REVIEW OF CURRENT DRIVE THEORY: SELECTED TOPICS

N. J. Fisch

Princeton Plasma Physics Laboratory
Princeton University, Princeton, NJ 08543

ABSTRACT

Two themes in current drive theory in tokamaks are reviewed, both relevant to the progression of tokamak experiments toward the reactor regime. First, we review our understanding of the physics of the tail electrons. These electrons are capable of carrying enormous rf-driven electric current, and, in the course of current-drive experiments worldwide not only has the current drive effect been demonstrated, but the underlying physical description of these tail electrons has been established. Second, anticipating the presence of the energetic alpha particles that result from D-T reactions in a reactor, we examine certain mechanisms through which these alpha particles can be used to facilitate current-drive.

KEYWORDS

Current drive; tokamaks; rf waves

1. INTRODUCTION

In recent years, there has been a great experimental effort in the area of rf driven currents. Almost every tokamak explores some aspect of noninductive currents. Nearly all of the research has been devoted to discovering and verifying methods of noninductive current drive for use in a reactor.

This review is highly selective; rather than reviewing the research in new methods of current drive, this article explores in some depth two themes: one, the physics of the tail electrons, and, two, current drive by waves in an alpha particle environment. The first theme is "backward-looking" — looking back upon the large theoretical and experimental effort in current drive by use of superthermal electrons, we find that we learned a great many important things that perhaps were not on our list of questions at the outset of this research effort. The second theme is "forward-looking" — looking ahead to the reactor regime, we ask where

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there might be opportunities to use the α-particles to benefit the current drive effect.

The discussion here draws primarily from the research of the author and his colleagues. For a more comprehensive review of the theory of current drive, and for a more complete list of references on the subject of current drive, see the review article by the author (Fisch, 1987).

2. Description of the Tail Electrons

The first theme is that in recent years a great deal has been learned about the physics of the tail electrons. Mechanisms of current drive that depend upon properties of these electrons have been discovered. New transport quantities, suitable for describing the unique physics associated with these electrons have been identified and calculated. Experimental evidence has corroborated in detail the theoretical description and the prediction of current drive mechanisms.

Moreover, and what has probably been overlooked or underappreciated in the excitement surrounding the finding of mechanisms of current drive, the very fact that these mechanisms could be predicted theoretically now gives us great confidence in the assumptions upon which the theory is based. To be specific, we now know that these fast electrons are governed by classical processes, at least so far as the parallel dynamics are concerned. This is not something that could be taken for granted or that we could have known for certain — prior to the extensive experimentation on the current-drive effect, it had been impossible to rule out the possibility that anomalous processes would dominate the dynamics of the fast, tail electrons.

The electron velocity distribution function f, in the vicinity of the fast electrons, can be described by writing $f = f_{\rm M}(1+\phi)$, where f_M is a Maxwellian distribution and ϕ describes a time-dependent perturbation brought about, for example, by rf heating. An excellent approximation is to treat ϕ as small in its contribution to the collision integral, so that f obeys the linearized Fokker-Planck equation. The evolution of ϕ may then be written as

$$f_{\rm M}\frac{\partial\phi}{\partial t} + \frac{q}{m}\mathbf{E}\cdot\frac{\partial}{\partial\mathbf{v}}f_{\rm M}\phi - C(\phi) = -\frac{\partial}{\partial\mathbf{v}}\cdot\mathbf{S} \tag{1}$$

where $\mathbf{v}$ is the electron velocity, C is a linearized collision term and $\mathbf{S}$ is an externally induced flux in velocity space. For example, in the case of an rf-induced flux that can be described well by quasilinear theory, we can write

$$\mathbf{S}(\mathbf{v}) = +\frac{q}{m}\mathbf{E}f_{\rm M} - D_{QL}(\mathbf{v})\cdot\frac{\partial}{\partial\mathbf{v}}f_{\rm M}(1+\phi), \tag{2}$$

where $D_{QL}(\mathbf{v})$ is the rf diffusion tensor. This diffusion tensor is proportional to the rf energy density excited in the plasma. Note that the first term in the expression for the flux is independent of both the perturbation and the rf excitation;

this term is responsible for the ohmic current. The second term is linear in the rf excitation.

