

Turbulent current drive

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Abstract. The Ohm's law is modified when turbulent processes are accounted for. Besides an hyper-resistivity, already well known, pinch terms appear in the electron momentum flux. Moreover it appears that turbulence is responsible for a source term in the Ohm's law, called here turbulent current drive. Two terms contribute to this source. The first term is a residual stress in the momentum flux, while the second contribution is an electro-motive force. A non zero average parallel wave number is needed to get a finite source term. Hence a symmetry breaking mechanism must be invoked, as for ion momentum transport. $E \times B$ shear flows and turbulence intensity gradients are shown to provide similar contributions. Moreover this source term has to compete with the collision friction term (resistivity). The effect is found to be significant for a large scale turbulence in spite of an unfavorable scaling with the ratio of the electron to ion mass. Turbulent current drive appears to be a weak effect in the plasma core, but could be substantial in the plasma edge where it may produce up to 10 % of the local current density.

1. Introduction

Turbulent acceleration was proposed as a possible mechanism for driving the intrinsic rotation observed in tokamaks [1, 2, 3]. This effect bears some similarity with turbulent heating, which was proposed as a potential heat source in magnetized plasmas [4, 5]. In practice, turbulent heating (resp. turbulent acceleration) corresponds to terms in the heat transport (resp. momentum) equation that do not depend on the temperature (current) or its gradient. When polarization terms are ignored, the sum over species of these sources vanishes, so that it corresponds in fact to heat and/or momentum transfer from some of the plasma constituents to others [5, 6, 7, 8, 9, 10, 11]. Net heating or acceleration is possible when radial fluxes at the edge are finite, due to finite Larmor radius effects and boundary conditions [10]. Thus a first question is whether turbulent acceleration, which appears as an effective electromotive force in the Ohm's law, is a significant effect. Another recent development comes from the calculation of various turbulent contributions to the ion momentum flux. Besides a turbulent viscosity, it has been shown that pinch terms contribute to the flux, and also a "residual stress" that is not proportional to the velocity or its gradient (see [12, 13] for overviews). Hence a second legitimate question is whether the residual stress contributes as a drive in the generalized Ohm's law.

The turbulent electro-motive force and the residual stress lead to a process that is called here "turbulent current drive". This modification of the Ohm's law was already mentioned in the past, and was dubbed "anomalous bootstrap current" [14], "turbulent dynamo" [15] or "source-like term in the Ohm's law" [16]. The terminology "anomalous bootstrap current" has



become somewhat misleading since another effect was recently proposed, called "turbulence-driven bootstrap current" [17]. It is related to turbulent velocity scattering in the phase space and detrapping via radial transport, hence different from the processes investigated in the present work. As for ion momentum transport, a non zero average parallel wavenumber is needed to get finite turbulent electromotive force and residual stress - this is a consequence of general symmetry considerations [13]. Hence a symmetry breaking mechanism must be invoked. The final result depends sensitively on the efficiency of the mechanism at work. Moreover the source term competes with the Spitzer resistivity, so that the final result depends strongly on the collisionality. The purpose of this work is to estimate the amplitude of turbulent current drive for two symmetry breaking mechanisms, namely the effect of a shear flow and the inhomogeneity of the turbulence intensity profile. The mismatch between the mode frequency and the electron transit frequency introduces an unfavorable scaling with the ratio of the electron to the ion mass. Nevertheless the effect is significant for large scale fluctuations. Symmetry breaking due to $E \times B$ shear flow and gradient of turbulence intensity are shown to provide comparable contributions. Turbulence current drive is found to be quite small in the core, but not in the edge where the level of fluctuations is large and the gradient lengths are small.

The paper is organized as follows. In section 2, local conservation equations for heat and momentum are derived from gyrokinetic equations. Explicit expressions of turbulent heating and momentum sources are detailed in section 3. Orders of magnitude are given and discussed. A conclusion follows.

