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Numerical Studies of Turbulence in LAPD

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Table of Contents

List of Figures

List of Tables

Acknowledgments

To be completed . . .

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CHAPTER 1

Introduction

CHAPTER 2

Turbulence and Instability

2.1 The Kolmogorov Paradigm of Turbulence

2.2 The Standard Plasma Paradigm of Linear Instability

2.3 Nonlinear Stability Effects

2.4 Nonlinear Stability Effects in Plasma Physics

CHAPTER 3

The Braginskii Fluid Model and LAPD

3.1 LAPD Suitability to the Braginskii Fluid Model

At a basic level, the state of a plasma is described by seven-dimensional distribution functions $f_j(\mathbf{x}, \mathbf{v}, t)$ for each species j. The behavior of the plasma is described by the system of kinetic equations (Boltzmann equations), which evolve the distribution functions forward in time:

$$\frac{\partial f_j}{\partial t} + \mathbf{v} \cdot \nabla f_j + \frac{e_j}{m_j}(\mathbf{E} + \mathbf{v} \times \mathbf{B}) \cdot \frac{\partial f_j}{\partial \mathbf{v}} = \left(\frac{\partial f_j}{\partial t}\right)_C . \tag{3.1}$$

$\left(\frac{\partial f_j}{\partial t}\right)_C$ is the change in the distribution function due to collisions. For plasmas, the collisions are Coulomb collisions, and the collision term takes the form of the Fokker-Planck operator. With this operator, Eq. 3.1 is called the Fokker-Planck equation. Now it is well known that the Fokker-Planck equation cannot be solved numerically for problems that require time intervals much larger than the electron-cyclotron time due to computational and time limitations. The phase space is just too large. Therefore, reduced equations, such as gyrokinetic, drift kinetic, or fluid equations have been derived to produce numerically tractable equations. These equations are all derived under certain physical assumptions such as strong guiding magnetic fields, small fluctuation levels, or slow spatial and/or time variations such that these different equations are best applied to different physical situations.

The equations that are arguably most suitable to describe waves and turbulence in LAPD (and fastest to solve numerically) are the fluid equations, specifically those derived by Braginskii [Bra65]. In deriving his equations, Braginskii approximates the solution as $f_j = f_j^0 + f_j^1$ where the zero-order piece f_j^0 is a Maxwellian and the first-order piece f_j^1 is a perturbation on the zero-order distribution function: $|f_j^1| \ll f_j^0$. The equations are then derived by taking moments of the Fokker-Planck equation to create coupled equations of the independent variables, n_j, $\mathbf{v_j}$, and T_j. Now certain requirements must hold to justify the Braginskii approximation, all of which have the flavor that macroscopic quantities must vary slowly in time and space. This is generally caused by strong relaxation processes such as collisions, which keep the distribution functions close to Maxwellians. In general, for the Braginskii equations to be applicable, processes of interest must occur on time intervals much greater than the collision time and quantities should vary slowly over distances traversed by the particles between collisions.

Specifically, the requirement that time variations must be slow can be written $\frac{d}{dt} \ll \nu$, where for electron drift wave turbulence, this is approximately $\omega_* \ll \nu_e$. Table 4.5, which displays typical LAPD operating parameters, shows that $\omega_*/\nu_e \sim 0.01$. The requirement that spatial quantities vary slowly compared to the collisional mean free path can be written simply for the direction parallel to the magnetic field as $\lambda_{ei} \sim \lambda_{ee} \ll L_\parallel$. For LAPD, $\lambda_{ei}/L_\parallel \sim 0.01$. For the direction perpendicular to the magnetic field, the same kind of relation $\lambda_{mfp} \ll L_\perp$ must also hold. However, due to the cyclotron motion of particles around the magnetic field, λ_{mfp} is really the larmor radius, unless the collisional mean free path is less than the larmor radius. For electrons, $\rho_e \ll \lambda_{ei}$ and $\rho_e/L_\perp \sim 10^{-4}$ where $L_\perp \sim 0.1m$. For the ions, the ion cyclotron frequency is close to the ion collision frequency, meaning that either the ion larmor radius or the ion mean free path may be used. Using the larmor radius, $\rho_i/L_\perp \sim 0.01$. Therefore, the collisionality is high enough and the machine dimensions are large enough so that the Braginskii fluid model should be applicable to LAPD.

