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PULSED LOWER-HYBRID WAVE PENETRATION IN REACTOR PLASMAS

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ABSTRACT

Providing lower-hybrid power in short, intense (GW) pulses allows enhanced wave penetration in reactor-grade plasmas. We examine nonlinear absorption, ray propagation, and parametric instability of the intense pulses. We find that simultaneously achieving good penetration while avoiding parametric instabilities is possible, but imposes restrictions on the peak power density, pulse duration, and/or r.f. spot shape. In particular, power launched in narrow strips, elongated along the field direction, is desired.

There is general agreement in the magnetic-fusion community as to the desirability, if not the feasibility, of a steady-state reactor. In a tokamak, achieving this goal requires the sustenace of a non-inductively driven toroidal current. Lower-hybrid (LH) waves are a very appealing choice as a current driver, as LH current drive has an impressive data base, and good theoretical and experimental current-drive efficiency. However, according to conventional wisdom, lower-hybrid current drive doesn't work in the middle of reactor-grade plasmas because accessibility restricts the phase velocity to be low enough that electron Landau damping is severe. But the power expended in Landau damping saturates (due to formation of a quasilinear or nonlinear plateau), allowing burnthrough with enough power. As the absorption formulas below show, the saturated power is in the GW range for a reactor-grade machine such as the International Thermonuclear Experimental Reactor (ITER); such power levels are prohibative if continuous but reasonable as peak levels in a pulsed scheme. Suitable sources of pulsed lower-hybrid power, such as induction-linac-driven relativistic klystrons, are presently under development.

We report here of investigations of several key issues for enhanced LH wave penetration. First is the basic issue of nonlinear absorption of the intense radiation. Because the absorption depends on the width of the parallel-wavenumber ($N_{\|}$) spectrum and the microwave spot shape in the interior of the plasma, ray tracing is important. A ray-tracing code also offers a convenient way of integrating the nonlinear opacity across realistic temperature and density profiles. Finally, because of the wave intensity, it is important to assess the role of parametric instabilities.

We consider first the absorption, which is calculated from the power expended in forming a nonlinear plateau. The result depends on the degree of relaxation of the plateau between successive (as seen by individual electrons) pulses. The limit of most interest (where there is the greatest advantage in pulsing for a fixed average power) is that where the plateau completely relaxes between successive pulses (but the relaxation time is long compared to the pulse duration). In this limit, Maxwellian electrons are incident on the microwave beam; the distribution of emerging electrons is altered only in the band of parallel velocities that can resonantly interact with the waves. Referring to Fig. 1, the power expended by the wave is that utilized in moving the initially Maxwellian electrons from (1) to (2) as they stream through the beam. Hence the rate of decay of the axially averaged (but temporally instantaneous) Poynting Flux $\langle S\rangle$ with path length s is:

$$\langle dS/ds\rangle \cong \left(nmv_t^3/4\pi^{1/2}L_{\|}\right)\exp(-\epsilon_1)H(\Delta\epsilon)$$

where $L_{\|}$ is the r.f. beam scale length along the magnetic field, ϵ_1 is the parallel energy normalized to the electron temperature T at the lower edge of plateau, $\Delta\epsilon$ is the full width of the plateau, $v_t \equiv (2T/m)^{1/2}$, and $H(x) \equiv \left(1-e^{-x}\right)\left(-1+x/2\right)+xe^{-x}$; note $H(x) \sim x^3/12$ for small x.

If the plateau cannot relax completely between pulses, the peak power required to form the plateau is not as big, since the incident f is partly flattened. The absorption is reduced

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approximately by the factor $1 + a\nu_{rep,eff}\tau_{c,eff}$, where $\tau_{c,eff}$ is the time for the plateau to relax, $a \sim 1$, and $\nu_{rep,eff}$ is the average repetition rate experienced by individaul electrons, $\nu_{rep,eff} \equiv \nu_{rep} A_b / A_\psi$ with ν_{rep} the actual pulse repetition rate, A_b the area spanned by electrons on a flux surface which pass through the r.f. beam during the pulse, and A_ψ the area of the flux surface. Hence, defining a damping decrement α by the relation $\langle dS/ds \rangle = -\alpha \langle S \rangle$, we obtain:

$$\alpha = \frac{n m v_t^3 H(\Delta\epsilon) \exp(-\epsilon_1)}{4\pi^{1/2} L_\parallel \langle S \rangle (1 + a\nu_{rep,eff}\tau_{c,eff})} \quad . \tag{1}$$

