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**PRINCETON UNIVERSITY
PLASMA PHYSICS LABORATORY**

**THEORY DIVISION
QUARTERLY REPORT***

July 1 - September 30, 1993

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I. Comparisons of Gyrofluid and Gyrokinetic Simulations

S. E. Parker, W. Dorland, R. A. Santoro, M. A. Beer, Q. P. Liu, W. W. Lee,
and G. W. Hammett

A. Overview

Successful work over the past decade on the formulation of gyrokinetic equations¹⁻³ and the more recent development of the gyrofluid equations⁴⁻⁶ provide two practical and complementary approaches for simulating low frequency gradient driven micro-turbulence in tokamak plasmas. The development of these model equations along with the astounding advances being made in high performance computing now make large scale numerical simulations an important tool for the improved understanding of anomalous transport. These advances have motivated a collaborative effort within the magnetic fusion community termed the “Numerical Tokamak Project” (NTP). Though still in its early stages, the NTP has generated interesting physics results, including the prediction of long wavelength spectral features, and the importance of self-generated poloidal shear flows in core Ion Temperature Gradient (ITG) mode simulations.⁷⁻¹¹ An important aspect of the NTP is checking the validity of the variety of models used, along with benchmarking codes during the development phase. Of equal importance is that comparisons between gyrofluid and gyrokinetic approaches help improve our understanding of plasma turbulence since much of the existing insights are based on fluid type models and analogies with Navier-Stokes turbulence. Past comparisons of the linearized gyrofluid equations^{5,6,8,12} show good agreement with linear gyrokinetic theory. More recently we have been making systematic nonlinear comparisons as well.^{13,14}

The ITG instability is used throughout this study. We will take the electrostatic limit and assume the electrons are adiabatic. The ITG mode has long been considered a plausible candidate for explaining the anomalous ion heat diffusivity in the core of tokamak plasmas, which is far above neoclassical predictions. In this study five different codes will be compared: two gyrokinetic δf particle codes (slab¹⁵ and toroidal¹¹), two gyrofluid codes (slab¹⁴ and toroidal^{8,12}), and a gyrokinetic toroidal Vlasov code.^{16,17} Comparisons will be made using a coherent two dimensional (2D) shearless slab model (where analytic comparisons can be made), a fully turbulent three dimensional (3D) sheared slab model, and finally a 3D toroidal model.

B. 2D Shearless Slab Comparisons

The first comparison is a simple “2D bounded sheared slab” model¹⁸ that greatly simplifies the comparison of the two simulations and also allows straightforward com-

Figure I.1: Saturation level of the electrostatic potential vs. $\theta L_n / \rho_i = k_{\parallel} L_n / (k_y \rho_i)$ for 2D shearless slab comparison. Gyrofluid and gyrokinetic simulation results and theoretical predictions.

parison with analytic theory. The system is bounded in x periodic in y and uniform in z . The magnetic field is uniform and has a small tilt in the y direction, $\hat{b} = \hat{z} + \theta \hat{y}$, providing a finite (and constant) $k_{\parallel} = \theta k_y$, and we neglect terms of order $(k_{\perp} \rho_i)^2$. We emphasize that this is a simplified coherent nonlinear test problem, and the turbulent 3D simulations are qualitatively quite different, with many modes excited. Fully turbulent 3D cases will be discussed below. Figure I.1 shows the nonlinear saturation level for the two simulations, with reasonable agreement at small $k_{\parallel}$. The gyrofluid calculation, however, overestimates the saturation level at large $k_{\parallel}$ —*e.g.*, 2.6 times too large at $k_{\parallel} = \theta(L_n / \rho_i)k_y = 0.24$.

C. Turbulent 3D Slab Comparisons

Next, we compare the codes using a turbulent 3D slab test problem including the effect of magnetic shear. The system is bounded in x and periodic in y and z . The radial boundary conditions were chosen to simplify the code comparison, but not on physical grounds. The system is large enough, however, making the transport time scale long enough to allow for a reasonable quasi-steady state. The local physical parameters were taken from a TFTR L-mode shot (termed the “NTP test case parameters”).^{13,14} The $k_y = 0$, $k_z = 0$ modes were suppressed for this comparison. Figure I.2 shows the k_y wavelength spectrum from the two simulations. The k_y

Figure I.2: Comparison of k_y wavelength spectrum for the turbulent 3D sheared test case.

spectrum is summed over all k_x and k_z and time averaged. They both peak in the same range $k_y \rho_i = 0.2$ and show little activity for $k_y \rho_i \gtrsim 0.8$. However, there is a 30–40% difference in the overall saturation level of Φ . Figure I.3 shows the ion heat diffusivity as a function of magnetic shear. The ion diffusivity is time averaged and measured after a quasi-steady state is reached. Details of the time histories can be found in Ref. 14.

