linear geometry where end losses dominate device performance.

In summary, a new magnetic confinement device has been built and tested. The principal advantages of this device are direct heating of the ions and high plasma β values (30%) since the confining field is surfacelike and the hot plasma is confined in the interior low-field region. The physical picture of enhanced ion current with the electron flow impeded by the cusp field and direct ion heating was confirmed with a particle simulation. The experimental results are a significant improvement over previous experiments. We feel that a suitably constructed TCX is possibly scalable to a reactor suitable for burning DT and advanced fuels.

We gratefully acknowledge the assistance of W. A. Peebles, H. Park, and D. Brower in helping to make the far-infrared laser and visible spectroscopic measurements.

This work was supported by the U. S. Department of Energy, Office of Fusion Energy, under Contracts No. DE-AM03-76SF00010 and No. PA DE-AT03-76ET53019.

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Radio-Frequency Current Drive in a Fusion-Producing Plasma

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A method for driving current in a fusion-producing plasma is proposed. Radio-frequency power is used to prohibit fusion-produced energetic particles from slowing down isotropically or to push them in a preferential direction. As a result, a net plasma current is generated whose efficiency is comparable to other current drive schemes.

PACS numbers: 52.50.Gj, 52.55.-s

Advantages of steady-state operation of fusion reactors have long been recognized. In some reactor concepts, such as tokamaks and pinches, the magnetic field produced by the current flowing in the plasma is an essential part of the confinement. The most commonly used method for driving plasma current in these experiments is to induce a toroidal electric field in the plasma by means of a time-varying magnetic flux. This method cannot operate in a steady-state fashion; therefore, alternative methods capable of driving plasma current continuously are desirable. Neutral beams,[1] heavy charged-particle beams,[2] and various rf waves[3-8] have been proposed for maintaining a steady-state plasma current by imparting to the electrons the momentum needed to compensate the resistive losses or by making the resistivity asymmetric.

In this paper, we propose to use the high-energy α particles produced by fusion reactions to sustain a steady-state current in plasma. The idea is to use rf power to prohibit the α particles from slowing down isotropically and to push the α particles in a preferential direction and thus form an α-particle beam. This α-particle beam will then transfer the momentum to electrons and sustain a current.

In a reactor, charged particles in the megaelectronvolt range (e.g., α particles) are produced along with neutrons by fusion reactions[9] using fuels like D and T and advanced fuels like D-D, D-^{3}He, etc. These α particles are born isotropically in velocity space and slow down via Coulomb interactions with background plasma particles. α particles interact predominantly with electrons until their energy is reduced to the critical value $\sim(m_\alpha/m_e)^{1/3}T_e$, where T_e is the electron temperature, and m_e and m_α are the masses of electrons and α particles, respectively. The α-particle distribution function remains isotropic during the slowing down, and there is no net current. In order to generate current, the α-particle distribution has to be made asymmetric by some means, e.g., rf power or particle loss.[10]

For simplicity, we assume that an α-particle beam is formed by injection of rf power. We would like to keep the α-particle beam energy above E_c, and thus minimize direct α-ion interaction. In steady state, the total forces on α particles and electrons are zero; then using the electron force balance relation one obtains an expression for current density:

$$J = n_\alpha Z_\alpha e(\bar{v}_\alpha - \bar{v}_e)(1 - Z_\alpha/Z_i). \tag{1}$$

We see that there is a net current (Ohkawa current)[1] in the plasma if the charge Z_i on the fuel ions is different from that on the α ions ($Z_\alpha = 2$). A more appropriate expression for the net current has been given by Start *et al.*[11]

As discussed before, energetic particles are born isotropically as a result of fusion reactions, and these particles interact primarily with plasma electrons. Injection of suitable rf waves traveling in a given direction would form, effectively, an α-particle beam. This we may identify as a state of induced anisotropy. Among the various rf waves considered for reactor plasma heating, the magnetosonic wave (fast wave)[12] seems particularly suitable for producing such a beam. The phase speed of such waves is of the order of Alfvén speed, which, in a reactor-type plasma, is roughly the same as the speed of α particles.