3.2 The Braginskii Equations

The Braginskii fluid equations are as follows: the continuity equation for species j, electrons or ions, is [Wes04, Bra65]

$$\frac{\partial n_j}{\partial t} = -\nabla \cdot (n_j \mathbf{v}_j). \tag{3.2}$$

The momentum balance equation is

$$n_j m_j \frac{d\mathbf{v}_j}{dt} = -\nabla p_j - \frac{\partial \Pi_{j\alpha\beta}}{\partial x_\beta} + n_j e_j (\mathbf{E} + \mathbf{v_j} \times \mathbf{B}) + \mathbf{R}_j. \tag{3.3}$$

$\Pi_{j\alpha\beta}$ is the stress tensor, which involves the products of viscosity coefficients and rate-of-strain tensor components. The viscosity coefficients are some of the several terms that are called transport coefficients. The transport coefficients are calculated by the Braginskii procedure in terms of n, $\mathbf{v}$, and T. $\mathbf{R}_j$, which involves several other transport coefficients, is the rate of collional momentum transfer. Note that $\mathbf{R}_i = -\mathbf{R}_e$ since collisions don't change the total plasma momentum. The momentum transfer from ions to electrons is given by

$$\mathbf{R}_e = -m_e n_e \nu_e (0.51 u_{\parallel e} + \mathbf{u}_{\perp e}) - 0.71 n_e \nabla_\parallel T_e - \frac{3}{2} \frac{n_e \nu_e}{\omega_{ce}} \mathbf{b} \times \nabla T_e \tag{3.4}$$

where $\mathbf{u} = \mathbf{v}_e - \mathbf{v}_i$ and ν_e is the electron collision frequency. $\mathbf{R}_e$ includes both the friction force and the thermal force, which, like the friction force, is due to electron-ion collisions, but originates from the temperature dependence of the collisionality. $\mathbf{R}_i$ is actually more than just the opposite of $\mathbf{R}_e$ because of the affect of collisions with neutrals, which is only important for the ions [PUC10]. So

$$\mathbf{R}_i = -\mathbf{R}_e - n_i m_i \nu_{in} \mathbf{v_i}. \tag{3.5}$$

The energy balance equation is

$$\frac{3}{2}n_j\frac{\partial T_j}{\partial t} = -n\mathbf{v}_j \cdot \nabla T_j - p_j\nabla \cdot \mathbf{v}_j - \nabla \cdot \mathbf{q}_j - \Pi_{j\alpha\beta}\frac{\partial v_j\alpha}{\partial x_\beta} + Q_j \tag{3.6}$$

where the term involving the stress tensor describes viscous heating. The electron heat flux (with more transport coefficients) is

$$q_e = n_eT_e\left(0.71u_\parallel + \frac{3\nu_e}{2\omega_{ce}}\mathbf{b}\times\mathbf{u}\right)+\frac{n_eT_e}{m_e\nu_e}\left(-3.16\nabla_\parallel T_e - \frac{4.66\nu_e^2}{\omega_{ce}^2}\nabla_\perp T_e - \frac{5\nu_e}{2\omega_{ce}}\mathbf{b}\times\nabla T_e\right) \tag{3.7}$$

where the first part of this expression constitutes convection, while the second part is conduction. The ion heat flux is

$$q_i = \frac{n_iT_i}{m_i\nu_i}\left(-3.9\nabla_\parallel T_i - \frac{2\nu_i^2}{\omega_{ci}^2}\nabla_\perp T_i - \frac{5\nu_i}{2\omega_{ci}}\mathbf{b}\times\nabla T_i\right). \tag{3.8}$$

