Zonal Flow as Pattern Formation

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Zonal flows — azimuthally symmetric, generally banded, shear flows — are spontaneously generated from turbulence and have been reported in atmospheric and laboratory plasma contexts. Recently, they have also been observed in astrophysical simulations. In magnetically confined plasmas, zonal flows are thought to play a crucial role in regulation of turbulence and turbulent transport. A greater understanding of zonal flow behavior is valuable for untangling a host of nonlinear processes in plasmas, including details of transitions between modes of low and high confinement.

Zonal flows remain incompletely understood, even regarding the basic question of the jet width. In the plasma literature, one finds modulational or secondary instability calculations of zonal flow generation, but these cannot provide information on a saturated state. Other theories typically make an assumption of long wavelength zonal flows and leave the zonal flow scale as an undetermined parameter. Within geophysical contexts, various authors have attempted to relate the jet width or spacing to length scales that emerge from the vorticity equation by heuristically balancing the magnitudes of the Rossby wave term and the nonlinear advection. Those scales include the Rhines scale and other, similar scales. A Rhines-like length scale is also obtained from arguments based on potential vorticity staircases. However, neither the heuristic Rhines estimates nor the paradigm of potential vorticity inversion and mixing generalize to more complex situations involving realistic plasma models. We are therefore motivated to seek a more systematic approach to determining the zonal flow width that may offer such a generalization.

A related topic is the merging of jets. Coalescence of two or more jets is ubiquitous in numerical simulations. The merging process occurs during the initial transient period before a statistically steady state is reached. It is clear that the merging is part of a dynamical process through which the zonal flow reaches its preferred length scale, but there has been little theoretical understanding of the merging phenomenon.

Our present work addresses these questions in the context of the stochastically forced generalized Hasegawa-Mima (GHM) equation for electrostatic potential, a model for magnetized plasma turbulence in the presence of a background density gradient. This model is mathematically similar to the barotropic vorticity equation on a β plane. Our analysis is related to several recent works that focused on that equation in the geophysical context. Importantly, numerical simulations of both models can display emergence of steady zonal flows. Because the GHM equation and the parameterizations of forcing and dissipation that we use are not realistic descriptions of plasma; however, the simplicity is an asset in understanding the qualitative behavior of these systems.

We study a statistical average of the flow. Statistical approaches enable one to gain physical insight by averaging away the details of the turbulent fluctuations and working with smoothly varying quantities. Sometimes, statistical turbulence theories strive for quantitative accuracy, which requires rather complicated methods. In contrast, our investigation is at a more basic level and concerns the fundamental nature of zonal flows interacting self-consistently with inhomogeneous turbulence.

We discover that from the statistical point of view, steady zonal flows emerge from homogeneous turbulence in a symmetry-breaking bifurcation. The bifurcation that generates these zonal flows obeys a classic amplitude equation, and therefore zonal flows can be understood as pattern formation. Two important results follow from the general properties of pattern-forming systems. First, the zonal flow wavelength is not unique. Indeed, in an idealized, infinite system, any wavelength within a certain continuous band corresponds to a solution. Second, of these wavelengths, only those within a smaller subband are linearly stable. Unstable wavelengths must evolve to reach a stable wavelength; this process manifests as merging jets.

Our basic model is the GHM equation in a uniform magnetic field in 2D,

\[
\partial_t w(x, y) + \mathbf{v} \cdot \nabla w - \kappa \partial_y \phi = \xi - \mu w - \nu (-1)^h \nabla^2 w, \tag{1}
\]

where \(\phi = (L_n/\rho_s) e \varphi/T_e \) is the normalized electrostatic
potential, $L_n$ is the density gradient scale length, $\rho_s$ is the sound radius, $T_e$ is the electron temperature, $w = \nabla^2 \phi - \hat{\alpha} \phi$ is the generalized vorticity and is related to ion gyrocenter density fluctuations $\delta n^G_t$ by $w = -(L_n/\rho_s) \delta n^G_t/n_0$ where $n_0$ is the background density, $\hat{\alpha}$ is an operator that is zero when acting on zonal flows and unity when acting on drift waves which respects within this 2D model the lack of adiabatic electron response to zonal flows, the magnetic field is in the $\hat{z}$ direction, $\mathbf{v} = \hat{z} \times \nabla \phi$ is the $\mathbf{E} \times \mathbf{B}$ velocity, $\mu$ is a constant frictional drag, $\nu$ is the viscosity with hyperviscosity factor $h$, $\xi$ is white-noise forcing, and $\kappa$ is related to the density scale length. Lengths are normalized to the sound radius $\rho_s$ and times are normalized to the drift wave period $\omega_s^{-1} = (L_n/\rho_s) \Omega_s^{-1}$. These normalizations and scalings are convenient to make $w, \phi$, and the active length and time scales of order unity, and additionally they allow us to set $\kappa = 1$.

The zonal flow behavior in numerical simulations of Eq. (1) is shown in Fig. 1(a). During the transient period, merging jets are observed, while in the late time, a statistically steady state is reached with stable unwaver- ing jets.

