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REPRODUCED FROM
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FOR TOKAMAK MICROTURBULENCE

By

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**Nonlinear Gyrokinetic Equations
for
Tokamak Microturbulence**

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Abstract

A nonlinear electrostatic gyrokinetic Vlasov equation, as well as Poisson equation, has been derived in a form suitable for particle simulation studies of tokamak microturbulence and associated anomalous transport. This work differs from the existing nonlinear gyrokinetic theories in toroidal geometry, since the present equations conserve energy while retaining the crucial linear and nonlinear polarization physics. In the derivation, the action-variational Lie perturbation method is utilized in order to preserve the Hamiltonian structure of the original Vlasov-Poisson system. Emphasis is placed on the dominant physics of the collective fluctuations in toroidal geometry, rather than on details of particle orbits.

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Nonlinear gyrokinetic equations formulated in terms of the perturbed distribution function¹ provide a useful theoretical basis for the various analytical drift wave turbulence models. For numerical studies utilizing particle simulation techniques, however, the gyrokinetic equations should be formulated in terms of the total distribution function.² Following the investigation by Lee, which demonstrated the importance of the linear and nonlinear polarization shielding in the Poisson equation and its impact on numerical schemes, Dubin et al.³ developed an energy-conserving set of gyrokinetic equations with Hamiltonian structure. A generalization to the electromagnetic (finite- β) case has been also achieved.⁴ The authors of Refs.3 and 4 considered slab geometry in order to isolate the crucial nonlinear physics issues.

An extension to realistic magnetic geometry in the electrostatic limit has also been attempted. Yang and Choi⁵ used the Hamiltonian method, which had been previously applied to the guiding center drift calculations,⁶⁻⁸ and derived a set of gyrokinetic Vlasov-Poisson equations. Although their derivation is mathematically correct, their equations do not conserve energy. This is due to the fact that they did not pursue the analysis of the gyrokinetic Vlasov equation to the order considered in Ref. 3, and, therefore, ignored the crucial nonlinear corrections to the effective potential. Even to the lowest order which has been published, their final equations are quite complex. Hagan and Frieman⁹ improved the gyrokinetic Vlasov equation of Ref. 5 by removing gyro-gauge-dependent terms and gave an illuminating discussion about the relations between the

formulations of Ref. 1 and Ref. 5. They unfortunately stopped the analysis at the lowest nontrivial order and did not consider the associated Poisson equation.

Despite the considerable progress made in Refs. 5 and 9, useful nonlinear gyrokinetic Vlasov-Poisson equations in toroidal geometry which contain all the relevant nonlinear physics are still needed. In this letter, we derive such equations via the phase-space Lagrangian variation method and the Lie perturbation theory.^{10,11} In doing so, the desirable Hamiltonian structure is automatically preserved. We adopt the usual gyrokinetic ordering of $\omega/\Omega_i \sim k_{\parallel}/k_{\perp} \sim e\phi/T_e \sim \rho_i/L_n = O(\epsilon)$ and $k_{\perp}\rho_i = O(1)$. Here, ω and Ω_i are the characteristic fluctuation frequency and the ion cyclotron frequency, respectively; $k_{\parallel}$ and $k_{\perp}$ are the components of the wave vector in the parallel and perpendicular direction with respect to the ambient magnetic field; ρ_i is the ion gyroradius; L_n is the density scale length; ϕ is the fluctuating electrostatic potential; and ϵ is a parameter indicating the smallness of ρ_i in comparison to L_n . This is not to be confused with another small parameter which describes the inhomogeneity of the equilibrium magnetic field: $\epsilon_B \sim \rho_i/L_B$, where L_B is the scale length of the magnetic field variation. For typical tokamak parameters, $\epsilon_B/\epsilon \sim L_n/L_B \sim a/R_0 \ll 1$, where a/R_0 is the inverse aspect ratio. We note that $v_{di}/v_{ti} \sim \epsilon_B$, where v_{di} is the drift speed of a thermal ion due to the magnetic gradient and curvature, and v_{ti} is the ion thermal speed. In this letter, we derive equations which are valid up to $O(\epsilon^2)$

and $O(\epsilon_B)$. The neglected terms are $O(\epsilon_B^2)$ and represent higher-order magnetic drifts.