The last transport coefficients are in the heating Q. The ion heating due to collisional heat exchange between ions and electrons is

$$Q_i = \frac{3m_e}{m_i}n_e\nu_e(T_e - T_i) \tag{3.9}$$

while the electron heating is

$$Q_e = -\mathbf{R}\cdot\mathbf{u} - Q_i. \tag{3.10}$$

The electron heat exchange involves an ohmic heating contribution ($\mathbf{R}\cdot\mathbf{u}$) that is absent from the ion heating because electrons colliding with ions transfer very little momentum to the ions.

3.3 The Vorticity Equation

Now the Braginskii equations in the previous section contain electric and magnetic fields which must be self-consistently determined by the charges and currents that are evolved by the equations. This is done with the inclusion of Maxwell's equations. Two of those equations are used to write the fields in terms of potentials:

$$\begin{aligned}\mathbf{E} &= -\nabla\phi - \frac{\partial \mathbf{A}}{\partial t} \\ \mathbf{B} &= \nabla \times \mathbf{A}.\end{aligned} \tag{3.11}$$

The next equation,

$$\nabla_\perp^2 \mathbf{A} = -\mu_0 \mathbf{j} \tag{3.12}$$

is used to relate the vector potential to the current, where the displacement current is neglected as is generally done in plasmas. The Poisson equation is not that useful for the main part of the plasma, in which the quasineutrality relation, $n_e = n_i \equiv n$, holds. The useful equation that can be used instead is the conservation of charge (or ambipolarity condition), $\nabla \cdot \mathbf{j} = 0$. The vorticity equation is derived from this conservation of charge equation.

The current is $\mathbf{j} = en(v_{\|i} - v_{\|e}) + en(\mathbf{v}_{\perp i} - \mathbf{v}_{\perp e})$. In LAPD, the parallel current is carried primarily by the fast streaming electrons, while the perpendicular current is primarily carried by the ions, which have larger Larmor radii. So the conservation of charge equation can be simplified to

$$\nabla_\| (n v_{\|e}) = \nabla_\perp \cdot (n \mathbf{v}_{\perp i}). \tag{3.13}$$

The perpendicular ion component of this equation is derived from Eq. 3.3 for the ions. Neglecting terms that have finite ion temperature (pressure and stress tensor), and solving for the ion velocity in the Lorentz force term, the perpendicular ion velocity has two terms [PUC10, SC03]:

$$\mathbf{v}_{\perp i} = \mathbf{v}_{pi} + \mathbf{v}_{\nu i} \tag{3.14}$$

where the the polarization velocity is $\mathbf{v}_{pi} = 1/\omega_{ci}\mathbf{b} \times (\partial_t + \mathbf{v}_i \cdot \nabla)\mathbf{v}_i$ and the Pedersen velocity is $\mathbf{v}_{\nu i} = \nu_{in}/\omega_{ci}\mathbf{b} \times \mathbf{v}_i$. All other terms have equal and opposite current contributions from the electron current. So,

$$\nabla_\parallel (n v_{\parallel e}) = 1/\omega_{ci} \nabla_\perp \cdot \left[n\mathbf{b} \times (\partial_t + \mathbf{v}_i \cdot \nabla + \nu_{in})\mathbf{v}_i \right]. \tag{3.15}$$

We now employ the approximation that $\nabla \cdot (n\mathbf{v}_i \sim \mathbf{v}_E \cdot \nabla n$ where $\mathbf{v}_E = \mathbf{E} \times \mathbf{B}/B^2 = -\nabla_\perp \phi \times \mathbf{B}/B^2$. Then,

$$\begin{aligned} \nabla_\parallel (n v_{\parallel e}) &= 1/\omega_{ci} \nabla_\perp \cdot \left[n\mathbf{b} \times (\partial_t + \mathbf{v}_E \cdot \nabla + \nu_{in})\mathbf{v}_E \right] \\ \nabla_\parallel (n v_{\parallel e}) &= -\frac{m_i}{B^2} \nabla_\perp \cdot \left[\times (\partial_t + \mathbf{v}_E \cdot \nabla + \nu_{in}) \nabla_\perp \phi \right]. \end{aligned} \tag{3.16}$$