We restrict ourselves to the quasilinear (QL) approximation of this system. To obtain the QL equations, we perform an eddy–mean decomposition, given by decom-

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ing jets.

The CE2 equations exhibit several important symme-

tries of translation and reflection, given by $\pi \rightarrow \pi + \delta \pi$, $(x, \pi) \rightarrow (-x, -\pi), (y, \pi) \rightarrow (-y, -\pi)$, and $(x, y) \rightarrow (-x, -y)$. Many studies of CE2 have been performed previously. Numerical simulations of CE2 also exhibit merging jets. For Eq. (2) there always exists a homogeneous equilib-rium, which arises from a simple balance between forcing and dissipation: $W = (2\mu + 2\nu D_h)^{-1} F, U = 0$. This equilibrium is stable in a certain regime of param-
eters. As a control parameter such as the friction $\mu$ is varied, this homogeneous state becomes zonostrophically unstable. Physically, zonostrophic instability occurs when dissipation is overcome by the mutually reinforcing processes of eddy tilting by zonal flows and production of Reynolds stress forces by tilted eddies. The instability eigenmode consists of perturbations spatially periodic in $\pi$ with zero real frequency, so that zonostrophic instability arises as a Type I instability of homogeneous turbulence. Zonostrophic instability within CE2 may be thought of as a variant of the modulational instability calculations of zonal flow generation.

Just beyond the instability threshold, a bifurcation analysis yields a perturbative amplitude equation for the bifurcating mode. This amplitude equation is con-
strained by the translation and reflection symmetries to take a universal form. The amplitude equation, sometimes referred to as the real Ginzburg-Landau equation, is

\[ \partial_t A(x, t) = A + \partial_x^2 A - |A|^2 A, \]  

(3)

where all coefficients have been rescaled to unity. Here, \( A \) is the complex, spatially varying amplitude of the eigenvector that is neutrally stable at the bifurcation point. Equation (3) also describes the bifurcation of Rayleigh-Bénard convection rolls, so the zonal flows are mathematically analogous to convection rolls. The derivation of Eq. (3) from Eq. (2) will be reported elsewhere.

The amplitude equation (3) is well understood. First, a steady-state solution exists for any wave number \(-1 < k < 1\) (to see this, observe that \( A = e^{i\alpha k} \) with \( |\alpha|^2 = 1 - k^2 \) is a solution). Second, only solutions with \( k^2 < 1/3 \) are linearly stable. This is demonstrated in Fig. 1(b), where an unstable solution that has been slightly perturbed undergoes merging behavior until a stable wave number is reached. These qualitative behaviors are also exhibited by the CE2 system, as we now show.

We proceed to find the steady-state solutions of Eq. (2). In the context of an infinite domain with no boundaries, these solutions are referred to as ideal states. Let \( q \) denote the basic zonal flow wave number of an ideal state. For a given \( q \), we solve the time-independent form of Eq. (2) directly. This approach is distinct from time integration of Eq. (2) to a steady state. Our procedure has two advantages, both related to the fact that ideal states exist for any \( q \) within a continuous band. First, we can specify precisely the \( q \) of the desired solution. Second, we can solve directly for all solutions, including unstable ones, rather than find only those which develop from time evolution.

An ideal state is represented as a Fourier-Galerkin series with coefficients to be determined. We expand as follows:

\[ U(x) = \sum_{p=-P}^{P} U_p e^{ipq\xi}, \]  

(4a)

\[ W(x, y | \xi) = \sum_{m=-M}^{M} \sum_{n=-N}^{N} \sum_{p=-P}^{P} W_{mnp} e^{i\alpha m x} e^{in y} e^{ipq\xi}. \]  

(4b)

While the periodicity in \( \xi \) is desired, the correlation function should decay in \( x \) and \( y \); periodicity in \( x \) and \( y \) is a consequence of using the convenient Fourier basis. Thus, \( a \) and \( b \), unlike \( q \), are numerical parameters. They represent the spectral resolution of the correlation function and should be small enough to obtain an accurate solution.

The CE2 symmetries allow us to seek a solution where \( U(x) = U(-x) \) and \( W(x, y | \xi) = W(-x, -y | \xi) = W(x, -y | -\xi) = W(-x, y | -\xi) \). These constraints,

along with reality conditions, force \( U_p \) to be real, \( U_p = U_{-p} \), and \( W_{mnp} = W_{-m,n,p} = W_{m,-n,p} = W_{m,n,-p} \).

