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A MEAN FIELD OHM'S LAW FOR COLLISIONLESS PLASMAS

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A MEAN FIELD OHM'S LAW FOR COLLISIONLESS PLASMAS

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A mean field Ohm's law valid for collisionless plasmas is derived kinetically. It is shown that contrary to conventional thinking, the resulting hyper-resistivity is significantly smaller than its fluid counterpart due to the fact that the turbulent decorrelation rate is linked to the rapid electron ballistic motion rather than the slower nonlinear mixing time. Moreover, the off-diagonal contributions to the parallel electron momentum flux are shown to result in Ohm's law renormalizations that dwarf the current diffusivity and break radial parity symmetry. Thus, the conventional wisdom of tearing and twisting parity solutions appears to be vitiated in the turbulent, collisionless regime.

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Resistive fluid models have served plasma physics as a useful paradigm for turbulence and anomalous transport. However, the utility of such models for describing present-generation experiments which operate at very low electron collisionalities is, at best, moot. Hence, there is clear motivation to extend the intuitively satisfying picture of resistive magnetohydrodynamics (MHD) into the collisionless domain. Generally speaking, the goal of such investigations is twofold: first, to understand the mechanism of electron momentum transport and its impact on current profile evolution, and second, to uncover possible fluctuation-driven dissipation mechanisms (i.e., turbulent electron viscosity or hyper-resistivity) which effectively replace resistivity as the agent of field-fluid decoupling in a collisionless plasma.

The concept of a turbulent electron viscosity or hyper-resistivity is, by now, well-established in the plasma literature. Such an effect was originally proposed by Schmidt and Yoshikawa¹ as a possible explanation for anomalous skin penetration in tokamaks. Later, an *ad hoc* anomalous electron viscosity term was also invoked by Kaw, Valeo, and Rutherford² as a means for accelerating tearing mode growth, and linking such accelerated growth to magnetic field-line stochasticity which, in turn, follows magnetic island overlap during a major disruption. The problem was reformulated kinetically by Jacobsen and Moses,³ who linked the anomalous electron viscosity to the velocity-integrated Rechester-Rosenbluth⁴ collisionless magnetic diffusion coefficient. Strauss⁵ and others⁶ recognized that a hyper-resistivity and associated counter-terms (negative resistivity) enter the picture as a consequence of the spectral transfer of mean square flux in reduced MHD turbulence. Subsequently, Bhattacharjee and Hameiri⁷ and Boozer⁸ derived the form of the hyper-resistivity from general considerations of selective dissipation (helicity vs. energy) with application to dynamo theory and Taylor relaxation. Craddock⁹ discussed the renormalization of the generalized Ohm's law, incorporating parallel density gradients. More recently, Itoh *et al.*¹⁰ have invoked the concept of a hyper-resistivity-driven MHD interchange-ballooning mode as the basis of a theory of *L*-mode confinement in tokamaks and stellarators (cf. also Connor¹¹). In that theory, the anomalous electron viscosity was linked to the random $\mathbf{E} \times \mathbf{B}$ convection of current fluctuations, i.e., it arises as a nonlinear electron inertia effect. However, as already mentioned, the relevance of a fluid prescription for collisionless tokamak discharges is highly questionable. Finally, Drake and Kleva¹² have applied the concept of anomalous electron diffusion to the problem of collisionless reconnection and the sawtooth crash. Thus, it can be seen that the question of hyper-resistivity is one of broad scope and importance, applicable to problems as diverse as anomalous tearing mode growth, reversed field pinch (RFP) relaxation, *L*-mode anomalous transport in tokamaks, and magnetic reconnection.

In this Letter, we present a systematic procedure for deriving a mean field (MF) Ohm's law for a collisionless plasma. The procedure is equally applicable to examining both mean profile evolution and test wave response. The approach, which is analogous to the procedure used for deriving the familiar "dynamo" equation in MF electrodynamics,¹³ is to calculate the modulations in the radial flux of parallel electron momentum which are induced by a test wave introduced into a background turbulent bath of am-

bient fluctuations. For brevity and simplicity, this analysis is limited to the case of an electrostatic turbulent background. As we shall discuss, the conclusions are generic and easily extended to the case of an electromagnetic bath. It will be seen that the simple fluid picture¹⁰ not only significantly overestimates the level of current transport, but also fails to capture off-diagonal “current pinch” terms which actually dominate the diagonal (hyper-resistive) current diffusion.