2. Energy and momentum transport equations

2.1. Magnetic equilibrium and gyrokinetic equation

The spatial coordinates are (χ, θ, φ) where χ is the poloidal magnetic flux normalized to 2π , θ is the poloidal angle and φ the toroidal angle. The Jacobian of this metric is $1/\mathbf{B} \cdot \nabla\theta$ where the unperturbed magnetic field is

$$\mathbf{B} = I\nabla\varphi + \nabla\varphi \times \nabla\chi \quad (1)$$

We consider the gyrocenter distribution function $\bar{F}(\mathbf{z})$ where $\mathbf{z} = (\chi, \theta, \varphi, v_{G\parallel}, \mu)$. Here $v_{G\parallel}$ is the gyrocenter parallel velocity and μ is the magnetic moment. To derive the conservation laws, it is convenient to use the collisionless gyrokinetic equation written in its conservative form [18, 19]

$$\partial_t \bar{F} + \frac{1}{B_{\parallel}^*} \nabla_{\mathbf{z}} \cdot (\dot{\mathbf{z}} B_{\parallel}^* \bar{F}) = 0 \quad (2)$$

where $\dot{\mathbf{z}} = d_t \mathbf{z}$ and $B_{\parallel}^* = B + \frac{m}{e} v_{G\parallel} \mathbf{b} \cdot (\nabla \times \mathbf{b})$ is the Jacobian of the gyrocenter transformation (m is the mass, e the charge and $\mathbf{b} = \mathbf{B}/B$ the unit vector along the field lines). For an electromagnetic turbulence, the gyrocenter equations of motion are

$$B_{\parallel}^* d_t \mathbf{x} = -B_{\parallel}^* [\bar{H}, \mathbf{x}] = v_{G\parallel} \mathbf{B}^* + \frac{1}{e} \mathbf{b} \times \nabla \Lambda \quad (3)$$

$$B_{\parallel}^* m d_t v_{G\parallel} = -B_{\parallel}^* [\bar{H}, v_{G\parallel}] = -\mathbf{B}^* \cdot \nabla \Lambda \quad (4)$$

where $\Lambda = \mu B + \bar{h}$,

$$\bar{h} = e(\bar{\phi} - v_{G\parallel} \bar{A}_{\parallel}), \quad (5)$$

is the gyroaverage perturbed hamiltonian, $\mathbf{B}^* = \mathbf{B} + \frac{m v_{G\parallel}}{e} \nabla \times \mathbf{b}$, and

$$\bar{H} = \frac{1}{2} m v_{G\parallel}^2 + \Lambda \quad (6)$$

is the hamiltonian, $\bar{\phi} = J \cdot \phi$ and $\bar{A}_{\parallel} = J \cdot A_{\parallel}$ are the perturbed gyro-averaged electric potential and parallel vector potential. The system is closed with a simplified Poisson equation

$$- \sum_{species} \nabla \cdot \left\{ \frac{\rho_{eq,s}}{B^2} \nabla_{\perp} \phi \right\} = e \sum_{species} n_G = \sum_{species} e \int \frac{2\pi}{m} B_{\parallel}^* d\mu dv_{G\parallel} J \cdot \bar{F} \quad (7)$$

and Ampère equation

$$-\nabla_{\perp}^2 A_{\parallel} + \sum_{species} \frac{\omega_{ps}^2}{c^2} A_{\parallel} = \mu_0 \sum_{species} e \int \frac{2\pi}{m} B_{\parallel}^* d\mu dv_{G\parallel} v_{G\parallel} J \cdot \bar{F} \quad (8)$$

where $\rho_{eq,s}$ is the unperturbed mass density, and ω_{ps} the plasma pulsation of a species 's'. We will use in the following the long wavelength limit of the gyroaverage operator

$$J \simeq 1 + \frac{1}{2} \nabla \cdot \left(\frac{m\mu}{e^2 B} \nabla_{\perp} \right) \quad (9)$$

This operator is consistent with the left hand side of the Poisson equation Eq.(7), and is self-adjoint.

2.2. Momentum and energy transport equation

These equations have been derived for the set of equations given above in refs [3, 20]. The local energy transport equation reads for each species

$$\partial_t E_K + \partial_{\chi} Q = W_E \quad (10)$$

where

$$E_K = \int d\tau^* \left(\frac{1}{2} m v_{G\parallel}^2 + \mu B \right) \bar{F} \quad (11)$$

is the kinetic energy, and $d\tau^* = B_{\parallel}^* 2\pi \frac{d\mu}{m} dv_{G\parallel} \frac{d\theta d\varphi}{\mathbf{B} \cdot \nabla \theta}$ is the phase space volume element in between two magnetic surfaces χ and $\chi + d\chi$. The heat flux Q is defined as

$$Q = \int d\tau^* \bar{H} v_G^{\chi} \bar{F} \quad (12)$$

and $v_G^{\chi} = \dot{\mathbf{z}} \cdot \nabla \chi$ is the radial projection of the gyrocenter velocity. The turbulent energy source is

