Kinetic Simulation of Steady States of Ion Temperature Gradient Driven Turbulence with Weak Collisionality

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(Received - Nov. 11, 2003)

NIFS-785

Nov. 2003

RESEARCH REPORT
NIFS Series
Kinetic simulation of steady states of ion temperature gradient driven turbulence with weak collisionality

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(Dated: November 14, 2003)

Statistically steady states of the ion temperature gradient driven turbulence with weak collisionality are investigated by means of a kinetic simulation based on the gyrokinetic equations. Even if the collision frequency is much lower than characteristic ones of the turbulence, it is indispensable to realizing a steady state of perturbed distribution function $\delta f$. In a low-collisionality limit, the low-order velocity-space moments of $\delta f$ as well as the ion heat transport flux agree with those in the quasi-steady state of the collisionless turbulence. A spectral analysis of $\delta f$ in the velocity-space clarifies the transfer and dissipation processes of the entropy variable associated with fluctuations, where the phase mixing, the $E \times B$ nonlinearity, and the finite collisionality are taken into account. A power-law scaling predicted by the theoretical analysis is also verified by the simulations in a subrange of the power spectrum which is free from the entropy production and the collisional dissipation.

Keywords: gyrokinetics, turbulence, transport, computer simulation

PACS numbers: 52.25.Fi, 52.35.Ra, 52.65.-y

I. INTRODUCTION

Turbulent transport in high-temperature plasmas has long been a key issue in the magnetic confinement fusion research [1], since it is considered as a main cause of the anomalous transport of particles and energy. Understanding the microscopic turbulence is important as the first step to prediction and control of the anomalous transport. Extensive simulation studies [2] on drift wave turbulence, such as the ion temperature gradient (ITG) mode, have revealed several important aspects of the turbulent transport in magnetically-confined plasmas, for example, the transport suppression by the self-generated zonal flow [3]. Nevertheless, saturation mechanism of the collisionless turbulence has been an open question. Since the collisionless gyrokinetic equation has the time-reversal symmetry, one needs to consider a coarse-grained form of the one-body velocity distribution function $f$ with small-scale fluctuations in order to define an irreversible transport process in collisionless turbulence.

It has been pointed out that, when a steady transport flux is observed in the collisionless turbulence, a quasi-steady state should be realized [4-6], where high-order velocity-space moments of the perturbed distribution function $\delta f$ continue to grow but the low-order ones are constant in average. Here, $\delta f \equiv f - f_{eq}$ is deviation from the equilibrium given by the Maxwellian velocity distribution $f_{eq}$. Existence of the quasi-steady state in the collisionless ITG turbulence has been confirmed by means of an Eulerian (so-called Vlasov) numerical simulation of the gyrokinetic equation [7]. The phase mixing generates fine-scale fluctuations of $\delta f$, and leads to continuous growth of the high-order moments as well as an entropy variable associated with fluctuations (that is defined by a square-integral of $\delta f$) of which the growth rate balances with the transport flux multiplied by a normalized ion temperature gradient (see section III A in more detail).

If the collisionless assumption has a practical meaning for considering the steady anomalous transport in an actual plasma with weak but finite collisionality, the quasi-steady state should be an idealization of a real statistically steady state in a weak-collisionality limit. Namely, for sufficiently low collision frequencies, statistical behaviors of the low-order moments such as the transport flux should agree with those in the collisionless turbulence. Our concern here is, thus, to find the collisionality dependence of the kinetic ITG turbulent transport, and to examine the above conjecture on the relation between the steady and quasi-steady states. Introduction of the finite collisionality allows the system to approach the real steady state. Even if the collision frequency is much lower than characteristic ones of the ITG modes, it definitely affects evolution of the system through dissipation of the fine-scale fluctuations of $\delta f$ in the velocity space.

In simulations of the ITG turbulence shown below, we employ a two-dimensional slab model without complication of interaction between the zonal flow and the turbulence, as has been done in our previous work [7]. The Eulerian kinetic simulation with the simplified model setting enables us not only to investigate fundamental processes in plasma turbulence, such as the $E \times B$ advection, the phase mixing and the collisional dissipation, but also to give a useful reference for construction of kinetic-fluid closure models [6, 8]. This is because collisionless fluid simulations of the steady turbulence transport are based on the conjecture on existence of the quasi-steady state of turbulence. A detailed comparison between the collisionless kinetic and fluid simulations of the slab ITG turbulence has recently been carried out [9], where the transport coefficient given by the fluid simulation with the nondissipative closure model is in good agreement with the kinetic results.
In the latter part of this paper, a velocity-space power spectrum of $\delta \nu$ represented by a quadratic form of the Hermite-polynomial expansion coefficients is investigated in analogy with the passive scalar convection in the homogeneous isotropic turbulence of a neutral fluid [10]. The spectral analysis elucidates the entropy transfer process from macro to micro scales in the phase space through the phase mixing and the $\mathbf{E} \times \mathbf{B}$ nonlinearity. The entropy variable is damped by collisions in a micro velocity scale. Similarly to the viscous-convective subrange in a power spectrum of the passive scalar, we identify a subrange in the power spectrum of $\delta \nu$ which is free from the entropy production and the collisional dissipation. A scaling-law for the spectrum of the entropy variable in this subrange will also be derived in this paper.