We begin our derivation with the nonlinear drift kinetic equation (NLDKE) for the fluctuating electron distribution function:

$$\left(\frac{\partial}{\partial t} + v_{\parallel} \hat{\mathbf{e}}_{\parallel} \cdot \nabla\right) \delta f_e + \delta \mathbf{v}_E \cdot \nabla (\delta f_e + F_{Me}) = \frac{e}{m_e} \delta E_{\parallel} \frac{\partial F_{Me}}{\partial v_{\parallel}}, \quad (1)$$

where $f_e = F_{Me} + \delta f_e$, $F_{Me} = (n_e/\pi^{1/2}v_{te}) \exp\{-[v_{\parallel} - v_0(r)]^2/v_{te}^2\}$ is a shifted electron Maxwellian, and $\delta \mathbf{v}_E = (c/B) \hat{\mathbf{e}}_{\parallel} \times \nabla \delta \phi$ is the $\mathbf{E} \times \mathbf{B}$ velocity. In writing Eq. (1), we have ignored electron-ion collisions which give rise to the usual resistivity in Ohm's law. Thus, the Ohm's law which we derive is valid for $\nu_{ei}/\omega_{te} \ll 1$, where ν_{ei} is the electron-ion collision frequency, and $\omega_{te} = k_{\parallel}v_{te}$ is the electron transit frequency. Note for example that for the Tokamak Fusion Test Reactor (TFTR), ν_{ei}/ω_{te} ranges from 2×10^{-4} in the plasma center to 0.3 at the plasma edge, so that the entire plasma is in the collisionless regime by this definition. Integrating over $v_{\parallel}$, we obtain:

$$\delta E_{\parallel} + \frac{1}{n_e e} \nabla_{\parallel} \delta p_e - \frac{4\pi}{\omega_{pe}^2} \frac{\partial \delta j_{\parallel}}{\partial t} - \nabla \cdot \delta \Gamma_e = 0, \quad (2)$$

where $\delta \Gamma_e = (4\pi/\omega_{pe}^2) \delta \langle \hat{j}_{\parallel} \hat{\mathbf{v}}_E \rangle$ is the modulation in the electron momentum flux induced by the test wave, the angular brackets denote a spectral average over the bath of short-wavelength background turbulence, and the fields associated with the background turbulence have been ‘hatted’ to distinguish them from the fields associated with the test wave (denoted with a ‘ δ ’). The first term on the lhs is the electrostatic field, $\delta E_{\parallel} = -\nabla_{\parallel} \delta \phi$, the second represents parallel compression, the third represents the influence of electron inertia, while the last term represents the induced modulation in radial transport. We now evaluate this last term using the nonadiabatic electron distribution function:

$$\begin{aligned} \Gamma_e &= -i \frac{4\pi e}{\omega_{pe}^2} \frac{c}{B} \int dv_{\parallel} v_{\parallel} \langle k_y \hat{\phi} \hat{f}_e \rangle \\ &= \frac{T_e}{e} \sum_k k_y \rho_e \frac{\omega_k}{\omega_{te}^2} \exp\left(-\frac{\omega_k^2}{\omega_{te}^2}\right) W_k \\ &\quad \times \left(\omega_k - k_{\parallel} v_0 - \left\{ 2 \left(\frac{\omega_k}{\omega_{te}} - \frac{v_0}{v_{te}} \right) \omega_{*j} + \omega_{*n} + \left[\left(\frac{\omega_k}{\omega_{te}} - \frac{v_0}{v_{te}} \right)^2 - \frac{1}{2} \right] \omega_{*T} \right\} \right), \quad (3) \end{aligned}$$

where $W_k = \langle |e \hat{\phi}_k / T_e|^2 \rangle$ is the background fluctuation intensity, $\omega_{*n} = k_y \rho_e v_{te} / L_n$, $\omega_{*T} = k_y \rho_e v_{te} / L_T$, $\omega_{*j} = k_y \rho_e dv_0/dr$, $L_n^{-1} = d \ln n_e / dr$, and $L_T^{-1} = d \ln T_e / dr$. The first two terms on the last line represent the Doppler-shifted frequency, while the remaining terms are the current, density, and temperature gradient-driven fluxes, respectively. For simplicity, we will assume a drift wave-like background bath. Thus, $\omega_k \sim \omega_{*n}$.

