CLM-P844

CLM-P844

The effect of energy loss on electron cyclotron current drive in tokamaks

R. O. Dendy and M. R. O'Brien

UK ATOMIC ENERGY AUTHORITY

Culham Laboratory

CLM-P844

The effect of energy loss on electron cyclotron current drive in tokamaks

R.O. Dendy and M.R. O'Brien

Culham Laboratory, Abingdon, Oxfordshire, OX14 3DB, U.K.

(EURATOM/United Kingdom Atomic Energy Authority Fusion Association)

Abstract

Two independent approaches, based on the inclusion of model loss terms in analytical Fisch-Boozer and numerical Fokker-Planck models, give good quantitative agreement on the reduction of electron cyclotron current drive (ECCD) efficiency by non-Coulomb energy loss. By fitting recent tokamak measurements of ECCD efficiency to these results, the characteristic timescale of non-Coulomb energy loss in velocity space is calculated.

(Submitted for publication in Physical Review Letters)

July, 1988

PACS Nos. 52.50.Gj 52.55.Fa 52.55.Pi

The creation of electric current in plasmas by the application of electromagnetic radiation whose frequency lies in the electron cyclotron range of frequencies (ECRF) was first suggested by Fisch and Boozer.[1] It is attractive, both for maintaining a bulk current in tokamaks, and for localised current drive for the control of specific instabilities.[2-4] Electron cyclotron current drive (ECCD) was first observed on the Culham Superconducting Levitron,[5] and preliminary observations of tokamak ECCD were obtained on TOSCA.[6] Detailed results on tokamak ECCD have only recently become available, from CLEO[7] at Culham, and WT-2[8] and WT-3[9] at Kyoto. In TOSCA and CLEO,[6,7] where the power density of the ECRF waves exceeded the Ohmic power density by an order of magnitude, the waves interacted mainly with thermal electrons. Both tokamaks displayed ECCD when the frequency ω of the waves matched the second harmonic of the electron cyclotron frequency ω_{ce} at the centre of the plasma. However, the measured ECCD efficiency on CLEO for $\omega/2\pi = 60$ GHz was three times smaller than that predicted by a Fokker-Planck code,[10] which nevertheless correctly predicted the absorption efficiency, and the scaling with toroidal field strength of the absorption and of the ECCD. No ECCD was observed on CLEO for the case $\omega/2\pi = 28$ GHz. These apparent differences between the measured and predicted ECCD efficiencies, which may[7] have been due to non-Coulomb energy loss mechanisms, motivate the present study. The Kyoto results, which involve the maintenance of an existing electron beam, rather than its creation by ECRF waves, are not considered further here. The ECCD mechanism[1] rests on the fact that the collisionality of an electron is reduced if its energy is increased by electron cyclotron heating. Such an electron loses parallel momentum at a slower rate than that which applied before its energy was increased. This enhanced persistence of parallel momentum for a selected, heated group of electrons, is responsible for the electric current. Clearly, any physical process which reduces the time for which a heated electron can contribute to the current reduces the efficiency of the current drive mechanism. Our aim is to quantify the extent to which ECCD efficiency is reduced as the magnitude of additional energy and momentum loss is increased. First, we consider a simple analytical model for the ECCD mechanism that extends the approach of Ref.1 by including loss terms. Next, we consider numerical solutions of a Fokker-Planck equation that includes a non-Coulomb energy loss term. The results from the numerical calculations lie close to a curve whose form was derived from our analytical model. This correspondence enables us to relate electron energy loss, through our model, to recent experimental measurements of ECCD.

