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The origin of fluctuations and cross-field transport in idealized magnetic confinement systems

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The study of plasma fluctuations and confinement in idealized systems such as octupoles and levitrons has contributed to the understanding of cross-field transport processes. The linear theory of plasma instabilities that cause fluctuations is well developed and can predict growth rates γ and wavelengths λ_x around lines of force. However, the theoretical prediction of cross-field transport coefficient $D_\perp$ is restricted to quasilinear estimates of upper bounds (for example, $D = \frac{1}{2}\gamma\lambda_x^2$) because of the complexity of the full nonlinear calculation. Such quasilinear estimates usually far exceed the measured values and are of limited worth.

A general view of the results from octupole and levitron experiments shows that under collisional conditions ($\lambda_{ei}/L < 0$) the diffusion coefficient, D, scales in the same way as classical collisional diffusion ($D \propto n/T_e^{\frac{1}{2}}B^2$). Agreement is closely approached in many cases, sometimes even in the presence of fluctuations. Under collisionless conditions ($\lambda_{ei}/L > 0$), Bohm diffusion scaling ($D \propto T_e/B$) is found in the few cases where the scaling law has been determined. This behaviour is consistent with the general scaling laws of Connor & Taylor (1977) but is not understood in detail. In addition there is evidence, both experimental and theoretical, that long-wavelength low-frequency electric fields (convection cells) can be generated nonlinearly from high-frequency fluctuations and can contribute to cross-field transport.

Introduction

Confinement of hot plasma by a magnetic field is essential for steady-state fusion. However, in early confinement experiments cross-field transport far exceeded the current predictions for plasmas in local thermal equilibrium (see Hinnov & Bishop 1966, for example). The observed confinement times were compatible with 'Bohm diffusion' (Bohm 1949) which has a diffusion coefficient defined as $D_B = T_e/16B$, where T_e is the electron temperature and B is the magnetic field strength. The reason for this poor confinement was not understood (for most plasmas of interest $D_B \gg D_{class}$, the rate due to collisions only). Although these plasmas were macroscopically stable they were not always microscopically stable, the distinction being between coherent motion of the bulk of the plasma and coherent motion restricted to small regions. To understand micro-instabilities and the mechanisms causing anomalous transport, several experiments were built which were designed to have 'ideal' plasma confinement properties. Some of the results from these experiments have been reviewed by Yoshikawa (1973), and Navratil & Post (1979). In this paper we discuss briefly micro-instabilities and cross-field transport and review results from these 'ideal' experiments including more recent work.

Ideal confinement systems

The 'ideal' system has toroidal geometry with no field lines entering or leaving the confinement region so that losses cannot occur along the field. In addition, by making the field axisymmetric, single particle confinement is assured through the conservation of canonical angular momentum. To avoid unnecessary current flow in the plasma, or plasma motion, the fields are generally steady state, that is, $\nabla \wedge \boldsymbol{E} = -\partial \boldsymbol{B}/\partial t = 0$. The effects of field errors and atomic processes such as recombination and excitation can be reduced to a low level in most experiments and are not considered further here.

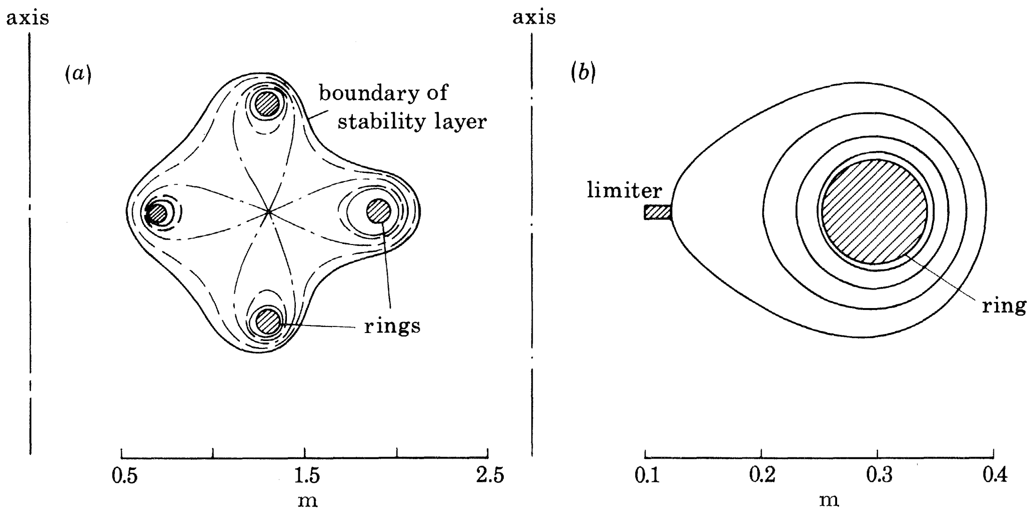

Figure 1. Schematic diagrams of (*a*) octupole and (*b*) levitron configurations. The shaded ring areas represent current-carrying hoops that are concentric about the axis. External coils are not shown.