We begin the analysis by briefly discussing the action-variational principle. A comprehensive discussion of the theory can be found in Refs.10 and 11. A single-particle Lagrangian written in terms of the canonical variables (q,p) is $L = p \cdot \dot{q} - h_C(q,p,t)$, where h_C is the usual canonical Hamiltonian. It can be generalized to an arbitrary phase space coordinate system $z = z(q,p,t)$ via the chain rule: $L = \vartheta_i \dot{z}^i - h$ ($i=1,\dots,6$; the summation convention is used), where ϑ_i and h are functions of z and t given by $\vartheta_i = p \cdot \partial q / \partial z^i$, $h = h_C - p \cdot \partial q / \partial t$. Now, the equation of motion can be derived from the variation of the integral $\int L dt$. The covariance of the variational formulation is manifested more clearly by considering the fundamental 1-form of Poincare-Cartan,

$$\vartheta = \vartheta_\mu dz^\mu = \vartheta_i dz^i - h dt, \quad (1)$$

where $z^0 = t$, $\vartheta_0 = -h$. Then, under the coordinate transform $Z^\mu = Z^\mu(z,t)$, we have

$$\vartheta = \vartheta_\mu dz^\mu = \Gamma_\mu dZ^\mu, \quad (2)$$

where $\Gamma_\mu = \vartheta_\sigma \partial z^\sigma / \partial Z^\mu$. That is, the fundamental 1-form will always have the same form for some Γ_μ in any phase-space coordinates. The variation of ϑ yields the Euler-Lagrange equation,

$$\hat{\omega}_{ij} dz^j / dt = \partial h / \partial z^i + \partial \vartheta_i / \partial t, \quad (3)$$

where $\hat{\omega}_{ij} = \partial \vartheta_j / \partial z^i - \partial \vartheta_i / \partial z^j$ are the components of the Lagrange tensor. It is important to note that the Euler-Lagrange equation is

invariant under the gauge transformation $\mathcal{J} \rightarrow \mathcal{J} + dS$, for any scalar function S on the extended phase space.

The canonical single-particle Hamiltonian in the presence of a perturbed electrostatic field in an arbitrary magnetic geometry is given by

$$h = (2m)^{-1} \{p - (e/c)A_0(q)\}^2 + e\phi(q,t).$$

Now, we write the fundamental 1-form (Lagrangian) in a noncanonical coordinate $z = (x,v)$ which are related to the canonical variables by $x = q$, $v = m^{-1} \{p - (e/c)A_0\}$. Then, according to Eqs.(1) and (2),

$$\mathcal{J} = \{(e/c)A_0(x) + mv\} \cdot dx - \{mv^2/2 + e\phi(x,t)\}dt, \quad (4)$$

The last term in Eq. (4) is the electrostatic perturbation. We can now eliminate the fast gyromotion of Eq.(4) through an averaging procedure. It is conceptually simpler and probably more physical to proceed through a number of steps systematically. First, we consider only the equilibrium part of $\mathcal{J}$, denoted by $\mathcal{J}_0$, and define the gyrocenter variables only in terms of the equilibrium magnetic field, such that the gyrophase dependence is removed from $\mathcal{J}_0$. Then, we consider the perturbed part together with the equilibrium part and find another transformation which removes the gyrophase dependence of the perturbed part. To the order which is considered in this paper, the former can be done straightforwardly, while the latter can be carried out most efficiently by the Lie transformation. The equilibrium part becomes

$$\mathcal{J}_0 = \{(e/c)A_0(x) + mv\} \cdot dx - (mv^2/2)dt. \quad (5)$$

Defining the gyrocenter variable as usual, we have

$$\mathbf{x} = \mathbf{R} + \{\mathbf{v}_\perp / \Omega(R)\} \mathbf{a}(R, \theta), \quad (6a)$$

$$M = v_\perp^2 / 2\Omega(R), \quad (6b)$$

$$U = \mathbf{v} \cdot \mathbf{b}, \quad (6c)$$

and

$$\theta = \tan^{-1}(\mathbf{v} \cdot \mathbf{e}_1 / \mathbf{v} \cdot \mathbf{e}_2), \quad (6d)$$

