Is Anomalous Transport "Diffusive"?

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ABSTRACT

It has often been assumed that the anomalous transport from saturated plasma instabilities is "diffusive" in the sense that the particle flux, Γ, the electron energy flux, q_e, and the ion energy flux, q_i, can be written in forms that are *linear* in the density gradient, dn/dr, the electron temperature gradient, dT_e/dr, and the ion temperature gradient dT_i/dr. In the simplest form, $\Gamma = -D_n^n(dn/dr)$, $q_e = -D_e^e n(dT_e/dr)$, and $q_i = -D_i^i n(dT_i/dr)$. A possible generalization of this is to include so-called "off-diagonal" terms, with $\Gamma = nV_n - D_n^n(dn/dr) - D_n^e(n/T_e)(dT_e/dr) - D_n^i(n/T_i)(dT_i/dr)$, with corresponding forms for the energy fluxes. Here, general results for the quasilinear particle and energy fluxes, resulting from tokamak linear microinstabilities, are evaluated to assess the relative importance of the diagonal and the off-diagonal terms. A further possible generalization is to include also contributions to the fluxes from higher powers of the gradients, specifically "quadratic" contributions proportional to $(dn/dr)^2$, $(dn/dr)(dT_e/dr)$, and so on. A procedure is described for evaluating the corresponding coefficients, and results are presented for illustrative, realistic tokamak cases. Qualitatively, it is found that the off-diagonal diffusion coefficients can be as big as the diagonal ones, and that the quadratic terms can be larger than the linear ones. The results thus strongly suggest that the commonly used "diffusive" approximation with only diagonal terms, $\Gamma = -D_n^n(dn/dr)$, and correspondingly for the energy fluxes, is not adequate in practice.

I. INTRODUCTION

Standard forms for transport equations and fluxes have been discussed recently by D. W. Ross.[1] Here, these equations will be extended in various ways and the coefficients will be evaluated explicitly for the quasilinear transport due to linear tokamak instabilities in realistic, illustrative cases. This will allow a test of the so-called "diffusive" approximation, where the quasilinear particle and energy fluxes of electrons and ions are linear in the density and temperature gradients. In this work, a pure hydrogenic plasma is considered so that the equilibrium densities satisfy $n_e = n_i = n$, and $dn_e/dr = dn_i/dr = dn/dr$, and the electron and ion quasilinear particle fluxes are equal on any chosen magnetic surface, $\Gamma_e = \Gamma_i = \Gamma$. Defining Q_j as the total energy flux for species j, the $(5/2)\Gamma T_j$ convective part is subtracted out, giving $q_j \equiv Q_j - (5/2)\Gamma T_j$. Alternative forms of the convective part are compared in Ref. 1, and it is concluded there that $(5/2)\Gamma T_j$ is the correct form. The simplest forms of the diffusion equations are

$$\Gamma = -D_n^n \frac{dn}{dr},$$
$$q_e = -D_e^e n \frac{dT_e}{dr}, \tag{1}$$
$$q_i = -D_i^i n \frac{dT_i}{dr}.$$

Here, only the so-called "diagonal" terms are kept. The generalization to include the so-called "off-diagonal" diffusion coefficients, as well as corresponding anomalous convection (pinch) velocities, V_n, V_e, and V_i, is

$$\Gamma = nV_n - D_n^n \frac{dn}{dr} - D_n^e \frac{n}{T_e}\frac{dT_e}{dr} - D_n^i \frac{n}{T_i}\frac{dT_i}{dr},$$
$$q_e = nT_eV_e - D_e^n T_e \frac{dn}{dr} - D_e^e n \frac{dT_e}{dr} - D_e^i n \frac{dT_i}{dr}, \tag{2}$$
$$q_i = nT_iV_i - D_i^n T_i \frac{dn}{dr} - D_i^e n \frac{dT_e}{dr} - D_i^i n \frac{dT_i}{dr}.$$

Equations (2) can be conveniently written in matrix form as

$$\begin{pmatrix} \Gamma/n \\ q_e/(nT_e) \\ q_i/(nT_i) \end{pmatrix} = \begin{pmatrix} V_n \\ V_e \\ V_i \end{pmatrix} - \begin{pmatrix} D_n^n & D_n^e & D_n^i \\ D_e^n & D_e^e & D_e^i \\ D_i^n & D_i^e & D_i^i \end{pmatrix} \begin{pmatrix} (dn/dr)/n \\ (dT_e/dr)/T_e \\ (dT_i/dr)/T_i \end{pmatrix} . \tag{3}$$

Normalizing this way, one finds that all of the diagonal and off-diagonal D_j^k (where j and $k = n$, e, and i) have the same units (m^2/sec, say), and so are directly comparable. For notational convenience, write $\hat{\Gamma} \equiv \Gamma/n$, $\hat{q}_j \equiv q_j/(nT_j)$, $\kappa_n \equiv -(1/n)(dn/dr)$, and $\kappa_j \equiv -(1/T_j)(dT_j/dr)$, so that Eq. (3) becomes

$$\begin{pmatrix} \hat{\Gamma} \\ \hat{q}_e \\ \hat{q}_i \end{pmatrix} = \begin{pmatrix} V_n \\ V_e \\ V_i \end{pmatrix} + \begin{pmatrix} D_n^n & D_n^e & D_n^i \\ D_e^n & D_e^e & D_e^i \\ D_i^n & D_i^e & D_i^i \end{pmatrix} \begin{pmatrix} \kappa_n \\ \kappa_e \\ \kappa_i \end{pmatrix} . \tag{4}$$