Let us recall the model for the collisional dynamics of current-carrying electrons that was given in Ref.1. The slowing-down rate due to Coulomb collisions of the current-carrying electrons is $\nu_E = \nu_o/u^3$. Here u is the dimensionless speed v/v_T, $v_T^2 = 2T_e/m_e$, T_e is the electron temperature in energy units, m_e is the electron mass, $\nu_o = n_e e^4 \ell n\Lambda \,/\, 4\pi\epsilon_o^2 m_e^2 v_T^3$, n_e is the electron number density, and $\ell n\Lambda$ is the Coulomb logarithm. The rate of loss of parallel momentum due to Coulomb collisions is $\nu_M = (2+Z_i)\nu_E$, where Z_i is the ion charge. If Coulomb collisions are the only mechanism by which energy and momentum are lost by the current-carrying electrons, we have

$$\frac{du}{dt} = -\nu_E u \quad , \quad \frac{du_\parallel}{dt} = -\nu_M u_\parallel . \tag{1}$$

It follows[1] that the parallel current remaining at time t is $j(t) = j_o[u(t)/u_o]^{2+Z_i}$, where u_o is the initial velocity of the current-carrying electrons. Then the contribution to the current of these electrons, integrated over all times, is the Fisch-Boozer current J_{FB}:

$$J = \int_o^{\infty} j(t)dt = -\int_{u_o}^{o} j(t)\frac{du}{\nu_E u} = \frac{1}{Z_i+5}\frac{j_o u_o^3}{\nu_o} \equiv J_{FB} . \qquad (2)$$

We shall construct our simple, amended model within this framework, and use J_{FB} to normalise the current that we derive.

It is well-known that the loss of energy in tokamak plasmas occurs at rates that exceed those calculated using Coulomb collision theory. However, the nature of the physical processes that lead to enhanced loss is not sufficiently well-understood for their functional dependence on particle velocity to be known. Thus, although it appears that Eq.(1) may require additional terms, the appropriate mathematical form of these terms is not clear. We therefore adopt a heuristic approach, and construct a simple variant of Eq.(1) that includes an enhanced loss rate. This model is intended to be analytically tractable, and to enable us to relate a dimensionless ECCD efficiency J/J_{FB} to a single dimensionless parameter that measures the enhanced loss of energy and momentum. Let us augment the Coulomb energy loss rate ν_E by an amount α_E which is independent of velocity. Then Eq.(1) is modified to

$$\frac{du}{dt} = -(\nu_E + \alpha_E)u \quad , \quad \frac{du_\parallel}{dt} = -(2+Z_i)(\nu_E + \alpha_E)u_\parallel \ . \tag{3}$$

From Eq.(3), following the approach of Eq.(2), we obtain

$$J = \int_0^\infty j(t)dt = -\int_{u_o}^{o} j(t) \frac{du}{(\nu_E + \alpha_E)u} \ . \tag{4}$$

We now denote the initial Coulomb energy loss rate of the current-carrying electrons by $\nu_{Eo} = \nu_o/u_o^3$ and define a dimensionless parameter λ, which is the ratio of the additional energy loss rate to ν_{Eo}: $\lambda = \alpha_E/\nu_{Eo}$. Setting $Z_i = 1$, Eq.(4) yields

$$J = J_{FB} \frac{2}{\lambda^2} [\lambda - \ell n(1+\lambda)] \ . \tag{5}$$

This dependence of J/J_{FB} on λ^{-1} gives the solid curve in Figs.1 and 2. The discrete points in Figs.1 and 2 represent the results of the following independent numerical approach to the problem.

Energy and momentum loss can be represented by writing the steady-state Fokker-Planck equation in the form

$$\frac{\partial f}{\partial t} = \left(\frac{\partial f}{\partial t}\right)_{waves} + \left(\frac{\partial f}{\partial t}\right)_{collisions} + \left(\frac{\partial f}{\partial t}\right)_{loss} = 0 \ . \tag{6}$$

Here $f(v_\perp, v_\parallel)$ is the electron velocity distribution function, with $v_\perp$

and $v_\parallel$ the components of velocity perpendicular and parallel to the magnetic field. The electron cyclotron current $J_{EC} \sim \int v_\parallel f(v_\perp, v_\parallel) d^3\underline{v}$ can be computed once Eq.(6) has been solved. We have used the following non-relativistic operators in the BANDIT code.[10] First, to represent the effect of electron cyclotron heating at the ℓth harmonic resonance, the wave diffusion operator is

