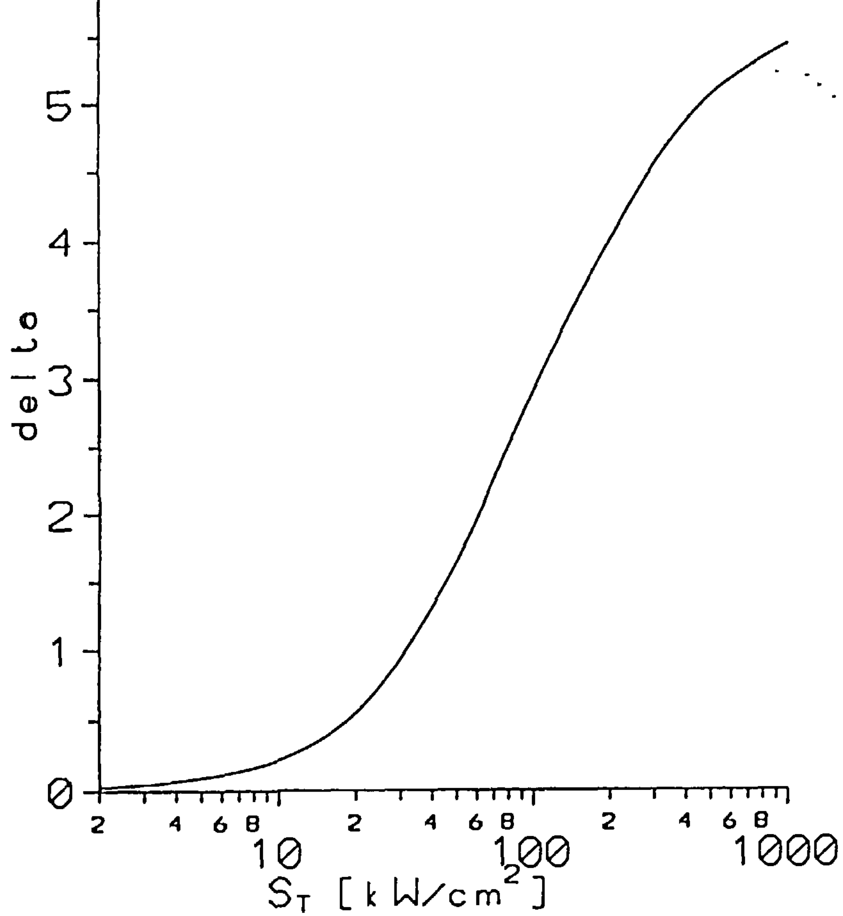
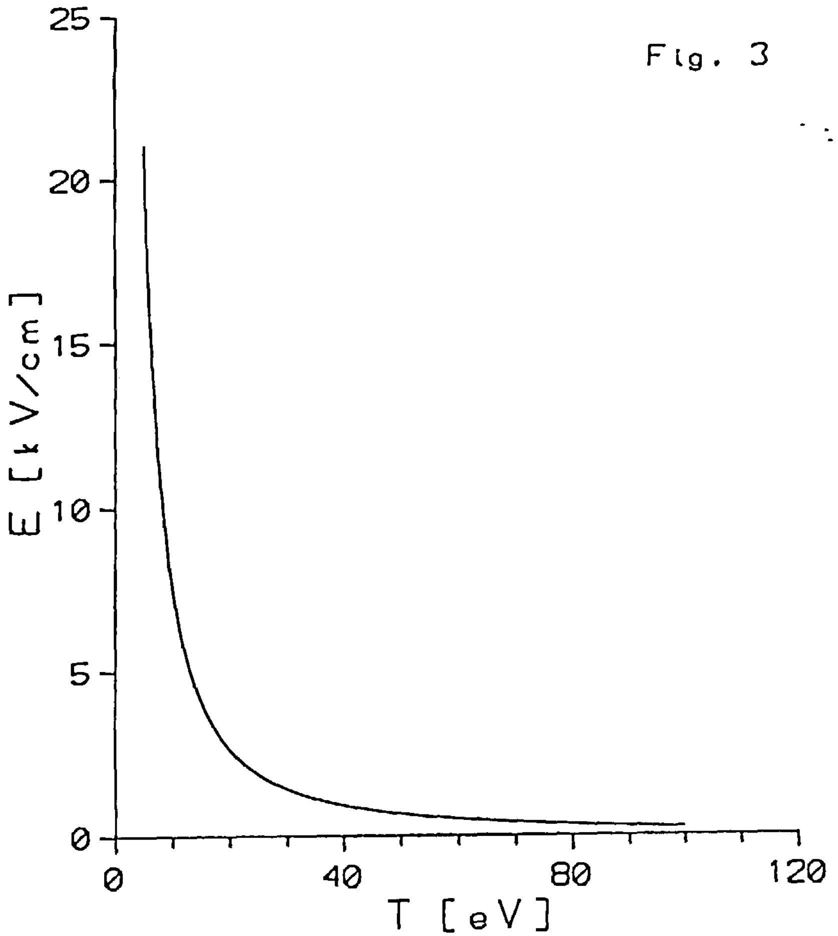