where $\mathbf{e}_1$ and $\mathbf{e}_2$ are arbitrary orthogonal unit vectors in the plane perpendicular to $\mathbf{b}$, $\mathbf{a} = \mathbf{e}_1 \cos \theta - \mathbf{e}_2 \sin \theta$, $\mathbf{c} = -\mathbf{e}_1 \sin \theta - \mathbf{e}_2 \cos \theta$, and $\Omega = eB_0(R)/mc$. At this point, we do not order terms in Eq. (6a) explicitly. This differs from the conventional guiding center drift calculations⁶⁻⁸ where the second term was explicitly ordered to be smaller than the first. By substituting Eq. (6) into Eq. (5) and keeping the terms up to $O(1)$, we get

$$\dot{\mathbf{x}}_0 = \{(\mathbf{e}/c)A_0(R) + mU\mathbf{b}(R)\} \cdot d\mathbf{R} + M d\theta - \{M\Omega(R) + mU^2/2\} dt. \quad (7)$$

It will be shown later that keeping the terms only up to $O(1)$ is consistent with our ordering and is sufficient for our purposes. Let us now introduce the electrostatic fluctuation ϕ and remove the gyrophase dependence from the resulting formulation. This can be done most efficiently via Lie transformation. The application of the Lie transformation is discussed in detail in Refs. 10 and 11. Here, we merely provide formulas which are pertinent to our derivation. The Lie coordinate transform can be formally written as

$$\mathbf{Z}^\mu = \mathbf{T} \mathbf{z}^\mu,$$

where $\mathbf{T}$ is, in general, a composition of a sequence of individual Lie transforms, $\mathbf{T} = \dots \mathbf{T}_3 \mathbf{T}_2 \mathbf{T}_1$.¹² An individual Lie transform $\mathbf{T}_n$ is the near-identity coordinate transformation, which can be written in

operator form as

$$T_n = \exp(\epsilon^n L_n). \quad (8)$$

Under T , a scalar transforms as $S = T^{-1}s$, and the operator L_n acts on scalars as $L_n f = g_n^\mu \partial f / \partial z^\mu$, where g_n^μ is the generator of the Lie transformation T_n , which satisfies $\partial Z^\mu / \partial \epsilon^n = g_n^\mu(Z)$. On the other hand, the 1-form $\mathcal{S} = \mathcal{S}_\mu dz^\mu$ transforms according to the rule

$$\Gamma = T^{-1} \mathcal{S} + dS, \quad (9)$$

where $\Gamma = \Gamma_\mu dZ^\mu$, and S represents a gauge transformation in phase space. The operator L_n acts on 1-forms as¹⁰

$$(L_n \mathcal{S})_\mu = g_n^\sigma (\partial \mathcal{S}_\mu / \partial z^\sigma - \partial \mathcal{S}_\sigma / \partial z^\mu). \quad (10)$$

Upon expanding the Γ , $\mathcal{S}$, and S of Eq. (9) in powers of ϵ , and expanding T_n in Eq.(8), we have

$$\Gamma_0 = \mathcal{S}_0 + dS_0. \quad (11a)$$

$$\Gamma_1 = \mathcal{S}_1 - L_1 \mathcal{S}_0 + dS_1. \quad (11b)$$

$$\Gamma_2 = \mathcal{S}_2 - L_1 \mathcal{S}_1 + \{(1/2)L_1^2 - L_2\} \mathcal{S}_0 + dS_2. \quad (11c)$$

The perturbation calculation consists of finding the transformation order by order by specifying S_n and g_n^μ in a way which simplifies Γ_μ .