There is a very powerful way of extracting from this equation important quantities that characterize particularly the salient properties of driven plasmas, such as the current drive efficiency (Fisch and Boozer, 1980), the runaway production rate (Fisch and Karney, 1985), the rf-enhanced conductivity (Fisch, 1986), and the rf enhancement to the radiation (Fisch, 1988). Suppose that the flux $\mathbf{S}$ is a given quantity, then equation (1) is a linear inhomogeneous partial differential equation for ϕ, and therefore may be solved by Green's function techniques. The power of this technique is that in many cases the exact distribution or, equivalently, the exact rf-induced flux need not be known exactly, but what is needed is the ratio of different plasma responses to the same excitation. For example, in calculating the current-drive efficiency, what is needed is the ratio of current to power dissipation, both linear functionals of the induced flux.

Thus, one can write the steady state current drive efficiency (in the absence of a dc electric field) as

$$\frac{J_{rf}}{P_d} = \frac{\mathbf{s}\cdot(\partial/\partial\mathbf{v})\,\chi(\mathbf{v})}{\mathbf{s}\cdot(\partial/\partial\mathbf{v})\,\epsilon(\mathbf{v})}, \tag{3}$$

where $\mathbf{s}$ is the unit vector in the direction of the rf-induced flux, $\chi(\mathbf{v})$ is the Green's function for the current drive, $\mathbf{v}$ is the velocity of the resonant electrons, *i.e.*, the vicinity in velocity space of the induced flux, and $\epsilon(\mathbf{v}) = mv^2/2$ is the kinetic energy of the resonant electrons.

In a similar fashion, one can define the so-called "hot" conductivity, $\sigma_H(P_d)$, which characterizes the increase in conductivity in an rf driven plasma. To be specific, write the total current as

$$J = \sigma_{sp}\,E + J_{rf} + \sigma_H(P_d)\,E, \tag{4a}$$

where the first term on the right is the ohmic current, the second term is the rf-driven current, and the third term, a function of the dissipated power and linear in the electric field E, is the enhancemnt to the ohmic current on account of the nonmaxwellian features caused by the rf dissipation in the plasma. Here, one can expand in the electric field the Green's function $\chi = \chi_0 + E\chi_1 + \cdots$, where χ_0 is the Green's function in the absence of the electric field, as in equation (3), and the hot conductivity takes the form

$$\sigma_H(P_d) = P_d\,\frac{\mathbf{s}\cdot(\partial/\partial\mathbf{v})\,\chi_1(\mathbf{v})}{\mathbf{s}\cdot(\partial/\partial\mathbf{v})\,\epsilon(\mathbf{v})}. \tag{4b}$$

As a third example of an important transport quantity linear in the power dissipated, consider the *incremental* runaway production rate can be written as

$$\dot{N}_R = \int d^3v\,\mathbf{S}\cdot\frac{\partial}{\partial\mathbf{v}}R(\mathbf{v}), \tag{5}$$

where $R(\mathbf{v})$ is the runaway probability of an electron with initial position in velocity space $\mathbf{v}$. This runaway probability is a well-defined quantity: it is the probability that an electron will first be accelerated to some arbitrarily high speed before being decelerated to some arbitrarily low speed — and this quantity has a precise definition in the dual limit. In Fig. 1, we show contours of $R(\mathbf{v})$ from Karney and Fisch (1986). Note that the runaway probability is 100% not only for electrons with large velocities in the direction of the electric force, but even for some electrons traveling initially at high speeds in the opposite direction.

The physics of fast electrons in rf-driven plasmas can be described by transport quantities such as $R(\mathbf{v})$ and $\chi(\mathbf{v})$. Note that these quanties are functions of velocity, *i.e.*, they depend on where in velocity space electrons are pushed by the rf waves. Thus, for example, the rf curent drive efficiency or the rf runaway probability increases with increasing velocity of the resonant electrons. Moreover, *differential* effects that might be the result of an rf-induced flux, such as the *incremental* runaway production or the *incremental* synchrotron radiation, depend also upon the direction of the flux.