Fig. 3



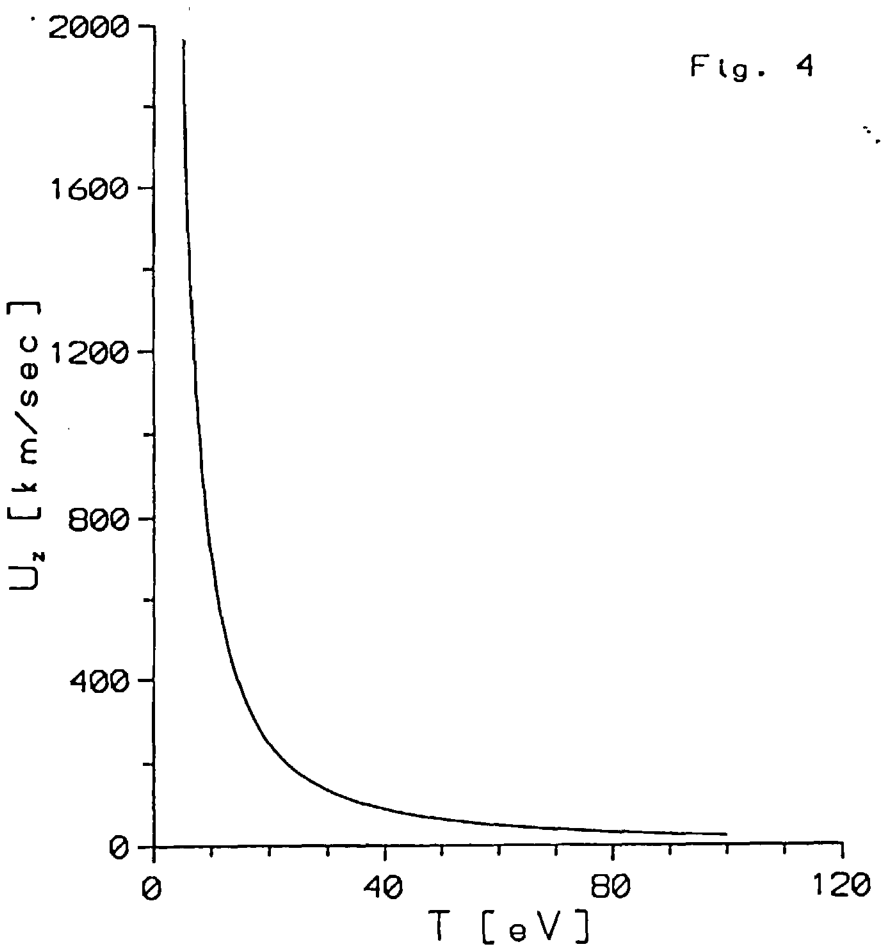


Fig. 4
2000
1600
1200
800
400
0
U_z [km/sec]
0
40
80
120
T [eV]