The overall level, as well as the scaling with magnetic shear, seem quite good.

D. 3D Toroidal Comparisons

Finally we make linear and nonlinear comparisons including the toroidal drifts and magnetic mirroring. The gyrofluid, gyrokinetic δf particle and gyrokinetic Vlasov codes simulate a small fraction of the tokamak plasma volume with the simulation domain being a flux tube using field line following coordinates.^{12,19}

There is fairly good linear agreement between the three different simulations.¹³ For the nonlinear toroidal comparisons, the $m = 0, n = 0$ mode was kept using $\delta n_e(m = 0, n = 0) = 0$, i.e. using the electron response $n_e = \Phi - \langle \Phi \rangle$.^{5,8,14,12} FLR effects in the Poisson equation, particularly the ion polarization density,² are very important for determining this component of Φ . Including this mode and the associated poloidal shear flows is important for achieving a saturated steady-state in the toroidal gyrofluid¹² and gyrokinetic field-line-following codes with adiabatic electrons.

The rms Φ and ion diffusivity are shown in Table I.1. Similar results (within a factor of two) have been reported by Waltz using a simpler gyrofluid model.⁹ Both

Figure I.3: Comparison of ion thermal diffusivity for the turbulent 3D sheared slab test case.

simulations exhibit a dominant $m = 0, n = 0$ mode. However, this mode is more pronounced in the gyrokinetic simulation, which is, in part, responsible for the higher gyrokinetic saturation level. In the future, the mechanisms for reducing the $m = 0, n = 0$ mode—*i.e.*, drift-orbit averaging—will be compared between the two models. Also, the effects of improved resolution, proper toroidal periodicity, and including the $m = 0, n \neq 0$ modes in the gyrokinetic simulation need to be explored further. This preliminary comparison shows agreement within a factor of two between the gyrofluid and gyrokinetic toroidal results, which is encouraging. However, we caution that this is only one test case (one set of physical parameters) and is not yet conclusive.

Codes	Linear ω	rms $e\Phi/T(L_n/\rho_i)$	$\chi_i L_n / (\rho_i^2 v_{ti})$
GK FLF	$-0.75 + 0.19i$	16.0	2.9
GF FLF	$-0.67 + 0.18i$	10.0	4.0
“Vlasov” FLF	$-0.61 + 0.18i$	—	—

Table I.1: Toroidal code comparison using the NTP parameters. Linear ω measured at $k_\theta \rho_i = 0.49$.

E. Discussion

Much attention to the details is necessary to get the agreement between gyrofluid and gyrokinetic simulations that was reported here. One must be careful to use the same boundary conditions and the same treatment of the $k_\theta = k_\parallel = 0$ component of Φ which corresponds to turbulence-generated sheared flows. One must also make sure that the two codes are consistent with each other—*e.g.*, are using the same factors of 2, π , etc. in all definitions. The results reported in Refs. 13 and 14 provide the first nonlinear tests of the gyrofluid models in fully turbulent 3-D systems. The results are generally encouraging, and indicate that models of phase-mixing, Landau-damping, and FLR effects in the gyrofluid equations⁵ may continue to work fairly well even in nonlinear turbulent systems. The slab codes have found excellent numerical agreement in the scaling of the turbulent heat diffusivity χ_i with magnetic shear (Fig. I.3). However, there are some puzzling areas where more work can still be done. There are $\sim 40\%$ differences in the rms Φ between the gyrofluid and gyrokinetic simulations, though similar qualitative scalings have been demonstrated. The gyrofluid saturation level was found to be higher than the gyrokinetic result in the coherent 2D slab test problem, whereas it was lower in the turbulent 3D slab case. In toroidal geometry, the field-line-following gyrofluid and gyrokinetic codes are within $\sim 50\%$ of each other in the turbulent χ_i and rms Φ , though there is more work which can be done to investigate sensitivity to small differences in filtering and boundary conditions.