$$\left(\frac{\partial f}{\partial t}\right)_{waves} = \frac{1}{v_\perp} \frac{\partial}{\partial v_\perp} \{ v_\perp^{2\ell+1} D_o \exp[-(v_\parallel - v_o)^2/w^2] \frac{\partial f}{\partial v_\perp} \} , \qquad (7)$$

where $D_o = 0.01 \, \nu_o v_T^2$. The mean parallel velocity of cyclotron resonant electrons corresponds to v_o; $w = 0.3 \, v_T$ gives a spread corresponding to a finite-width wavenumber spectrum. Non-linear effects are insignificant at this small value of D_o. Electron trapping is not included in our treatment, which corresponds to ECCD close to the tokamak magnetic axis, as on CLEO.[7] The collision operator in Eq.(6) is an extension of that given in Ref.10, and includes two physical processes: pitch-angle scattering of electrons on massive ions ($Z_{eff} = 1$), and electron-electron collisions which are calculated from a Legendre decomposition of the distribution function, truncated to four terms. The zeroth term is taken to be Maxwellian so that, as usual in Fokker-Planck approaches to ECCD,[11,12] energy is not conserved by our collision operator. This appears to be the lesser of two unavoidable problems. Alternatively, we could use a full collision operator that conserves energy. In this case, for fixed wave power, the energy content of the steady-state distribution would depend on the magnitude of the loss term. This would complicate interpretation of the results more than the alternative that we have chosen, since we would be considering ECCD efficiency for plasmas with different temperatures. Furthermore, we aim at consistency with existing theory in the limit $(\partial f/\partial t)_{loss} \rightarrow 0$. To achieve a steady-state solution of Eq.(6) in this case, the energy input from the waves must be balanced by energy loss, which is introduced through the collision operator.[11,12] This approach appears reasonable provided the non-Coulomb energy loss timescale greatly exceeds the characteristic collisional timescale. Finally, the energy loss operator in Eq.(6) is taken to be

$$\left(\frac{\partial f}{\partial t}\right)_{loss} = \frac{1}{v^2}\frac{\partial}{\partial v}\left(\frac{v^3 f}{2\tau_L}\right) \tag{8}$$

This causes the energy content of the distribution to decay with characteristic loss time τ_L, which we assume to be independent of the velocity. The same operator was employed in Ref. 13. We shall see that the parameter τ_L plays a role analogous to that of α_E^{-1} in the analytical approach.

We have solved Eq.(6) using the code described in Ref.10 with a mesh of 150 speed points $(0 \leqslant v \leqslant 6v_T)$ and 100 pitch angle points $(0 \leqslant \Theta \leqslant \pi)$. In Fig.1, points representing the electron cyclotron current J_{EC} obtained for fundamental resonance (ℓ=1) are plotted against τ_L, with v_o as an additional parameter. Results for second harmonic resonance (ℓ=2) are shown in Fig.2. J_{EC} is normalised with respect to J_∞, the electron cyclotron current obtained from Eq.(6) when $\tau_L \rightarrow \infty$, so that by Eq.(8) $(\partial f/\partial t)_{loss}$ is zero. We note that the present code,[10] like previous studies,[11,12] gives good agreement between J_∞ and J_{FB} defined by Eq.(2). A dimensionless variable is obtained from τ_L by normalising to τ_R, the Coulomb collision timescale of the electrons which carry most of the current. These electrons are identified using a hybrid of the Fisch-Boozer and Fokker-Planck approaches. The Fisch-Boozer[1] ECCD efficiency $J_{FB}/P \sim v_{\parallel} v$, and the power P absorbed is $v^2(\partial f/\partial t)_{waves}$. Then, using Eq.(7), we find the value of $v_{\perp}$ which maximises the integrand in

$$J \sim \int v_{\parallel} (v_{\perp}^2 + v_{\parallel}^2)^{3/2} \left(\frac{\partial f}{\partial t}\right)_{waves} v_{\perp} dv_{\perp} dv_{\parallel} \tag{9}$$

when $v_{\parallel} = v_o$. For example, when $v_o \simeq v_T$, as in the CLEO ECCD experiments,[7] the resulting values of $v_{\perp}$ for $\ell=1$ and $\ell=2$ are $1.35v_T$ and $1.69v_T$ respectively. Thus, both components of velocity play an important role in determining τ_R. Figs.1 and 2 contain results from both analytical (solid curve) and numerical (discrete points) calculations. These have been plotted with the pairs $(J/J_{FB}, J_{EC}/J_{\infty})$ and $(\lambda^{-1}, \tau_L/\tau_R)$ sharing the same units on the same axes.