Equation (4) can be regarded as a "Taylor series" for the fluxes, $\hat{\Gamma}$, $\hat{q}_e$, and $\hat{q}_i$ in powers of the gradients, κ_n, κ_e, and κ_i, truncated at first order. The extension of Eq. (4) to second powers in the gradients is

$$\begin{aligned} \begin{pmatrix} \hat{\Gamma} \\ \hat{q}_e \\ \hat{q}_i \end{pmatrix} &= \begin{pmatrix} V_n \\ V_e \\ V_i \end{pmatrix} + \begin{pmatrix} D_n^n & D_n^e & D_n^i \\ D_e^n & D_e^e & D_e^i \\ D_i^n & D_i^e & D_i^i \end{pmatrix} \begin{pmatrix} \kappa_n \\ \kappa_e \\ \kappa_i \end{pmatrix} \\ &+ \begin{pmatrix} E_n^{nn} & E_n^{ee} & E_n^{ii} \\ E_e^{nn} & E_e^{ee} & E_e^{ii} \\ E_i^{nn} & E_i^{ee} & E_i^{ii} \end{pmatrix} \begin{pmatrix} \kappa_n^2 \\ \kappa_e^2 \\ \kappa_i^2 \end{pmatrix} + \begin{pmatrix} F_n^{ne} & F_n^{ni} & F_n^{ei} \\ F_e^{ne} & F_e^{ni} & F_e^{ei} \\ F_i^{ne} & F_i^{ni} & F_i^{ei} \end{pmatrix} \begin{pmatrix} \kappa_n \kappa_e \\ \kappa_n \kappa_i \\ \kappa_e \kappa_i \end{pmatrix} . \end{aligned} \tag{5}$$

Equation (4) is a good approximation to Eq. (5) if the quadratic E_j^{kk} and F_j^{kl} contributions to the fluxes are small compared to the linear D_j^k contributions, and Eq. (1) is a good approximation to Eq. (4) if the off-diagonal D_j^k are small compared to the diagonal D_j^j and the V_j are small compared to the D_j^k contributions. It is these conditions which will be investigated here for realistic tokamak instabilities.

A comprehensive electromagnetic kinetic eigenfrequency eigenfunction code[2] for high toroidal mode number tokamak instabilities such as the trapped-electron-η_i mode and the MHD ballooning mode is employed here. This code also calculates the quasilinear fluxes of particles and energy for each plasma species using formulas which are direct generalizations of those derived in Ref. 3. The eigenmode equations, the solution methods employed, and the quasilinear flux expressions are all presented in detail in Ref. 2, so this material will not be repeated here. However, this calculation can be described briefly as constituting a solution of the linearized gyrokinetic equation in toroidal geometry without any frequency expansions, using a model collision operator. It employs the ballooning formalism at lowest order in $1/n$ (toroidal mode number), so the calculation is local to a single, chosen magnetic surface. Thus, the calculation includes transit frequency harmonics (Landau damping, *etc.*) for circulating particles of each species, bounce frequency harmonics for trapped particles of each species, magnetic drift frequency effects, and full finite Larmor radius effects. Input data, such as radial density and temperature profiles, for the calculation of the results presented here come either from experimental measurements or from transport code results. The MHD equilibrium used in the instability calculation is computed numerically from the corresponding pressure and safety factor profiles, or else an algebraic model MHD equilibrium is used.

The quasilinear fluxes, while being only bilinear in the fluctuation amplitudes, can in general be nonlinear in the density and temperature gradients. Given a set of gradients, κ_n, κ_e, and κ_i, the code calculates the quasilinear particle and energy fluxes, $\hat{\Gamma}$, $\hat{q}_e$, and $\hat{q}_i$, proportional to the square of the unknown saturation level ϕ_0 for the perturbed electrostatic potential of the mode. The results to be presented here will allow us to assess the strength of this nonlinearity in the gradients.

The specific procedure for doing this is described in Sec. II. Numerical results for realistic tokamak cases are presented in Sec. III for the linear truncation in Eq. (4) and in Sec. IV for the quadratic truncation in Eq. (5). Discussion and conclusions are given in Sec. V.

II. Gradient Variation

For practical computation, start with a base parameter set, denoted by a zero superscript, and compute $\hat{\Gamma}^0$, $\hat{q}_e^0$, and $\hat{q}_i^0$ for $\kappa_n = \kappa_n^0$, $\kappa_e = \kappa_e^0$, and $\kappa_i = \kappa_i^0$. Here, κ_n^0, κ_e^0, and κ_i^0 will represent the actual experimental or transport code values on a chosen magnetic surface. Then, define a series of parameter sets in which κ_n, κ_e, and κ_i are varied slightly up or down by amounts $\Delta\kappa_n$, $\Delta\kappa_e$, and $\Delta\kappa_i$, respectively, in different combinations, from the base parameter set, with $\hat{\Gamma}^m$, $\hat{q}_e^m$, and $\hat{q}_i^m$ being recomputed for each of these sets of gradient values, where the superscript m identifies the parameter set. The κ_n, κ_e, and κ_i values for the needed parameter sets are all listed in Table I.

