

# Density Troughs in the Ionosphere Sustained by Transport Barriers

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## Key Points:

- Density troughs formed in the Earth's ionosphere are likely to develop fast gradient instabilities but typically live longer than expected
- We argue that while driving turbulent plasma transport along the density gradient, the drift-wave instability also generates macroscopic (zonal) flows across the gradient and magnetic field
- Zonal flows efficiently self-regulate the instability-transport relation, essentially suppressing the two

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## Abstract

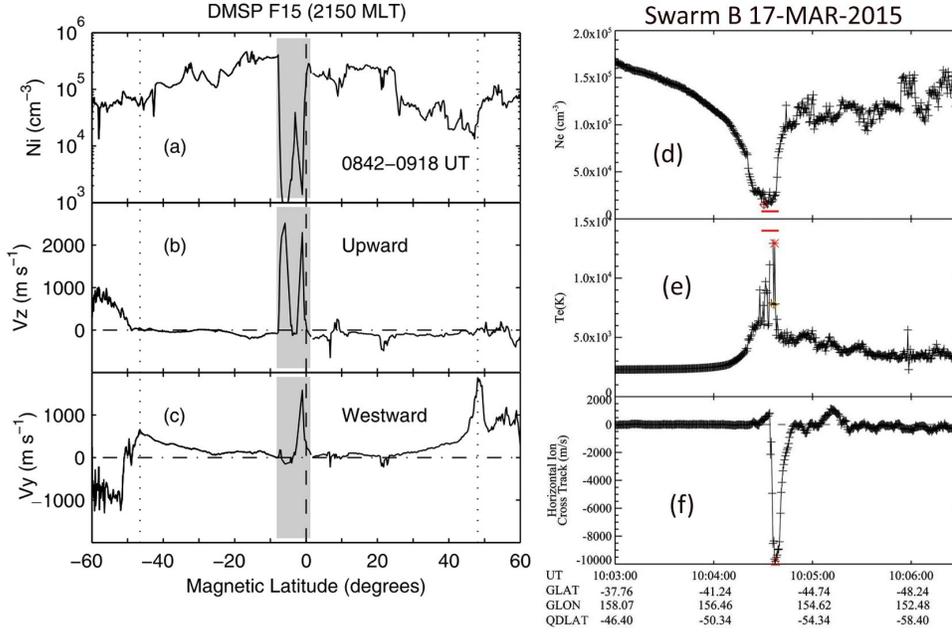
This study explores the relaxation and sustainability of density irregularities and plasma flows in the Earth's ionosphere. To do this, we use a modified model of drift-wave turbulence known as the Hasegawa-Wakatani model. Similar to turbulent processes in laboratory plasmas, we explore a powerful mechanism that can reduce the turbulent plasma transport. This mechanism is associated with the creation of 'zonal flows,' cutting across the gradients of particle density and magnetic fields. They work effectively to minimize particles' random movement and reduce the turbulence causing this movement. The zonal flows create transport barriers in areas where the density gradient is steepest and where drift waves grow most vigorously. The transport barriers significantly delay the refill of low-density regions with surrounding plasma. They also lead to changes in the electric potential of the plasma and influence the movement of ions in the direction of the magnetic field. Our research investigates how these zonal flows are generated and how effectively they sustain density irregularities. We also examine the spectral characteristics of turbulence in and around these barriers.

## 1 Introduction

During active geomagnetic conditions, midlatitude and subauroral ionosphere just equatorward of the auroral zone often becomes disturbed by a number of phenomena that could adversely affect modern technology. In particular, the disturbed subauroral convection is dominated by enhanced, westward,  $\mathbf{V} = \mathbf{V}_W$ , plasma streams: narrow subauroral ion drifts (SAID) near midnight and broad subauroral polarization streams (SAPS) on the duskside, containing deep-density depletions, manifested in troughs elongated in the magnetic field direction (e.g., Foster & Burke, 2002; E. V. Mishin, 2013). Plasma density irregularities inside subauroral irregular plasma density troughs (e.g., Mishin et al., 2003; Sinevich et al., 2022) may induce scintillations in radio waves transmitted by GNSS or communication satellites, thus affecting the quality of received signals (e.g., Basu et al., 2001; Ledvina et al., 2002; Nishimura et al., 2021; Pradipta et al., 2023). While the SAID and SAPS onsets are close to that of magnetospheric substorms, e.g., (E. Mishin et al., 2017), the phenomena typically last a few more hours well into the substorm recovery (e.g., He et al., 2017). Various scenarios for the SAID and SAPS generation have been critically ascertained by E. V. Mishin (2023).