Finally, we are ready to consider the $\mathcal{S}_0$ given by Eq. (7). For simplicity, we take $dS_0 = 0$ and $g_n^t = 0$. The latter condition is imposed since we do not transform time. From Eq. (4), we have

$$\mathcal{S}_1 = -e\phi(R + \mathbf{a} \cdot \mathbf{v}_\perp / \Omega) dt. \quad (12)$$

Then, Eq. (11b) yields

$$\Gamma_1 = dS_1 - e\phi dt + g_1^\theta dM - g_1^M d\theta + m(g_1^R \cdot b dU - g_1^U \cdot b \cdot dR) \\ + (e/c)g_1^R \times B^* \cdot dR + (\Omega g_1^M + mUg_1^U + Mg_1^R \cdot \nabla \Omega)dt, \quad (13)$$

where $g_n^R = (g_n^{R1}, g_n^{R2}, g_n^{R3})$, $B^* = \nabla \times A_0 + (mc/e) U \nabla \times b$.

In Eq. (13), we can choose dS^1 and g_1^μ such that all the $\Gamma_{1\mu}$ vanish except for Γ_{1t} . Also, by requiring that there is no θ -dependent term in Γ_{1t} , we get

$$\Gamma_{1t} = -e\langle\phi\rangle. \quad (14)$$

For the second order,

$$\Gamma_2 = dS_2 - (1/2)g_1^i \partial_i e \tilde{\phi} dt + g_1^i \partial_i e \phi dt \\ + g_2^\theta dM - g_2^M d\theta + m(g_2^R \cdot b dU - g_2^U \cdot b \cdot dR) \\ + (e/c)g_2^R \times B^* \cdot dR + (\Omega g_2^M + mUg_2^U + Mg_2^R \cdot \nabla \Omega)dt. \quad (15a)$$

By requiring the same constraints on $\Gamma_{2\mu}$, we have

$$\Gamma_{2t} = e^2/2\Omega [\partial/\partial M \langle \tilde{\phi}^2 \rangle + 1/\Omega \langle \nabla \tilde{\phi} \cdot b \times \nabla \tilde{\phi} \rangle], \quad (15b)$$

where $\langle\phi\rangle = 1/2\pi \oint d\theta \phi(R+p, t)$, $\tilde{\phi} = \phi - \langle\phi\rangle$, and $\tilde{\phi} = \int \theta \tilde{\phi} d\theta$.

Upon adding Eqs. (14) and (15) together, we obtain the fundamental 1-form which describes the particle motion in gyrocenter variables in the presence of a perturbed electrostatic field

$$\Gamma = ((e/c)\overline{A_0}(\overline{R}) + m\overline{U}b(\overline{R})) \cdot d\overline{R} + \overline{M}d\overline{\theta} - \{m\overline{U}^2/2 + \overline{M}\Omega(\overline{R}) + e\psi\}dt, \quad (16)$$

where $\psi = \langle\phi\rangle - e/2\Omega[\partial/\partial M \langle \tilde{\phi}^2 \rangle + 1/\Omega \langle \nabla \tilde{\phi} \cdot b \times \nabla \tilde{\phi} \rangle]$ is the renormalized potential. We use $\overline{}$ to indicate the gyrophase-averaged coordinates. The nonvanishing components of the Lagrange tensor defined in Eq. (3) are

$$\overline{\omega_{RiRj}} = (e/c)\epsilon_{ijk}B_k^*, \quad (17a)$$

$$\overline{\omega_{RU}} = -mb, \quad (17b)$$

$$\hat{\omega}_t \bar{M} = \Omega + e \partial \Psi / \partial \bar{M}, \quad (17c)$$

$$\hat{\omega}_t \bar{U} = m \bar{U}, \quad (17d)$$

$$\hat{\omega}_t \bar{R} = M \bar{\nabla} \Omega + e \bar{\nabla} \Psi, \quad (17e)$$

$$\hat{\omega}_{\theta \bar{M}} = -1. \quad (17f)$$

The nontrivial components of the Euler-Lagrange equation, Eq. (3), yield

$$d\bar{M}/dt = 0, \quad (18a)$$

$$d\bar{\theta}/dt = \Omega + e \partial \Psi / \partial \bar{M}, \quad (18b)$$

$$b \cdot d\bar{R}/dt = \bar{U}, \quad (18c)$$

$$- \{ (e/c) B + m \bar{U} \bar{\nabla} \times b \} \times d\bar{R}/dt - m b d\bar{U}/dt = \bar{M} \bar{\nabla} \Omega + e \bar{\nabla} \Psi. \quad (18d)$$