Only results for parameter sets 0 to 3 are actually needed to evaluate all of the D_j^k in Eq. (4). Using standard formulas[4] for the numerical evaluation of derivatives, it is found that

$$
\begin{aligned}
D_n^n &= \frac{\hat{\Gamma}^1 - \hat{\Gamma}^0}{\Delta\kappa_n}, \quad D_n^e = \frac{\hat{\Gamma}^2 - \hat{\Gamma}^0}{\Delta\kappa_e}, \quad D_n^i = \frac{\hat{\Gamma}^3 - \hat{\Gamma}^0}{\Delta\kappa_i},\\
D_e^n &= \frac{\hat{q}_e^1 - \hat{q}_e^0}{\Delta\kappa_n}, \quad D_e^e = \frac{\hat{q}_e^2 - \hat{q}_e^0}{\Delta\kappa_e}, \quad D_e^i = \frac{\hat{q}_e^3 - \hat{q}_e^0}{\Delta\kappa_i},\\
D_i^n &= \frac{\hat{q}_i^1 - \hat{q}_i^0}{\Delta\kappa_n}, \quad D_i^e = \frac{\hat{q}_i^2 - \hat{q}_i^0}{\Delta\kappa_e}, \quad D_i^i = \frac{\hat{q}_i^3 - \hat{q}_i^0}{\Delta\kappa_i}
\end{aligned}
\tag{6}
$$

to lowest order in the $\Delta\kappa_j/\kappa_j^0$. Once the D_j^k have been evaluated, then the anomalous convection velocities are given, from Eq. (4), as

$$
\begin{aligned}
V_n &= \hat{\Gamma}^0 - D_n^n\kappa_n^0 - D_n^e\kappa_e^0 - D_n^i\kappa_i^0,\\
V_e &= \hat{q}_e^0 - D_e^n\kappa_n^0 - D_e^e\kappa_e^0 - D_e^i\kappa_i^0,\\
V_i &= \hat{q}_i^0 - D_i^n\kappa_n^0 - D_i^e\kappa_e^0 - D_i^i\kappa_i^0.
\end{aligned}
\tag{7}
$$

In evaluating the D_j^k in Eq. (5), results for parameter sets 0 to 6 can also be used, and the same results are necessary to evaluate the E_j^{kk}. Results for all of parameter sets 7 to

18 are needed to evaluate the F_j^{kl} in Eq. (5). Again using standard formulas[4] for numerical evaluation of derivatives, it is found that

$$
\begin{aligned}
D_n^n &= \frac{\hat{\Gamma}^1 - \hat{\Gamma}^4}{2\Delta\kappa_n}, \quad D_n^e = \frac{\hat{\Gamma}^2 - \hat{\Gamma}^5}{2\Delta\kappa_e}, \quad D_n^i = \frac{\hat{\Gamma}^3 - \hat{\Gamma}^6}{2\Delta\kappa_i}, \\
D_e^n &= \frac{\hat{q}_e^1 - \hat{q}_e^4}{2\Delta\kappa_n}, \quad D_e^e = \frac{\hat{q}_e^2 - \hat{q}_e^5}{2\Delta\kappa_e}, \quad D_e^i = \frac{\hat{q}_e^3 - \hat{q}_e^6}{2\Delta\kappa_i}, \\
D_i^n &= \frac{\hat{q}_i^1 - \hat{q}_i^4}{2\Delta\kappa_n}, \quad D_i^e = \frac{\hat{q}_i^2 - \hat{q}_i^5}{2\Delta\kappa_e}, \quad D_i^i = \frac{\hat{q}_i^3 - \hat{q}_i^6}{2\Delta\kappa_i},
\end{aligned}
\tag{8}
$$

and

$$
\begin{aligned}
E_n^{nn} &= \frac{\hat{\Gamma}^1 + \hat{\Gamma}^4 - 2\hat{\Gamma}^0}{2\Delta\kappa_n^2}, \quad E_n^{ee} = \frac{\hat{\Gamma}^2 + \hat{\Gamma}^5 - 2\hat{\Gamma}^0}{2\Delta\kappa_e^2}, \quad E_n^{ii} = \frac{\hat{\Gamma}^3 + \hat{\Gamma}^6 - 2\hat{\Gamma}^0}{2\Delta\kappa_i^2}, \\
E_e^{nn} &= \frac{\hat{q}_e^1 + \hat{q}_e^4 - 2\hat{q}_e^0}{2\Delta\kappa_n^2}, \quad E_e^{ee} = \frac{\hat{q}_e^2 + \hat{q}_e^5 - 2\hat{q}_e^0}{2\Delta\kappa_e^2}, \quad E_e^{ii} = \frac{\hat{q}_e^3 + \hat{q}_e^6 - 2\hat{q}_e^0}{2\Delta\kappa_i^2}, \\
E_i^{nn} &= \frac{\hat{q}_i^1 + \hat{q}_i^4 - 2\hat{q}_i^0}{2\Delta\kappa_n^2}, \quad E_i^{ee} = \frac{\hat{q}_i^2 + \hat{q}_i^5 - 2\hat{q}_i^0}{2\Delta\kappa_e^2}, \quad E_i^{ii} = \frac{\hat{q}_i^3 + \hat{q}_i^6 - 2\hat{q}_i^0}{2\Delta\kappa_i^2},
\end{aligned}
\tag{9}
$$