Note that this description is, in principle, far more detailed than descriptions of bulk transport quantities, since there is the detailed dependence on velocity space location. For example, the Spitzer conductivity is an integrated quantity, the integration being caried out over the full velocity distribution, but the "hot conductivity" is a function of where and how the rf is absorbed in velocity space. Similarly, the Dreicer runaway velocity is just one number, whereas the runaway probability function $R(\mathbf{v})$ assigns to every velocity space location $\mathbf{v}$ a unique runaway probability. It is just this detail that often proves extremely useful; for example, the detailed description of the current drive efffect points to the most efficient waves to employ, and the details of the runaway production rate allow us to deduce (see below) bounds on confinement times.

These transport quantities enable a detailed description of the rf-driven plasma, but how can they be employed without knowing in detail the velocity space location of the rf-induced flux? In fact, in many problems, the important features of the flux can be deduced through other considerations, making the Green's function approach more powerful yet. The direction of the flux is given by the nature of the rf waves; for example, for electrostatic waves the flux is entirely in the parallel direction. Moreover, the resonance condition very much localizes $\mathbf{s}$ in velocity space too. Therefore, the flux is often known up to a multiplicative constant. This multiplicative constant is the same, however, for all quantities of interest, so the *ratio* of quantities of interest, such as the ratio of current generated to power dissipated, is determined.

The localization of $\mathbf{s}$ in velocity space can be seen, for example, in the case of current drive by lower hybrid waves (Fisch, 1978). These waves damp on superthermal electrons, namely those with parallel velocities about 4 or 5 times