The close coincidence between the results of the independent analytical and numerical models leads to the following conclusions. First, we may use Eq.(5), with the identification of variables introduced in Figs. 1 and 2, to fit the results of the Fokker-Planck calculations. Since $J_{FB} \simeq J_{\infty}$, the results of Eq.(6) are described by

$$J_{EC} = J_{FB}\, 2(\tau_L/\tau_R)^2 \left[\tau_R/\tau_L - \ell n(1 + \tau_R/\tau_L)\right] \tag{10}$$

We note that τ_R is a function of v_o, and that the choice of normalisation τ_L/τ_R is central to the coincidence of numerical results for different values of v_o. Second, the analytical description is based on the dynamics of a single group of current-carrying electrons, whereas the Fokker-Planck description is global in velocity space. The coincidence of the results indicates the power of the Fisch-Boozer[1] approach which, extended here to include energy loss, remains in agreement with the collective Fokker-Planck treatment. Third, the analytical and numerical results agree most closely for the highest value of v_o considered. This is the case for which the resonant electrons are least collisional, so that their coupling to the bulk plasma is weakest and their treatment as a separate entity, as in the analytical approach, is most appropriate. Fourth, we may use Eq.(10) to obtain information about the value of τ_L/τ_R from experimental measurements of ECCD efficiency. Recall that, in Eq.(8), τ_L is a measure of non-Coulomb energy loss and τ_R is the Coulomb collision time of the current-carrying electrons. In the CLEO tokamak second harmonic ECCD experiments,[7] the parallel velocity v_o at which peak current drive occurs is close to v_T. Then for source frequencies of 60 GHz ($T_e \simeq 1.25$ keV, $n_e \simeq 6\times10^{18}m^{-3}$) and 28 GHz($T_e \simeq 1.25$ keV, $n_e \simeq 3\times10^{18}m^{-3}$), the procedure indicated at Eq.(9) gives $\tau_R = 0.85$ ms and 1.7 ms respectively. At 60 GHz, it was found that the ECCD efficiency was three times smaller than that predicted from J_∞ by the Fokker-Planck code. From Fig. 2, this suggests $\tau_L/\tau_R \simeq 0.3$, and hence $\tau_L \simeq 0.25$ ms. The global energy confinement time τ_E for this discharge, obtained from energy balance, is 1ms. Now τ_E is determined by the transport properties of the entire plasma, whereas in our model τ_L relates primarily to the part of the plasma in which ECCD is concentrated. Thus, it is reasonable that τ_L is found to be less than τ_E, but of similar magnitude. The absence of detectable ECCD at 28 GHz suggests, by Fig.2, that in this case $\tau_L/\tau_R \leqslant 0.1$. Then $\tau_L \lesssim 0.17$ ms, which is significantly smaller than the value for 60 GHz. The decline in τ_L is matched by a decline in τ_E, which is 0.5ms for the 28 GHz case.

The Fisch-Boozer analytical approach to collisionality, and its extension in this paper, are not restricted to any particular wave damping scheme. The model of Ref.1 has been used successfully, without a loss term, to explain the efficiencies observed in lower hybrid current drive (LHCD) experiments which involve electrons with $v_\parallel >> v_T$.[14] In these experiments, the global confinement time τ_E is substantially less than τ_R.[14,15] The fact that a model without a loss term successfully predicts the LHCD efficiently suggests that $\tau_L/\tau_R \gtrsim 1$. It follows that $\tau_L >> \tau_E$, so that the electrons that carry the current lose energy much more slowly than the bulk plasma, in agreement with previous experimental and theoretical results.[16-18]