Two ways of achieving macroscopic stability have been used. The first relies on the stabilizing properties of an average magnetic well (Rosenbluth & Longmire 1957) in the m.h.d. limit. The appropriate average is defined by the function $U = \oint \mathrm{d}l/B$ for a closed-line system and the requirement for stability for a low pressure plasma is $\nabla p \cdot \nabla U > 0$, where p is the plasma pressure. The toroidal octupole, which has four current-carrying concentric rings within the plasma (figure 1*a*), provides a large magnetic-well region. The second method uses the stabilizing properties of shear (Johnson *et al.* 1958). The combination of the poloidal field of a single current-carrying ring and a toroidal field provides a system of closed nested flux surfaces with strong shear between the field lines in each surface. This configuration is illustrated in figure 1*b*. One other system to which we refer is the linear quadrupole (Phillips *et al.* 1978) with two straight current-carrying conductors within the plasma. This configuration has a shallow magnetic well but does not have the absolute confinement properties of a toroidal system.

Classical diffusion in ideal systems

The minimum rate for cross-field transport of energy and particles in a plasma is simply due to Coulomb collisions. Many authors have evaluated the collisional transport coefficients, and a review of the subject for plasmas in toroidal geometry has been given by Hinton & Hazeltine (1976). In slab geometry the particle flux for singly charged ions is

$$\Gamma = a_e^2 \nu_{ei}[-\nabla n(1+T_i/T_e) - n\nabla T_i/T_e + 0.5n\nabla T_e/T_e], \tag{1}$$

where a_e is the electron cyclotron orbit radius, ν_{ei} is the electron–ion collision frequency, n is the electron density and T_e and T_i are the electron and ion temperatures. By analogy with Fick's law, $\Gamma = -D\nabla n$, the classical diffusion coefficient is

$$D_{class} = a_e^2 \nu_{ei} \propto n/T_e^{\frac{1}{2}} B^2. \tag{2}$$

In toroidal geometry, additional transport occurs owing to the non-uniformity of the toroidal field: in this case the transport is referred to as neo-classical diffusion. As well as these collisional effects, the thermal fluctuation spectrum can give a significant additional cross-field flux known as 'convective diffusion' (Taylor & McNamara 1971; Okuda & Dawson 1973). This can dominate diffusion for low density plasmas in strong magnetic fields which have a two-dimensional character. All of the classical processes give loss rates that are acceptably small for steady-state fusion: experimentally, higher rates are observed and the way in which these can be produced by micro-instabilities is considered next.

Theory of fluctuations and transport

(a) *General considerations*

The linear theory of plasma instabilities is well developed but there is no complete theory relating particle and energy transport to the existence of any particular instability. Fluctuations growing above the thermal level require a source of free energy (Fowler 1965). The principal source is the plasma confinement itself and it is this free energy that is degraded during collisional diffusion. The growth of fluctuations coupled to this source therefore leads to an outflow of plasma, and we shall determine this from quasilinear theory in § 2 (*b*). However, the theory is incomplete. It provides no ultimate sink for the energy, consequently a continuous outflow requires continuous fluctuation growth. In contrast, experiment shows that statistically stationary fluctuations may lead to plasma loss. Presumably the energy input to fluctuations is balanced by energy loss to a suitable sink. In some geometries fluctuations may propagate out of the system, otherwise the energy loss must be due to nonlinear effects transferring energy to other modes, which in turn are damped or propagate out of the system. These effects can reduce the growth rate, which, in the limit, may become zero to give a state of nonlinear marginal stability where there is no energy flow and no plasma loss except for residual collisional effects. There is no *a priori* connection between the existence of fluctuations and continuous plasma loss, and a complete transport theory requires a detailed knowledge of the pathways and rates of energy and particle flow. This knowledge is not available for any confined plasma system, even though the basic nonlinear processes are known and their rates have been measured and compared with theory (Franklin 1977). We are therefore forced to rely on order-of-magnitude estimates of transport: some estimates of upper bounds are discussed in § 2 (*c*).

(b) *Linear and quasilinear theory*

Many instabilities are described in the literature (see, for example, the review by Tang (1978)) but here we refer only to the 'drift' and 'flute' instabilities. We shall treat the drift wave as an example and show how its instability leads to plasma loss. We work in 'slab geometry' and use

Cartesian coordinates with z-axis parallel to the magnetic field $\boldsymbol{B}$ and x-axis parallel to the gradient of mean density n_0. The equations appropriate to low frequencies ($\omega < \omega_{ei}$) and the electrostatic limit are given below.

Ion continuity: $$\frac{\partial \tilde{n}_i}{\partial t} + \tilde{v}_E \frac{dn_0}{dx} = 0; \tag{3}$$

plasma drift speed: $$\tilde{v}_E = -\frac{\nabla\tilde{\phi} \wedge \boldsymbol{B}}{B^2} \equiv -\frac{1}{B}\frac{\partial \tilde{\phi}}{\partial y}; \tag{4}$$

electron equilibrium parallel to $\boldsymbol{B}$: $$\tilde{n}_e = n_0 e\tilde{\phi}/T_e; \tag{5}$$

and quasi neutrality: $$\tilde{n}_e = \tilde{n}_i = \tilde{n}, \quad \text{say.} \tag{6}$$

A tilde marks the fluctuating quantities, which, after being Laplace-transformed in time and Fourier-transformed in space, become the complex amplitudes $\tilde{n}_k$, $\tilde{\phi}_k$ of a mode with wavevector $\boldsymbol{k}$ and frequency ω. From these equations we obtain

$$\omega \tilde{n}_k = k_y \tilde{\phi}_k B^{-1} dn_0/dx, \tag{7}$$

$$\tilde{n}_k = n_0 e\tilde{\phi}_k/T_e, \tag{8}$$

and eliminating $\tilde{n}_k$ and $\tilde{\phi}_k$ gives the dispersion relation

$$\omega = \omega_{*e} \equiv \frac{k_y T_e}{eBn_0}\frac{dn_0}{dx}. \tag{9}$$

Equation (9) shows that ω is real, which indicates neutral stability.