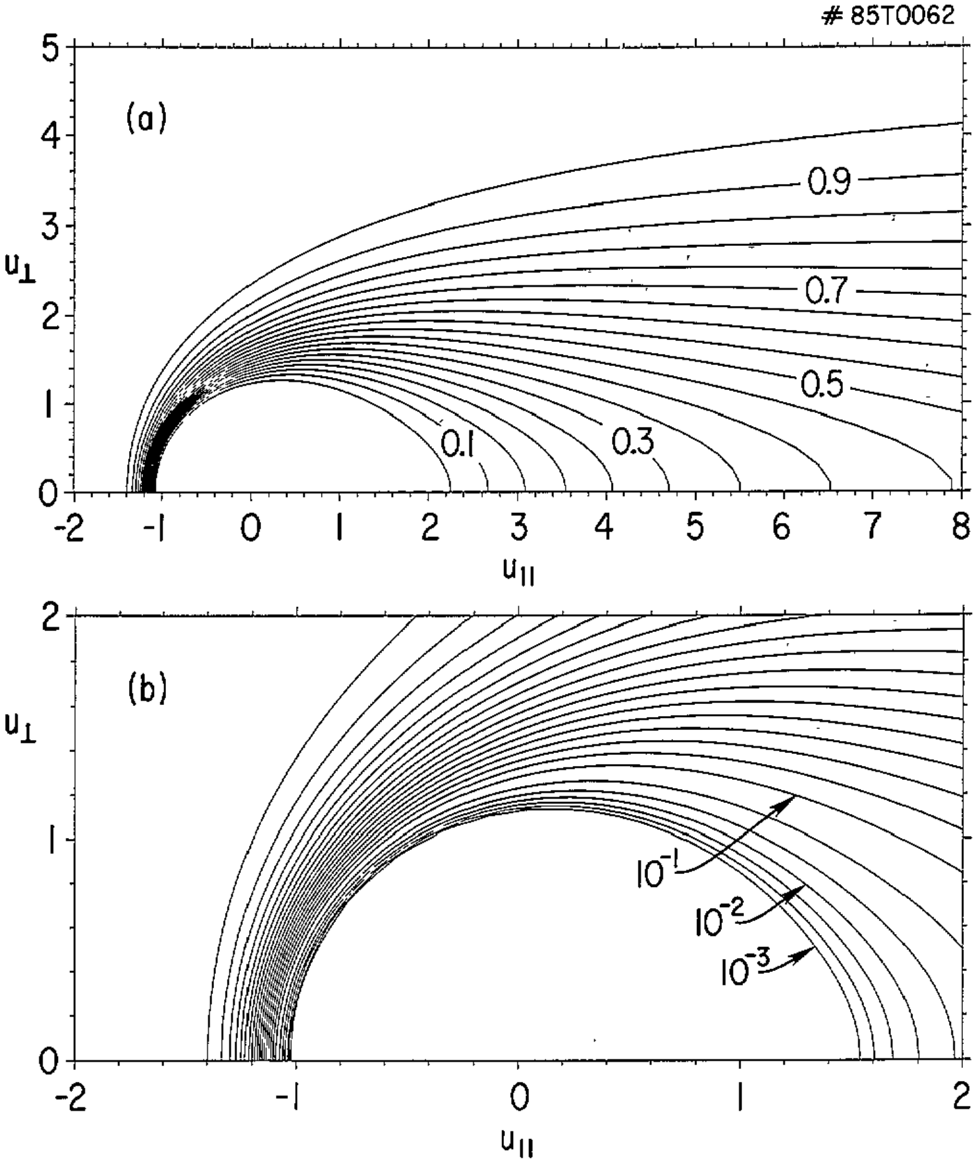

Fig. 1. Runaway probability $R(\mathbf{u})$, where $\mathbf{u} \equiv \mathbf{v}/v_R$ is a velocity normalized to a runaway speed. Parts (a) and (b) show R on two different scales. In (a), the contours are equally spaced at intervals of 0.05. In (b), the lowest seven contours are geometrically spaced at intervals of $10^{1/3}$, between 10^{-3} and 10^{-1}; the remaining contours are spaced equally at intervals of 0.05 as in (a) (Karney and Fisch, 1986).

the thermal velocity. The parallel velocity is given, in fact, through the resonance condition, but the perpendicular velocity, except in highly unusual cases, may be deduced too. For an rf spectrum that is not very broad, the electron distribution remains roughly Maxwellian in the perpendicular velocity direction, even in the presence of intense rf power. This means that $v_{\parallel} \gg v_{\perp}$ for resonant electrons, which is sufficient localization to calculate accurately most quantities of interest, such as the current drive efficiency. Even when the rf spectrum is broad enough so that the amount of current generated may be difficult to surmise in the absence of a full Fokker-Planck calculation, the rf-induced fluxes might still be sufficiently localized for purposes of calculating the current drive efficiency.

3. Evidence for the Classical Description

Before the intensive experimental investigation of the current drive effect through rf absorption by fast electrons, it was assumed but unproved that the superthermal electrons could be described by classical collision theory, in other words, it remained to question whether some anomalous collective effect, possibly involving a collisionless instability associated with an asymmetric velocity distribution function, or possibly through Cerenkov or other collective emission, might dominate the slowing down and scattering of an enhanced tail of superthermal electrons in a tokamak.

The fact that Spitzer resistivity appeared to be confirmed in other experiments was only proof that the integrated resistivity was dominated by classical collision processes; however, for effects that might depend largely on the dynamics of a relatively small percentage of the electron population, the integrated result is not sufficient evidence.

It is, in fact, the lower hybrid current drive and current ramp-up experiments that now rule out any such surprises. The theoretical predictions concerning the current-drive effect have now been documented in detail and in large experiments — as much as 2 MA of current have now been driven by waves in the JET and JT-60 tokamaks. Since the current drive effect itself relies upon a detailed description of the electron collisions, the evidence that documents the effect also documents in detail the dynamics of the fast electrons, showing that these electrons indeed are dominated by the classical collision processes.

The most detailed evidence comes from the PLT (Princeton Large Torus) series of current-drive and ramp-up experiments (Jobes *et al.*, 1985). In the ramp-up experiments, in which the current rises due to the rf waves, as much as 40% of the rf power was converted into poloidal field energy. The fact that the energy conversion can be so efficient, something that is consistent with the theory, is strong evidence for the model. These experiments spanned several parameter regimes, leading to different physics regimes too, including that of steady-state current drive, ramp-up of the current, and even the unsuccessful sustainment of the current.

In Fig. 2, we show the experimental data from PLT as plotted by Karney *et al.* (1985), where an attempt was made to check the theory of the electron dynamics without making many assumptions concerning the details of either the theory of wave propagation or wave damping. This was accomplished by comparing dimensionless quantities, each of which depended upon the wave being absorbed. Over 250 shots were tabulated, and with formally only two free adjustable parameters (but related parameters, so essentially only one free adjustable parameter), the fit to the theoretical prediction was remarkable. The choice of these two free parameters was made to avoid resolving the so-called "spectral gap problem," namely, the observation that the lower hybrid waves appear to damp on electrons even when the launched spectrum of waves appears to contain parallel phase velocities so high as to be unlikely to interact with very many electrons. There are several possibilities that may account for this effect, including the possibility that the phase velocity of the absorbed spectrum of waves is slower, through an upshift in wavenumber, than is the phase velocity of the launched spectrum of waves. In the analysis of the PLT data, the upshift in wavenumber and the amount of absorption were the formal adjustable parameters. The two parameters are related, since greater upshift results in greater absorption.