We obtain a system of nonlinear algebraic equations for the coefficients \( U_p, W_{mnp} \) by substituting the Galerkin series into Eq. (2) and projecting onto the basis functions. To demonstrate the projection for Eq. (2a), let \( \phi_{mnp} = e^{i\alpha m x} e^{in y} e^{ipq\xi} \). We project Eq. (2a) onto \( \phi_{rst} \) by operating with

\[ \left( \frac{2\pi}{a} \frac{2\pi}{b} \frac{2\pi}{q} \right)^{-1} \int_{-\pi/a}^{\pi/a} dx \int_{-\pi/b}^{\pi/b} dy \int_{-\pi/q}^{\pi/q} d\xi \phi_{rst}^*. \]  

(5)

For instance, the term \((U_+ - U_-)\partial_t W \) projects to \( I_{\xi \text{stp}^\prime mnp}^{(1)} U_p W_{mnp} \), where repeated indices are summed over, \( I_{\xi \text{stp}^\prime mnp}^{(1)} = im \delta_{m,r} \delta_{\rho^\prime+p-t,0} (\sigma_+ - \sigma_-) \), \( \sigma_\pm = \text{sign}(\alpha_\pm /b) \), and \( \alpha_\pm = nb \pm sp^\prime q/2 \). The other terms of Eq. (2a), as well as Eq. (2b), are handled similarly. In total, we generate as many equations as there are coefficients.

The system of nonlinear algebraic equations is solved with a Newton’s method. Figure 2 shows the zonal flow amplitude coefficients \( U_p \) as functions of \( q \) at \( \mu = 0.21 \) (\( R_\beta = 1.48 \)) and \( \mu = 0.19 \) (\( R_\beta = 1.51 \)). In the unshaded region, ideal states are stable. The vertical lines correspond to various instabilities which separate the regions (see Fig. 3).
our infinite-domain results are modified merely by the domain. When periodic boundary conditions are used, observed in QL simulations.

is consistent with the dominant zonal flow wavenumber from independent QL simulations (crosses). The stationary ideal states vanish to the left of $D$. Here, $a = 0.06$, $b = 0.08$, $M = 20$, $N = 33$, $P = 5$, and other parameters are given in the text. $\gamma$ is varied by changing $\mu$ while holding other parameters fixed.

expanded as a Bloch state:

$$
\delta W(x, y | \xi, t) = e^{i\sigma t} e^{iQ\xi} \sum_{mnp} \delta W_{mnp} e^{im\xi} e^{in\eta} e^{ip\sigma \xi},
$$

$$
\delta U(\xi, t) = e^{i\sigma t} e^{iQ\xi} \sum_{p} \delta U_{p} e^{ip\sigma \xi},
$$

where $Q$ is the Bloch wave number and can be taken to lie within the first Brillouin zone $-q/2 < Q < q/2$. We do not use a $Q_{x}$ or $Q_{y}$ because as previously mentioned the periodicity in $x$ and $y$ is artificial. The perturbation equations are projected onto the basis functions in the same way as in the ideal state calculation. This projection results in a linear system at each $Q$ for the coefficients $\delta W_{mnp}$ and $\delta U_{p}$; this determines an eigenvalue problem for $\sigma$. The equilibrium is unstable if for any $Q$ there are any eigenvalues with Re$\sigma > 0$.

The stability diagram is shown in Fig. 3. As the control parameter we adopt $\gamma = \varepsilon^{1/4} \beta^{1/2} \mu^{-5/4}$, an important dimensionless parameter controlling the zonal flow dynamics. To vary $\gamma$, we change $\mu$ and hold other parameters fixed at their previous values. The stable ideal states exist inside of the marginal stability curve marked $S$. Near the threshold, marginal stability is governed by the Eckhaus instability, a long-wavelength universal instability. Farther from threshold, the instability transitions into new, nonuniversal instabilities; details will be reported elsewhere. The zonal flows are spontaneously generated for $\gamma > 6.53$. For $\gamma > 6.53$, the stability curve is consistent with the dominant zonal flow wavenumber observed in QL simulations.

Numerical simulations typically occur within a finite domain. When periodic boundary conditions are used, our infinite-domain results are modified merely by the discretization of wave numbers. This affects not only the possible equilibria, but also any perturbations and hence the stability boundaries too.

For a time-evolving system, the exact $q$ that is ultimately chosen within the stability balloon results from a dynamical process and is not addressed in a systematic way by the present study.

While the CE2 equations exhibit spontaneously generated zonal flows, it is true that they neglect many physical effects. An important piece of physics missing from the CE2 equations is the nonlinear eddy self-interaction, which clearly cannot be ignored in general. Furthermore, the CE2 equations involve one-time correlation functions rather than the more general two-time functions. The lack of time-history information means that most of the effects of wave propagation are discarded. At least one particular instance of the qualitative failure of CE2 has been noted.

Yet, the basic mathematical structure of the theory presented here arises only from symmetry arguments and general properties of the zonostrophic instability. If one were to include the important physics neglected in CE2, those general symmetries and properties should remain intact. Therefore, we expect our qualitative conclusions to likewise remain valid.

In summary, by analyzing a second-order statistical model of an ensemble of interacting zonal flows and turbulence, we have shown that zonal flows constitute pattern formation amid a turbulent bath. This continues previous work to provide a firm analytic understanding of zonal flow generation and equilibrium within CE2. We calculated the stability diagram of steady zonal jets and explained the merging of jets as a means of attaining a stable wave number. In general, the use of statistically averaged equations, perhaps with more sophisticated closures, and the pattern formation methodology provide a path forward for further systematic investigations of zonal flows and their interactions with turbulence.

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