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II. Recent Advances in Gyrokinetic Theory

John A. Krommes

A. Introduction

The nonlinear gyrokinetic (GK) equation (GKE)¹⁻³ is now routinely studied numerically by large particle simulation codes⁴ and other techniques⁵; it is an important tool for the study of low-frequency microinstabilities and anomalous transport in tokamaks. Gyrokinetic research continues to benefit from interactions between theory and (numerical) experiment. In this article I summarize some of the recent advances in GK theory.

Nonlinear gyrokinetics was pioneered by Frieman and Chen.¹ Lee² proposed a modified derivation that produced an equation in characteristic form that was particularly well suited to particle simulation; his work initiated an ambitious program in GK simulation that continues to this day. An important theoretical advance was made by Dubin et al.,³ who showed how to use modern Hamiltonian techniques to systematically derive a nonlinear GKE that was guaranteed to obey Liouville's theorem. The Hamiltonian procedure was refined and employed by later workers⁶⁻⁸ to derive GKEs for general magnetic geometries and finite- β fluctuations.

B. The gyrokinetic fluctuation-dissipation theorem

The Hamiltonian nature of the GKE has profound implications, especially when coupled with the particle-simulation technique. For example, the fluctuation-dissipation theorem (FDT) can be employed to predict the thermal-equilibrium noise in a GK simulation. Of course, one is ultimately interested in nonequilibrium phenomena, but knowledge of the equilibrium noise provides (1) a constraint on the minimum number of particles necessary for an effective turbulence simulation (the noise must be smaller than the expected saturation level of the turbulence), and (2) a nontrivial check on the fidelity of the code (the predicted and measured equilibrium wavenumber spectra should agree).

The equilibrium spectrum for electrostatic GK plasma was first predicted some years ago⁹ and later compared successfully with the simulations. That work introduced the concept of the *gyrokinetic vacuum*, a medium with a large dielectric constant $\epsilon_{\perp}(\mathbf{k}) \approx \omega_{pi}^2/\omega_{ci}^2 \gg 1$. This concept was extensively elaborated in Ref. 10. In brief, a GK plasma is a collection of gyrocenters placed into the GK vacuum, and is analogous to a collection of actual charged particles placed into the true vacuum (whose dielectric constant is unity). There are important similarities and differences between the two models. The collections of both particles and gyrocenters are dielectric media; both will polarize in response to an externally applied potential, but they will do so in different ways: the dielectric response function $\mathcal{D}_P(\mathbf{k}, \omega)$ of the true plasma is different from the analogous function $\mathcal{D}_G(\mathbf{k}, \omega)$ of the GK plasma—and both differ from the relative GK permittivity $\epsilon_{\perp}(\mathbf{k})$. This stems from the fact that the motions of the particles and gyrocenters differ. For constant magnetic field $\mathbf{B}$, the true particles have four types of motion: motion along $\mathbf{B}$, gyration about $\mathbf{B}$, $\mathbf{E} \times \mathbf{B}$ drift, and polarization drift. The gyrocenters have only two types of motion: parallel motion, and $\mathbf{E} \times \mathbf{B}$ drift. Of course, gyration has been averaged away in GKs. Notably, the gyrocenters do *not* undergo polarization drift; that effect has been subsumed into $\epsilon_{\perp}$ —i.e., into the properties of the GK vacuum. This separation of the $\mathbf{E} \times \mathbf{B}$ - and polarization-drift effects is one of the most important properties of the GK plasma, as it substantially expedites numerical solutions of the equations.

The electrostatic GK wavenumber spectrum is

$$\frac{1}{8\pi} \langle |\delta \mathbf{E}_\perp|^2 \rangle(\mathbf{k}) \approx \frac{1}{\varepsilon_\perp(\mathbf{k})} \frac{T/2}{1 + k_\perp^2 \rho_s^2}. \quad (\text{II.1})$$

This bears a marked resemblance to the well-known many-particle thermal spectrum and demonstrates, through the factor of $\varepsilon_\perp^{-1}$, that GK discreteness-induced fluctuations are small.⁹

With recent emphasis on finite- β effects, it became important to generalize Eq. (II.1) to the weakly electromagnetic regime. This requires a matrix form of the FDT, which introduces various complications and subtleties. The calculations are most conveniently done by introducing a covariant formalism¹¹ in which φ and $A_\parallel$ form the components of a two-vector. Short surveys of the calculations can be found in Refs. 12 and 13. The details are fully presented in Ref. 10; the results are tabulated in Table IV of that reference.