However, the drift wave can be driven unstable by dissipative mechanisms that cause a phase shift between $\tilde{n}$ and $\tilde{\phi}$. For instance, electron–ion collisions drive the collisional drift instability and inverse electron Landau-damping drives the collisionless drift instability. The transfer of electrons from trapped to passing orbits by means of collisions leads to the dissipative trapped-electron instability. We define the collisionality of a plasma by the ratio λ_{ei}/L, where λ_{ei} is the mean free path for electron–ion collisions and L is a typical scale length of the magnetic field taken parallel to B. The collisionality is proportional to T_e^2/n and collisional effects are expected to be important if $\lambda_{ei}/L < 1$. Trapped-particle effects are expected when $\omega_b/\nu_{eff} > 1$, where ω_b is the particle bounce frequency and ν_{eff} the effective collision frequency for particle de-trapping.

The flute instability, defined by $k_z = 0$, is made unstable by adverse field-line curvature rather than dissipation. It is observed in multipoles, while in a levitron high shear and low plasma pressure convert it into the electrostatic 'resistive-g' instability.

The outward flux of plasma due to fluctuations is given by the covariance of density and radial velocity:

$$\Gamma_x = \langle \tilde{n}\tilde{v}_E \rangle \equiv -B^{-1}\langle \tilde{n}\,\partial\tilde{\phi}/\partial y \rangle.$$

Fourier analysis shows that each mode contributes to Γ_x only if $\tilde{n}_k$ and $\tilde{\phi}_k$ are out of phase. From (7) this implies that ω is complex: if $\omega_r = \omega - i\gamma$, the phase shift is $\alpha = \arctan \gamma/\omega_r$ and for $\gamma/\omega_r \ll 1$, we have

$$\Gamma_x = \frac{1}{2}\frac{k_y \tilde{n}_k \tilde{\phi}_k}{B}\sin\alpha \approx \frac{1}{2}\frac{k_y \tilde{n}_k \tilde{\phi}_k}{B}\frac{\gamma}{\omega_r}, \tag{10}$$

which is to be summed over all k to give the total flux. Equation (10) illustrates the close connection between growth and plasma loss. Stringer (1977) derives a similar result more rigorously and shows clearly that for self-consistency $\tilde{n}_k$ and $\tilde{\phi}_k$ must be increasing with time. However, experimentally, phase shifts and plasma loss often co-exist with statistically stationary amplitudes of $\tilde{n}_k$ and $\tilde{\phi}_k$, and this requires an energy sink as discussed previously.

(c) *Upper bounds on diffusion rates*

As the amplitudes increase, a succession of possible nonlinear effects must be considered, each of which may be capable of limiting the growth. If the limit is due to mode–mode coupling or related effects, the flux will be given by (10), but no general statement about the amplitudes can be made. On the other hand, if the limit is due to nonlinear reduction of the growth rate it is much easier to write down plausible expressions for the limiting amplitudes, but for consistency the flux must now fall below that given by (10). Nevertheless, an upper bound can be constructed by using the linear growth rate and the limiting amplitudes.

We write (10) by using (9) as

$$\Gamma_x = \tfrac{1}{2}\gamma\left(\frac{\tilde{\phi}_k^2}{T_e^2}r_n^2\right)\frac{\mathrm{d}n_0}{\mathrm{d}x} \tag{11}$$

(Stringer 1977), where $r_n = n_0(\mathrm{d}n_0/\mathrm{d}x)^{-1}$ is the density gradient scale length. If it is assumed that growth ceases when the instantaneous density gradient changes sign, then we may write

$$\tilde{n}_k/n_0 \lesssim \lambda_x/r_n, \tag{12}$$

where λ_x is the radial scale of the mode amplitude, and (8), (11) and (12) give

$$\Gamma_x \leqslant \tfrac{1}{2}\gamma\lambda_x^2\,\mathrm{d}n_0/\mathrm{d}x. \tag{13}$$

By Fick's law, this gives a diffusion coefficient $D_{\mathrm{q.l.}} \approx \frac{1}{2}\gamma\lambda_x^2$ (Kadomtsev 1965). A similar result was obtained by Dupree (1968) by considering wave-particle scattering effects.

For the upper bound of (12) to be physically reasonable it must not exceed the limit set by the available free energy. This was calculated by Fowler (1965). Strictly, since thermodynamic equilibrium implies an unconfined plasma of uniform density, the entire thermal energy of a confined plasma is free energy in this sense. To obtain more significant limits we must restrict the possible motions. For example, fluctuations of short range $\lambda_x \ll r_n$ can convert only the free energy within that range. This gives a bound on fluctuation energy of

$$W \lesssim \tfrac{1}{6}n_0\,T_e\lambda_x^2/r_n^2.$$

For drift waves, the corresponding density fluctuation level is

$$\tilde{n}/n_0 \lesssim (k_y a_i)^{-1}\lambda_x/r_n,$$

where a_i is the ion Larmor radius. Thus the condition given by (12) is physically acceptable in the usual case $k_y a_i < 1$, and the flux given by (13) is usually used for comparison with experiment.