Further information on the characteristic timescales of ECCD can be obtained from time-dependent Fokker-Planck calculations. In Fig.3, the electric current is shown as a function of time t, again normalised to τ_R, for two cases with $v_o = 1.5\ v_T$. The upper curve has τ_L infinite, corresponding to the absence of an explicit loss term, and is calculated for a distribution that is Maxwellian at $t=0$. The lower curve has $\tau_L/\tau_R = 0.75$, and is calculated for a distribution which at $t=0$ is a solution of Eq.(6) with $(\partial f/\partial t)_{waves}$ set to zero. We note that the curves are initially very close, but diverge for $t \gtrsim 0.2\tau_R$, and that the steady-state current in the lower curve is established by $t \simeq \tau_R \simeq \tau_L$. Physically, the first point reflects the fact that for initial times

$t \ll \tau_L$, the loss term has no effect on ECCD, which is determined by the balance of wave diffusion with collisions alone. The second point shows that once $t \simeq \tau_L$, equilibrium between wave diffusion and loss terms has been reached, and the system evolves no further.

In conclusion, we have studied the effect of energy loss on ECCD using two independent approaches. The analytical model of Eq.(3), which is a heuristic extension of Fisch-Boozer[1] theory, leads to Eq.(5). This is shown as the solid curve in Figs.1 and 2. Numerical solution of the Fokker-Planck equation, Eq.(6), with a simplistic loss term defined by Eq.(8), produces the discrete points in Figs. 1 and 2. The good agreement between the two approaches suggests that the results are not strongly model-dependent, and shows that the results of the Fokker-Planck calculations follow the simple formula of Eq.(10). Thus, Eq.(3) appears to be a good simple representation of the physics contained in the global Fokker-Planck approach. Our results enable us to use experimental measurements of ECCD efficiency to obtain characteristic timescales of non-Coulomb energy loss which appear consistent with experiment. The results quantify the importance, for efficient ECCD, of selecting resonant electrons that possess optimal energy retention characteristics. They indicate that where ECCD efficiencies below the Fisch-Boozer value have been measured, this may be a consequence of relatively poor confinement. Conversely, the projected ECCD efficiency in large tokamaks with better confinement remains good.

References

[1]N.J, Fisch and A.H. Boozer, Phys. Rev. Lett. 45, 720 (1980).

[2]A.H. Reiman, Phys. Fluids 26. 1338 (1983).

[3]R.O. Dendy, Plasma Phys. 27, 1243 (1985).

[4]D.W. Ignat and P.H. Rutherford, Proceedings of the Course and Workshop on Applications of rf waves to Tokamak Plasmas (Monotypia Franchi, Perugia, 1985) Vol. 2, p.525.

[5]D.F.H. Start et al., Phys. Rev. Lett. 48 (1982) 624.

[6]M.W. Alcock et al., Plasma Physics and Controlled Nuclear Fusion Research (Proc. 9th Int. Conf. Baltimore, 1982) (IAEA, Vienna, 1983), Vol.2, p.51

[7]B. Lloyd et al., Culham Laboratory Preprint CLM-P830, accepted for publication in Nucl. Fusion.

[8]A. Ando et al., Phys. Rev. Lett. 56, 2180 (1986).

[9]H. Tanaka et al., Phys. Rev. Lett. 60, 1033 (1988).

[10]M. O'Brien et al., Nucl. Fusion 26, 1625 (1986).

[11]C.F.F. Karney and N.J. Fisch, Nucl Fusion 21, 1549 (1981).

[12]J.G. Cordey et al., Plasma Phys. 24, 73 (1982).

[13]V.V. Alikaev and V.L. Vdovin, Sov. J. Plasma Phys. 9, 538 (1983).

[14]C.F.F. Karney et al., Phys. Rev. A32, 2554 (1985).

[15]F.C. Jobes et al., Phys. Rev. Lett. 55, 1295 (1985).

[16]G. Fussmann et al., Phys. Rev. Lett. 47, 1004 (1981).

[17]T.K. Chu et al., Nucl. Fusion 26, 1319 (1986).

[18]H.E. Mynick and J.D. Strachan, Phys. Fluids 24, 695 (1981).