This paper is organized as follows. Our simulation model is described in Sec. II. Simulation results are presented in Sec. III, where the statistically steady state of the weakly collisional turbulence is discussed in Sec. III A. The collision frequency dependence of the transport flux and a convergence check to mode truncation in the wave number space are given in Secs. III B and C, respectively. The spectral analysis of the distribution function is described in Sec. IV. The obtained results are summarized in Sec. V.

II. SIMULATION MODEL

Our simulation model considered here is the same as in the previous work on the collisionless ITG turbulence except for an ion-ion collision term. We consider a periodic two-dimensional slab configuration with translational symmetry in the $z$-direction, where the uniform magnetic field is set in the $y$-$z$ plane such that $\mathbf{B} = B(z + \theta y)$ with $\theta \ll 1$. Equations numerically solved here are derived from a $v_l$-integral of the gyrokinetic equations [11] by neglecting the parallel nonlinear term and by assuming $\delta v_{k}(y, v_{l}) = \tilde{f}_{k}(y, v_{l})F_{M}(v_{l})$, where $F_{M}$ denotes the Maxwellian velocity distribution. We also assume constant density and temperature gradients of the background ions in the $z$-direction with much larger scale lengths $L_{n} = -d \ln n_{i}/dz$ and $L_{T} = -d \ln T_{i}/dz$ than the fluctuation wave lengths. Therefore, we arrive at the following equations represented in the wave number space $k \equiv (k_{y}, k_{z})$ as

$$\partial_{t} \tilde{f}_{k} + i \nu v_{l} k_{y} \tilde{f}_{k} + \sum_{k' \neq k} (k'_{y} k'_{z} - k_{y} k_{z}) \Psi_{k'} \tilde{f}_{k'} =$$

$$-i k_{y} \Psi_{k} \left[ 1 + (v_{l}^{2} - 1 - k^{2}) \eta_{i}/2 + \Theta_{v_{l}} \right] F_{M}(v_{l}) + G_{t}(\tilde{f}_{k})$$

and

$$\left[ 1 - \Gamma_{o}(k^{2}) \right] \phi_{k} = e^{-k^{2}/2} \int \tilde{f}_{k}(v_{l})dv_{l} - \bar{n}_{e,k},$$

where the electric potential $\phi_{k}$ is related to $\Psi_{k}$ by $\Psi_{k} = e^{-k^{2}/2} \phi_{k}$ with $k^{2} = k_{y}^{2} + k_{z}^{2}$. The background electron temperature $T_{e} = 1$ and the adiabatic electron response are also assumed, such that $\bar{n}_{e,k} = \phi_{k}$ for $k_{y} \neq 0$. Here, for comparison between the collisionless and weakly collisional turbulence, we consider a limiting case with no zonal flow component of $k_{y} = 0$ by fixing $\tilde{f}_{k_{y}=0} = \phi_{k_{y}=0} = 0$, as has been done in our previous simulation [7] with the aim of simulating a large transport level observed in a toroidal geometry. In the slab configuration, otherwise, the turbulence is too severely suppressed by the zonal flow to cause a mean transport [7]. The assumption of no zonal flow also enables us to examine a finite-collisionality effect on turbulence without complication of the collisional damping of the zonal flow and its interplay with turbulence [12]. In addition, we neglect $k_{z} = 0$ modes of $\tilde{f}_{k}$ and $\phi_{k}$, since they are included in the background part with constant density and temperature gradients in the $z$-direction [13].

Equations (1) and (2) are normalized as follows: $\nu = x'/\nu_{i}$, $y = y'/\nu_{i}$, $v = v'/\nu_{i}$, $t = t_{0}/\nu_{i}$, $\bar{f} = f_{0}L_{n}/\nu_{i}$, and $\phi = e^{f_{0}/2} \phi_{k_{z}=0}$, where $v_{l}$, $\nu_{i}$, $\Gamma_{o}(k^{2})$, $\Omega_{l}$, $\nu_{i}$, $\eta_{i}$, and $I_{T}$ are the ion thermal velocity, the ion thermal gyro-radius, the background plasma density, the elementary charge, and the ion background ion temperature ($T_{i} = m_{i} \nu_{i}^{2}$; $m_{i}$ means the ion mass), respectively. $\nu$ means a dimensional quantity. $\eta_{i}$ is defined as $\eta_{i} = \theta_{n} L_{n}/\nu_{i}$. $\eta_{i}$ and $\Gamma_{o}(k^{2})$ are given by $\eta_{i} = L_{n}/L_{T}$ and $\Gamma_{o}(k^{2}) = \exp(-k^{2})I_{o}(k^{2})$, respectively. $I_{o}(z)$ is the 0-th modified Bessel function of $z$.