What were the implications of these experiments? The immediate interpretation concerns the parallel velocity space dynamics of fast electrons. The dependencies of the electron dynamics upon density, electric field, and parallel phase velocity must, to make such a fit, be as predicted classically, leaving little room for anomalous effects. Note that this is a far deeper statement on the nature of the electron dynamics than could be offered merely through a detailed study of Spitzer resistivity, which is an integrated quantity.

Consistent results were obtained on other tokamaks such as the MIT Alcator C tokamak (Porkolab, 1985), and the Asdex tokamak (Leuterer *et al.*, 1985).

4. Confinement Bounds From Current Ramp-up Experiments

The current drive and current ramp-up experiments can be used to provide both upper and lower bounds to the confinement time of the fast electrons. The fact that the current drive and current ramp-up effects were observed at all indicates immediately that the fast electrons must be reasonably well confined, at least on the order of their collision time. What has not been appreciated is that these experiments also give an upper bound to the confinement time of these electrons. The upper bound may be deduced from the high efficiency of the energy conversion.

The upper bound arises because, during ramp-up, there is an electric field that

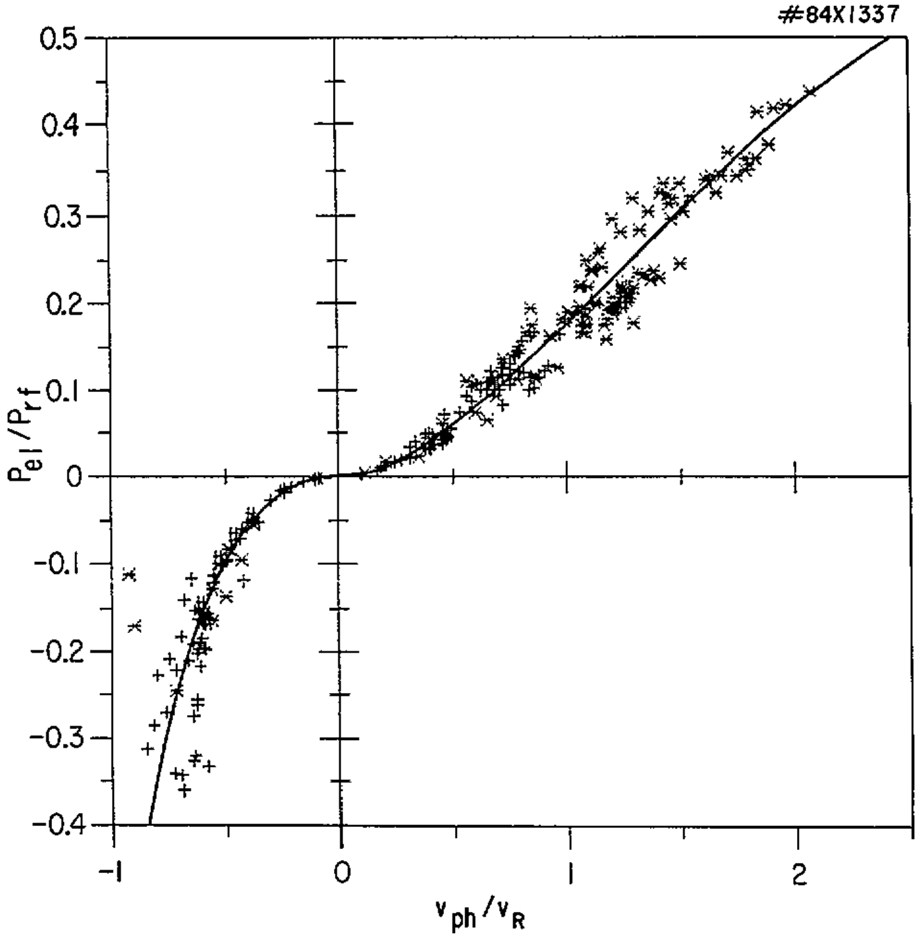

Fig. 2. P_{el}/P_{rf} vs. v_{ph}/v_R for 273 PLT shots. Here, v_{ph}/v_R is a normalized wave phase velocity, and P_{el}/P_{rf} is a normalized measure of ramp-up efficiency. The rf power P_{rf} varied from 0 to 300 kW, the density varied from 1.5×10^{12} to 6×10^{12} cm^{-3}, and the plasma current varied from 150 to 400 kA. Three wavedguide phasings were used, 60° (∗), 90° (+), and 135° (#). (Karney, Fisch, and Jobes, 1985).

opposes the current drive. If there are runaway electrons accelerated by this electric field — the so-called "backward" runaways, because they flow opposite to the current carriers — the ramp-up is hampered. If there are substantial backward runaways that are confined indefinitely, the plasma behaves essentially as a superconductor and it resists a change in the current.

To calculate the upper bound, note that, apart from the problem of backward runaways, the conversion of wave energy to poloidal field energy during ramp-up is efficient only if the rf works primarily against an opposing electric field rather than against collisions. This implies that the electric field must be strong, which, however, is the limit in which we might expect substantial backward runaway electron production. The runaway electrons, in undergoing constant acceleration by the dc electric field, convert the field energy back to electron kinetic energy, thus limiting the conversion efficiency.