$$W_E = - \int d\tau^* \bar{h} \partial_t \bar{F} \quad (13)$$

In a similar way, the toroidal momentum conservation equation reads

$$\partial_t \mathcal{L}_{\varphi} + \partial_{\chi} \Pi_{\varphi}^{\chi} + \partial_{\chi} T_{\varphi}^{\chi} = \mathcal{J}^{\chi} \quad (14)$$

where

$$\begin{aligned} \mathcal{L}_{\varphi} &= m \int d\tau^* u_{\varphi} \bar{F} \\ \Pi_{\varphi}^{\chi} &= m \int d\tau^* \bar{F} u_{\varphi} v_G^{\chi} \\ T_{\varphi}^{\chi} &= \int^{\chi} d\chi \int d\tau^* \bar{F} \partial_{\varphi} \bar{h} \\ \mathcal{J}^{\chi} &= e \int d\tau^* v_G^{\chi} \bar{F} \end{aligned} \quad (15)$$

Here $u_\varphi = \frac{I}{B} v_{G\parallel}$ is the gyrocenter toroidal angular velocity at requested order. An explicit expression of the polarization term T_φ^χ is given in [20, 3]. In most cases, it can be ignored. The local angular momentum conservation equation can then be recast as

$$\partial_t \mathcal{L}_\varphi + \partial_\chi \Pi_\varphi^\chi = R_0 M_\parallel \quad (16)$$

where the local source of toroidal momentum source is

$$M_\parallel = - \int d\tau^* \bar{F} \frac{I}{BR_0} \nabla_\parallel \bar{h} \quad (17)$$

where the parallel gradient is defined as

$$\nabla_\parallel = \frac{1}{B} \mathbf{B} \cdot \nabla = \frac{I}{BR^2} \left(\partial_\varphi + \frac{1}{q} \partial_\theta \right) \quad (18)$$

and R is the major radius (R_0 is the major radius on axis).

3. Explicit expressions of fluxes and sources

The particle conservation equation and the equations Eqs.(10,16) can be recast as follows [3]

$$\partial_t N + \partial_\chi \Gamma_N^\chi = 0 \quad (19)$$

$$Nm \partial_t V_\parallel + \partial_\chi \Gamma_V^\chi = M_\parallel \quad (20)$$

$$\frac{3}{2} N \partial_t T + \partial_\chi \Gamma_T^\chi = W_T \quad (21)$$

where N is the density, V_\parallel is the parallel fluid velocity and T the temperature. The quasilinear expressions of the fluxes and sources are [3]

$$\begin{pmatrix} W_T \\ M_\parallel \\ \Gamma_N^\chi \\ \Gamma_V^\chi \\ \Gamma_T^\chi \end{pmatrix} = \frac{T}{e} \sum_{\mathbf{n}} \int_{-\infty}^{+\infty} d\omega' \int d\tau^* F_{eq} \left| \frac{\bar{h}_{\mathbf{n}}}{T} \right|^2 \frac{\Delta\omega_{\mathbf{n}}}{(\omega' - \omega_{\mathbf{n}})^2 + \Delta\omega_{\mathbf{n}}^2} n (\omega' - \mathbf{n} \cdot \boldsymbol{\Omega}_*) \delta(\omega' - \mathbf{n} \cdot \boldsymbol{\Omega}) \begin{pmatrix} e (\omega' - k_\parallel V_\parallel) / n \\ ek_\parallel / n \\ 1 \\ mv_{G\parallel} \\ \mathcal{E} - \frac{3}{2} T \end{pmatrix} \quad (22)$$

where F_{eq} is the unperturbed distribution function (typically a Maxwellian), \mathcal{E} the particle kinetic energy, n the toroidal wavenumber. The triplets of integers \mathbf{n} refer to the Fourier development of various fields when using angle variables. Here $\omega' = \omega - n\omega_E$, and $\mathbf{n} \cdot \boldsymbol{\Omega}$ is the resonant frequency calculated in the frame of reference where the radial electric field vanishes, i.e. $\mathbf{n} \cdot \boldsymbol{\Omega} = n\omega_D + k_\parallel v_{G\parallel}$ for passing particles, and $\mathbf{n} \cdot \boldsymbol{\Omega} = n\omega_{Dt}$ for trapped particles, where ω_{Dt} is the precession frequency. The pulsation $\mathbf{n} \cdot \boldsymbol{\Omega}_*$ is an energy dependent diamagnetic pulsation, which is detailed in the next section. From now on the calculations are performed in the frame of reference where ω_E locally vanishes. This does not forbid a finite velocity shear $\partial_\chi \omega_E$, which is a source of acceleration.