Experimental results

(a) *Classical diffusion*

In both octupoles and levitrons classical diffusion rates have been achieved. In the experiments of Ohkawa *et al.* (1971) in the GA-octupole, classical scaling of the particle confinement time, τ_p, with both n and B was observed. The absolute value was two-thirds of the classical value but this discrepancy was within the experimental error. A similar result was obtained by Navratil & Post

(1977) in the Wisconsin octupole. In both these experiments only poloidal field was used. With toroidal field added, particle loss rates were consistent with neoclassical scaling (Ohkawa *et al.* 1972). Finally, Tamano *et al.* (1973) made experiments at the low plasma densities and high magnetic field strengths required for convective diffusion to be dominant. The expected scaling with T_e, n, B and ion mass was confirmed but the absolute value of the cross-field flux was 3.5 times that expected. Since the theory was derived for a slab geometry this discrepancy is possibly not significant.

TABLE 1. PRINCIPAL RESULTS ON INSTABILITIES IN OCTUPOLES

laboratory	collisionality régime	wave character	instability	transport	reference
general atomic, San Diego	collisionless, $\omega_{\mathrm{b.e.}} > \nu_{\mathrm{eff}}$	$\omega_{\mathrm{R}} = 0.3\omega_{*\mathrm{e}}$ $v_{\mathrm{phase}} \| v_{*\mathrm{e}}$ $\omega_{\mathrm{R}} \approx \nu_{\mathrm{eff}}$	dissipative trapped electron	Bohm diffusion scaling with T_e and B_θ	Prater *et al.* (1978)
University of Wisconsin	collisionless $\omega_{\mathrm{b.i.}} > \nu_{\mathrm{eff}}$	(i) $v_{\mathrm{phase}} \| v_{*\mathrm{e}}$ (ii) $v_{\mathrm{phase}} \| v_{*\mathrm{i}}$	dissipative trapped ion collisionless trapped ion	enhanced loss but scaling not reported	Drake (1977)
University of Wisconsin	collisional	$v_{\mathrm{phase}} \| v_{*\mathrm{e}}$ $k_\perp a_i \approx 0.3$	resistive drift	classical diffusion	Navratil & Post (1977)

(b) *Fluctuations and enhanced losses in octupoles*

Fluctuations and enhanced transport have also been observed in octupole experiments, and in several cases the micro-instability thought to be responsible has been identified. The main results and parameters are given in table 1. In the experiments of Prater *et al.* (1978) the scaling of the diffusion coefficient was $D \propto T_e/B_\theta$, similar to Bohm scaling but with the poloidal field alone playing a part. No satisfactory explanation for this behaviour was given. In the experiment of Navratil & Post (1977) the loss rate was close to the classical value even though large-amplitude fluctuations were present ($\Delta n/n \lesssim 0.2$). This result would be consistent with (10) if the phase angle, α, between $\tilde{n}$ and $\tilde{\phi}$ were close to zero. However, the measured value of sin α was about 0.1 giving a value for Γ_x some one thousand times larger than that observed. This phase measurement was not corrected for the effect of temperature fluctuations: such fluctuations can affect the observed $\tilde{\phi}$ through sheath potential fluctuations (Ellis & Motley 1974), which illustrates one of the major difficulties of deducing the flux from probe measurements.

(c) *Fluctuations and enhanced losses in levitrons*

In the F.M.1 experiment (a levitron) particle confinement time was found to scale as $\tau_p \propto T_e^{\frac{3}{4}}/nB^2$ for $T_e \lesssim 2$ eV, as shown in figure 2*a* (Sinnis *et al.* 1972). This is a similar scaling to that predicted for classical transport, but the absolute value of the loss rate was about five times too large. The extra loss rate was thought to be due to resistive-drift waves (Okabayashi & Arunasalam 1977). For $T_e \gtrsim 2$ eV, τ_p decreases as T_e^{-1} which again suggests Bohm scaling. By applying localized heating Ejima & Okabayashi (1975) measured the conductivity $K_\perp$ and diffusion coefficient $D_\perp$ as a function of T_e in this higher temperature régime. The results are shown in figure 2*b* where, for comparison, $K_\perp$ has been defined so that it has the same dimensions as $D_\perp$. The fluctuation amplitude in these measurements was larger for $\eta_e > 0$ than for $\eta_e < 0$, where $\eta_e = \nabla \ln T_e/\nabla \ln n$, which suggests the trapped-electron instability. The quasilinear estimate of the ratio $K_\perp/D_\perp$ was calculated for this instability and found to be half the observed value but estimates for the absolute values of $K_\perp$ and $D_\perp$ were not given.