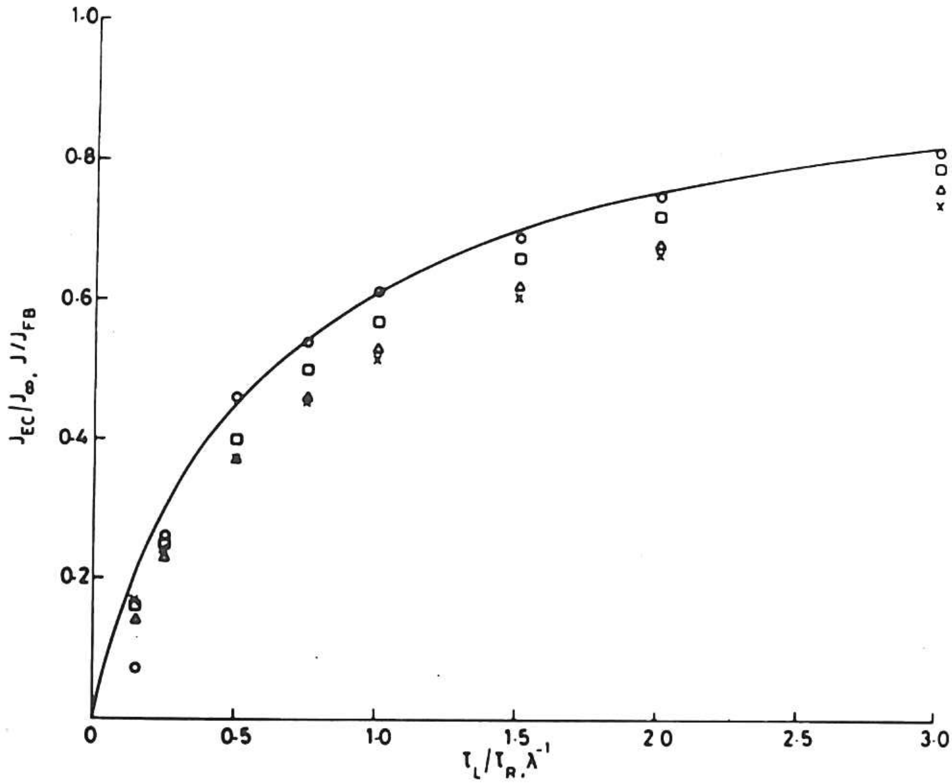

Fig. 1 Current drive efficiency at the fundamental electron cyclotron resonance. Solid curve: analytical model. Discrete points: Fokker-Planck calculations, with resonant parallel velocity $v_0 = v_T$(x), $1.5v_T(\Delta)$, $2v_T(\Box)$, $2.5v_T(0)$.