In the limit where the plateau relaxes negligibly between pulses, the absorption of *average* power is the same as for continuously applied power; it can be obtained, for example, by calculating the (time-averaged) current and dividing by the current-drive efficiency. That result is reproduced by Eq. (1) in the limit of large $\nu_{rep,eff}\tau_{c,eff}$, provided that one eliminates S in favor of its time-averaged value.

The absorption depends sensitively on the plateau energy ϵ_1 and on the plateau width $\Delta\epsilon$. An upper bound on ϵ_1 is set by the accessibility condition $N_\parallel \gtrsim N_a$, where $N_a = (1 + x^2 - \omega_{pi}^2/\omega^2)^{1/2} + x$ and $x = \omega_{pe}/\omega_{ce}$, the ratio of the electron plasma and cyclotron frequencies. Thus there is strong dependence on n/B^2, where n is density and B is the magnetic-field strength. The plateau width, which in the weak-r.f. limit depends only on the spread in $N_\parallel$, now has a lower bound set by single-pass acceleration in the intense r.f. electric field,

$$(\Delta u/\bar{u})_E \cong \left(S_\perp(\mathrm{GW/m^2})/N_\parallel\right)^{1/4} A^{-1/8} (G/\mathcal{E}_\parallel)^{1/2} (10\,\mathrm{GHz}/f)^{1/2} \tag{2}$$

where f is the wave frequency, A is the ion mass in AMU, and $G \equiv x^{1/2}/(1 - x^{-2})^{1/4}$.

Using (2) in (1), we find the penetration length $L_p \equiv S_\perp / S\alpha$ to vary as $S^{1/4}$ for fixed u_1, and to be of the order of a meter for infrequent, 1 GW/m^2 pulses with $L_\parallel \sim 5$ m and nominal ITER parameters ($B = 5.3$ T, $n = 0.7 \times 10^{20}$ m^{-3}, $T = 25$ keV) (compared to about 10 cm. for linear absorption).

The absorption coefficient (1) accounts only for resonant interaction of electrons with the electric field. Additionally, however, the acceleration of the resonant electrons leads to a charge separation which persists for of order of an electron toroidal transit time, and hence on the pulse timescale, a d.c. electric field which accelerates the electron distribution function in the opposite direction from the resonant acceleration. This effect enhances penetration, but has not yet been adequately quantified.

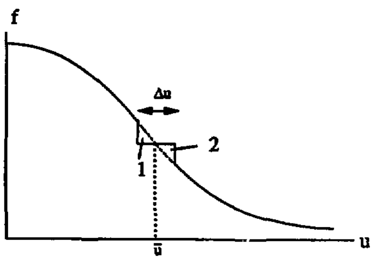

FIG. 1. Distribution function incident on (dashed) and emerging from (solid) r.f. beam vs. parallel velocity

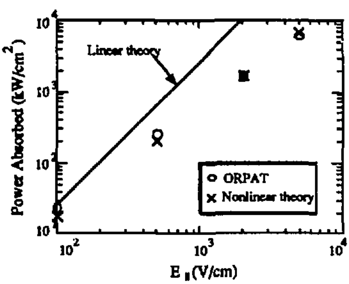

FIG. 2. r.f. power absorpbed for $N_\parallel = 1.7$, $f = 8$ GHz, $T_e = 25$ keV, $L_\parallel = 75$ cm, and $n = 8 \times 10^{19}$ m^{-3}

We have compared the absorption predicted by Eq. (1) with results from a multi-particle orbit code ORPAT. Here, Maxwellian electrons are followed for a single pass through a monochromatic, single $N_{\|}$ beam. Fig. 2 illustrates the good agreement with the nonlinear theory, as well as the reduction in absorption compared to the predictions of linear theory (weak r.f. result).