The parallel advection term on the left hand side of Eq. (1) contributes to generation of fine-scale fluctuations of $\tilde{f}_{k}$ in the velocity space, that is, the phase mixing. The instability drive is contained in the first term on the right-hand side of Eq. (1). The second term on the right-hand side denotes the ion-ion collision term for which we employ the Lenard-Bernstein model collision operator,

$$G_{t}(\tilde{f}_{k}) = \nu \partial_{v_{l}} \left[ \theta_{v_{l}} + \eta_{i} \right] \bar{f}_{k}(v_{l}),$$

with the collision frequency $\nu$ normalized by $v_{l}/L_{n}$. The collision operator in Eq. (3) makes $\tilde{f}_{k}$ approach $F_{M}$ preserving the mass. Although the momentum and the energy are not conserved, it doesn't cause a significant influence on the results in the present study.

Velocity-space derivatives in the collision term are calculated in the velocity wave number space $k$ into which $\tilde{f}_{k}(v_{l})$ is Fourier-transformed from the $v_{l}$ space discretized by a uniform grid in a range of $-v_{\text{max}} \leq v_{l} \leq v_{\text{max}}$ with $v_{\text{max}} = 10\nu_{i}$. Then, they are transformed back to the $v_{l}$ space. $\tilde{f}_{k}$ is fixed to zero at $v_{l} = \pm v_{\text{max}}$, since the fluctuation amplitude at the velocity-space boundary is negligibly small. For the collisionless case, we set $v_{\text{max}} = 5\nu_{i}$. We have employed sufficient resolution of the velocity space in accordance to the magnitude of $\nu$; finer grid spacing for $v_{l}$ is necessary for a smaller value of $\nu$. Numerical time integration is carried out by the
fourth-order Runge-Kutta-Gill method with careful convergence checks to the time step $\Delta t$ so as to keep enough accuracy, while a nondissipative time-integration scheme is employed for the collisionless simulation \cite{14, 15}. The minimum and maximum values of the wave number are set to $k_{\text{min}} = 0.1$ and $k_{\text{max}} = 3.2$, respectively, for both of the $k_x$- and $k_y$-directions with the 3/2-rule for de-aliasing in the spectral method. Results of a convergence check to $k_{\text{max}}$ are given in Sec. III C.

III. SIMULATION RESULTS

A. Steady state of weakly collisional turbulence

In the system described above, we note a balance equation of entropy variable defined by a functional, $\delta S = \sum_k \int d\nu \left| \frac{\bar{f}_k}{F_M} \right|^2 / 2$, that is,

$$\frac{d}{dt} (\delta S + W) = \eta Q_i + D,$$

which is derived from Eq.(1) by multiplying $\bar{f}_k/F_M$ (where asterisk denotes complex conjugate) and taking the velocity-space integral and summation over $k$. Here, $Q_i$, $W$, and $D$ are defined as the perpendicular ion heat flux $Q_i = \sum_k \int d\nu \left| \frac{\bar{f}_k}{F_M} \right|^2 \nu f_k / 2$, the potential energy $W = \sum_k \int d\nu \left[ 1 + (T_e/T_i)(1 - \Gamma_0) \right] \bar{f}_k^2 / 2$, and the collisional dissipation $D = \sum_k \int d\nu \fint_k \frac{\bar{f}_k}{F_M} = -\nu \sum_k \int d\nu \left[ \partial_\nu \bar{f}_k + \eta \nu \frac{\bar{f}_k^2}{F_M} \right] < 0$, respectively. It is also remarked that $\delta S$ is rewritten as $\delta S = S_M - S_m$ within the second order for $\delta f = f - F_M$, where $S_M = -\int d^2v F_M \ln f$ and $S_m = -\int d^2v f \ln f$ represent macroscopic and microscopic entropy per unit volume, respectively. ($\cdots$) means ensemble average. Here, $\delta S$ corresponds to the opposite sign of the excess entropy defined by Glansdorff and Prigogine \cite{16}.

In the collisionless system with $\nu = 0$, as we have shown in Ref.\cite{7} for the no zonal flow case, the quasisteady state characteristic by monotonical increase of $\delta S$ is realized in turbulence while keeping $W$ and $Q_i$ constant in average, that is,

$$\frac{d\delta S}{dt} = \eta Q_i,$$

where $\cdots$ indicates time averaging on a certain period longer than a characteristic time of the turbulence. On the other hand, even if $\nu$ is much smaller than inverse of the characteristic time of instabilities, introduction of the collision term may lead to statistically steady turbulence, where not only low-order moments but also the distribution function itself are statistically steady. Therefore, in the case with finite collisionality, it is expected that $\frac{d\delta S}{dt} \approx 0$ and