The discovery (MacDonald et al., 2018) of peculiar subauroral arcs – Strong Thermal Emission Velocity Enhancement (STEVE) and the picket fence – inside extremely high speed and electron temperature SAID channels with deep troughs instigated a series of systematic observational efforts, both spacecraft- and ground-based (see a comprehensive review by (Nishimura et al., 2023), complemented with theoretical modeling (e.g., Liang et al., 2021, 2022; E. V. Mishin & Streltsov, 2024). The observations testify to intrinsic mechanisms of sustainability, self-organization, and ultimate disruption of SAID. Indeed, the density and temperature gradients such as at the SAID density walls and strong  $\mathbf{E} \times \mathbf{B}$  shear flows have long been recognized as potentially efficient drivers of several instabilities (Tsunoda, 1988). Should they develop without a significant back-reaction from the SAID to sustain its structure, the unstable waves would likely grow to the point of SAID disruption.

An example of ionospheric structures pertinent to the present study and in which steep density gradients form is equatorial spread F (ESF) plasma bubbles. They were studied theoretically in (e.g., Ossakow & Chaturvedi, 1978; Ott, 1978). The ESF bubbles are frequently accompanied by very fast ion drifts (Huang et al., 2007), as shown in one observation example in Fig.1 from (Huang et al., 2007). In this particular case, the DMSP spacecraft crosses the ESF plasma bubble just south of the magnetic equator, detecting the plasma density drops by more than two orders of magnitude. The vertical velocity has two peaks inside the bubble at 2000 m/s. The westward velocity has



**Figure 1.** (left column) Adopted from (Huang et al., 2007) with authors’ permission including the following caption: Latitudinal distribution of the ionospheric ion density and velocity measured by the (a) DMSF F15 satellite between 0842 and 0918 UT on 29 October 2003.  $V_z$  is the vertical velocity component, and  $V_y$  is the horizontal velocity component. The shaded regions indicate deep density depletion. The vertical dotted lines indicate the location of the ion horizontal velocity peak in the SAPS region. (right) Adopted from (Martinis et al., 2022) with authors’ permission: Plasma parameters measured along the Swarm B trajectory in the Northern Hemisphere mapped to 425 km in the Southern Hemisphere. From top to bottom: electron density ( $N_e$ ), electron temperature ( $T_e$ ), and horizontal plasma drift velocity ( $V_{hor}$ ).

66 only one strong peak inside the bubble that reaches 1500 m/s. This flow morphology indicates that the two jets inside the bubble are not collinear. *For generality, we also show*  
 67 *the SAID channel detected by the Swarm-B satellite at about 500 km altitude during a*  
 68 *bright STEVE arc event (Martinis et al., 2022). Such  $\delta T_e - \delta n_e - \delta V_{hor}$  variation across*  
 69 *the channel is typical for strong SAID events. Frequently, the upward vertical flow ve-*  
 70 *locity is also enhanced inside SAID channels.*  
 71

72 One may distinguish between two scenarios of a turbulent relaxation of a steep density  
 73 gradient after the driver that creates this gradient is switched off. For the sake of  
 74 argument, let us consider an isolated plasma layer with strong density gradients across  
 75 the magnetic field between two broader plasma regions with constant densities. The first  
 76 scenario can be termed quasi-linear. It could occur when the underlying instability wears  
 77 out its source. For example, the density profile or velocity shear can be flattened by a  
 78 turbulent density and momentum transport, respectively. Unstable small-scale flow ir-  
 79 regularities, comparable with the ion-acoustic gyroradius, drive the transport. Under the  
 80 second scenario, instability takes a more violent route by forming large coherent struc-  
 81 tures, which cut across the flow. This avenue of plasma transport is often associated with  
 82 strong eddies, blobs, streamers, or even avalanches (e.g., Carreras et al., 1998; Zweben  
 83 et al., 2003; Garcia et al., 2005; Dippolito et al., 2011). The avalanches promptly trans-  
 84 port material and heat across the main flow. These transport regimes relax the density

85 and temperature gradients much faster than in the quasi-linear scenario. For example,  
 86 the coherent flow disruption can occur in a matter of turnover time of large eddies. This  
 87 second possibility is often characterized as an avalanche transport regime. The first, rel-  
 88 atively benign, transport mechanism is often quantified by the Bohm diffusivity,  $D_B =$   
 89  $cT_i/16eB$ . Moreover, it can be further reduced to the so-called gyro-Bohm regime,  $D_{gB} \sim$   
 90  $D_B \rho_s / L_n$ . Here  $\rho_s = \sqrt{T_e / M \omega_{ci}^{-1}}$  is the ion-acoustic gyro-radius,  $\omega_{ci}$  is the ion gyro-  
 91 frequency, and  $L_n^{-1}$  is the characteristic density gradient.