In the Culham levitron it is found that, near the edges of the plasma, waves occur with frequencies above 30 kHz and with $k_y a_i \approx 0.3$, close to the value of 0.5 predicted for drift waves of maximal growth. Further in, there are low frequency coherent modes with $k_y a_i < 0.1$ (Ashby 1977). Profiles of n, T_e and amplitude of the low frequency modes are shown in figure 3*a* for medium

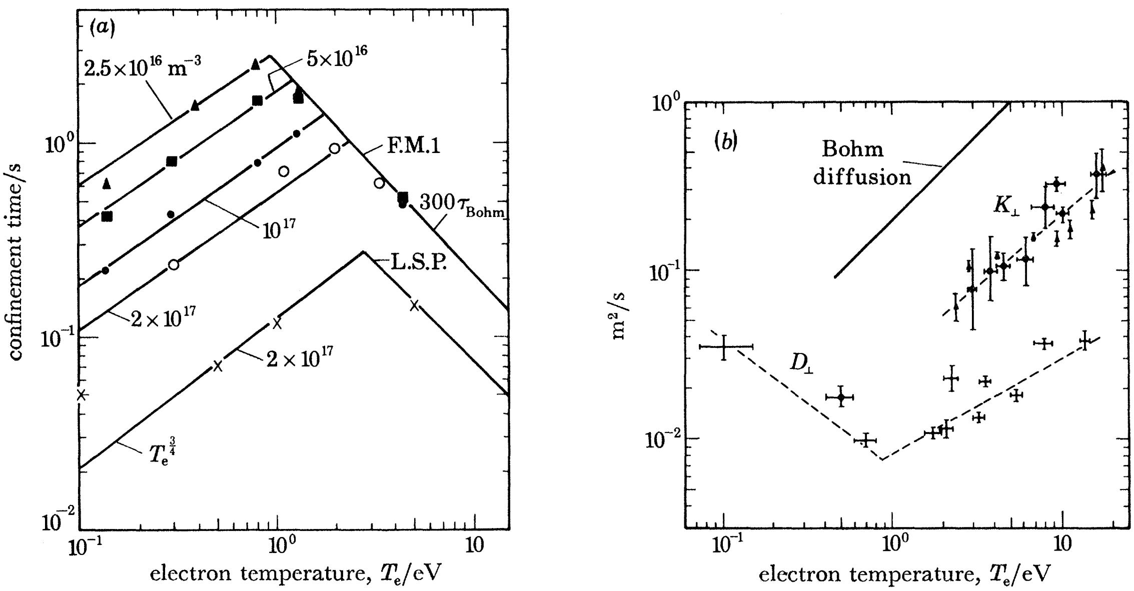

FIGURE 2. (*a*) Variation of confinement time with electron temperature in the F.M.1 and L.S.P. experiments (Sinnis *et al.* 1972). (*b*) Variation of thermal conductivity and diffusion coefficients ($K_\perp$ and $D_\perp$) with electron temperature in the F.M.1 experiment (Ejima & Okabayashi 1975).

shear and collisional conditions (Ainsworth *et al.* 1978). The frequency spectrum of the modes is shown in figure 3*b* as a function of the position on the profile for the same data, the darker regions representing higher amplitude on a logarithmic scale. The modes are centred on the integer rational surfaces whose positions are indicated along the abscissa. There is a well defined frequency observed in the laboratory frame in each case (Cordey *et al.* 1979). The toroidal wavelength λ_ϕ for the main component is equal to the pitch length (L_p), i.e. the poloidal mode number $m = 1$, and $k_\parallel = 0$ near the centre of each mode. The modes appear only where the field curvature is destabilizing, i.e. where $\nabla B \cdot \nabla p > 0$. In addition, for low values of B_ϕ/B_θ, where B_θ is the toroidal field, fluctuations have large amplitude with $\tilde{n}/n$ of the order of 0.1, but as B_ϕ/B_ϕ is raised, thereby increasing the shear, the amplitude falls and the modes are stabilized. These properties fit those expected for the resistive-g/ion-temperature gradient instability treated in the electrostatic limit by Cordey *et al.* (1980).

The ratio of the local diffusion coefficient to D_{class} is shown in figure 4*a* as a function of the average shear length $\bar{L}_s$ (note that shear strength is inversely related to $\bar{L}_s$). The diffusion approaches the classical value at sufficiently strong shear. The values of $D_{\text{q.l.}}$ obtained from the theory are compared with the observed D as a function of radius in figure 4*b*. They are much larger than required to explain the experimental losses over most of the profile. Scaling of $D_{\text{q.l.}}$ from the theory is predicted to be similar to the classical scaling, but with an added shear dependence. This is consistent with the measured scaling, and the instability could provide an alternative to the resistive-drift instability as an explanation for the pseudo-classical scaling seen in F.M.1 (see figure 2*a*).