$$\eta Q_i \approx -D > 0.$$

In order to examine effects of the finite collisionality, we have performed several simulations for different $\nu$'s. Throughout the simulation runs shown below, we set $\eta_i = 10$ and $\Theta = 2.5$. For these parameters, in the collisionless case, the angular frequency $\omega_\nu$ and the linear growth rate $\omega_\nu$ of the most unstable mode with $k_x = 0.1$ and $k_y = 0.3$ are $\omega_\nu = -0.957$ and $\omega_\nu = 7.73 \times 10^{-2}$, respectively, of which changes due to finite values of $\nu$ are negligible in the present parameter range. $\delta S$ given by the weakly collisional simulation of the ITG turbulence for $\nu = 1.25 \times 10^{-4}$ is plotted as a function of time in Fig.1, where the collisionless simulation result is also shown as a reference. Time-evolutions of $\delta S$ for different values of $\nu$ which is changed from $1 \times 10^{-6}$ to $8 \times 10^{-3}$ are similar to that for $\nu = 1.25 \times 10^{-4}$ in Fig.1. In the quasisteady state of the collisionless turbulence, one finds the monotonical increase of $\delta S$, while the potential energy $W$ and the heat flux $Q_i$ are saturated \cite{7}. In the finite collisionality case, however, the growth of $\delta S$ ceases in the turbulence as well as $W$ and $Q_i$. It means that not only the low-order but also the high-order moments of $\delta f$ are statistically steady in the weakly collisional case, since $\delta S$ is also represented by the sum of squares of the velocity-space moments, $\delta S = \sum_k \delta S_k$ [see Eq.(8) for definition of $\delta S_k$] \cite{6}. Time-histories of each term in Eq.(4), $\delta f(\delta S)/dt$, $\delta f/W$, $\delta f/Q_i$, and $D$ for $\nu = 1.25 \times 10^{-4}$, are plotted in Fig.2, where data are running averaged for a time period of $\tau = 50$. One can see that the collisional dissipation $D$ balances with the mean transport, that is, $\eta_i Q_i \approx -D$, while $\delta S/dt \approx \delta W/dt \approx 0$. This also means that the weakly collisional turbulence is in the real statistically steady state. The high accuracy in calculation of the entropy balance is achieved by the sufficient velocity-space resolution of the Eulerian kinetic simulation.
B. Collision frequency dependence of transport

Collision frequency dependence of the ion heat transport coefficient, \( \chi_i \equiv Q_i/\eta_i \), is summarized in Fig. 3 where the time-average is taken from \( t = 1000 \) to 3000. According to the value of \( \nu \), the time step \( \Delta t \) is changed from 1/80 to 1/320 so that the numerical error in Eq.(4) should be much smaller than \( \eta_iQ_i \) and \( D \). The error bars are estimated from the standard deviation of running-averaged \( \chi_i \) for a time period \( \tau = 10 \). In a range of \( 1.25 \times 10^{-4} < \nu < 8 \times 10^{-3} \), \( \chi_i \) has a logarithmic dependence on \( \nu \). The \( \nu \)-dependence of \( \chi_i \) becomes quite weak for lower collision frequencies (\( \nu < 1.25 \times 10^{-4} \)), where \( \chi_i \) approaches a level of the collisionless one shown by a horizontal dashed line in Fig.3, that is, \( \chi_i \approx 0.36p_i^iT_i/L_n \).

From the results shown in Figs.1 to 3, it is summarized that, if \( \nu \) is small enough, then, the collision term does not influence the low-order moments of \( \delta f \) as well as the transport coefficient \( \chi_i \), while it is indispensable to realizing the statistically steady turbulence through damping of the high-order ones generated by the phase mixing. These facts agree with a concept that the quasisteady state is regarded as an idealization of the real steady state in the weak-collisionality limit [6].

C. Convergence check to mode truncation

We have also carried out simulation runs with different values of the maximum wave number \( k_{\text{max}} \) in mode truncation such as \( k_{\text{max}} = 1.6, 3.2, 4.8, 6.4, \) and 12.8, for \( \nu = 1.25 \times 10^{-4} \). The results are summarized in Fig.4 in terms of the ion heat transport coefficient \( \chi_i \). The numerical simulation is carried out up to \( t = 3000 \), and the time-averaging period is the same as those in Fig. 3. However, the simulation with \( k_{\text{max}} = 12.8 \) could only be run up to \( t = 1200 \), because it requires a large amount of computational cost. The time averaging is, thus, taken from \( t = 1000 \) to 1200, although it doesn't affect the results. One can see good convergence of \( \chi_i \) for \( k_{\text{max}} \geq 3.2 \), while \( \chi_i \) for \( k_{\text{max}} = 1.6 \) is about 30% larger than the others. One of the reasons is that, in case of \( k_{\text{max}} = 1.6 \), the potential amplitude at \( k = k_{\text{max}} \) is not sufficiently damped by the finite Larmor radius (FLR) effect that is represented as \( \exp(-k^2/2) \) in the definitions of \( \phi_k \) and \( W_k \). Convergence of the simulation results to \( k_{\text{max}} \) is also discussed in the next section.