For the magnetosonic wave, we determine the momentum transferred to α particles through Landau damping and transit-time damping processes using the quasilinear equation for particle distribution function:

$$\frac{\partial f_\alpha}{\partial t} = \pi\left(\frac{e_\alpha}{m_\alpha}\right)^2 \lim_{V\to 0} \frac{V}{(2\pi)^3} \frac{\partial}{\partial v_\parallel} \int d^3k\, \delta(k_\parallel v_\parallel - \omega) \left| E_\parallel - \frac{i m_\alpha k_\parallel v_\perp{}^2 B_\parallel}{2 e_\alpha B_0}\right|^2 \frac{\partial}{\partial v_\parallel} f_\alpha . \tag{2}$$

The symbols ω, $k_\parallel$, and $B_\parallel$ represent wave frequency, parallel wave number, parallel wave electric field, and parallel wave magnetic field, respectively. In Eq. (2), we have integrated over all the modes in volume V that are excited externally, in order to obtain the necessary current drive. We shall only consider circulating α particles, and discuss the effect of the magnetically trapped particles later. We start with the α-particle distribution function given by

$$f_\alpha = n_\alpha \delta(v - v_\alpha)/4\pi v_\alpha{}^2. \tag{3}$$

Here, n_α is the density per unit volume, and v_α is the velocity of α particles at birth.

For Landau damping, interaction of the parallel component of the wave electric field with α particles whose parallel velocities are equal to the wave phase velocity leads to transfer of momentum and energy to the particles. For transit-time damping, it is the $\mu_\alpha \nabla B_\parallel$ term that transfers momentum to the particle, μ_α being the magnetic moment.

The rate of momentum input into α particles per unit volume is

$$\dot{P}_{\alpha\parallel} = 2\pi m_\alpha \frac{\partial}{\partial t} \int v^3\, dv \cos\theta \sin\theta\, d\theta f_\alpha, \tag{4}$$

where θ is the angle between the direction of the velocity vector and that of the toroidal magnetic field. Combining Eqs. (2) and (3) with Eq. (4), one obtains

$$\dot{P}_{\alpha\parallel} = 2\pi m_\alpha \int v^3\, dv \cos\theta \sin\theta\, d\theta \left(\cos\theta \frac{\partial}{\partial v} - \frac{1}{v\sin\theta}\frac{\partial}{\partial\theta}\right)\pi\left(\frac{e_\alpha}{m_\alpha}\right)^2 \frac{V}{(2\pi)^3}$$
$$\times \int \frac{d\vec{k}}{k_\parallel} \left| E_\parallel - \frac{i m_\alpha k_\parallel v_\perp{}^2 B_\parallel}{2 e_\alpha B_0}\right|^2 \delta\left(v\cos\theta - \frac{\omega}{k_\parallel}\right)\left(\cos\theta\frac{\partial}{\partial v} - \frac{1}{v\sin\theta}\frac{\partial}{\partial\theta}\right)\frac{n_\alpha \delta(v - v_\alpha)}{4\pi v_\alpha{}^2}. \tag{5}$$

The current thus generated may be obtained as

$$j_\alpha = (e_\alpha/m_\alpha)\dot{P}_{\alpha\parallel}\tau_s\,, \tag{6}$$

where τ_s is the characteristic time for slowing down of α particles on electrons. After some algebra, it can be shown that

$$j_\alpha = \frac{3n_\alpha e_\alpha^3\tau_s}{8\pi^2 m_\alpha^2 v_\alpha^2}\lim_{V\to 0} V\int\frac{d\vec{k}}{|k_\parallel|}|E_\parallel|^2\left(\frac{\omega}{k_\parallel v_\alpha}\right)^3 \tag{7a}$$