Fig. 5

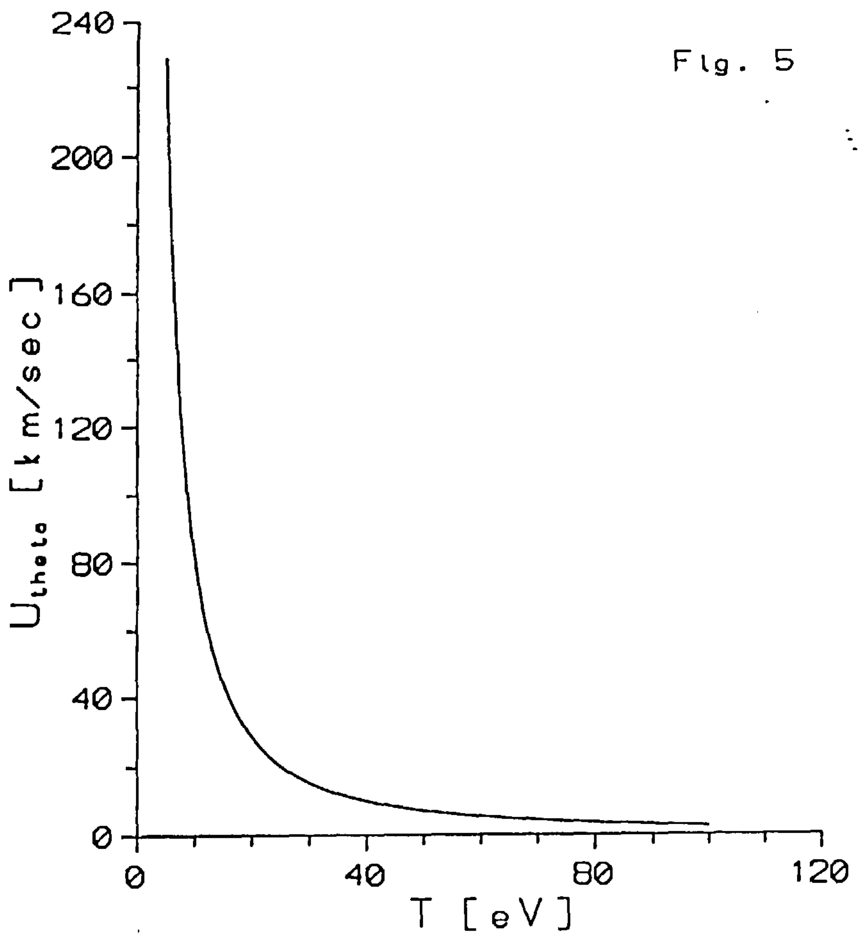


Fig. 6