Even a small number of runaways ruins the ramp-up effect. Every resonant electron supports the current ramp-up for a slowing down time, but, in the event that it becomes a backwards runaway, works against it, with even greater efficacy, for a confinement time. Therefore, a necessary assumption in explaining the observed data was that these runaway electrons could not be long confined.

First, let us calculate the lower bound on the comfinement time of superthermal electrons; the current carriers (in the forward direction) must be confined in order to achieve the current drive effect. The current drive effect relies upon the longer slowing down time of electrons when propelled to higher velocities, so that the full current drive effect is achieved in about one collision time, τ_c, of these fast electrons. Thus, if the effect is not to be destroyed or substantially diminished by the finite confinement time $\tau^f_{\rm conf}$ of the fast electrons, we must have

$$\tau_c \ll \tau^f_{\rm conf}. \tag{6}$$

To get a rough upper bound, we require that the countercurrent, $J_{\rm back}$ be much less than the forward current $J_{\rm for}$, *i.e.*,

$$J_{\rm back} = en_R v_R \ll J_{\rm for} = en_f v_f, \tag{7}$$

where n_R and v_R are the backward runaway density and velocity respectively, and n_f and v_f are the forward current carrier density and velocity respectively. (Here, v_R is a fluid velocity or average velocity of the runaway distribution, not to be confused with v_R in Fig. 1, which is proportional to the Dreicer runaway speed.) Equation (7) is equivalent also to the statement that the rf power that flows into the magnetic field energy during ramp-up of the current is less than the power that flows into the kinetic energy of the backward runaways.

Now, whenever a forward current carrier is created by the rf, there is a probability $R(v)$ that it will become a backward runaway; hence, the rate of production of backward ruanaways is related directly to the rate of promotion of bulk electrons to the tail of the distribution, *i.e.*

$$\dot{n}_R = R(v)\dot{n}_f. \tag{8}$$

Here, we assume that any production of runaways comes from the rf-generated superthermal tail of the electron distribution, rather than from the bulk of the distribution as might happen in very strong electric fields (the so-called "slide-away" effect).

Let us approximate the forward current carrier velocity v_f and the backward runaway velocity v_R as fixed; in practice, forward current carriers spend most of their lifetime (before slowing down) at about the initial, resonant parallel velocity, which in the case of lower hybrid current drive is several times the thermal velocity, so one can take $v_f \simeq 4v_T$. On the other hand, if the confinement time of backward runaways $\tau^R_{\rm conf}$ is long, the backward runaway velocity v_R quickly becomes relativistic, so that $v_R \simeq c$, where c is the velocity of light.