when Landau damping dominates, and

$$j_\alpha = \frac{e_\alpha n_\alpha v_\alpha^2\tau_s}{16\pi}\lim_{V\to 0} V\int\frac{d\vec{k}}{|k_\parallel|}k_\parallel^2\left|\frac{B_\parallel}{B_0}\right|^2\left\{15\left(\frac{\omega}{k_\parallel v_\alpha}\right)^9+\left(\frac{\omega}{k_\parallel v_\alpha}\right)^7-2\left(\frac{\omega}{k_\parallel v_\alpha}\right)-\left(\frac{\omega}{k_\parallel v_\alpha}\right)^3-13\left(\frac{\omega}{k_\parallel v_\alpha}\right)^5\right\} \tag{7b}$$

when transit-time damping prevails. With knowledge of j_α, the density of the return current in the plasma can be obtained.

An estimate of the energy input $\dot{W}_\alpha$ can also be obtained in a manner similar to that of $\dot{P}_{\alpha\parallel}$. After some algebra, one can show that

$$\dot{W}_\alpha = \pi m_\alpha\frac{\partial}{\partial t}\int v^2\,dv\sin\theta\,d\theta\,v^2 f_\alpha = \frac{\pi m_\alpha n_\alpha}{2v_a}\left(\frac{e_\alpha}{m_\alpha}\right)^2\lim_{V\to\infty}\frac{V}{(2\pi)^3}\int\frac{d\vec{k}}{|k_\parallel|}|E_\parallel|^2\left(\frac{\omega}{k_\parallel v_\alpha}\right)^2 \tag{8a}$$

when Landau damping dominates, and

$$\dot{W}_\alpha = \frac{m_\alpha n_\alpha v_\alpha^3}{8\pi^2}\lim_{V\to\infty} V\int\frac{d\vec{k}}{|k_\parallel|}k_\parallel^2\left|\frac{B_\parallel}{B_0}\right|^2\left(\frac{\omega}{k_\parallel v_\alpha}\right)^2\left(\frac{\omega}{k_\parallel v_\alpha}+\frac{3}{4}\frac{\omega^4}{k^2 v^4}-\frac{7}{4}\right) \tag{8b}$$

when transit-time damping prevails. The corresponding electron return current is then given by

$$j_{\rm tot} = j_\alpha(1 - Z_\alpha/Z_{\rm eff})\,. \tag{9}$$

A quantity of practical interest is the efficiency of the current-drive scheme, viz., the current generated per unit of power:

$$I/P = \int j_{\rm tot}\,dS\Big/\int \dot{W}_\alpha\,dV\,. \tag{10}$$

By combining Eqs. (8), (9), and (18), and assuming $v_{\rm ph} = \omega/k_\parallel \gtrsim v_\alpha$, one can show that the above ratio is proportional to $T^{3/2}/nV_{\rm ph}$ and an approximate numerical estimate may be obtained (considering absorption only through transit-time effect) as $I/P \sim 0.5$ A/W, where we have assumed plasma parameters similar to that of an INTOR-type tokamak reactor:

$$r = 1.5\text{ m},\quad R = 5\text{ m},\quad V \approx 200\text{ m}^3,$$
$$\overline{T} = 20\text{ keV},\quad \overline{n}_e = 10^{20}\text{ m}^{-3},$$
$$B = 5\text{ T with } Z_{\rm eff} = 1.5\,. \tag{11}$$

The above estimate for current drive compares very favorably with other methods of current drive, viz., energetic neutral beam injection or other rf drive techniques.