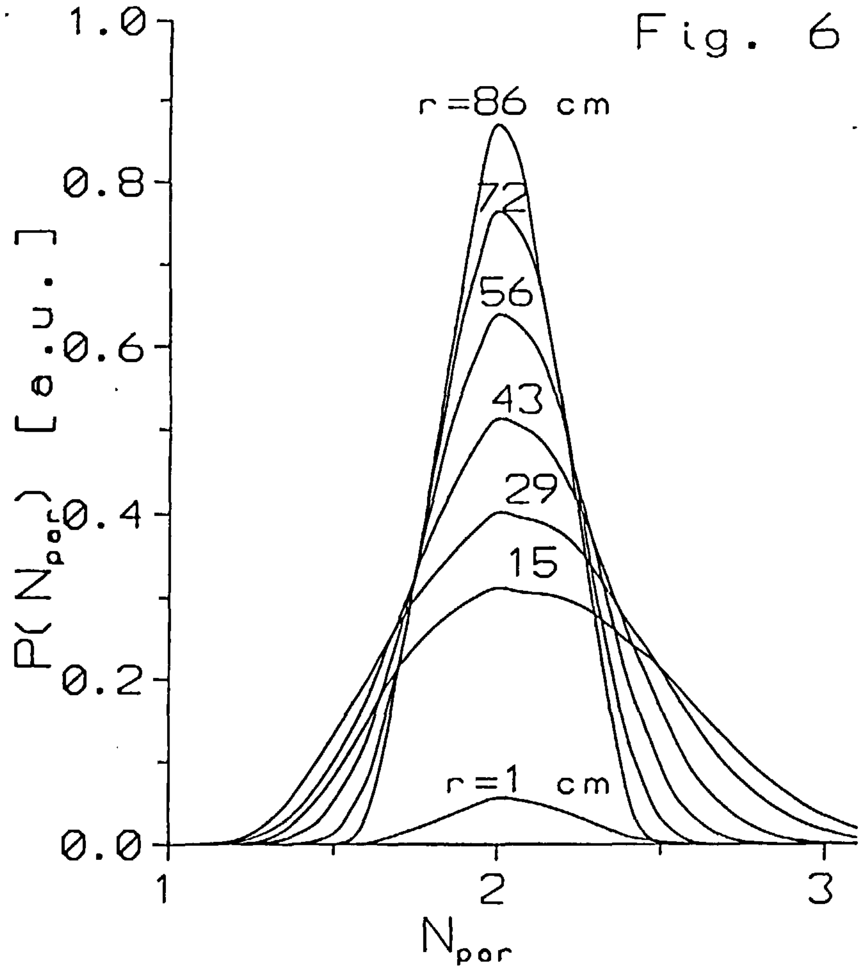


Fig. 7