3.1. Practical expressions for fluxes and sources

Let us now consider one species with mass m_s and charge e_s (for electrons $e_s = -e$). For a Maxwellian equilibrium distribution function, sources and fluxes are given by the expressions Eq.(22), which can be recast as

$$\begin{pmatrix} \frac{W_T}{N_s e_s} \\ \frac{M_{\parallel}}{N_s e_s} \\ \frac{\Gamma_N^x}{N_s} \\ \frac{\Gamma_T^x}{N_s T_s} \\ \frac{\Gamma_V^x}{N_s m_s v_{Ts}} \end{pmatrix} = -\frac{T_s}{e_s} \sum_{\mathbf{n}} \int_{-\infty}^{+\infty} d\omega \int_0^{+\infty} du e^{-u} \int_{-\infty}^{+\infty} \frac{d\zeta}{\sqrt{2\pi}} e^{-\frac{\zeta^2}{2}} n \left| \frac{\bar{h}_{\mathbf{n}}}{T} \right|^2 D_{\mathbf{n}\omega} \frac{\Delta\omega_{\mathbf{n}}}{(\omega - \omega_{\mathbf{n}})^2 + \Delta\omega_{\mathbf{n}}^2} \delta(R_{\mathbf{n}\omega}) \begin{pmatrix} \frac{\omega}{n} \\ \frac{k_{\parallel}}{n} \\ 1 \\ u + \frac{\zeta^2}{2} - \frac{3}{2} \\ \varepsilon_t \zeta \end{pmatrix}$$

where the denominator $D_{\mathbf{n}\omega}$ is

$$D_{\mathbf{n}\omega} = \omega + n \frac{T_s}{e_s} \partial_{\chi} \Xi_s - \varepsilon_t k_{\parallel} v_{Ts} \mathcal{V}_{s\parallel} \quad (23)$$

with

$$\partial_{\chi} \Xi_s = \partial_{\chi} \ln N_s + \left(u + \frac{\zeta^2}{2} - \frac{3}{2} \right) \partial_{\chi} \ln T_s + \varepsilon_t \zeta \partial_{\chi} \mathcal{V}_{s\parallel} \quad (24)$$

The resonance term $R_{\mathbf{n}\omega}$ is

$$R_{\mathbf{n}\omega} = \omega - k_{\parallel} v_{Ts} (\mathcal{V}_{s\parallel} + \zeta) - n\omega_{Ds} \left(u + (\zeta + \mathcal{V}_{s\parallel})^2 \right) \quad (25)$$

for passing particles and

$$R_{\mathbf{n}\omega} = \omega - n\omega_{Ds} u \quad (26)$$

for trapped particles, where $\mathcal{V}_{s\parallel} = V_{\parallel s}/v_{Ts}$ is the parallel fluid velocity normalized to the species thermal velocity v_{Ts} . The variable $u = \mu B/T_s$ and $\zeta = v_{G\parallel}/v_{Ts}$ are respectively the normalized perpendicular energy and parallel velocity. For trapped (resp. passing) particles, $\varepsilon_t = 0$ (resp. $\varepsilon_t = 1$), i.e. trapped particles do not carry parallel momentum. The curvature frequency is defined as $\omega_{Ds} = \frac{q}{r} \frac{T_s}{e_s B R}$, the thermal velocity is $v_{Ts} = \sqrt{T_s/m_s}$, and the Larmor radius is $\rho_s = \frac{m_s v_{Ts}}{e_s B}$. The expression of the precession frequency that is used here is valid only for deeply trapped particles, and is given to illustrate the structure of fluxes. It should be replaced by its exact expression for accurate calculations. The structure of fluxes and sources given above is close to the one derived in [16], using similar techniques.

3.2. Orders of magnitude

Current is carried by passing electrons. We further simplify the problem by ignoring trapped electrons and keeping transit resonances only. In that case the angle variables are simply the cyclotron angle, and the poloidal and toroidal angles. The force acting on electrons is

$$M_{\parallel e} = \frac{N_e T_e}{R_0} \frac{1}{\sqrt{2\pi}} \sum_{\mathbf{k}} \exp \left\{ -\frac{1}{2} \left(\frac{\omega_{\mathbf{k}}}{\omega_{te}} \right)^2 \right\} \left| \frac{\bar{h}_{\mathbf{k}}}{T_e} \right|^2 k_{\parallel} R_0 \frac{\omega_{\mathbf{k}} - \omega_{*e}}{|\omega_{te}|} \quad (27)$$

where the components of \mathbf{k} are the poloidal and toroidal wavenumbers (m, n) , $\omega_{*e} = n \frac{T_e}{e} \partial_{\chi} \Xi_e|_{\zeta=\omega_n/\omega_{te}}$, and $\omega_{te} = k_{\parallel} v_{Te}$ (n is the toroidal wavenumber).