It may be pointed out here that the efficiency is inversely proportional to $v_{\rm ph}$ which we have taken to be of the order of $\overline{v}_\alpha$. Since the factors which dictate the minimum value of $\overline{v}_\alpha$ (e.g., current strength, α-particle density, electron temperature, etc.) vary from device to device, no attempt has been made here to get maximum efficiency. Instead, $\overline{v}_\alpha$ has been taken to be the velocity corresponding to an average α energy of 1.9 MeV. We may also note that we have considered only the case of magnetosonic wave injection for two important reasons: The phase velocity along the field line is about the Alfvén speed, which is close to the α-particle speed for reactor-type parameters; furthermore, this wave can propagate deep into the plasma over a wide range of frequencies and thus provide the necessary push on the α particles needed for current generation.

In this frequency range, we expect magnetosonic waves also to interact with electrons through Landau damping and transit-time damping. The parallel electric field in the magnetosonic wave is related to the wave magnetic field by the relation $E_z \approx -(V_e^2/c\Omega_e)(\partial B_z/\partial z)$ (see, e.g., Ref. 12). It is particularly interesting to note there is a partial cancellation between Landau damping and transit-time terms for the electrons.[13] When we use the above relationship between E_z and B_z of the wave in order to obtain the contribution from α particles, Landau damping and transit-time items now add, with the primary contribution coming from the transit-time term.

When we consider the case of wave absorption through the transit-time damping mechanism, it can be shown that absorption by α particles is at least of the same order (if not higher) as compared to that by the electrons. This can be seen by comparing the electron absorption rate given by

$$\dot{W}_e = \frac{n_e m_e v_e^2}{4}\,\omega\pi^{1/2}\left(\frac{\omega^2}{k_\parallel^2 v_e^2}\right)\exp\left(-\frac{\omega^2}{2k_\parallel^2 v_e^2}\right)\left|\frac{B_\parallel}{B}\right|^2 \tag{12}$$

with that given in Eq. (8b). To calculate the electron absorption rate, we have taken a Maxwellian distribution for the electrons and assumed a monochromatic fast-wave spectrum, for the sake of simplicity. The magnetosonic waves may be generated as cavity modes and, with appropriate choice of frequency and wave number, one may obtain a spectrum for which the parallel electric field is much weaker than the magnetic field. Under such circumstances, the primary mode of energy absorption from waves is through transit-time damping. A numerical calculation has been carried out by Bhadra *et al.*[14] to elucidate the interaction of magnetosonic cavity modes with both α particles and electrons in a reactorlike plasma. With appropriate choices of frequency and wave number, a set of cavity modes have been obtained which generate currents both through direct interaction with electrons and through interaction with energetic α particles. The computed I/p ratios for α-particle–driven current are found to be a factor of 3 to 4 larger than that obtained from direct electron interaction. Considering the fact that the wave absorption rates by the two species are, at least, comparable, it is thus expected that the resultant current generated would be primarily due to the α-particle effect. Moreover, the current generated by the direct interaction of the electrons with the waves is additive to that due to the α particles for $Z_{\text{eff}} < Z_\alpha$.

In the above, we have not taken into account the effect of magnetically trapped electrons, either on the "return" current or concerning absorption of magnetosonic waves. In the light of present-day tokamak experiments, where electron confinement has been consistently found to be anomalous, it is not unreasonable to expect that the behavior of electrons in the tokamak plasmas cannot be modeled through the neoclassical theory of trapped electrons. Thus, the question regarding the nature of the effect of magnetically trapped electrons in tokamak plasmas remains open.

Also, in the above calculations, we have used a δ-function distribution for α particles (with some average energy of about 1.9 MeV), for the sake of simplicity. To be more realistic, one needs to use a distribution function which shows the effect of Coulomb slowing down, viz., $f_\alpha \to (v^3 + v_c^3)^{-3}$. The usage of such a distribution function brings in analytical difficulties without changing the basic physical nature of the current-drive mechanism proposed. However, such a distribution function is expected to require a wider band for the spectrum of fast waves in order to enhance the efficiency of this method.