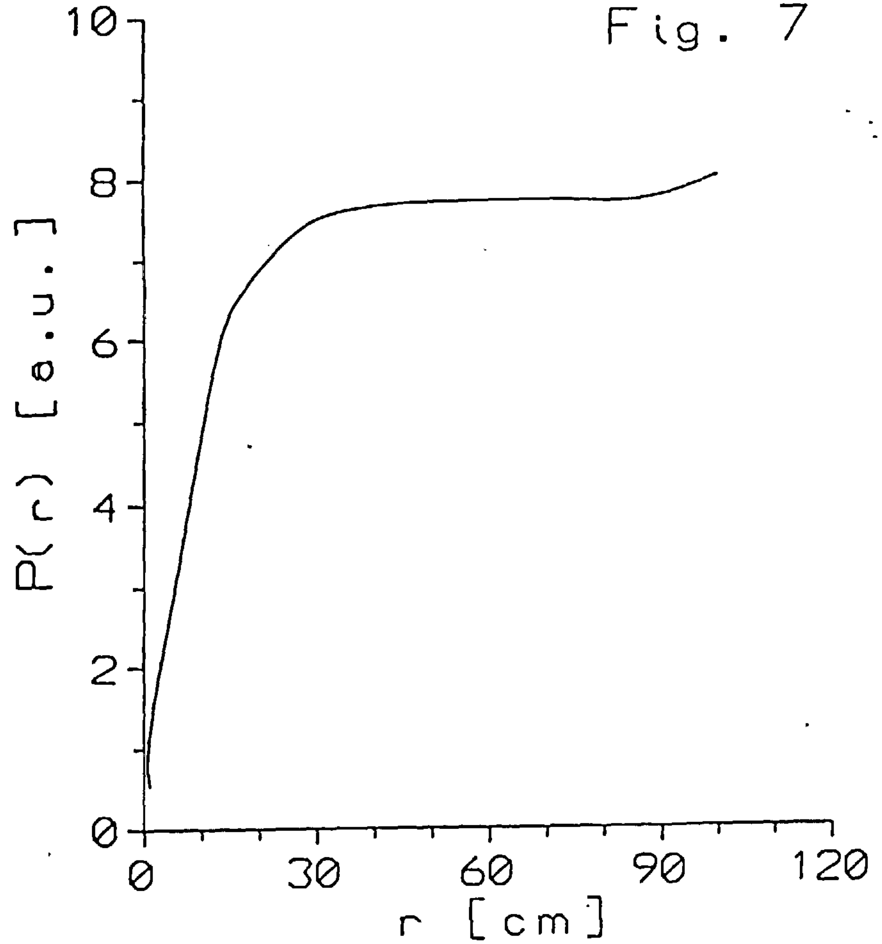
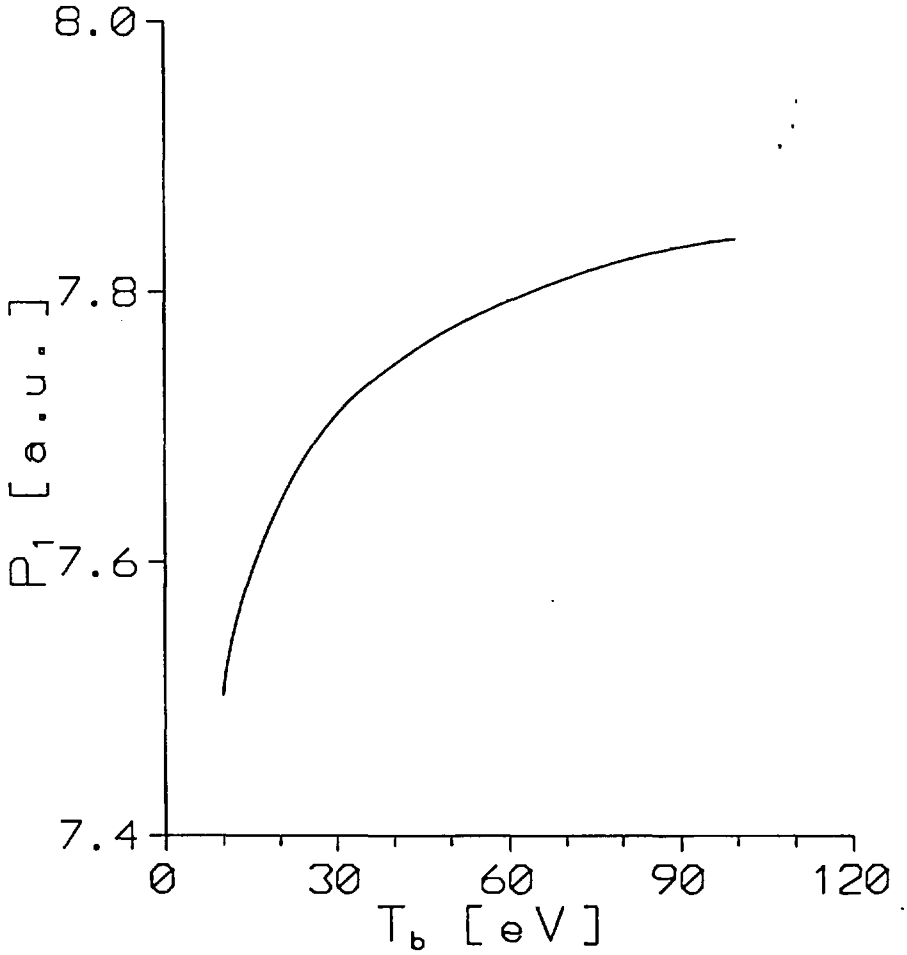