and

$$
\begin{aligned}
F_n^{ne} &= \frac{\hat{\Gamma}^7 - \hat{\Gamma}^8 - \hat{\Gamma}^9 + \hat{\Gamma}^{10}}{4\Delta\kappa_n\Delta\kappa_e}, \quad F_n^{ni} = \frac{\hat{\Gamma}^{11} - \hat{\Gamma}^{12} - \hat{\Gamma}^{13} + \hat{\Gamma}^{14}}{4\Delta\kappa_n\Delta\kappa_i}, \quad F_n^{ei} = \frac{\hat{\Gamma}^{15} - \hat{\Gamma}^{16} - \hat{\Gamma}^{17} + \hat{\Gamma}^{18}}{4\Delta\kappa_e\Delta\kappa_i}, \\
F_e^{ne} &= \frac{\hat{q}_e^7 - \hat{q}_e^8 - \hat{q}_e^9 + \hat{q}_e^{10}}{4\Delta\kappa_n\Delta\kappa_e}, \quad F_e^{ni} = \frac{\hat{q}_e^{11} - \hat{q}_e^{12} - \hat{q}_e^{13} + \hat{q}_e^{14}}{4\Delta\kappa_n\Delta\kappa_i}, \quad F_e^{ei} = \frac{\hat{q}_e^{15} - \hat{q}_e^{16} - \hat{q}_e^{17} + \hat{q}_e^{18}}{4\Delta\kappa_e\Delta\kappa_i}, \\
F_i^{ne} &= \frac{\hat{q}_i^7 - \hat{q}_i^8 - \hat{q}_i^9 + \hat{q}_i^{10}}{4\Delta\kappa_n\Delta\kappa_e}, \quad F_i^{ni} = \frac{\hat{q}_i^{11} - \hat{q}_i^{12} - \hat{q}_i^{13} + \hat{q}_i^{14}}{4\Delta\kappa_n\Delta\kappa_i}, \quad F_i^{ei} = \frac{\hat{q}_i^{15} - \hat{q}_i^{16} - \hat{q}_i^{17} + \hat{q}_i^{18}}{4\Delta\kappa_e\Delta\kappa_i},
\end{aligned}
\tag{10}
$$

to lowest order in the $\Delta\kappa_j/\kappa_j^0$. From Eq. (5), the anomalous convection velocities are then given as

$$
\begin{aligned}
V_n &= \hat{\Gamma}^0 - D_n^n\kappa_n^0 - D_n^e\kappa_e^0 - D_n^i\kappa_i^0 - E_n^{nn}(\kappa_n^0)^2 - E_n^{ee}(\kappa_e^0)^2 - E_n^{ii}(\kappa_i^0)^2 - F_n^{ne}\kappa_n^0\kappa_e^0 - F_n^{ni}\kappa_n^0\kappa_i^0 - F_n^{ei}\kappa_e^0\kappa_i^0, \\
V_e &= \hat{q}_e^0 - D_e^n\kappa_n^0 - D_e^e\kappa_e^0 - D_e^i\kappa_i^0 - E_e^{nn}(\kappa_n^0)^2 - E_e^{ee}(\kappa_e^0)^2 - E_e^{ii}(\kappa_i^0)^2 - F_e^{ne}\kappa_n^0\kappa_e^0 - F_e^{ni}\kappa_n^0\kappa_i^0 - F_e^{ei}\kappa_e^0\kappa_i^0, \\
V_i &= \hat{q}_i^0 - D_i^n\kappa_n^0 - D_i^e\kappa_e^0 - D_i^i\kappa_i^0 - E_i^{nn}(\kappa_n^0)^2 - E_i^{ee}(\kappa_e^0)^2 - E_i^{ii}(\kappa_i^0)^2 - F_i^{ne}\kappa_n^0\kappa_e^0 - F_i^{ni}\kappa_n^0\kappa_i^0 - F_i^{ei}\kappa_e^0\kappa_i^0.
\end{aligned}
\tag{11}
$$