IV. SPECTRAL ANALYSIS OF THE DISTRIBUTION FUNCTION IN THE VELOCITY SPACE

A. Theoretical framework

The parallel advection term in Eq.(1) generates large-scale fluctuations of the perturbed distribution function
\( \delta f \) in the velocity space (phase mixing). The small-scale components of \( \delta f \) are effectively damped by the collision term with the second derivative in \( \eta \). Here, we investigate a power spectrum of the velocity distribution function \( \delta f \). For the spectral analysis of \( \delta f \), it is meaningful to pay attention to the transfer of the entropy variable \( \delta S \) from macro to micro velocity scales. Using the basic equations, Eqs. (1) and (2), we obtain

\[
\frac{d}{dt} \left[ \delta S_n + \delta_n \frac{1}{2} \sum_k |\Phi_k|^2 \left( 1 - \Gamma_0(b_k) \right) \right] = J_{n-1/2} - J_{n+1/2} + \delta_n \eta Q_n - 2\nu n \delta S_n, \tag{7}
\]

where

\[
\delta S_n = \sum_k \delta S_{k,n} = \sum_k \frac{1}{2} n |\tilde{f}_{k,n}|^2, \tag{8}
\]

\[
J_{n-1/2} = \sum_k \Theta k_n n \text{Im}(\tilde{f}_{k,n} \tilde{f}_{k,n-1}), \tag{9}
\]

\[
J_{n+1/2} = \sum_k \Theta k_n n \text{Im}(\tilde{f}_{k,n} \tilde{f}_{k,n+1}), \tag{10}
\]

and \( \delta_{n,m} = \delta(n = m, 0; n \neq m) \). Here \( \tilde{f}_{k,n} \) \( (n = 0, 1, 2, \ldots) \) are defined as coefficients in the Hermite-polynomial expansion of \( f_k \),

\[
\tilde{f}_k(t) = \sum_{n=0}^{\infty} \tilde{f}_{k,n} H_n(\eta) H_n(\eta). \tag{11}
\]

In the steady state, the left-hand side of Eq.(7) vanishes. We see that \( J_{n-1/2} (J_{n+1/2}) \) represents the entropy transfer from the \( (n-1) \)th \((n \)th) to the \( n \)th \((n+1) \)th Hermite-polynomial portion. The third and fourth terms on the right-hand side of Eq.(7) represents the entropy production due to the downward turbulent heat flux in the temperature gradient and the collisional dissipation, respectively. It is important to note that, in the range \( n \geq 3 \), the entropy production rigorously disappears, which is the reason why the Hermite-polynomial expansion is employed here. A clear cutoff of the entropy production like this never occurs if we use the Fourier expansion in terms of \( \exp(i\eta \theta) \) \((-\infty < \theta < \infty)\) as basis functions.

Note that

\[
H_n(\eta) e^{-\eta^2/2} = \frac{i^n}{\sqrt{2\pi}} \int_{-\infty}^{\infty} d\eta e^{i\eta^2/2} \eta^n, \tag{12}
\]

where the function \( e^{-\eta^2/2} \eta^n \) of \( \eta \) has the maximum absolute value at \( \eta = \pm \sqrt{n} \) and is expanded around it as

\[
e^{-\eta^2/2} \eta^n \approx e^{-n/2} \eta^n [1 - (1 \pm \sqrt{n})^2]. \tag{13}
\]

Thus, the main contribution to the nth component of the Hermite expansion is from the Fourier components with \( |l/\sqrt{n} + 1| < 1/\sqrt{n} \) so that we may use the relation \( n \approx l^2 \) for \( n \gg 1 \). The inverse of Eq.(12) is given by

\[
e^{-l^2/2} \approx \sum_{n=0}^{\infty} \frac{H_n(x)}{n!} \frac{x^n}{n!}, \tag{14}
\]

The phase mixing process described in Eq.(1) causes the factor \( \exp(\gamma(t) \eta) \) in the distribution function, where \( \gamma(t) \) satisfies \( d\gamma(t)/dt = -\Theta k_v \). Then, the sign of \( \gamma(t) \) is opposite to that of \( k_v \) for large \( t \) so that we write \( \gamma(t) \approx -(k_v/k_v)(\sqrt{n(t)} \) with the order \( n(t) \) of the Hermite-polynomial expansion as a function of \( t \).