C. Generalized weighting schemes for particle simulation

The algorithms used for GK particle simulation have undergone considerable evolution. Currently the most popular one is due to Parker and Lee (PL).¹⁴ It is a “non-linearly weighted δf method.” In a δf method, only the fluctuating distribution δf is simulated; the form of the “background” distribution f_0 is inserted analytically (it is often taken to be a Maxwellian). In the derivation of PL, the perturbed distribution was written as

$$\delta f(\mathbf{z}, t) \approx \delta \hat{f}(\mathbf{z}, t) \doteq \frac{1}{\bar{n}_s} \sum_{i=1}^{N_s} \tilde{w}_i(t) \delta(\mathbf{z} - \tilde{\mathbf{z}}_i(t)). \quad (\text{II.2})$$

This “weighted Klimontovich” form introduces a particle weight $w(t)$ that is appropriately evolved along with the phase-space coordinates of the particles. The equation used by PL was

$$\frac{d}{dt} w_i(t) = -(1 - w) \mathbf{V}_E \cdot \nabla \ln f_0, \quad (\text{II.3})$$

where the gradient on the right-hand side introduces the density and temperature scale lengths. The factor of w on the right-hand side generalizes the earlier partially linearized algorithm of Dimits and Lee.¹⁵

In many, although not all cases, the PL algorithm is correct. In their original derivation, several points of confusion arose: (1) the same distribution function was variously represented as either smooth or singular; and (2) inhomogeneous sources were treated inconsistently. Recently¹⁶ the derivation has been generalized¹⁷ and placed on a sounder basis by reworking the formalism in terms of a phase-space sampling technique (see also Ref. 18). In this approach a careful distinction is drawn between smooth functions of phase space and Klimontovich representations. The

fundamental weighting algorithm is best derived without reference to Klimontovich functions. However, it was also shown that a Klimontovich function can be usefully employed in the derivation of an analytic formula for the sampling noise in a δf simulation. The calculation¹⁷ of that noise is the first analytic attempt for δf methods.

One interesting byproduct of the new approach is that the role and proper treatment of sources becomes much clearer. Indeed, it was found that if the original kinetic equation to be solved is $df/dt = S$, then the most general weighting algorithm involves two independent weights w and p that obey

$$\frac{d}{dt}w_i(t) = -D_i(t)p_i(t) + \frac{S_i(t)}{F_i(t)}, \quad \frac{d}{dt}p_i(t) = D_i(t)p_i(t), \quad (\text{II.4})$$

where

$$D(t) \doteq \frac{D}{Dt} \ln f_0 + \text{div} \mathbf{V}, \quad (\text{II.5})$$

$\mathbf{V}$ is the phase-space flow velocity, which may be compressible, and F is the probability distribution of the marker trajectories. For $S = 0$, one can eliminate p to find $p = 1 - w$, thereby recovering the algorithm of PL.

D. The entropy paradox and steady-state transport

I now turn to issues related to the analytic prediction of the steady-state fluxes $\bar{\Gamma}$. Statistical closure theories are by now relatively well understood,¹⁹ although their solutions for practical GK problems require much further work (currently in progress²⁰). Another approach that has been employed^{21,22} is the use of steady-state balance equations. In particular, if one introduces a small-amplitude approximation to the information-theoretic entropy,

$$\bar{S} \approx -V^{-1} \int dz \frac{1}{2} \delta f^2(\mathbf{z}, t) / f_0(\mathbf{z}, t). \quad (\text{II.6})$$

one finds that

$$\frac{d}{dt} \bar{S} = -\bar{\Gamma} \kappa + \text{other terms}. \quad (\text{II.7})$$

In certain models of completely collisionless gyrokinetics, the “other terms” have been argued to be negligible. This leads to an apparent relationship between the rate of entropy production and the flux that has been numerically verified in some cases.²²

However, such a balance raises two immediate questions: (1) How does one interpret the minus sign in Eq. (II.7)? (2) How can one be in steady state if the entropy is changing in time? Collectively, one may call these issues the “entropy paradox.”