Fluctuation measurements indicate that the spectral turbulence intensity $I_{\mathbf{k}} = |\bar{h}_{\mathbf{k}}/T_e|^2$ is maximum at low wavenumbers $k_{\perp} \rho_i = 0.1$, and decays as k_{\perp}^{-3} for $k_{\perp} \rho_i \leq 1$, faster beyond $k_{\perp} \rho_i = 1$ [21] (k_{\perp} is the perpendicular wavenumber and ρ_i the thermal ion gyroradius). For a frequency of the order of the diamagnetic frequency, the ratio $\frac{\omega_{\mathbf{k}}}{\omega_{te}}$ behaves as $\frac{q R_0}{s L_p} \frac{\rho_i}{w_k} \sqrt{\frac{m_e}{m_i}}$, where q is the safety factor, s the magnetic shear, w_k the mode width, L_p a gradient length (for simplicity all gradient lengths are assumed to be comparable). It appears that in spite of an unfavorable mass scaling, the dominant contribution comes from low wave number fluctuations, at long as quasi-linear theory is valid. An order of magnitude of the average wavenumber due to $E \times B$ shear flow is [22]

$$\bar{k}_{\parallel} R_0 = -2 \frac{q}{s} k_{\theta}^2 \rho_i^2 \frac{1}{\omega_{\mathbf{k}}} \left(1 - \frac{\omega_{\mathbf{k}}}{n \omega_{De}} \right) \frac{dV_{E\theta}}{dr} \quad (28)$$

here $\gamma_E = \frac{dV_{E\theta}}{dr}$ is the flow shear rate. An order of magnitude of the flow shear rate at low parallel Mach number is $\gamma_E \simeq \rho_i \frac{v_{Ti}}{L_p^2}$. Hence the mean parallel wave number is $\bar{k}_{\parallel} R_0 \simeq \frac{q}{s} k_{\theta} \rho_i \frac{\rho_i}{L_p}$. Another mechanism comes from the gradient of turbulence intensity [23]. In that case $\bar{k}_{\parallel} R_0 \simeq s k_{\theta} \frac{w_k^2}{L_I}$, where $\frac{1}{L_I} = \frac{d \ln(I_{\mathbf{k}})}{dr}$. Hence both effects leads to the same scaling for the mean parallel wave numbers in the gyroBohm framework $w_k \equiv \frac{1}{k_{\theta}} \equiv \rho_i$, i.e. $\bar{k}_{\parallel} R_0 \equiv \frac{\rho_i}{L_p}$. Assuming isotropic spectra $w_k \simeq \frac{1}{k_{\theta}}$, it turns out that the effect of turbulence intensity gradient is stronger at low wave numbers. The turbulent force is then approximated by

$$M_{\parallel e} \simeq \frac{N_e T_e}{L_p} \frac{q \rho_i}{L_I} \sqrt{\frac{m_e}{m_i}} I_{turb} \quad (29)$$

where I_{turb} is the turbulence intensity at the spectrum peak. The effective electric field

$$E_{eff} \simeq \frac{M_{\parallel e}}{N_e e} \simeq \frac{T_e}{e R_0} \sqrt{\frac{m_e}{m_i}} \frac{q R_0}{L_p} \frac{\rho_i}{L_I} I_{turb} \quad (30)$$

can then be compared to the inductive field, of the order of $0.1V/m$. It appears that E_{eff} is usually small in the core plasma, i.e. of the order of $10^{-3}V/m$ or less. However, this is not true in the plasma edge, where the level of fluctuations is large, and the gradient scales are small. We choose the following parameters at the pedestal height $N_e = 1.10^{19} m^{-3}$, $T_e = 100eV$, $R_0 = 1m$, $a = 0.3m$, $L_p = L_I = 0.01m$, $B = 1T$, $q = 3$, $\sqrt{\frac{m_i}{m_e}} = 60$, $I_{turb} \simeq \left| \frac{\delta N}{N_e} \right|^2 \simeq 10^{-3}$. One then finds $\frac{T_e}{e R_0} \simeq 10^2 V/m$, $\sqrt{\frac{m_e}{m_i}} \frac{q R_0}{L_p} \simeq 5$, and $\frac{\rho_i}{L_I} \simeq 0.072$. This leads to an effective electromotive field of the order of $3.610^{-2}V/m$, i.e. can be considered as significant. Another way to estimate the turbulence contribution to the current is to compare the driven current to the bootstrap current. The local bootstrap current in the edge scales as $J_{boot} \simeq 5 \sqrt{\frac{R}{a}} q \frac{N_e T_e}{B L_p}$, while