Next, defining the vorticity as $\varpi \equiv \nabla_\perp \cdot (n \nabla_\perp \phi)$, the vorticity equation reads,

$$\frac{\partial \varpi}{\partial t} = -\mathbf{v}_E \cdot \nabla_\perp \varpi - \nabla_\perp \mathbf{v}_E : \nabla_\perp (n \nabla_\perp \phi) - \frac{B^2}{m_i} \nabla_\parallel (n v_{\parallel e}) - \nu_{in} \varpi. \tag{3.17}$$

Finally, the term with the tensor product can be rewritten in a different form [PUC10]:

$$\frac{\partial \varpi}{\partial t} = -\mathbf{v}_E \cdot \nabla_\perp \varpi + \frac{1}{2} (\mathbf{b} \times \nabla_\perp n) \cdot \nabla_\perp \mathbf{v}_E^2 - \frac{B^2}{m_i} \nabla_\parallel (n v_{\parallel e}) - \nu_{in} \varpi. \tag{3.18}$$

3.4 Minimizing the Equation Set for LAPD Parameters

3.4.1 The Reduced Equations

The continuity equations 3.2 for electrons and ions do not have to both be used due to the quasineutrality condition $n_e = n_i \equiv n$. So, if one focuses on the electron continuity equation, then,

$$\frac{\partial n}{\partial t} = -\nabla \cdot (n\mathbf{v}_e). \tag{3.19}$$

Now, $\mathbf{v}_e = \mathbf{v}_{\perp e} + v_{\parallel e}$, where $\mathbf{v}_{\perp e} = \mathbf{v}_E + \mathbf{v}_{de} + \mathbf{v}_{pe}$, the terms being the $\mathbf{E} \times \mathbf{B}$ velocity, the diamagnetic velocity, and the polarization velocity. To a good approximation, $\nabla \cdot (n\mathbf{v}_{\perp e}) = \mathbf{v}_E \cdot \nabla n$ [PUC10, SC03]. So, the continuity equation reads

$$\frac{\partial n}{\partial t} = -\mathbf{v}_E \cdot \nabla n - \nabla_{\parallel}(nv_{\parallel e}). \tag{3.20}$$

Next, the momentum equations 3.3, of which there are six (three for electron velocity components and three for ion velocity components) are reduced to two here. The first is the vorticity equation 3.18, in which we used the perpendicular momentum equations to derive it. The second is the equation for the parallel electron momentum. We neglect the parallel ion momentum equation since $v_{\parallel e} \gg v_{\parallel i}$ for LAPD. The electron parallel momentum equation is then

$$nm_e \frac{\partial v_{\parallel e}}{\partial t} = -nm_e \mathbf{v}_E \cdot \nabla v_{\parallel e} - \nabla_{\parallel} p_e - enE_{\parallel} - 0.71n\nabla_{\parallel} T_e - 0.51 m_e n \nu_e v_{\parallel e}, \tag{3.21}$$

where the viscous terms have been neglected. The conservation of energy equations 3.6 are left. Since the ion temperature in LAPD is very low ($T_i \leq 1$ eV), the ion energy equation is neglected. The electron energy equation is [SC03]

$$\frac{3}{2}n\frac{\partial T_e}{\partial t} = -\frac{3}{2}n\mathbf{v}_E\cdot\nabla T_e - p_e\nabla_\| v_{\|e} + 0.71 T_e\nabla\cdot(nv_{\|e}) + \nabla_\|(\kappa_{\|e}\nabla_\| T_e) + 0.51 m_e n\nu_e v_{\|e}^2 - 3\frac{m_e}{m_i}n\nu_e T_e, \tag{3.22}$$

where $\kappa_{\|e} = 3.16\frac{nT_e}{m_e\nu e}$.