For \( n \geq 3 \) in the steady state, we find from Eq.(7) that

\[
-2\nu n \delta S_n = J_{n+1/2} - J_{n-1/2} \approx \frac{dJ_n}{dt}, \tag{15}
\]

where \( n \) is treated as a continuous variable and the finite difference is approximated by the derivative. The ratio of \( J_n \) to \( \delta S_n \) is written as

\[
\frac{J_n}{\delta S_n} = \frac{(n + 1/2)! \sum_k \Theta k_n n \text{Im}(\tilde{f}_{k,n-1/2} \tilde{f}_{k,n+1/2})}{\frac{1}{2} n! \sum_k |\tilde{f}_{k,n}|^2} \approx 2\Theta \sqrt{n} \sum_k |k_v| |\tilde{f}_{k,n}|^2 \sum_k |\tilde{f}_{k,n}|^2 \approx 2\Theta \sqrt{n} |k_v|, \tag{16}
\]

where \( (n + 1/2)!/n! = \Gamma(n + 3/2)/\Gamma(n + 1) \approx \sqrt{n} \) for large \( n \) and the averaging operator \( \langle f \rangle = \sum_k |f_{k,n}|^2 / \sum_k |f_{k,n}|^2 \) are used. In Eq.(16), we put \( \tilde{f}_{k,n-1/2} \tilde{f}_{k,n+1/2} \approx -(k_v/k_v) \tilde{f}_{k,n} \) by assuming \( n \)-dependence of the phase of \( \tilde{f}_{k,n} \) to be described mainly by Eq.(14) with \( l \approx -(k_v/k_v) \sqrt{n} \).

Now, let us examine the role of \( E \times B \) convection term. In analogy with the study by Batchelor on the spectrum of the passive scalar for wavelength smaller than the Kolmogorov scale in the large Prandtl number case [10], we postulate that, for large \( n \), \( \tilde{f}_{k,n} \) varies so rapidly that \( E \times B \) flow acting on \( \tilde{f}_{k,n} \) is regarded as a steady one which is statistically independent of \( \tilde{f}_{k,n} \). Then, the strain of the steady flow is considered to cause the exponential growth of the wave number of the convected variable. Thus, writing \( k_v(t) \propto e^{\gamma t} \) and using \( \sqrt{n} = -(k_v/k_v) \gamma t \approx \Theta \gamma t |k_v|, \) we have \( \Theta |k_v| \approx \gamma \sqrt{n} \) which is substituted into Eq.(16) to yield

\[
J_n/\delta S_n = 2\gamma n. \tag{17}
\]

Substituting Eqs.(17) into (15) and integrating it with respect to \( n \), we obtain

\[
\delta S_n = \sigma \frac{2}{\gamma n} \exp \left( \frac{-\nu n}{\gamma} \right), \tag{18}
\]

where \( \sigma \) coincides with the entropy production rate as shown by the constraint derived from Eq.(7),

\[
\sigma = 2\nu \int_0^{\infty} \eta \delta S_n \eta n \approx 2\nu \sum_n n \delta S_n = \eta Q_n. \tag{19}
\]
Thus, in the range where neither entropy production nor collisional dissipation occurs \(1 < n < \gamma/\nu\), we expect the power-law of \(\delta S_n \propto 1/n\) with \(J_n = \sigma = \text{const.}\), which is analogous to the passive scalar spectrum and its power transfer in the viscous-convective subrange.

In the above analytical treatment, \(\langle |k_\parallel| \rangle_n \propto \sqrt{n}\) increases indefinitely with \(n\). However, in the numerical simulation performed in the present study, there exists the upper limit of \(|k_\parallel|\). Even if the potential amplitude is sufficiently damped at the maximum wave number in the simulation, still \(f_{\parallel,n}\) for large \(|k_\parallel|\) and large \(n\) is continuously produced by the combination of the \(E \times B\) convection and the phase mixing process. Therefore, saturation of \(\langle |k_\parallel| \rangle_n\), with increasing \(n\) is anticipated due to the upper limit of \(|k_\parallel|\). In this case, taking \(\Theta(\langle |k_\parallel| \rangle_n) = \gamma M\) as independent of \(n\), we obtain from Eqs. (15) and (16),

\[
\frac{J_n}{\delta S_n} = 2\gamma M \sqrt{n},
\]

and

\[
\delta S_n = \frac{\sigma}{2\gamma M \sqrt{n}} \exp \left(-\frac{2 \nu n^{3/2}}{3 \gamma M}\right),
\]

where \(\sigma\) is the same as given by Eq. (19).

B. Application to the simulation results

Let us compare the theoretical analysis given above with the simulation results. For different collision frequencies \(\nu\), we have plotted \(\delta S_n\) and \(J_{n-1/2}/\sigma\) in Figs. 5 and 6, respectively. The maximum wave number is \(k_{\text{max}} = 3.2\). One can see that, \(\delta S_n\) at \(n = 1, 2,\) and \(3\) has almost the same values for different \(\nu\), and is followed by a power-law profile but with \(\nu\)-dependent exponents \(\delta S_n \propto n^{-\alpha}\) where \(\alpha(\nu) > 0\). Then, \(\delta S_n\) exponentially decays in the dissipation range, where the collision term is dominant, of which a typical value of \(n\) is scaled as \(n_0 \propto \nu^{-2/3}\) in consistent with Eq. (21). Thus, \(k_{\text{max}}\) affects the spectra for high \(n\)’s.