Equation (18d) can be put into more physically transparent form by taking projections. Taking the cross product of Eq. (18d) with b , we derive the equations for the gyrocenter drift,

$$d\bar{R}/dt = 1/B^* \{ \bar{U} B^* + b \times (c/e) \bar{M} \bar{\nabla} \Omega + b \times c \bar{\nabla} \Psi \}, \quad (19)$$

where $B^* = B + (mc/e) \bar{U} b \cdot \bar{\nabla} \times b$. Using the vector identity $B^* = B^* b + (mc \bar{U}/e) b \times b \cdot \bar{\nabla} b$, we can write Eq.(19) as

$$d\bar{R}/dt = \bar{U} b + (1/B^*) b \times \{ (c/e) \bar{M} \bar{\nabla} \Omega + (mc \bar{U}^2/e) b \cdot \bar{\nabla} b + c \bar{\nabla} \Psi \}. \quad (20)$$

Also, taking the dot product of Eq. (18d) and B^* , we obtain the parallel acceleration equation,

$$\begin{aligned} m d\bar{U}/dt &= -(1/B^*) B^* \cdot (\bar{M} \bar{\nabla} \Omega + e \bar{\nabla} \Psi) \\ &= - (b + (mc \bar{U}/e B^*) b \times b \cdot \bar{\nabla} b) \cdot (\bar{M} \bar{\nabla} \Omega + e \bar{\nabla} \Psi). \end{aligned} \quad (21)$$

Note that the familiar curvature and ∇B drifts of $O(\epsilon_B)$ are included in Eq.(20). Also, a modification to the parallel acceleration is included in Eq. (21). Although it is possible to find higher-order geometric corrections to the drifts, Eqs. (20) and (21) already include the most relevant physics associated with particle drifts in toroidal

geometry for practical purposes. In a number of previous publications,^{5,9} B^* has been approximated by B , ignoring the parallel velocity correction. Although those equations are also valid to the same order, our equations preserve the Hamiltonian structure exactly.

Now, following a procedure similar to that of Ref.3, it is straightforward to obtain an energy-conserving set of gyrokinetic Vlasov-Poisson equations in general geometry. They are

$$\partial F / \partial t + d\bar{U} / dt \cdot \partial F / \partial \bar{U} + d\bar{R} / dt \cdot \partial F / \partial \bar{R} = 0, \quad (22a)$$

$$\nabla^2 \phi(x, t) = -4\pi e \left[\int (F(\bar{Z}) + e / \Omega_i \{ \tilde{\psi} \partial F / \partial \bar{M} + 1 / \Omega_i \nabla \tilde{\phi} \cdot b \times \nabla F \}) s(\bar{X} - x + \bar{\rho}) d^6 \bar{Z} - n_e \right], \quad (22b)$$

where the term inside { } on the right-hand side of Eq. (22b) is the polarization density,^{2,3} and n_e can be straightforwardly obtained from the nonlinear electron drift-kinetic equation. The invariant energy for this system is³

$$\begin{aligned} E = & \int F_i m_i (U^2 / 2 + MB) d^6 Z + (1/2) \int f_e m_e v_z^2 d^6 z \\ & + (8\pi)^{-1} \int |sE|^2 d^3 x \\ & + e^2 / 2 \Omega_i \int [\partial / \partial M \langle \tilde{\phi}^2 \rangle + 1 / \Omega_i \langle \nabla \tilde{\phi} \cdot b \times \nabla \tilde{\phi} \rangle] F_i d^6 Z = \text{const.} \end{aligned} \quad (23)$$

It is instructive to compare our results with those of Ref. 5. First of all, our system of equations is much simpler algebraically than those of Ref. 5. Moreover, our equations conserve energy while those of Ref. 5 do not. This property is quite important for particle simulation purposes. We also note that the Poisson equation and energy conservation were not considered in Ref. 9. In conclusion, Eqs. (20-22) contain most of the relevant linear and nonlinear physics for describing microturbulence in tokamaks. For example, it is

straightforward to recover the governing equations of the 'toroidal' drift wave^{1,13} from our equations. The ample physics contents and the relative simplicity of our equations should be very useful for simulation applications.

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