With the velocities fixed, the forward and backward currents change according to

$$\begin{aligned} \frac{d}{dt}J_{\rm back} &= e\dot{n}_R v_R - J_{\rm back}/\tau^R_{\rm conf} \\ \frac{d}{dt}J_{\rm for} &= e\dot{n}_f v_f - J_{\rm for}/\tau_c, \end{aligned} \tag{9}$$

where, we assume, the forward currents are largely destroyed by collisions (since, as per equation (5), $\tau_c \ll \tau^f_{\rm conf}$), but the backward runaway current, being more energetic by far and hence less collisional, is destroyed when the runaways leave the tokamak. For the PLT experiment, the typical ramp-up times are larger than collision or confinement times, so one can approximate $d/dt \simeq 0$ above, so that we can write

$$\begin{aligned} J_{\rm back} &= \tau^R_{\rm conf} e\dot{n}_R v_R \\ J_{\rm for} &= \tau_c e\dot{n}_f v_f. \end{aligned} \tag{10}$$

Now, using equations (8) and (10), we can rewrite equation (7) as

$$\frac{J_{\rm back}}{J_{\rm for}} = \frac{\tau^R_{\rm conf}}{\tau_c}\frac{v_R}{v_f}R(v) \ll 1. \tag{11}$$

This represents an upper bound to the confinement time of the backward runaway electrons.

Let us make the rough assumption that the confinement times of all fast superthermal electrons are the same, whether we are considering the tail current carriers or the backward runaways, *i.e.*, $\tau^R_{\rm conf} = \tau^f_{\rm conf} \equiv \tau_{\rm conf}$. Then, the following rough bounds on the confinement time of the fast electrons can be deduced from the PLT data:

$$\frac{v_f}{v_R}\frac{\tau_c}{R({\rm v})} > \tau_{\rm conf} > \tau_c, \tag{12}$$

where $\tau_c = \tau_c(\mathbf{v})$ is the slowing down time of an electron initially resonant at velocity $\mathbf{v}$. The second inequality states merely that electrons are confined long enough to be slowed down classically, something we believe to occur because of the remarkable fit of the theory to the data. The first inequality states that the effect of runaway electrons must be less important than the effect of resonant electrons.

Typical slowing down times in these experiments were on the order of ten milliseconds and typical runaway probabilities in the resonant region were on the order of 0.1. If we take $v_f/v_R \simeq 1/5$, we can deduce from equation (12) a confinement time in PLT of between one and two τ_c, or about fifteen milliseconds.

5. Tapping Free Energy in α-Particles

The second theme of this paper is that, in order to improve greatly on the efficiency of lower hybrid and possibly other kinds of current drive, the α-particles in a reactor might be put to use. Because of the substantial theoretical understanding and experimental database of current drive by lower hybrid waves, this method is the first focus of attention in extrapolating current drive to a reactor. At the same time, because our understanding of current drive by lower hybrid waves is fairly complete, unless some new physics is taken into account, such as the interaction with α-particles, it would be unlikely that there would be found in a reactor an efficiency substantially different from the presently predicted efficiency. This section summarizes research by Fisch and Rax (1992b) on possibilities in using α-particles.

It had been thought that in the most favorable wave regimes in a tokamak reactor, the current-drive efficiency may be reduced greatly because these waves tend also to be absorbed by energetic α-particles (Wong and Ono, 1984). The calculation by Wong and Ono assumes, however, an infinite homogeneous plasma. It is possible, however, that, since the α-particles do tend to concentrate near the tokamak center, there may be free expansion energy that may be tapped. Then the α-particles need not damp the wave, and may even under certain circumstances amplify the wave. The wave amplification by α-particles can be accompanied, of course, by wave damping by electrons, leading to very efficient current drive.

The α-particles may be supposed to have a reasonably steep spatial gradient. The spatial gradient arises naturally from the fusion production rate; this rate is proportional to the square of the ion density, and, near the operating point of a D-T tokamak reactor, also would be about proportional to the square of the ion temperature. Whatever spatial peaking accompanies the plasma pressure is therefore exaggerated in the α-particle density profile.

With waves, there is the opportunity to diffuse the α-particles along a direction in energy-configuration space. In other words, even if the α-particles are monotonically decreasing in energy on each magnetic surface, but do exhibit a large spatial gradient, then there may be more energetic α-particles in the high density central region than there are low energy α-particles towards the tokamak periphery. Under the influence of suitable waves, the α-particles would tend to diffuse in energy-configuration space to the less energetic peripheral region. A wave that would accomplish this is an electrostatic wave, such as the lower hybrid wave, with substantial poloidal momentum.