Profiles of \(\delta S_n\) obtained by simulations depend on \(\nu\) as well as \(k_{\text{max}}\) and seem to be described by mixture of Eqs. (18) and (21). Certainly, for sufficiently small \(\nu\) \((\leq 1.25 \times 10^{-4})\), we have confirmed \(0.5 < \alpha < 1\). If the collisional dissipation on the left-hand side of Eq. (15) is neglected, one finds that \(J_{n-1/2} \propto \text{const.}\). In correspondence, for \(\nu \leq 1.25 \times 10^{-4}\), there exists a flat profile of \(J_{n-1/2}\) where neither the entropy production nor the collisional dissipation is seen. For larger \(\nu\), the subrange with constant \(J_{n-1/2}\) ends at relatively small \(n\), which means that the entropy production region on the low \(n\) side and the dissipation range are not separated well. Thus, lower-order moments and transport are influenced by \(\nu\). Therefore, it results in the spectrum of \(\delta S_n\) with \(\alpha \geq 1\) for the relatively large values of \(\nu\) \((\nu > 1.25 \times 10^{-4})\).

It is also noteworthy that the logarithmic dependence of \(\chi_1\) on \(\nu\) is observed in Fig. 4 for the same parameter range of \(\nu\) as that discussed here.

\[
\begin{align*}
\nu=2 \times 10^{-3} & \quad \nu=10^{-3}/32 \\
\nu=10^{-3}/2 & \quad \nu=10^{-3}/128 \\
\nu=10^{-3}/8 & \quad \nu=10^{-3}/512
\end{align*}
\]

Order of Hermitian n

\[
\begin{align*}
\delta S_n & \propto 1/n \\
& \propto 1/\sqrt{n}
\end{align*}
\]

FIG. 5: Spectrum of \(\delta S_n\) versus order of the Hermite polynomial expansion \(n\) for different collision frequencies \(\nu\). The maximum wave number is \(k_{\text{max}} = 3.2\). Upper and lower dotted straight lines represent \(\propto 1/\sqrt{n}\) and \(\propto 1/n\), respectively.

\[
\begin{align*}
\nu=2 \times 10^{-3} & \quad \nu=10^{-3}/32 \\
\nu=10^{-3}/2 & \quad \nu=10^{-3}/128 \\
\nu=10^{-3}/8 & \quad \nu=10^{-3}/512
\end{align*}
\]

Order of Hermitian n

\[
\begin{align*}
J_{n-1/2} & \propto \sigma \\
& \propto \nu^{-2/3}
\end{align*}
\]

FIG. 6: Plots of \(J_{n-1/2}\) defined in Eq. (9) normalized by \(\sigma\) for different collision frequencies \(\nu\) with \(k_{\text{max}} = 3.2\).

The entropy variable \(\delta S_{k,n}\) before taking summation over \(k\) in Eq. (7) is transferred in the \((k,n)\)-space by the phase mixing as well as the \(E \times B\) advection. The latter effect can be found in profiles of \(\delta S_{k,n} \equiv \sum_k \delta S_{k,n}\), of which cross-sectional plots are shown in Fig. 7 for the case with \(\nu = 1.25 \times 10^{-4}\) and \(k_{\text{max}} = 12.8\). A steeply-peaked profile of \(\delta S_{k,n}\) at around \(k_0 = 0\) in \(n < 100\) (as shown by the upper three lines in Fig. 7) disappears with increasing \(n\). This is caused by the growth of \(\langle |k_\parallel| \rangle\) due to the strain of the \(E \times B\) flow. Then, the whole profile
gradually broadens, until the saturation of $\langle |k_y| \rangle$ occurs due to the finite $k_{\text{max}}$.

The growth and saturation of the spectrum-averaged wave number $\langle |k_y| \rangle$ defined in Eq. (16) is more clearly recognized in Fig. 8. For $k_{\text{max}} = 12.8$, $\langle |k_y| \rangle$ grows nearly in proportional to $\sqrt{n}$ from $n = 3$ to $\sim 100$, which agrees with the estimate of $\Theta |k_y| \approx \gamma \sqrt{n}$ used for derivation of Eq. (17). As $n$ increases, then, the growth of the wave number slows down due to the upper limit $k_{\text{max}}$. A similar evolution of $\langle |k_y| \rangle$ is also found for smaller $k_{\text{max}}$, although the slowdown of the wave number growth is observed at lower $n$. Even if $k_{\text{max}}$ is small, thus, the values of $\delta S_n$ at $n = 1$, $2$, and $3$ are unchanged, except for the case of $k_{\text{max}} = 1.6$ where $\langle |k_y| \rangle$ increases slower than $\sqrt{n}$.

The above result is consistent with the convergence check of the transport coefficient $\chi_1$ to $k_{\text{max}}$ given in Sec. III C. Therefore, we can conclude that the obtained $\chi_1$ as well as the entropy transfer process on the macro velocity scale but in a small $|k|$ region is insensitive to the maximum wave number of $k_{\text{max}} \geq 3.2$.