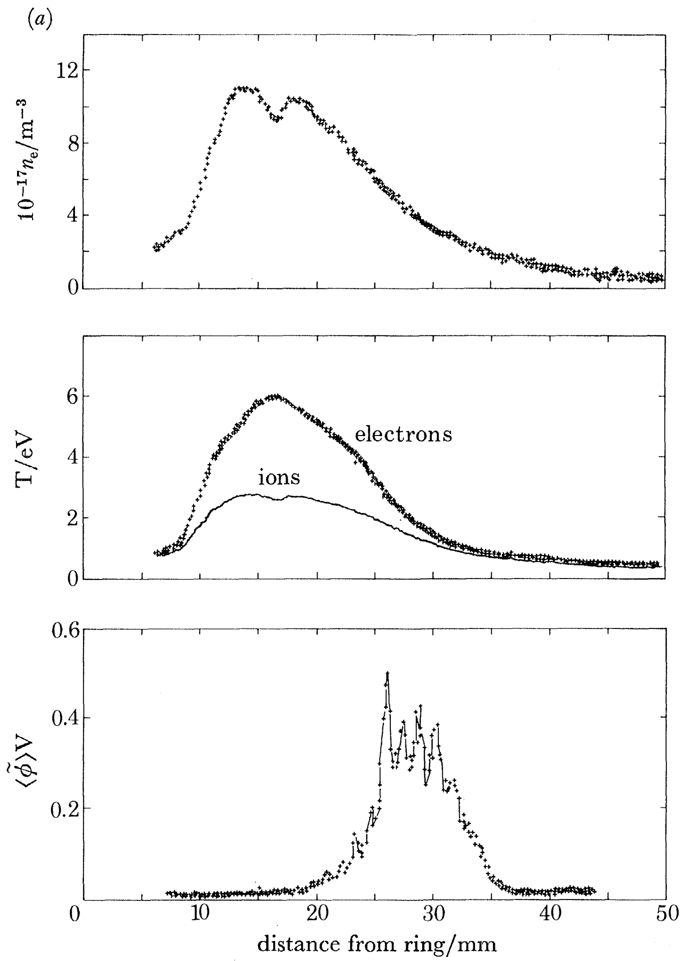

FIGURE 3*a*. Electron density, temperature and fluctuation amplitude profiles in the Culham levitron. (Ainsworth *et al.* 1978.)

It is interesting to note that fluctuations are also observed which have very low frequencies (v.l.f.) of the order of a few hertz and $\lambda_\phi > L_p$: these might be convection-cell fields. Ashby (1979) showed by using a triple correlation technique (Kim & Powers 1979) that such v.l.f. fluctuations can be due to nonlinear coupling between drift waves. Cheng & Okuda (1978) have shown that in a computational model also, convection-cell fields can be generated by nonlinear coupling between degenerate drift-modes and that these can give a dominant contribution to the cross-field transport.

(*d*) *Experiments in a linear quadrupole*

A direct link between fluctuations and particle diffusion has been made in a linear quadrupole experiment described by Phillips *et al.* (1978). The main contribution to the cross-field diffusion was due to fluctuations with a frequency near 1 kHz (Williamson & Rusbridge 1979) which are attributed to the 'shallow-well' flute instability (Hastie & Taylor 1971). The particle flux

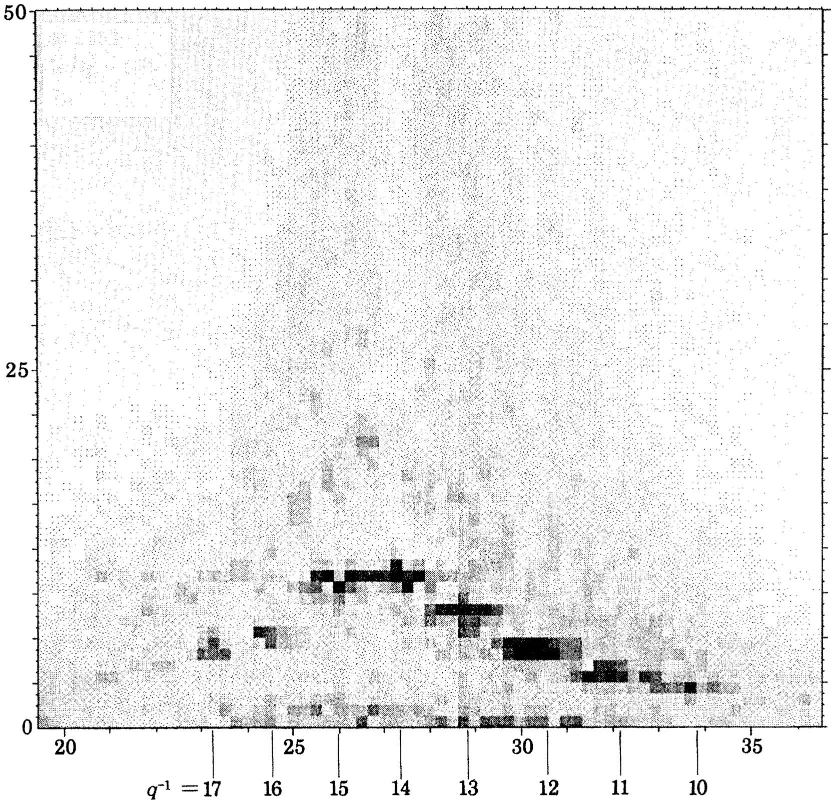

FIGURE 3*b*. Fourier spectra of the fluctuations shown in (*a*) across the profile. Darker shades refer to larger amplitude with logarithmic grading between shades. (Figure 3*b* was supplied by W. H. W. Fletcher and E. M. Jones.)