and since the electron velocity shift is sonic at best, $v_0/v_{te} \sim (m_e/m_i)^{1/2} \ll 1$, and also $\omega_k/\omega_{te} \sim (R/L_n) k_y \rho_e \ll 1$. *En passant*, we note that our conclusions would be unaffected if we had assumed a drift Alfvén-like bath, since then, $\omega_k/\omega_{te} = v_A/v_{te} = [(m_e/m_i)(T_i/T_e)]^{1/2} \beta^{-1/2} \ll 1$. Using parameters in the core of TFTR, for example, ω_k/ω_{te} is of order 0.05 for electrostatic microturbulence and 0.1 for Alfvénic turbulence. The electron current flux then reduces to

$$\Gamma_e \simeq \frac{T_e}{e} \sum_k k_y \rho_e \frac{\omega_k}{\omega_{te}^2} \left[\omega_k - k_{\parallel} v_0 + 2 \left(\frac{\omega_k}{\omega_{te}} - \frac{v_0}{v_{te}} \right) \omega_{*j} + \omega_{*n} - \frac{\omega_{*T}}{2} \right] W_k. \quad (4)$$

If we now consider the modulation of this flux by the test wave-induced inhomogeneity, we find

$$\begin{aligned} \delta \Gamma_e \simeq \frac{T_e}{e} \sum_k k_y \rho_e \frac{\omega_k}{\omega_{te}^2} & \left(\left[\omega_k - k_{\parallel} v_0 + 2 \left(\frac{\omega_k}{\omega_{te}} - \frac{v_0}{v_{te}} \right) \omega_{*j} + \omega_{*n} - \frac{\omega_{*T}}{2} \right] \delta W_k \right. \\ & \left. + k_y \rho_e v_{te} \left[2 \left(\frac{\omega_k}{\omega_{te}} - \frac{v_0}{v_{te}} \right) \nabla_x \frac{\delta j_{\parallel}}{j_0} + \nabla_x \frac{\delta n_e}{n_e} - \frac{1}{2} \nabla_x \frac{\delta T_e}{T_e} \right] W_k \right), \end{aligned} \quad (5)$$

where $j_0 = en_e v_{te}$. In other words, the test fluctuation induces modulations in the current, density, temperature, and the (background) turbulent spectrum. In order to evaluate the latter, we consider the random refraction of $(\mathbf{k}, \omega_k)$ drift wave packets¹⁴ propagating in a medium modulated by a long-wavelength, low-frequency $[(\mathbf{q}, \Omega_q) \ll (\mathbf{k}, \omega_k)]$ test wave. The wave packet (bath constituent) propagation is most conveniently described by the radial spectral width-averaged wave kinetic equation:

$$\frac{\partial \delta N_k}{\partial t} + v_{gk,y} \frac{\partial \delta N_k}{\partial y} - \frac{\partial \delta \omega_k}{\partial y} \frac{\partial N_{k0}}{\partial k_y} = 0, \quad (6)$$

where $N_{k0}(\mathbf{x}, \mathbf{k}, t) = k_{\perp}^2 \rho_s^2 [\partial(\omega_k \epsilon_k)/\partial \omega_k] W_k = k_{\perp}^2 \rho_s^2 (1 + k_{\perp}^2 \rho_s^2) W_k/2$ is the pseudo-action density, $\mathbf{v}_{gk} = -(\omega_k \rho_s, \mathbf{k}_{\perp} + \mathbf{v}_*/2) [2\rho_s/(1 + k_{\perp}^2 \rho_s^2)]$ is the drift wave group velocity ($\mathbf{v}_* = c_s \hat{\mathbf{e}}_{\parallel} \times \nabla \ln n_e$), $\omega_k = \omega_{k0} + \delta \omega_k$, $\omega_{k0} = \omega_{*n}/(1 + k_{\perp}^2 \rho_s^2)$, and