3.4.2 The Electrostatic Justification

CHAPTER 4

LAPD Simulation Details

4.1 The Equations

4.2 Finite Difference Schemes

4.3 Sources

4.4 Boundary Conditions

4.5 Profiles and Parameters

Species	^{4}He
Z	1
n	2.5×10^{18} m^{-3}
T_e	5 eV
T_i	$\lesssim 1$ eV
B_0	0.1 T
$L_\parallel$	17 m
a	0.4 m
λ_D	10^{-5} m
ω_{ci}	2.4×10^6 rad/s
ω_{ce}	1.8×10^{10} rad/s
ρ_e	5.3×10^{-5} m
ρ_i	$\sim 1 \times 10^{-3}$ m
ρ_s	5×10^{-3} m
v_{te}	9.4×10^5 m/s
c_s	1.1×10^4 m/s
v_A	7×10^5 m/s
β	5×10^{-4}
m_e/m_i	$1.4 \times 10{-}4$
$\ln\Lambda$	11
ν_e	7.2×10^6 Hz
λ_{ei}	0.13 m
ν_i	$\sim 10^6$ Hz
ν_{in}	1.2×10^3 Hz
$\kappa_\parallel^e$	9.8×10^{23} eV/m^2 s
η_0^i	$\sim 10^{12}$ eV s/m^3
ω_*	$\sim 5 \times 10^4$ rad/s

Table 4.1: Typical LAPD parameters

CHAPTER 5

Linear Instabilities

5.1 Drift Waves

5.2 Conducting Wall Mode

In this section, we consider the linear instability caused by a plasma bounded by two conducting walls on the boundaries where the magnetic field lines terminate (the axial boundaries). The instability is dependent upon Bohm sheath boundary conditions described in the following subsection. A point to note is that these boundary conditions are not necessarily the correct ones for LAPD. In tokamaks, the scrape-off-layer (SOL) can be characterized by a number of different regimes such as the sheath-limited, conduction-limited, or detached divertor regimes [Sta00]. The primary factor that controls which regime the SOL is in is the dimensionless collisionality $\nu^* \sim L/\lambda$ where L is the SOL length and λ is the electron or ion collision length. For LAPD, $\nu^* \sim 100$, which would put it in the detached regime. The Bohm sheath boundary condition is derived for low collisionality (the sheath-limited regime), so LAPD probably does not contain such a boundary condition. Yet, it is still academically instructive to apply such a boundary condition to LAPD because it creates a new linear instability, which can be used to test the robustness of LAPD's nonlinear instability.

5.2.1 The Bohm Sheath Boundary Condition

It is known that to good approximation, a plasma bounded by a wall can be divided into two regions: the main plasma and the Debye sheath [Sta00]. The Debye sheath is a small region adjacent to the wall, generally several Debye lengths long. It has a net positive charge ($n_i > n_e$) that shields the negative charge on the wall and serves to deflect some of the electrons that flow into the sheath. The sheath does not completely shield the negative wall, however, and a small electric field penetrates into the main plasma (the ambipolar field), which mostly serves to accelerate the cold ions toward the wall, and slightly retard the electrons before entering the sheath. In the main plasma, the quasineutrality relation holds ($n_i = n_e$).