However, there are some practical questions as to what to include in the variation process so as to pick up various indirect effects. Specifically, in the fully electromagnetic version of the computer code, as dn/dr, dT_e/dr, and dT_i/dr are varied, n, T_e, T_i, B_0, β, and the toroidal mode number are held fixed. Also, the changed converged values of the eigenfrequency ω are used, which can depend on all of the gradients in a nonlinear fashion, and which can thus introduce nonlinear dependence on the gradients into the fluxes. A difficult question is what to do about the undetermined saturation level ϕ_0 for the perturbed electrostatic potential that enters the quasilinear flux formulas. It is decided here to take $|e\phi_0/T_e| = |C/k_\theta r_p|$, a mixing-length type assumption, where r_p is the *total* pressure gradient scale length, so that all three of the changing equilibrium gradients enter the saturation condition. Here, C is kept constant during the gradient variation, but is chosen at the end to set the final D_e^e (often called χ_e) to an experimentally typical value, say 1 m^2/sec. This is an ad-hoc saturation criterion, in the absence of a generally accepted nonlinear theory for these modes, but it should be adequate for the present purpose of investigating the nature of the gradient dependence of the fluxes.

III. Diagonal and Off-Diagonal Diffusion Coefficients

In this and the next section, results will be presented for several cases. The first corresponds to a Tokamak Fusion Test Reactor (TFTR) "supershot,"[5] (shot 35782).[6] The radial profiles of n, T_e, and T_i come from experimental measurements, and the MHD equilibrium is computed numerically with the corresponding pressure and q profiles. Here, the MHD equilibrium is held fixed as the local gradients are varied. The instability calculation is carried out for the trapped-electron-η_i mode on a chosen magnetic surface with $r/a \simeq 1/2$. The local parameters on this surface are: $q = 1.54$, $(r/q)(dq/dr) = 1.18$, $n_e = n_i = n = 2.91 \times 10^{19}\mathrm{m}^{-3}$, $r_n \equiv -(d\ln n/dr)^{-1} = 0.345$ m, $T_e = 4.12$ keV, $r_{Te} = 0.242$ m, $\eta_e = 1.43$, $T_i = 7.17$ keV, $r_{Ti} = 0.152$ m, $\eta_i = 2.27$, $T_i/T_e = 1.74$, $r = 0.415$ m, $a = 0.792$ m, $R_0 = 2.45$ m, $B_0 = 47.2$ kG, $\beta = 1.05\%$, toroidal mode number $= 55$, $k_\theta\rho_i = 0.495$, and $\nu_e = \nu_i = 0$

(collisionless). Then, for the trapped-electron-η_i mode on this surface, the eigenfrequency is $\omega/\omega_{*e} = -0.749 - 0.471i$, or $\omega = (2.78 - 1.75i) \times 10^5 \mathrm{sec}^{-1}$, for the eigenmode with the fewest nodes along the field line for the perturbed potentials, which balloon strongly to the outside of the torus. The corresponding basic gradient parameter values are $\kappa_n = 2.90\ \mathrm{m}^{-1}$, $\kappa_e = 4.14\ \mathrm{m}^{-1}$, and $\kappa_i = 6.57\ \mathrm{m}^{-1}$. In this case, $\Delta\kappa_n$, $\Delta\kappa_e$, and $\Delta\kappa_i$ are chosen to be 1% of the corresponding basic gradient values. Then, computing $\hat{\Gamma}^m$, $\hat{q}_e^m$, and $\hat{q}_i^m$ for parameter sets $m = 0$ to 3, as defined in Table I, and substituting the results in Eqs. (6) and (7), Eq. (4) is evaluated for this case as

$$\begin{aligned}\begin{pmatrix}\hat{\Gamma}\\ \hat{q}_e\\ \hat{q}_i\end{pmatrix} &= \begin{pmatrix}-7.3\\ 1.3\\ -16.0\end{pmatrix}\frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix}1.74 & 0.04 & 0.74\\ -0.93 & 1.00 & -0.44\\ 0.56 & 1.31 & 2.60\end{pmatrix}\frac{\mathrm{m}^2}{\mathrm{sec}}\begin{pmatrix}\kappa_n\\ \kappa_e\\ \kappa_i\end{pmatrix}\\ &= \begin{pmatrix}-7.3\\ 1.3\\ -16.0\end{pmatrix}\frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix}10.1\\ -1.5\\ 24.1\end{pmatrix}\frac{\mathrm{m}}{\mathrm{sec}} = \begin{pmatrix}2.8\\ -0.2\\ 8.1\end{pmatrix}\frac{\mathrm{m}}{\mathrm{sec}},\end{aligned} \tag{12}$$

with $C = 0.134$. Qualitatively, V_n and V_i are inward while V_e is slightly outward, and each is somewhat smaller than the corresponding D_j^k matrix contribution. Thus, there are substantial partial cancellations, leaving net outward fluxes for $\hat{\Gamma}$ and $\hat{q}_i$, and a slight net inward flux for $\hat{q}_e$. Of course, if the $(5/2)\Gamma T_j$ convective terms are included, the total energy fluxes are both outward: $\hat{Q}_e \equiv \hat{q}_e + (5/2)\hat{\Gamma} = 6.7$ m/sec and $\hat{Q}_i \equiv \hat{q}_i + (5/2)\hat{\Gamma} = 15.0$ m/sec with this normalization. It should be remarked that, physically, the V_j do not have independent existences in the sense of allowing nonzero fluxes in the limit of zero gradients. Rather, they reflect the finite thresholds on the gradients to have instability, with the fluxes rising steeply above these thresholds. Below the thresholds, the mode is damped and the corresponding fluxes are just zero. This implies that the response to a perturbation of the gradients would be larger than estimated from the steady-state fluxes on the basis of Eq. (1), where there are no anomalous convection velocities. In the D_j^k term in Eq. (12), it is seen that all of the diagonal and off-diagonal elements are of the same order, except for $D_n^e = 0.04$.

and that D_e^n and D_e^i are negative. In this case, then, Eq. (1) is a poor approximation to Eq. (4).