The nonlinear estimates (1) and (2) of the wave absorption and plateau broadening have been incorporated in a toroidal ray tracing code[1] to study the wave penetration problem using realistic plasma profiles in tokamak geometry. The ray tracing calculation[1] has been modified for noncircular equilibria and combined with an MHD equilibrium solver.[2] A rate equation for the wave energy flux is integrated along each ray path and has the form $d\langle S\rangle/dt = -2\gamma\langle S\rangle$, where $2\gamma = \alpha\times$ group velocity.

The density and pressure profiles are taken to be $n_e(\psi) = n_e(0)(1-\psi_n)$ and $p(\psi) = p(0)(1-\psi_n)^2$, where the normalized flux is $\psi_n \equiv (\psi-\psi_0)/(\psi_a-\psi_0)$, $\psi_0 \equiv \psi(0)$, $\psi_a \equiv \psi(a)$, and $T_e(\psi)$, $T_i(\psi)$ are found by taking $p(\psi) = n_e(\psi)[T_e(\psi)+T_i(\psi)]$. The other plasma parameters used were minor radius $a = 1.8$ m, major radius $R_0 = 5.5$ m, elongation $\kappa = 2.22$, triangularity $\delta = 0.61$, $B_0 = 5.3$ T, plasma current $I_p = 18$ MA, effective charge $Z_{eff} = 1.5$, $n_e(0) = 8\times10^{19}$ m^{-3}, and $T_e(0) = T_i(0) \simeq 33$ keV. The r.f. parameters were $f_0 = 8$ GHz, $N_{\|}^0 = 2.0$, $\Delta N_{\|}/N_{\|}^0 = (\Delta u/u)_0 = 0.05$, $\nu_{rep} = 5$ kHz, power $P_{LH} = (1-50)$ GW, and waveguide-array area $A_B = L_{\|}h = 1$ m^2; h is the height normal to field lines. A plot of $N_{\|}$ vs. ψ for five rays launched from a vertical position of $y_0 = +0.5$ m, for $N_{\|}^o = (1.9, 1.95, 2.0, 2.05, 2.10)$ is shown in Fig. 3. The results indicate that rays which start at the plasma edge ($\psi = 0$) with a spacing in $N_{\|}$ of 0.05 tend to keep this spacing on their first pass into the plasma, but spread substantialy following their closest approach to the magnetic axis. Similar results are found for $N_{\|}(\psi)$ plots made at other vertical positions, $y_0 = (+0.25, 0, -0.25, -0.5)$ m. This property makes it easier to ensure that waves will penetrate but not overshoot. Note that the nonlinear plateau width (2) is $\simeq 0.1-0.13$; there is thus no need to launch a $N_{\|}$ spectrum much narrower. Spatial broadening of the initial waveguide area (r.f. beam) has been found to be minimal. In fact, the r.f. beam area has been found to decrease slightly from 1.0 m^2 to 0.8 m^2 during the first pass of the r.f. beam into the plasma. The projection of the r.f. beam on each flux surface was found by tracing up to 100 "marker" rays emanating from the waveguide at the plasma surface.

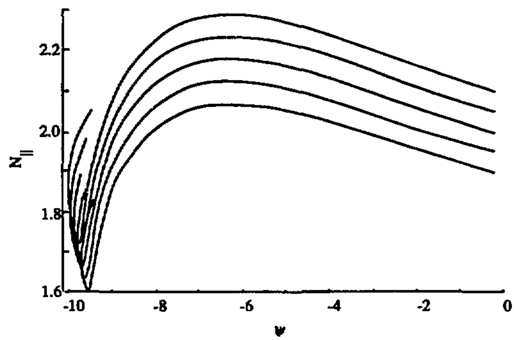

FIG. 3. Parallel wavenumber vs. flux surface

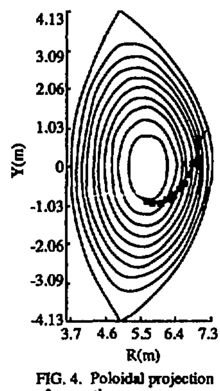

FIG. 4. Poloidal projection of ray path

The best penetration is achieved using long, narrow waveguides. An example of this is shown in Fig. 4 for the parameters given above with $P_{LH} = 1$ GW, $h = 0.1$ m, $L_{\|} = 10$ m and vertical launch point $y_0 = 1.0$ m. Each tick mark along the ray path represents a 10%