$$\delta \omega_k(\mathbf{q}) = \frac{k_{\perp}^2 \rho_s^2}{1 + k_{\perp}^2 \rho_s^2} \mathbf{k}_{\perp} \cdot \delta \mathbf{v}_E = -i k_y \rho_e v_{te} \frac{k_{\perp}^2 \rho_s^2}{1 + k_{\perp}^2 \rho_s^2} \nabla_x \frac{e \delta \phi}{T_e} \quad (7)$$

is the modulation induced in the background fluctuation oscillation frequency by the test wave. Fourier-analyzing the wave kinetic equation, we obtain the modulation of the action spectrum as background waves propagate at their group velocity through the inhomogeneous medium:

$$\delta W_k(\mathbf{q}) = 2\pi k_y \rho_e v_{te} \frac{k_{\perp}^2 \rho_s^2}{(1 + k_{\perp}^2 \rho_s^2)^2} \delta(\Omega_q - q_y v_{gk,y}) q_y \frac{\partial N_{k0}}{\partial k_y} \nabla_x \frac{e \delta \phi}{T_e}. \quad (8)$$

Substituting Eq. (8) into Eq. (5), and plugging back into Eq. (2), we finally obtain the Fourier-analyzed Ohm's law:

$$i \left(-q_{\parallel} v_{te} + \nabla_x \cdot D_W \nabla_x \right) \frac{e \delta \phi_q}{T_e} - i q_{\parallel} v_{te} \frac{\delta p_{\parallel q}}{p_e} + \left(-i \Omega_q + \nabla_x \cdot D_J \nabla_x \right) \frac{\delta j_{\parallel q}}{j_0} + \nabla_x \cdot D_n \nabla_x \left(\frac{\delta n_q}{n_e} - \frac{\delta T_q}{2T_e} \right) = 0. \quad (9)$$

where δT_q is the fluctuating parallel temperature, and

$$D_J = 2 v_{te}^2 \sum_k k_y^2 \rho_e^2 \frac{\omega_k^2}{\omega_{te}^3} \left(1 - \frac{k_{\parallel} v_0}{\omega_k}\right) W_k \quad (10a)$$

$$D_n = v_{te}^2 \sum_k k_y^2 \rho_e^2 \frac{\omega_k}{\omega_{te}^2} W_k \quad (10b)$$

$$D_W(q_y, \Omega_q) = 2\pi v_{te}^2 \sum_k k_y^3 \rho_e^3 \frac{\omega_k}{\omega_{te}^2} \frac{k_{\perp}^2 \rho_s^2}{(1 + k_{\perp}^2 \rho_s^2)^2} \delta(\Omega_q - q_y v_{gk,y}) \\ \times \left[\omega_k - k_{\parallel} v_0 + 2 \left(\frac{\omega_k}{\omega_{te}} - \frac{v_0}{v_{te}} \right) \omega_{*j} + \omega_{*n} - \frac{\omega_{*T}}{2} \right] q_y \frac{\partial N_{k0}}{\partial k_y} \quad (10c)$$

are the hyper-resistivity (current diffusivity), particle diffusivity, and spectral diffusivity, respectively.

Equation (9) is the main result of this Letter and can be seen to have several important and surprising implications for predictions of transport of parallel electron momentum and anomalous electron dissipation. First, the collisionless hyper-resistivity is a factor $(\omega_k/\omega_{te})^3 \ll 1$ smaller than its naive fluid counterpart.¹⁰ The origin of this reduction factor becomes clear when we rewrite the hyper-resistivity in the form $D_J = \sum_k \langle |\delta \mathbf{v}_k|^2 \rangle \tau_{ck} F(\omega_k/\omega_{te})$, where $\delta \mathbf{v}_k$ is the turbulent velocity field, τ_{ck} is the decorrelation time, and $F(\omega_k/\omega_{te})$ is a factor related to parallel electron compression. The crucial point is the form of τ_{ck} and $F(\omega_k/\omega_{te})$. For collisionless electron transport, $\tau_{ck}^{-1} = \omega_{te}$ as opposed to $\Delta \omega_k \sim \omega_k \ll \omega_{te}$ in the simple fluid picture,¹⁰ i.e., the decorrelation rate is linked to the (rapid) electron transit frequency rather than the (slower) nonlinear mixing time. Furthermore, $F(\omega_k/\omega_{te}) = (\omega_k/\omega_{te})^2 \ll 1$ in the collisionless regime, reflecting the tendency of parallel electron compressibility to render the electron response quasi-adiabatic, thus drastically reducing current diffusion.