The same procedure has been carried out for three other cases. These will now be described more briefly. The next two cases have input parameters taken from runs of the BALDUR transport code[7] for a now superseded design of the Compact Ignition Tokamak (CIT) with major radius $R_0 = 1.6$ m. Specifically, on the chosen magnetic surface with $r/a \simeq 0.5$, at the BALDUR-predicted value of local $\beta = 8.34\%$, both the trapped-electron-η_i mode and the kinetically calculated MHD ballooning mode are unstable. Both of these modes will be considered as separate cases. Also, for these cases, the so-called "s-α" model MHD equilibrium[8] is employed instead of a numerical MHD equilibrium, with the Shafranov shift parameter $\alpha \propto dp/dr$ varying as the gradients are changed. In this case, the trapped-electron-η_i mode at $k_\theta \rho_i = 0.75$ has eigenfrequency $\omega/\omega_{*e} = -0.0299 + 0.127i$, and the MHD ballooning mode at $k_\theta \rho_i = 0.40$ has eigenfrequency $\omega/\omega_{*e} = -1.76 + 0.0273i$. The basic parameter set has $\kappa_n^0 = 1.70\ \mathrm{m}^{-1}$, $\kappa_e^0 = 0.684\ \mathrm{m}^{-1}$, and $\kappa_i^0 = 0.917\ \mathrm{m}^{-1}$, with the $\Delta\kappa_j$'s chosen to be 1% of the corresponding κ_j^0's. Also, for this case, $\eta_e = 0.403$ and $\eta_i = 0.541$. Carrying through the same procedure for the trapped-electron-η_i mode for this case results in

$$\begin{aligned}\begin{pmatrix}\dot{\Gamma}\\ \dot{q}_e\\ \dot{q}_i\end{pmatrix} &= \begin{pmatrix}-3.7\\ 2.5\\ 4.2\end{pmatrix}\frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix}3.01 & 0.67 & 0.20\\ -3.20 & 1.00 & 0.25\\ -4.57 & -0.73 & 1.30\end{pmatrix}\frac{\mathrm{m}^2}{\mathrm{sec}}\begin{pmatrix}\kappa_n\\ \kappa_e\\ \kappa_i\end{pmatrix}\\ &= \begin{pmatrix}-3.7\\ 2.5\\ 4.2\end{pmatrix}\frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix}5.7\\ -4.5\\ -7.1\end{pmatrix}\frac{\mathrm{m}}{\mathrm{sec}} = \begin{pmatrix}2.0\\ -2.0\\ -2.9\end{pmatrix}\frac{\mathrm{m}}{\mathrm{sec}},\end{aligned} \tag{13}$$

for $C = 0.214$. Here, V_n is inward while V_e and V_i are outward. All of the diagonal and off-diagonal D_j^k are comparable in magnitude, and the diagonal elements are all positive, while several of the off-diagonal elements are negative. Again, there is substantial cancellation between the V_j convection terms and the D_j^k diffusion terms, leading to net outward $\dot{\Gamma}$ and to net inward $\dot{q}_e$ and $\dot{q}_i$. However, again the total energy fluxes are both outward, with $\dot{Q}_e = 3.1$ m/sec and $\dot{Q}_i = 2.3$ m/sec, in the normalization of Eq. (13).

The corresponding result for the MHD ballooning mode is

$$\begin{aligned}
\begin{pmatrix} \hat{\Gamma} \\ \hat{q}_e \\ \hat{q}_i \end{pmatrix} &= \begin{pmatrix} -19.4 \\ 7.7 \\ -0.1 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix} 7.43 & 4.11 & 5.73 \\ -3.76 & 1.00 & -2.55 \\ 1.36 & 21.26 & -11.22 \end{pmatrix} \frac{\mathrm{m}^2}{\mathrm{sec}} \begin{pmatrix} \kappa_n \\ \kappa_e \\ \kappa_i \end{pmatrix} \\
&= \begin{pmatrix} -19.4 \\ 7.7 \\ -0.1 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix} 20.7 \\ -8.0 \\ 6.6 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}} = \begin{pmatrix} 1.2 \\ -0.3 \\ 6.4 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}}.
\end{aligned} \tag{14}$$

for $C = 0.849$. For this instability, V_n is inward, while V_e is outward and V_i is much smaller. Again, the off-diagonal D_j^k are of comparable or larger magnitude than the diagonal elements. In this case, D_i^i, D_e^n, and D_e^i are negative, and the others positive. Here, there is substantial cancellation between the convective V_n and V_e and the corresponding D_j^k contributions, but not for V_i. The net $\hat{\Gamma}$ and $\hat{q}_i$ are outward, while the net $\hat{q}_e$ is slightly inward. However, the total energy fluxes are again both outward, with $\dot{Q}_e = 2.8$ m/sec and $\dot{Q}_i = 9.6$ m/sec with the normalization of Eq. (14).