The well-known Bohm criterion along with other considerations restricts the ions to move into the sheath entrance at the sound speed $c_s = \sqrt{T_e/m_i}$. We consider here the case where there is no external biasing; in other words, the end plates are electrically isolated and floating. The wall can be set to an arbitrary potential, say $\phi_w = 0$, while the potential at the sheath entrance is then the positive floating potential ϕ_{sf}. This potential difference across the sheath reflects slow electrons that enter the sheath. The electrons approximately maintain a cutoff Maxwellian velocity distrubution throughout the sheath, and at the wall, their velocity is retarded by a Boltzmann factor due to the floating potential. In total, the current to the wall is [BCR93]

$$J_{\|} = en \left[c_s - \frac{(T_e/m_e)^{1/2}}{2\sqrt{\pi}} exp^{(-\frac{e\phi_{sf}}{T_e})} \right]. \tag{5.1}$$

Note that this is not only the current to the wall, but also the current going into the sheath edge, as long as current isn't escaping radially and there isn't an ionization source within the sheath. Furthermore, since the wall is electrically isolated, the equilibrium current at the wall vanishes. This sets the value for the floating potential to be $\phi_{sf} = \Lambda T_e/e$ with $\Lambda = ln(\frac{1}{2\sqrt{\pi}}\sqrt{\frac{m_i}{m_e}})$. Note that T_e is a function of radius, necessitating that ϕ_{sf} is also a function of radius. Thus, a radial equilibrium temperature gradient produces a radial equilibrium electric field. It is noted that $J_{\|}$ need not vanish on every field line since the end plates are conducting and charges can move around on the plate, however, the vanishing equilibrium current is generally a fair approximation [BCR93].

On the other hand, the fluctuating component of the current is allowed to vary between field lines. The first order fluctuating component is obtained by linearizing Eq. 5.1, giving the result:

$$\tilde{J}_{\|} = eN_0 c_{s0} \left[\frac{e\tilde{\phi}}{T_{e0}} - \Lambda \frac{\tilde{T}_e}{T_{e0}} \right]. \tag{5.2}$$

This expression for the current sets the fluctuating axial boundary condition of the plasma and is often called the Bohm Sheath boundary condition. This current condition holds both at the wall and at the sheath entrance. So rather than taking the simulation domain all the way to the wall, simulations often end at the sheath entrance and employ this analytically derived boundary condition to the boundaries of the main plasma. Then one doesn't have to worry about the small spatial scales and the non-quasineutrality of the sheath. The corresponding boundary conditions for the other fluid variables such as the density and temperature have recently been derived by Loizu et al. [LRH12].

The conducting wall mode instability in the case considered here is purely an electron temperature gradient instability, although other types of gradients can cause it [BCR93]. Electron temperature fluctuations are advected by electrostatic potential fluctuations and feed off the equilibrium electron temperature gradient as in the case of the thermal drift waves. However, in contrast to the thermal drift waves, the coupling between the temperature and potential fluctuations comes through the sheath boundary condition rather than through the adiabatic response.

5.2.2 Bohm Sheath Boundary Implementation

[XRD93]

CHAPTER 6

The Nature of LAPD Turbulence

6.1 A Visual Examination

6.2 A Statistical Examination

CHAPTER 7

Energy Dynamics Formalism

7.1 Total Energy and Dynamics

7.2 Spectral Energy Dynamics

CHAPTER 8

Nonlinear Instability for the Periodic Simulation

8.1 The Energy Spectra

8.2 Energy Dynamics Result

8.3 n=0 Suppression

CHAPTER 9

Energy Dynamics for the Non-periodic Simulations

9.1 The Importance of Axial Boundary Conditions

9.2 Fourier Decomposing Non-periodic Functions

9.3 Energy Dynamics Results

9.4 Linear vs Nonlinear Structure Correlation

CHAPTER 10

Finite Mean Flow Simulations

10.1 The LAPD Biasing Experiment

10.2 New Linear Instabilities

10.3 Statistical Comparisons to Experiment

10.4 Energy Dynamics Results

CHAPTER 11

Conclusion

APPENDIX A

The BOUT++ Code

A.1 The Object-Oriented Fluid Framework

A.2 Explicit Finite Differences

A.3 The Physics Inputs

APPENDIX B

Grid Convergence

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