The effect of such a wave, say traveling in the y-direction (poloidal) but causing diffusion in the x-direction (radial), can be captured in a one dimensional diffusion equation (Fisch and Rax, 1992a). By integrating the quasilinear diffusion equation for the α-particles over parallel velocity space, one can obtain

$$\frac{\partial}{\partial \tau}F(\epsilon, X, t) = \frac{\partial}{\partial \epsilon}\epsilon F + \left(\frac{\partial}{\partial \epsilon} + \frac{\partial}{\partial X}\right)\frac{D(\epsilon, X)}{\sqrt{\epsilon - \epsilon_w}}\left(\frac{\partial}{\partial \epsilon} + \frac{\partial}{\partial X}\right)F + \frac{S}{\sqrt{1-\epsilon}}, \quad (13)$$

where F is the α-particle distribution integrated over parallel velocity, and where we have used the normalized variables $\epsilon = v_\perp^2/v_\alpha^2$, $X = x/L$, $\tau = 2\nu t$, and $\epsilon_w = (\omega/k_y v_\alpha)^2$; and where $\nu = 16\sqrt{2\pi m_e}e^4 n_e \ln\Lambda/3T_e^{3/2}m_\alpha$ is the slowing down rate of α-particles on electrons, where v_α is the α-particle birth speed, and where $L = v_\alpha^2 k_y/2\omega\Omega_\alpha$. Here, ω is the wave frequency, and k_y is the wave poloidal wavenumber. The first term on the right represents slowing down on electrons, the dominant collisional relaxation for energetic α-particles. The third term is related to the α-particle source, where $S(x) = \dot{N}(x)/4\pi v_\alpha^2 \nu$, and where $\dot{N}$ is the α-particle production rate per unit length. The second term gives the quasilinear diffusion of the α-particles, where the diffusion coefficient may be written as

$$D(\epsilon, X) = \begin{cases} (\omega/\nu\epsilon_w^3)(v_{\rm osc}/v_\alpha)^2 & \text{if } \epsilon > \epsilon_w \text{ and } 0 \le X \le A; \\ 0 & \text{otherwise,} \end{cases} \quad (14)$$

where $v_{\rm osc} = 2eE/m\omega$ would be the oscillating velocity of an α-particle in the electrostatic field of strength E, and where A is the normalized radial extent of the lower hybrid power.

By solving this diffusion equation, it is possible to quantify the utility of the density gradient (Fisch and Rax, 1992b). It turns out that the α-particle-energy can be extracted, especially if the scheme is combined with methods of selective ash removal (Mynick, 1992).

The consequences of tapping α-particle energy, and using it for current drive, are very favorable for tokamak operation. Channeling the α-particle free energy to lower hybrid waves traveling in one toroidal direction to drive current is very efficient, because this energy is immediately where it is needed, avoiding power conversion inefficiencies. These conversion inefficiencies would be substantial, because the fusion heat is first converted to low grade heat within the reactor wall, then this heat is converted to electricity, then the electricity powers a microwave generator, and that generator finally injects waves back into the plasma. A second advantage is that, since the free energy is at the expense of expansion, unwanted helium ash is removed from the center. A third advantage is plasma stability; since the α-particles are depleted in energy, less α-particle energy is available to fuel unwanted instabilities.

Note that the effect can be large; if 5% of the fusion power output in a D-T reactor is to be recirculated for lower hybrid current drive, and, if 10–20% of the α-particle power, which is 20% of the fusion power, can be tapped by utilizing the waves, then the circulating power can be held to only 1–3% of the power output, effectively enhancing the efficiency of the current drive by a factor of 1.7–5! The cost of a reactor is very sensitive to the extent of recirculated power, so the savings here could be important. In a D-He^3 reactor, in which the reaction products are protons, lower hybrid waves could similarly tap the energetic proton power, but with even greater effect, since the energetic protons dominate the total fusion power.

Although there is potential for extracting the α-particle free energy using lower hybrid waves, this scheme is, at this stage, highly speculative. Although waves that might accomplish the effect obey the dispersion relation, it remains to be shown, even theoretically, that such waves can be launched efficiently from the plasma periphery, and then can penetrate the α-particle gradient in a reactor.

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