It should be noted that *simple* fluid theory is intrinsically incapable of recovering this fundamentally kinetic result. Indeed, it might be expected that proper incorporation of parallel dynamics would recover the reductive factors of ω_k/ω_{te} in the kinetic result, since the origin of the latter is in parallel dynamics. In fact, as we will presently show, parallel compression only partially recovers the missing factors. Consider the renormalization of the simple, two-field system of equations:

$$\frac{\partial \delta j_{\parallel}}{\partial t} - e v_{te}^2 \nabla_{\parallel} \delta n_e + \nabla \cdot \left(\frac{c}{B} \delta j_{\parallel} \hat{\mathbf{e}}_{\parallel} \times \nabla \delta \phi \right) = 0, \quad (11)$$

$$\frac{\partial \delta n_e}{\partial t} + \frac{1}{e} \nabla_{\parallel} \delta j_{\parallel} = 0. \quad (12)$$

Equation (11) is the current evolution equation, where the second term represents parallel compression, while the last term is the turbulent $\mathbf{E} \times \mathbf{B}$ advection of the current, and Eq. (12) is the continuity equation. Our aim in adopting this reduced description is to keep only those terms that enter explicitly into the derivation of the fluid hyper-resistivity. Other terms have been dropped in the interest of keeping the argument as transparent

as possible. To renormalize the nonlinear term, namely $(ic/B) \sum_k \hat{\mathbf{e}}_{\parallel} \cdot \mathbf{k} \times \mathbf{q} \delta\phi_k \delta j_{\parallel q}^{(2)}$, we solve for the beat fields from

$$\begin{aligned} (\omega'' + i\Delta\omega'') \delta j_{\parallel q}^{(2)} + ev_{te}^2 q_{\parallel}'' \delta n_{eq}^{(2)} &= i \frac{c}{B} \hat{\mathbf{e}}_{\parallel} \cdot \mathbf{k} \times \mathbf{q} \delta\phi_k \delta j_{\parallel q}, \\ -(\omega'' + i\Delta\omega'') \delta n_{eq}^{(2)} + \frac{1}{e} q_{\parallel}'' \delta j_{\parallel q}^{(2)} &= 0, \end{aligned}$$

to obtain

$$\begin{aligned} \delta j_{\parallel q}^{(2)} &= i \frac{c}{B} \left(\omega'' + i\Delta\omega'' + \frac{\omega_{te}''^2}{\omega'' + i\Delta\omega''} \right)^{-1} \hat{\mathbf{e}}_{\parallel} \cdot \mathbf{k} \times \mathbf{q} \delta\phi_k \delta j_{\parallel q}, \\ &\simeq i \frac{c}{B} \frac{\omega''}{\omega_{te}''^2} \hat{\mathbf{e}}_{\parallel} \cdot \mathbf{k} \times \mathbf{q} \delta\phi_k \delta j_{\parallel q}, \end{aligned}$$

where we have assumed $\omega_{te}''^2 \gg \omega'' \sim \Delta\omega''$. The renormalized current equation in Fourier space is thus:

$$-iq_{\parallel} v_{te} \frac{\delta n_q}{n_e} + \left(-i\Omega_q + \nabla_x \cdot D_J^{fluid} \nabla_x \right) \frac{\delta j_{\parallel q}}{j_0} = 0, \quad (13)$$

where

$$D_J^{fluid} = v_{te}^2 \sum_k k_y^2 \rho_e^2 \frac{\omega_k}{\omega_{te}^2} W_k. \quad (14)$$

Thus, even accounting for parallel compression in the simple fluid picture, $D_J^{fluid}/D_J \simeq \omega_{te}/\omega_k \gg 1$. Note that in the analysis presented in Ref. 10, which ignores parallel compression even in the fluid picture, the hyper-resistivity is off by an extra two powers of ω_k/ω_{te} .