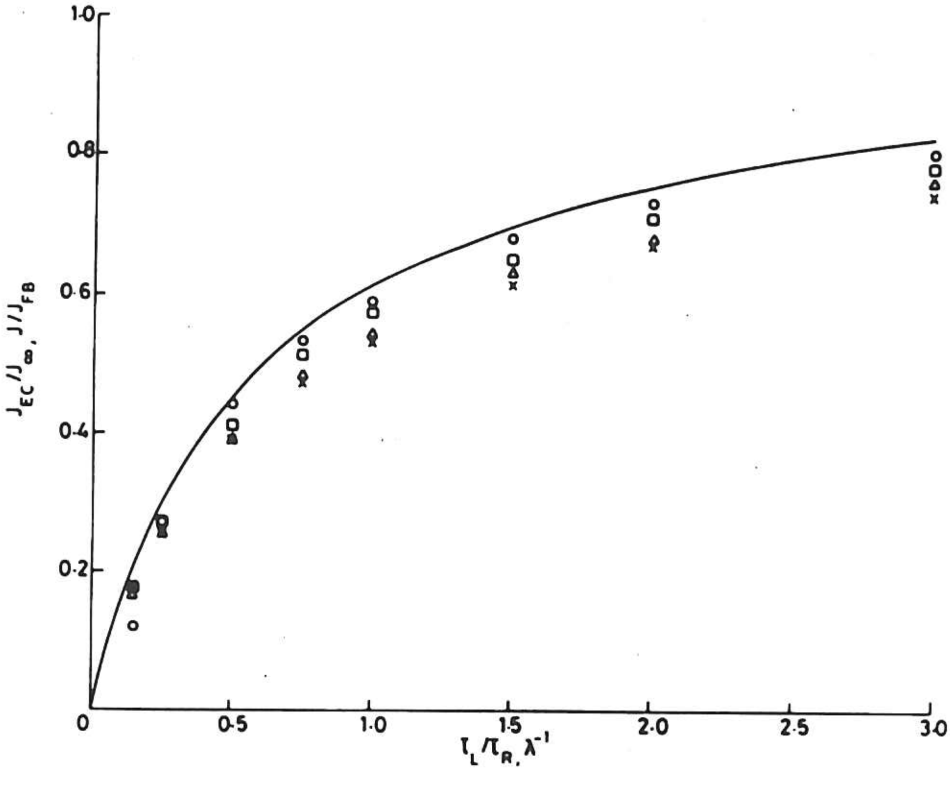

Fig. 2 Current drive efficiency at the second harmonic electron cyclotron resonance. Solid curve: analytical model. Discrete points: Fokker-Planck calculations, with resonant parallel velocity $v_0 = v_T$(x), $1.5v_T(\Delta)$, $2v_T(\Box)$, $2.5v_T(0)$.

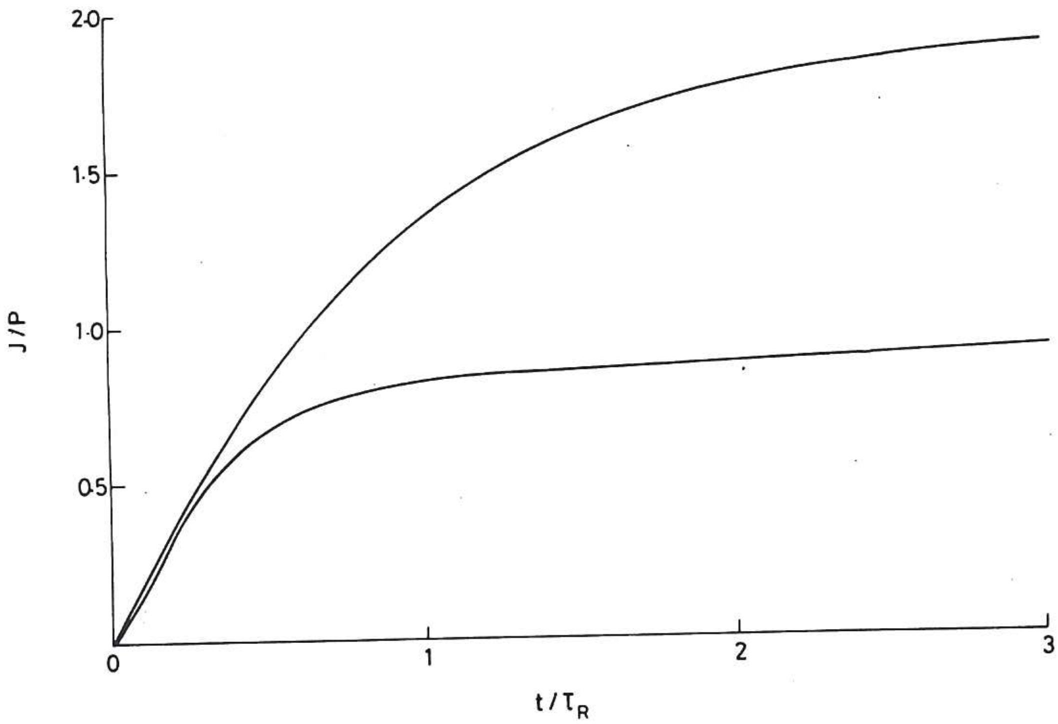

Fig. 3 Current rise, using a time-dependent Fokker-Planck code. Upper curve: τ_L infinite. Lower curve, $\tau_L = 0.75\,\tau_R$. In both cases, $v_0 = 1.5\,v_T$.