It is meaningful to estimate $\gamma$ and $\gamma_M$ from the simulation result. Let us suppose $\langle |k_y| \rangle \sim 0.2 \sqrt{n}$ according to Fig. 8, and substitute it into $\Theta |k_y| \approx \gamma \sqrt{n}$. Thus, we find $\gamma \sim 0.5$ for $\Theta = 2.5$. In the dissipation range of $n \gtrsim 10^3$, however, $\langle |k_y| \rangle \sim 6$ for $k_{\text{max}} = 12.8$ as shown in Fig. 8, and thus, $\gamma_M = \Theta |\langle k_y \rangle| \sim 15$. The spectrum $\delta S_n$ obtained by the simulation with $k_{\text{max}} = 12.8$ is, then, compared with those given by the theoretical analysis with $\gamma = 36$ in Fig. 9, where the solid, dashed, and dotted lines indicate the simulation result, Eq. (18) with $\gamma = 0.5$, and Eq. (21) with $\gamma_M = 15$, respectively. One can see that the simulation result is nearly proportional to that of Eq. (18) in a range of $3 \leq n \leq 10^4$. In this range $n \leq 10^3$, the constant factor $\gamma / 2 \gamma_M$ in Eq. (18) gives smaller values of $\delta S_n$ than that obtained by the simulation. The reason of this is understood as follows. We should recall that the constant factor is derived from the constraint $2\nu \int n \delta S_n \, dn = \sigma$ where $\sigma$ is evaluated from the simulation result. Then, in the range $n \lesssim 10^3$, the spectrum $\delta S_n$ in Eq. (18) needs to be smaller than that in the simulation in order to satisfy the constraint because, for $n \gtrsim 10^3$, the latter spectrum is smaller than the former due to the effect of finite $k_{\text{max}}$. However, in the dissipation range of $n \gtrsim 10^3$, the spectrum found in the simulation is well fitted by Eq. (21). The above results show that the transfer process of the entropy variable observed by the Eulerian kinetic simulations of the slab IGT turbulence can be well described by combining the analytical expressions in Eqs. (18) and (21) in Sec. IV A.

It is expected that, if we can employ an sufficiently high value of $k_{\text{max}}$, the simulation will reproduce the spectrum $\delta S_n$ in Eq. (18) for the whole range $n \geq 3$. 
V. CONCLUDING REMARKS

We have carried out Eulerian kinetic simulations of the slab ITG turbulence with weak collisionality, where we have employed the gyrokinetic equation (integrated for $v_i$) with the Lenard-Bernstein model collision operator and the quasi-neutrality condition. Introduction of finite collisionality enables us to find the real statistically steady state of turbulence, where not only the turbulence energy and the transport flux but also the entropy variable $\delta S = \int \delta f^2 / 2F_i dt$ are constant in average. Then, the ion heat transport flux $Q_i$ multiplied by $\eta_i$ is balanced with the collisional dissipation. It is contrast to the quasisteady state of the collisionless turbulence, where $d(\delta S)/dt$ balances with $\eta_i Q_i$ but with constant energy [7]. A parameter survey for the collision frequency $\nu$ shows the logarithmic dependence of $\chi_i$ on relatively large values of $\nu$. For sufficiently low collision frequency, however, the transport coefficient $\chi_i$ approaches a value in the collisionless case, which means that the low-order velocity-space moments of the distribution function in the quasisteady state of the collisionless turbulence agree with those in the real steady state of weakly collisional one. It is also confirmed that $\chi_i$ has a good convergence to the maximum wave number $k_{max} \geq 3.2$ employed in the simulations.

We have also made the spectral analysis of the distribution function in the velocity space by means of the Hermite-polynomial expansion with the Maxwellian weight function. The entropy variable $\delta S_m$ defined by a power spectrum of the distribution function has almost the same values at $n = 1, 2, and 3 even for different values of $\nu$, where $n$ denotes the order of the Hermite-polynomial expansion. This is consistent with the $\nu$-dependence of $\chi_i$ for the low collisionality case described above as well as the conjecture by Krommes and Hu such as ‘flux determines dissipation’ [4]. The entropy variable produced on the low $n$ side is transferred in the $(k, n)$-space by the phase mixing and the $E \times B$ nonlinearity. Then, it is dissipated by the collision term on a high $n$ side. Our theoretical analysis of the spectrum well describes how $\delta S_m$ ($n \gg 3$) depends on $n, \nu$, and $k_{max}$. In analogy with the passive scalar convection for wave length smaller than the Kolmogorov scale in the large Prandtl number case, the scaling-law, $\delta S_m \propto 1/n$, is found in the subrange of the $n$-space where neither entropy production nor collisional dissipation occurs.

Acknowledgments

Part of this work was completed when one of the authors (T.-H.W.) was an exchange scientist at the Institute for Fusion Studies, the University of Texas at Austin, under the Joint Institute for Fusion Theory, U.S.-Japan Fusion Cooperation Program. This work is partially supported by grants-in-aid of Ministry of Education, Culture, Sports, Science and Technology (No.12860497 and 14780387).


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