A second finding of our work is that the mean field Ohm's law contains off-diagonal, spatial parity-breaking contributions associated with test wave density and temperature gradients [last set of terms in Eq. (9)] *in addition to* the diagonal current diffusion. These contributions arise from test wave-induced modulations of the *off-diagonal* pieces of the radial flux of parallel electron momentum driven by density and temperature gradients, respectively. Such contributions may be physically thought of as an effective rms diamagnetic velocity, as they induce nonlinear frequency shifts determined by the background turbulence. Although all gradients (i.e., available free energy sources) couple through stresses to drive the large-scale current, they each have different parities. Thus, the conventional wisdom of tearing and twisting parity solutions of the MHD Ohm's law appears to be vitiated in the turbulent, collisionless regime. Significantly, the hyper-resistivity is subdominant compared to both the particle and thermally-driven "current pinches" by a factor $\omega_k/\omega_{te} \ll 1$. Finally, we note that the spectral diffusivity, D_W , is primarily reactive rather than dissipative and, in this sense, is fundamentally different from the other two coefficients. Taken together, it can be seen that all these results significantly complicate expectations based on naive fluid models.¹⁰

A few other points bear mention. First, since ballistic electron motion does not affect the radial transport of perpendicular heat, radial transport of *parallel* electron heat and momentum is a factor $(\omega_k/\omega_{te})^2 \ll 1$ smaller than the radial transport of *perpendicular*

heat. Second, as mentioned earlier, magnetic fluctuations will alter these results only by taking $\hat{\phi} \rightarrow \hat{\phi} - (v_{\parallel}/c) \hat{A}_{\parallel}$, and consequently introducing a contribution proportional to $\langle \hat{b}_r \hat{p}_{\parallel} \rangle$ in Γ_e . It can readily be seen then that the magnetic contributions are reduced by a factor $\omega_k/\omega_{te} \ll 1$ relative to the electrostatic contributions. Also, while this calculation has been quasilinear in its methodology, it should be noted that considerations of more elaborate self-consistency constraints¹⁵ will act to couple radial transport to ion dissipation, thereby reducing it even further. Third, the present calculation has focused exclusively on *circulating* electrons, since it is this class of particles which supports the parallel current. However, the effect of turbulence on bootstrap processes, which requires consideration of *trapped* electrons, is beyond the scope of the present calculation. Fourth, the effects discussed here are related to the spatial *transport* of current and hence, should not be confused as being in conflict with experimental observations of neoclassical resistivity. In other words, although total helicity is conserved, it nonetheless can get redistributed.¹⁶ By the same token, our analysis pertains only to the manner in which the plasma current density gets locally *redistributed*; the *total* plasma current remains unchanged. Finally, it should be reemphasized that the calculation here refers to the *self-consistent* transport of parallel current. The problem of fast electron transport in RFP's¹⁷ is adequately described by *test-particle* transport models.^{3,4}

We conclude with a brief comment on recent work on hyper-resistive ballooning modes,¹⁰ which claims to demonstrate the possibility of a self-sustaining state of D_J -driven turbulence. Our analysis has revealed that ignoring ballistic electron motion, as done in that work, is valid only when $\tau_{ck}\omega_{te} \sim \omega_{te}/\omega_k < 1$. As pointed out earlier, this is never satisfied for drift-Alfvénic turbulence for experimentally-relevant parameters. For a drift wave-like background, this condition requires that the mode spatial scale satisfy $\Delta x < (L_s/L_n)\rho_e$. Assuming $T_e \sim 10$ keV, $B = 5$ T, and $L_s/L_n \simeq 10$ (a conservatively large estimate for typical L -mode plasmas, the focus of Ref. 10), this amounts to $\Delta x < 5 \times 10^{-2}$ cm. But then, such highly-localized modes do not overlap with each other, and hence cause very feeble transport. Hence, the assumptions made in Ref. 10 appear to be internally inconsistent.

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