Fig. 8



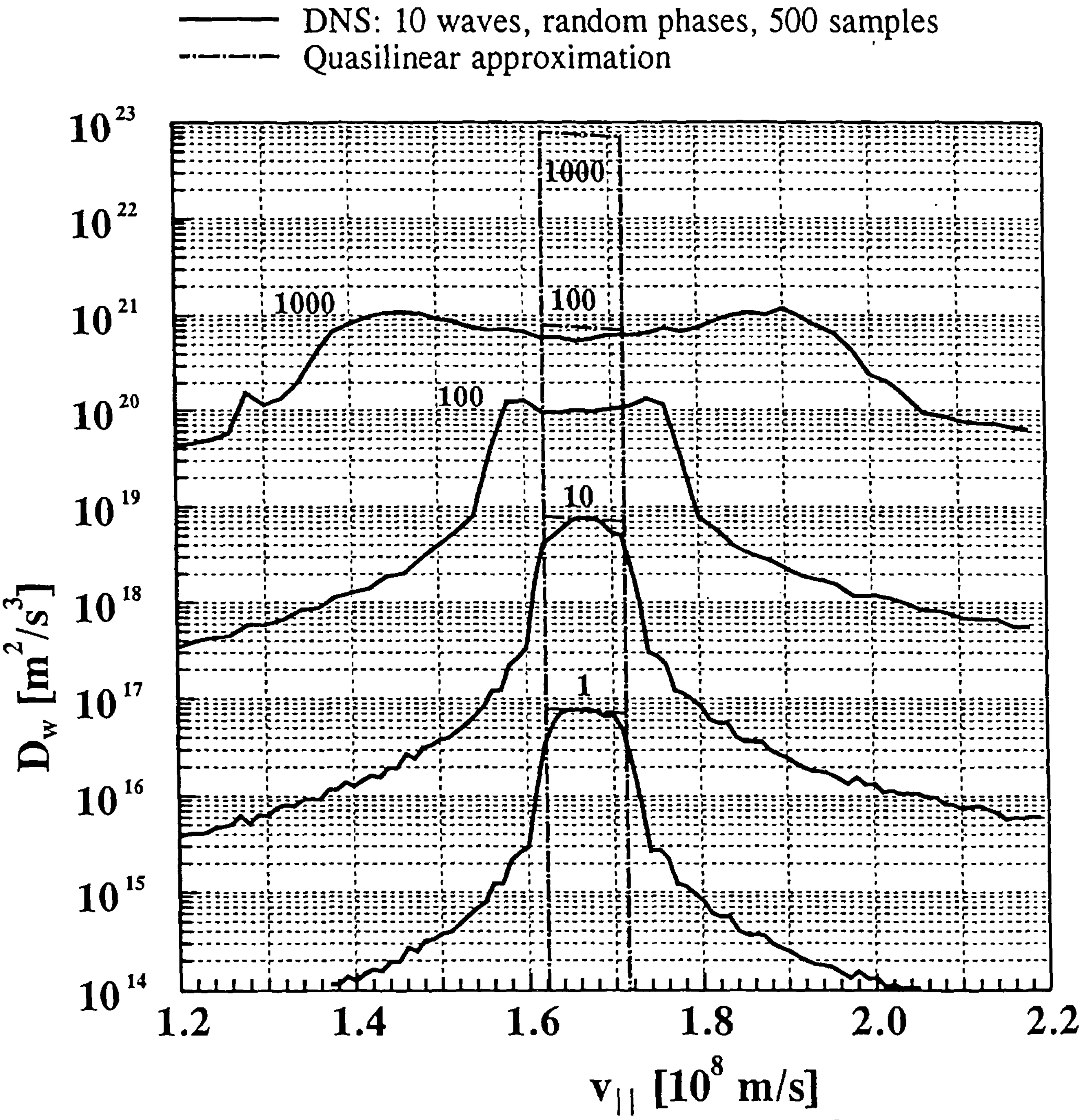

DNS: 10 waves, random phases, 500 samples
Quasilinear approximation
1000
100
10
1
10^{23}
10^{22}
10^{21}
10^{20}
10^{19}
10^{18}
10^{17}
10^{16}
10^{15}
10^{14}
D_w [m^2/s^3]
1.2
1.4
1.6
1.8
2.0
2.2
$v_{||}$ [10^8 m/s]

Fig. 9