The last case to be considered here is a TFTR α-storage mode case[9] (here used without the α-particles), again with the "s-α" model MHD equilibrium, for the trapped-electron-η_i mode with local $\beta = 1.21\%$, $\eta_e = 0.888$, and $\eta_i = 0.713$. The basic parameter set, from the BALDUR transport code results, has $\kappa_n^0 = 2.51$ m^{-1}, $\kappa_e^0 = 2.23$ m^{-1}, and $\kappa_i^0 = 1.79$ m^{-1}. In this case, $\Delta\kappa_n$ is chosen as 10% of κ_n^0, $\Delta\kappa_e$ as 1% of κ_e^0, and $\Delta\kappa_i$ as 2.5% of κ_i^0. Carrying through the same procedure results in

$$\begin{aligned}
\begin{pmatrix} \hat{\Gamma} \\ \hat{q}_e \\ \hat{q}_i \end{pmatrix} &= \begin{pmatrix} -2.9 \\ 0.9 \\ 2.9 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix} 1.88 & 0.02 & -0.38 \\ -1.32 & 1.00 & -0.09 \\ -2.83 & 0.02 & 1.44 \end{pmatrix} \frac{\mathrm{m}^2}{\mathrm{sec}} \begin{pmatrix} \kappa_n \\ \kappa_e \\ \kappa_i \end{pmatrix} \\
&= \begin{pmatrix} -2.9 \\ 0.9 \\ 2.9 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix} 4.1 \\ -1.2 \\ -4.5 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}} = \begin{pmatrix} 1.2 \\ -0.3 \\ -1.6 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}},
\end{aligned} \tag{15}$$

for $C = 0.139$. For this case, V_n is inward, while V_e and V_i are outward. Again, most of the off-diagonal D_j^k are of comparable or larger magnitude than the diagonal elements. In this case again, several of the D_j^k are negative. Also, there is substantial cancellation between the convective V_j and the corresponding D_j^k contributions. The net $\dot{\Gamma}$ is outward, while the net $\dot{q}_e$ and $\dot{q}_i$ are inward. However, the total energy fluxes are again both outward, with $\dot{Q}_e = 2.6$ m/sec and $\dot{Q}_i = 1.3$ m/sec with the normalization of Eq. (15).

IV. Linear and Quadratic Gradient Dependences

Using the same TFTR supershot case employed in the previous section, the coefficients in Eq. (5) are evaluated as described in Sec. II. Recomputing the $\dot{\Gamma}^m$, $\dot{q}_e^m$, and $\dot{q}_i^m$ for parameter sets $m = 0$ to 18 as specified in Table I, and then evaluating the D_j^k, E_j^{kk}, and F_j^{kl} according to Eqs. (8) to (11), Eq. (5) in this case becomes

$$\begin{aligned}
\begin{pmatrix} \dot{\Gamma} \\ \dot{q}_e \\ \dot{q}_i \end{pmatrix} &= \begin{pmatrix} -200.3 \\ 95.3 \\ 139.6 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix} 1.68 & 0.04 & 0.71 \\ -0.87 & 1.00 & -0.41 \\ 0.31 & 1.04 & 2.48 \end{pmatrix} \frac{\mathrm{m}^2}{\mathrm{sec}} \begin{pmatrix} \kappa_n \\ \kappa_e \\ \kappa_i \end{pmatrix} \\
&+ \begin{pmatrix} 2.14 & 0 & 0.42 \\ -2.14 & 0 & -0.42 \\ 8.56 & 6.31 & 1.87 \end{pmatrix} \frac{\mathrm{m}^3}{\mathrm{sec}} \begin{pmatrix} \kappa_n^2 \\ \kappa_e^2 \\ \kappa_i^2 \end{pmatrix} \\
&+ \begin{pmatrix} 0 & 8.50 & -0.17 \\ -1.13 & -3.78 & 0.99 \\ -0.38 & -23.84 & 1.65 \end{pmatrix} \frac{\mathrm{m}^3}{\mathrm{sec}} \begin{pmatrix} \kappa_n \kappa_e \\ \kappa_n \kappa_i \\ \kappa_e \kappa_i \end{pmatrix} \\
&= \begin{pmatrix} -200.3 \\ 95.3 \\ 139.6 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix} 9.7 \\ -1.1 \\ 21.5 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix} 36.0 \\ -36.0 \\ 260.7 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}} + \begin{pmatrix} 157.3 \\ -58.4 \\ -413.7 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}} \\
&= \begin{pmatrix} 2.7 \\ -0.2 \\ 8.1 \end{pmatrix} \frac{\mathrm{m}}{\mathrm{sec}},
\end{aligned} \tag{16}$$