decrease in the wave power due to nonlinear damping. The maximum penetration of the ray was to $\psi(x_p) \simeq -10.0$, where $n_e(x_p) \simeq 7.1 \times 10^{19}$ m^{-3} and $T_e(x_p) \simeq 31$ keV (here the magnetic axis corresponds to $\psi \simeq -11.0$). Rays launched from the midplane or below did not get as close to the axis. While a wide launcher results in superior wave penetration, it may be difficult to install in a tokamak. Therefore, we considered a more conventional r.f. beam with $L_{\parallel} = 1$ m, $h = 1$ m and $P_{LH} = 1$ GW (simple port size in ITER). Then, the maximum penetration distance of the beam was only $\psi(x_p) \simeq -5.0$, $n_e(x_p) \simeq 3.5 \times 10^{19}$ m^{-3}, and $T_e(x_p) \simeq 20$ keV. Finally, we note that as P_{LH} is increased from 1 GW to 50 GW with fixed waveguide dimensions of $L_{\parallel} = 1$ m and $h = 1$ m, the maximum penetration distance increased from $\psi(x_p) \simeq -5.0$ to $\psi(x_p) \simeq -7.0[n_e(x_p) = (3.5 \rightarrow 5.0) \times 10^{19}$ m$^-3$ and $T_e(x_p) = (20 - 24.5)$ keV]. The saturation in x_p with increasing r.f. power is mostly due to an increase in the nonlinear broadening of the plateau width (Δu).

Parametric instabilities may occur during high power LH pulse injection, especially in the plasma edge region. This may, in turn, restrict the allowed maximum peak power density, and/or maximum pulse length. A detailed theoretical and numerical study of parametric instabilities has been carried out using the previously developed theory and code.[3,4] The results may be summarized as follows. The dominant instability in the edge region ($n_e(a) \gtrsim 3 \times 10^{19}$ m^{-3}, $T_e \gtrsim 1$ keV) is a dissipative ion cyclotron quasi-mode, coupled with a lower-hybrid side-band, for incident power densities $S_{\perp} = P_{LH}/A < 10$ GW/m^2. The instability is driven by both $\mathbf{E}_{\perp} \times \mathbf{B}_0$ drift and parallel ($\mathbf{E}_{\parallel}$) drift of electrons. Typical growth rates are $\gamma/\omega_0 \gtrsim 0.02$ for power densities $S_{\perp} \gtrsim 10$ GW/m^2, depending on edge parameters. The lower density and temperature regions ($n_e \sim 10^{19}$ m^{-3}, $T_i \sim T_e \sim 100$ eV) are more susceptible to instability, and growth rates of the order of $\gamma/\omega_0 \sim 0.01$ may be achieved with power densities as low as $S_{\perp} \sim 1$ GW/m^2. However, the growth rates decrease rapidly as the temperature and density rise. The growth may be limited by convective losses which limit spatial growth and/or limited pulse length. To limit growth to acceptable values, it may be necessary to cluster waveguide arrays to axial widths of $\Delta z \gtrsim 30$ cm, or heights normal to field lines of $\Delta y \gtrsim 10$ cm, or pulse lengths $\Delta t \gtrsim 10$ ns. Limiting the height is the best choice for good penetration. The waveguide clusters may have to be stacked after suitable spatial separation depending upon pulse length and rep-rate. High rep-rate (50 kHz) may be required to achieve the necessary average power for reactor applications.

In conclusion, we find that the use of pulsed power significantly enhances lower-hybrid wave penetration, and is consistent with avoidance of severe parametric instabilities for suitable choices of launcher shape (narrow strips, elongated along field lines, are best), pulse duration and size. Unresolved issues include nonlinear filamentation, quantification of the charge-separation effect, and the feasibility of long, skinny launching structures.

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