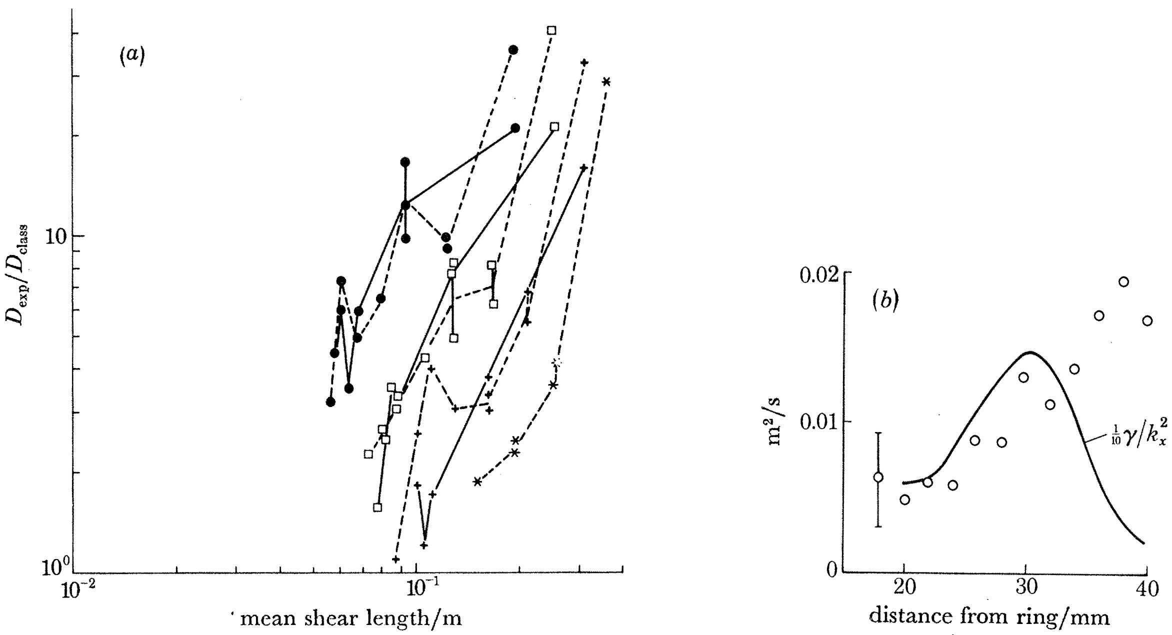

FIGURE 4. (*a*) Ratio of experimental to classical diffusion coefficients, D_{exp}/D_{class}, as a function of average shear length for ring currents of 180 kA (——) and 110 kA (– – –), and the following distances from the ring: ★, 20 mm; +, 25 mm, □, 30 mm; ●, 35 mm. (*b*) Profile of D_{exp} compared with the corresponding predictions of γ/k_x^2 from the resistive-*g* instability theory of Cordey *et al.* (1980).

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deduced from measurement of $\tilde{n}$, $\tilde{\phi}$, k_y and phase shift was within a factor of three of that obtained from the rate of plasma decay. The upper bound for the diffusion coefficient given by (13) was evaluated by using γ from the theory of the instability and λ_x from the experimental observations; it exceeded the diffusion coefficient calculated from the particle flux by a factor of fifty.

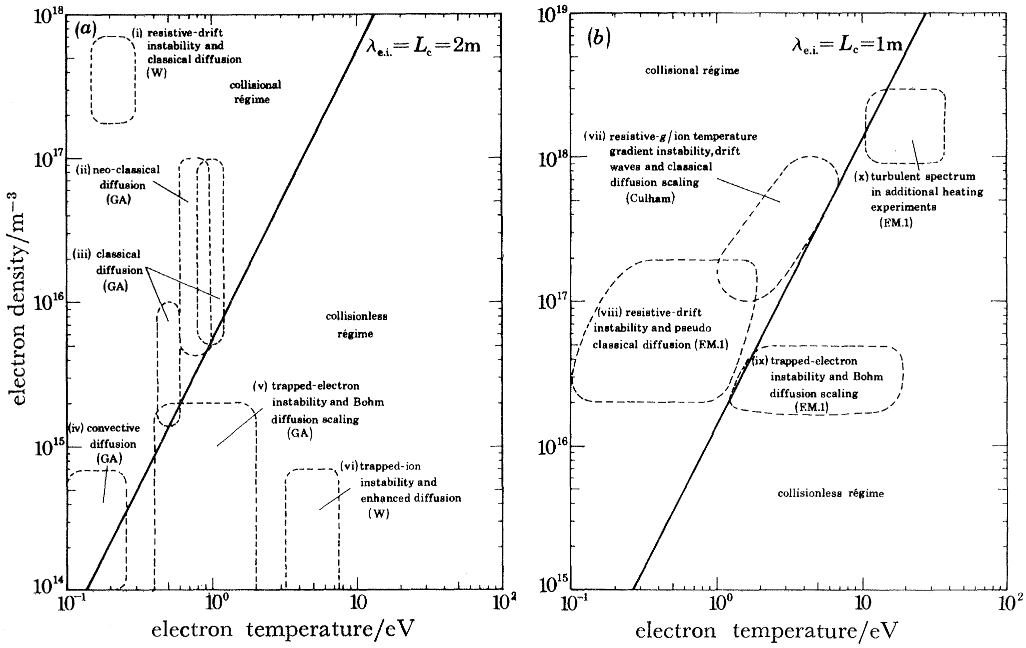

Figure 5. Operating régimes for fluctuation and confinement measurements made in (*a*) octupoles with connection length $L_c \approx 2$ m and (*b*) levitrons ($L_c \approx 1$ m) as a function of electron temperature and density. References are as follows: (i) Navratil & Post (1977); (ii) Ohkawa *et al.* (1974); (iii) Ohkawa *et al.* (1971); (iv) Tamano *et al.* (1973); (v) Prater *et al.* (1978); (vi) Drake (1977); (vii) Cordey *et al.* (1979); (viii) Sinnis *et al.* (1972); (ix) Ejima & Okabayashi (1975); (x) Okabayashi & Arunasalam (1977).