for $C = 0.134$. With this normalization, $\dot{Q}_e = 6.9$ m/sec and $\dot{Q}_e = 15.0$ m/sec. In these results, we see that the quadratic (in the gradients) E_j^{kk} and F_j^{kl} term contributions are an order of magnitude larger than the linear D_j^k term contributions! What this means is that the "Taylor series" in powers of the gradients is in fact diverging instead of converging. The V_j's are large, giving near cancellations with the E_j^{kk} and F_j^{kl} term contributions, and leaving much smaller net fluxes. Thus, approximation of Eq. (5) by Eq. (4), let alone by Eq. (1), is seen to be quite unacceptable. In fact, the nonlinearity in the gradients appears from Eq. (16) to be sufficiently strong that the only acceptable course is to compute the fluxes from the quasilinear expressions to all orders in the gradients, without ever expanding in powers of the gradients.

V. Conclusions

A comprehensive electromagnetic kinetic eigenfrequency eigenfunction code[2] for high toroidal mode number tokamak instabilities such as the trapped-electron-η_i mode and the MHD ballooning mode also computes the quasilinear particle and energy fluxes for each plasma species. By varying separately the equilibrium density gradient, electron temperature gradient, and ion temperature gradient from an initial set of values, and recalculating the particle flux, the electron energy flux, and the ion energy flux for each new set of values, it has been shown to be possible to evaluate both the diagonal and off-diagonal anomalous diffusion coefficients corresponding to a particular linear instability, as well as the corresponding anomalous convection (pinch) velocity. The results for realistic tokamak cases presented in Sec. III indicate that: a) The anomalous convection velocities for particles and electron energy and ion energy are large enough to give substantial partial cancellation with the diffusion term, but are not big enough to reverse the sign of the flux. b) The diagonal and off-diagonal diffusion coefficients are generally all of the same order, though some of the off-diagonal elements can be small or negative, depending on the element and the case considered.

The formalism has also been extended here to examine the dependence of the fluxes on second powers of the gradients. A procedure for evaluating the corresponding transport coefficients is described in Sec. II. Results were presented for a realistic TFTR supershot case in Sec. IV, which indicate that the quadratic (in the gradients) contributions to the fluxes are much larger than the linear contributions! In other words, the "Taylor series" for the fluxes in powers of the gradients is in fact diverging. This implies that commonly used approximations for the transport equations, such as Eq. (4), where the fluxes are linear in the gradients, are not acceptable. In this sense, anomalous transport is, in fact, not "diffusive."

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TABLE I. Parameter sets for gradients. Here, $\kappa_n \equiv -(1/n)(dn/dr)$, $\kappa_e \equiv -(1/T_e)(dT_e/dr)$, and $\kappa_i \equiv -(1/T_i)(dT_i/dr)$.

parameter set m	κ_n	κ_e	κ_i
0	κ_n^0	κ_e^0	κ_i^0
1	$\kappa_n^0 + \Delta\kappa_n$	κ_e^0	κ_i^0
2	κ_n^0	$\kappa_e^0 + \Delta\kappa_e$	κ_i^0
3	κ_n^0	κ_e^0	$\kappa_i^0 + \Delta\kappa_i$
4	$\kappa_n^0 - \Delta\kappa_n$	κ_e^0	κ_i^0
5	κ_n^0	$\kappa_e^0 - \Delta\kappa_e$	κ_i^0
6	κ_n^0	κ_e^0	$\kappa_i^0 - \Delta\kappa_i$
7	$\kappa_n^0 + \Delta\kappa_n$	$\kappa_e^0 + \Delta\kappa_e$	κ_i^0
8	$\kappa_n^0 + \Delta\kappa_n$	$\kappa_e^0 - \Delta\kappa_e$	κ_i^0
9	$\kappa_n^0 - \Delta\kappa_n$	$\kappa_e^0 + \Delta\kappa_e$	κ_i^0
10	$\kappa_n^0 - \Delta\kappa_n$	$\kappa_e^0 - \Delta\kappa_e$	κ_i^0
11	$\kappa_n^0 + \Delta\kappa_n$	κ_e^0	$\kappa_i^0 + \Delta\kappa_i$
12	$\kappa_n^0 + \Delta\kappa_n$	κ_e^0	$\kappa_i^0 - \Delta\kappa_i$
13	$\kappa_n^0 - \Delta\kappa_n$	κ_e^0	$\kappa_i^0 + \Delta\kappa_i$
14	$\kappa_n^0 - \Delta\kappa_n$	κ_e^0	$\kappa_i^0 - \Delta\kappa_i$
15	κ_n^0	$\kappa_e^0 + \Delta\kappa_e$	$\kappa_i^0 + \Delta\kappa_i$
16	κ_n^0	$\kappa_e^0 + \Delta\kappa_e$	$\kappa_i^0 - \Delta\kappa_i$
17	κ_n^0	$\kappa_e^0 - \Delta\kappa_e$	$\kappa_i^0 + \Delta\kappa_i$
18	κ_n^0	$\kappa_e^0 - \Delta\kappa_e$	$\kappa_i^0 - \Delta\kappa_i$

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