A recent analysis²³ addresses the entropy paradox from first principles and shows how to resolve it. An exactly soluble model problem illustrates the general discussion.

The minus sign can be understood by noting that the information-theoretic entropy is defined relative to a reference state. When the reference state is taken to be the background distribution f_0 , $\bar{S}$ is maximized in that state, and any process that generates fluctuations will reduce $\bar{S}$. The question has been discussed thoroughly by Jaynes.²⁴ The issue of steady state is more problematical. Strictly speaking, there can be no steady state in the absence of (collisional) dissipation, which I parametrize here by μ . It is shown in Ref. 23 that one must be very cautious about the interchange of the limits $t \rightarrow \infty$, then $\mu \rightarrow 0$. When μ is first set to zero, one finds a time-evolving entropy; however, for small but finite μ one finds true steady states in which $\dot{\bar{S}} = 0$ and the flux is in balance with the “other terms,” which include collisional dissipation. As discussed in detail in Ref. 23, this does *not* mean that the flux must scale with μ ; in fact, typically fluxes depend only weakly on μ , if at all.

E. Summary

In summary, gyrokinetic theory is healthy and continues to evolve. Theoretically, use has been made of modern Hamiltonian methods,³ the fluctuation–dissipation theorem,¹⁰ phase-space sampling methods and Klimontovich kinetic theory,¹⁷ and insights from nonlinear statistical dynamics regarding the nature of forced, dissipative steady states.²³ The practical consequences have been systematically derived nonlinear gyrokinetic equations that (1) obey Liouville’s theorem; (2) are readily simulated by low-noise, nonlinearly weighted algorithms; and (3) predict steady-state fluxes that are being analyzed by a variety of standard techniques drawn from statistical turbulence theory and nonlinear dynamics. This analysis is far from finished. A complete understanding of the physical implications of nonlinear gyrokinetics will require the cooperative efforts of many people over a long period of time. However, it is comforting to know that the foundations of the gyrokinetic method are firm.

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III. Theory Publication List for July 1993 – September 1993

Cheng, C.Z. and Q. Qian, *Theory of Ballooning Mirror Instabilities for Anisotropic Pressure Plasmas in the Magnetosphere*, submitted to J. of Geophys. Res. and PPPL Report 2935 (September, 1993).

Fu, G.Y., C.Z. Cheng, and K.L. Wong, *Stability of the Toroidicity-Induced Alfvén Eigenmode in Axisymmetric Toroidal Equilibria*, Phys. Fluids B 5, 4040 (1993) and PPPL Report 2938 (September, 1993).

Schivell, J., D.A. Monticello, N. Fisch, and J.-M. Rax, *Modeling of MeV Alpha Particle Energy Transfer to Lower Hybrid Waves*, submitted to Plasma Phys. and Control. Fus. and PPPL Report 2943 (October, 1993).

Yamada, M., F. Levinton, J. Manickam, Y. Nagayama, and N. Pomphrey, *Experimental Study of Magnetic Reconnection During a Sawtooth Crash in a High Temperature Tokamak Plasma*, submitted to Phys. Rev. Letts. and PPPL Report 2952 (January 1994).

IV. Theory Visitors for July 93 – September 1993

Chen, Liu, University of California at Irvine, Dept. of Physics, Irvine CA, September 13 – 23, 1993.

Drozdov, Vladimir, Keldysh Institute of Applied Mathematics, Moscow, Russia, September 15 – June 15, 1994.

Gorelenkov, Nikolai, Troitsk Moscow Region, Russia, September 1 - June 1994.

Kim, Chang-Bae, Soong Sil University, Seoul, Korea, August 2, 1993.

Maurice, Sylvestre, Universite Sabatier, Observatoire Midi-Pyrenees, Laboratoire D'Astrophysique de Toulouse, Toulouse, France, July 11 – August 9, and August 30 – September 3, 1993.

Penn, Greg, MIT, Cambridge, MA, July 13 – August 13, 1994.

Rayevski, Dimitri, Courant Institute, New York University, July 8, September 22, 23, 1993.

Salas, Angel, CIEMAT, Madrid, Spain, June 17 – August 13, 1993.

Tessarotto, Massimo, University of Trieste, Dept. of Scientific Mathematics, Trieste, Italy, September 1 – October 20, 1993.

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