Summary

A summary of confinement experiments in ideal systems is shown in figures 5*a*, *b* on an n, T_e-diagram. A line is drawn on each diagram separating collisional and collisionless régimes. In the octupoles, the transport is classical in the collisional régime whereas in the collisionless régime Bohm scaling is found in the one experiment in which the scaling is determined. In the levitrons, the transport scales classically in the collisional régime also, but only approaches the correct absolute value at high shear. In the collisionless régime again Bohm scaling of the transport is found in the one experiment in which the scaling is determined. This pattern of behaviour is compatible with the general scaling laws for plasma confinement derived by Connor & Taylor (1977) but this does not in any sense explain the transport. The diffusion scaling predicted for the resistive-g mode by quasilineary theory does agree with that observed, but the theory always over-estimates the transport and must be considered as having limited usefulness. This over-estimation of the transport conforms, of course, with the interpretation of the limits as upper bounds only.

There is evidence, both experimental and theoretical, that large-scale convection cells, not considered by the quasilinear theory, can be generated by micro-instabilities through nonlinear processes and can contribute to cross-field transport.

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We conclude that the experiments in ideal confinement systems have shown that classical loss rates can be achieved and have provided useful information on micro-instabilities and the scaling of enhanced losses. Detailed understanding of enhanced transport is lacking, however, and only crude comparisons with theory have been possible.

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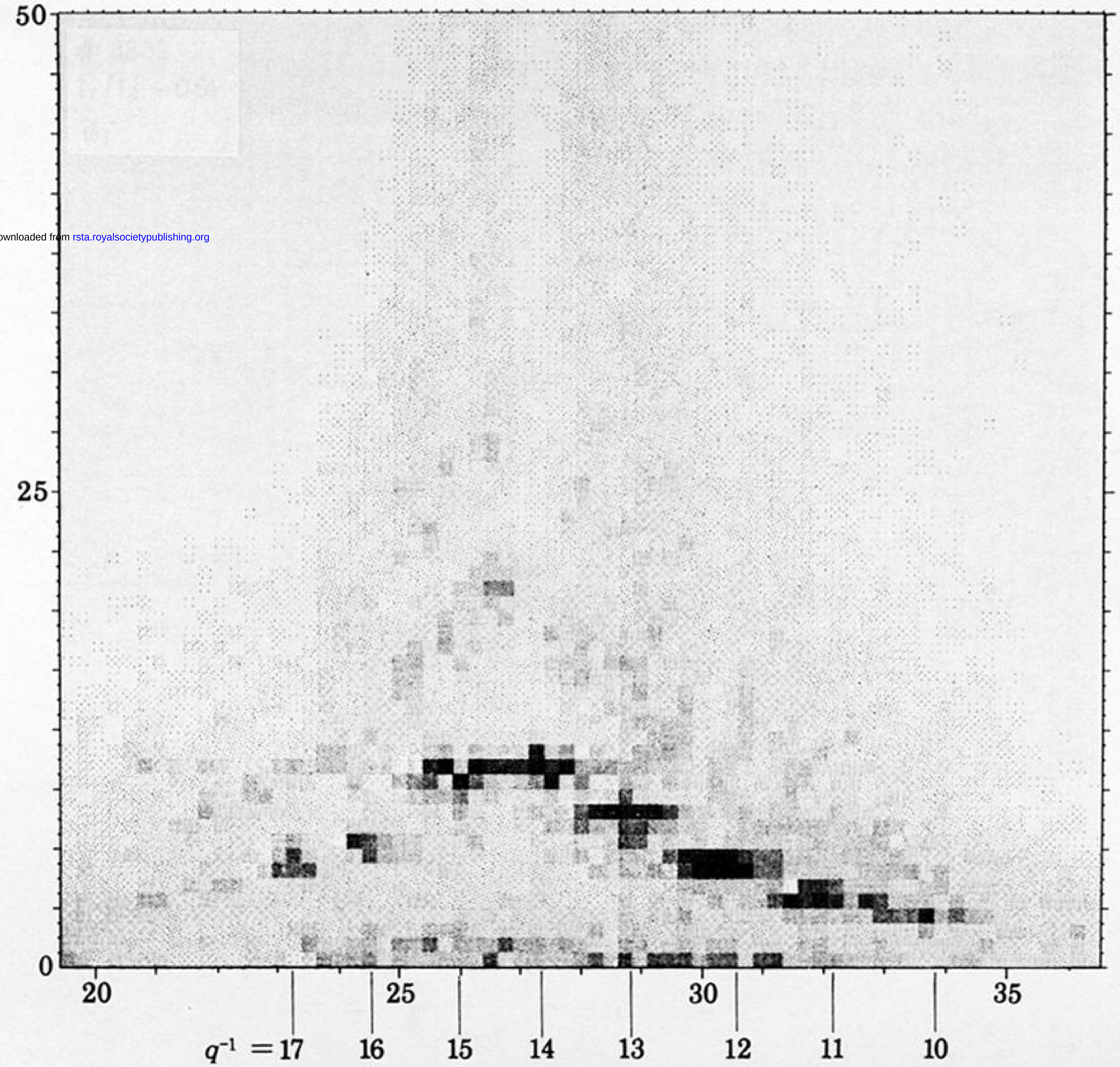

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GURE 3*b*. Fourier spectra of the fluctuations shown in (*a*) across the profile. Darker shades refer to larger amplitude with logarithmic grading between shades. (Figure 3*b* was supplied by W. H. W. Fletcher and E. M. Jones.)