In this paper, a method for maintaining a steady-state current in a fusing plasma has been proposed. Since the method employs an rf wave that is being considered as a likely candidate for reactor heating and uses α particles born in fusion reaction, such a scheme fits naturally into the steady-state fusion reactor operation. This method is not only applicable to a D-T burning fusion reactor, but also appears suitable for current drive in fusion reactors burning advanced fuels.

The processes related to the method proposed here have been observed in experimental conditions. On the adiabatic-toroidal-compressor tokamak, rf waves have been found to interact strongly with high-energy ion beams and to have a clamping effect on the beam.[15] Experimental observation of the Ohkawa current has also been carried out in some recent experiments on the DITE tokamak.[16]

This work was supported by the U. S. Department of Energy under Contract No. DE-AT03-76ET51011.

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Spectrum of Small-Scale Density Fluctuations in Tokamaks

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The power spectrum of plasma density fluctuations in the range of frequencies of drift waves has been investigated in the Princeton Large Torus tokamak. The results suggest that the observed fluctuations evolve to a strong nonlinear state, and therefore they emphasize the need for a complete nonlinear theory of drift waves in tokamaks to assess their effects on plasma transport.

PACS numbers: 52.25.Gj, 52.25.Lp, 52.55.Gb

The tokamak is the most successful of all the magnetic confinement schemes. Nevertheless, this type of configuration is still plagued by anomalous heat losses which remain one of its most critical issues.

It has been known for a long time that the turbulence produced by drift-wave instabilities can adversely affect the transport of plasma. More recently Horton and Estes[1] have shown that the confinement of electrons in several Ohmically heated tokamaks is consistent with the predictions of the quasilinear theory of drift waves. Nevertheless the role played by these instabilities in the transport of plasma is not yet completely understood.

In this Letter, new data on the spectrum of density fluctuations in the range of frequencies of drift waves are presented. Their implications are that a complete nonlinear interaction of drift waves is essential for assessing their effects on the confinement of plasma in tokamaks.

Drift waves are excited by the free energy stored in the plasma by macroscopic inhomogeneities. According to the linear theory of drift waves in tokamaks,[2] the range of frequencies is $\omega/\omega_*^e \lesssim 1$, where $\omega_*^e = k_\theta c T_e/eBL_n$ (with k_θ the poloidal wave number and $L_n = |d\ln n/dr|^{-1}$ the density scale length). The upper limit, $\omega \approx \omega_*^e$, is reached by electron drift waves with long wavelengths (i.e., $k_\theta \rho_i \ll 1$, with ρ_i the ion Larmor radius) that propagate along the electron diamagnetic direction. The ion inertia and the finite ion Larmor radius come into play when $k_\theta\rho_i \gtrsim 1$ with the result of decreasing the value of $|\omega|$. As the ion temperature profile becomes steeper than the density profile (i.e., $\eta_i = d\ln T_i/d\ln n > \eta_0 \simeq 1-2$) the sign of ω is reversed, and we have drift waves which propagate along the ion diamagnetic direction. Various experimental observations[3-5] have indicated that a small-scale turbulence exists in tokamaks in a range of frequencies around ω_*^e, with wavelengths $\lambda \gtrsim \rho_i$ and amplitude $\langle(\tilde{n})^2\rangle^{1/2} \simeq \langle n\rangle/(kL_n)$. One intriguing result has been the broad frequency spectra of the observed fluctuations whose spectral width $\delta\omega = (2-4)\omega_*^e$. If this is a real feature of the microturbulence of tokamaks it means that the plasma fluctuations induced by drift instabilities evolve very rapidly to a nonlinear state. Unfortunately the poor spatial resolution of those measurements could not rule out the possibility that the observed broadenings were simply produced by radial variations of plasma parameters. To clarify this crucial point, the power spectrum of density fluctuations in the range of frequencies of drift waves has been investigated in the Princeton Large Torus tokamak with scattering of microwaves.

An array of antennas, all located in the same poloidal plane, was used for launching an electromagnetic wave and for collecting the waves scat-