the turbulence driven current can be estimated by balancing the turbulent force $M_{\parallel e}$ with the collisional electron/ion friction force $\nu_{ei} \frac{m_e}{e} J_{turb}$ (ν_{ei} is the electron/ion collision frequency). An estimate of the turbulence driven current is

$$J_{turb} = \frac{e}{m_e \nu_{ei}} \frac{N_e T_e}{L_p} \frac{q \rho_i}{L_I} \sqrt{\frac{m_e}{m_i}} I_{turb} \quad (31)$$

so that the ratio of the turbulent to bootstrap current is

$$\frac{J_{turb}}{J_{boot}} = \frac{1}{5} \sqrt{\frac{a}{R_0}} \frac{v_{Te}}{L_I \nu_{ei}} I_{turb} \quad (32)$$

Using the parameters above, it is found that $J_{turb}/J_{boot} \simeq 0.11$, i.e. 10%.

Another turbulence contribution comes from the flux of momentum Γ_V^χ , whose expression is given by Eq.(22). Using a radial coordinate close to a minor radius and neglecting the particle flux, the momentum flux is $\Gamma_V = N_e m_e \langle \tilde{v}_{Er} \tilde{V}_{\parallel} \rangle$. The structure of this flux is close to the one found for ions in the context of plasma "intrinsic rotation" (see overviews [12, 13] and references therein). It can be written as

$$\Gamma_V = -N_e m_e D_V \partial_r V_{\parallel} + N_e m_e v V_{\parallel} + \Pi_{\parallel} \quad (33)$$

where D_V is a viscosity, v is the pinch velocity and Π_{\parallel} the residual stress. The electron viscosity D_V can also be interpreted as an hyper-resistivity since the electron parallel velocity is a good proxy of the current. This is not by far a new idea [24, 25]. The value of the turbulent viscosity is still subject to debate. It is probably in the range of $[0.1 - 1] m^2/s^{-1}$. The corresponding diffusion time is to be compared with the electron collision time. Its contribution to the Ohm's law is negligible for the current evolution at large scale, but could matter for some specific problems like MHD inertial layers. A second contribution comes from pinch terms. The expression of v is quite complex, as several effects contribute. However perpendicular compressibility often provides the dominant contribution, of the order of $v \simeq D_V/R$. If the viscosity is substantial, this effect cannot be ignored either for small scale dynamics. Finally a third contribution comes from the residual stress. An order of magnitude is

$$\Pi_{\parallel} \simeq N_e T_e \frac{L_p}{R_0} \sum_{\mathbf{k}} i k_{\parallel} R_0 \omega_{*e} \tau_{ck} \left| \frac{\bar{h}_{\mathbf{k}}}{T_e} \right|^2 \quad (34)$$

where $\tau_{ck} = 1/\Delta\omega_{\mathbf{k}}$ is a correlation time and $\Delta\omega_{\mathbf{k}} \simeq \omega_{te}$ for passing electrons. It appears that the divergence of this flux is of the same order as the source term Eq.(29).

4. Conclusion

In summary, turbulence processes may modify substantially the Ohm's law. It is already known that turbulence produces a significant hyper-resistivity. However it also appears that it is responsible for pinch terms in the momentum flux. Moreover turbulence produces a source term in the Ohm's law that is not proportional to the current or its gradient. Two terms contribute to this source. The first term is the residual stress in the momentum flux. The second one is an electromotive force, that competes with the electron-ion collision friction (Spitzer resistivity). A symmetry breaking mechanism is needed to provide a finite average parallel wave number that produces a non zero current drive. $E \times B$ shear flows and radial gradient of turbulence intensity give similar orders of magnitude. The effect is significant for a large scale turbulence in spite of an unfavorable scaling with the ratio of the electron to ion masses. Turbulent current drive appears to be a weak effect in the plasma core, but could be substantial in the plasma edge where the level of fluctuations is large and the gradient lengths are small.

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