92 However, flattening the *SAID may take more than a few hours* (e.g., He et al., 2017).  
 93 If the Kelvin-Helmholtz instability were advancing in the SAID shear flow, large eddies  
 94 would mix the density and temperature much earlier. In addressing this dilemma, it is  
 95 helpful to turn to its analogy with magnetic plasma confinement in laboratory devices,  
 96 such as tokamaks and stellarators, discussed in the above references. We have learned  
 97 that there is a third scenario in which the transport-driving turbulence exhibits an in-  
 98 verse cascade of unstable drift wave (DW) perturbations across the magnetic field that  
 99 self-organize in alternating streams comprising a shear flow. These streams, termed zonal  
 100 flows (ZF), are directed across the magnetic field and the instability-driving density/temperature  
 101 gradients. The flows react back to the wave instability and associated particle transport  
 102 in two ways. First, they stretch DW eddies by imposing a shearing motion on them. This  
 103 shortens the scale of the eddies across the ZF, thus leading to turbulence damping. How-  
 104 ever, the damping does not erase the DWs completely, as they maintain the ZF itself.  
 105 Even if a nearly complete turbulence suppression happens temporarily, the instability  
 106 resumes, and the process may evolve in a limit-cycle oscillation regime (e.g., Malkov et  
 107 al., 2001; E.-J. Kim & Diamond, 2003; Diamond et al., 2005; Malkov & Diamond, 2009;  
 108 Estrada et al., 2011; Schmitz et al., 2012; Dam et al., 2013).

109 Second, ZFs affect the plasma transport itself. It comes from the suppression, as  
 110 mentioned earlier, of the transport-driving DW eddies and their stretching across the den-  
 111 sity gradient. It follows that an efficient suppression of turbulent transport and preser-  
 112 vation of a plasma trough may be achieved by nonlinear feedback from the instability  
 113 to maintain the ZF. Mechanisms of self-organization and self-sustainability are well-known  
 114 in magnetic fusion and geophysical fluid dynamics research areas. Magnetically isolated  
 115 plasma structures supported by strong ExB shear flows have been studied for decades.  
 116 Their generation and control became an instrument of choice for the particle and heat  
 117 flux suppression in the magnetic fusion devices after the discovery of the so-called H-mode  
 118 plasma confinement in the ASDEX Tokamak (Wagner et al., 1982). Here, 'H' stands for  
 119 the "high" confinement regime that emerges from the L-mode (low) confinement regime  
 120 characterized by an enhanced turbulent transport across the field (see Diamond et al.,  
 121 2005; Hahn & Diamond, 2018; Groebner & Saarelma, 2023, for review). The transport  
 122 driving DW turbulence is generated by confined plasma's temperature and density gra-  
 123 dients, thus severely restricting the confinement quality and duration. The transition to  
 124 the H-mode occurs when Reynolds stresses, built up by the DW eddies, generate the above-  
 125 discussed ZFs.

126 Significant differences exist between the transition mechanisms to the H-mode in  
 127 magnetic fusion devices, resulting in steep density, temperature gradients, and ionospheric  
 128 troughs. However, building up Reynolds stresses in turbulence is generic. Its role in main-  
 129 taining shear flows along the walls of the troughs needs to be investigated once gradient-  
 130 driven instabilities are present in such systems. This mechanism enhances the trough life-  
 131 time and controls the spectrum of turbulence. It emerges from a delicate balance between  
 132 the turbulence energy input at short scales from the instability, an inverse cascade ac-  
 133 celerating the flow, and a forward cascade where the unused fluctuation energy dissipates.  
 134 These processes are coordinated in such a way as to leave the turbulence level at a bare  
 135 minimum, thus also minimizing the transport. Unsurprisingly, ZFs are not robust, but  
 136 they are not exceptional either. Once established, they typically remain stable and self-  
 137 sustaining within certain boundaries in the parameter space. Upon crossing these bound-

aries, a strong Kelvin-Helmholtz instability may set in, resulting in large eddies that disrupt the flow. Overall, this process would constitute a tertiary instability of the initial plasma configuration (Rogers et al., 2000; Numata et al., 2007). The first two are the DW instability driven by the density and pressure gradients, followed by the secondary instability constituting the amplification of the ZF.

The objectives of the present study are (i) to extend our understanding of the mechanisms whereby significant density and temperature contrasts are maintained in the subauroral and other flows in the ionosphere, (ii) to investigate their morphology and establish parameter requirements for their generation by possible magnetospheric drivers, and (iii) to identify possible common physics behind various flows in the ionosphere. Special attention will be devoted to the spectra of turbulence resulting from the instability and structure formation.

We shall primarily pursue a numerical approach in this investigation. Specifically, we will employ a system of fluid equations similar to those used in describing the DW-ZF turbulence, as, e.g., the Hasegawa-Wakatani (HW) system (Hasegawa & Wakatani, 1987) in magnetic fusion energy research and geostrophic flows. Together with its earlier subset, known as the Charney-Hasegawa-Mima model, the HW model has been broadly used for more than four decades in studying DW turbulence in plasmas and Rossby waves in atmospheric and oceanic dynamics.

At first, however, we provide a modified system to better suit the ionospheric conditions and lay the ground for future model development. In addition to the electrostatic potential and particle density operative for the zonal flow generation in the HW model, we include the parallel flow velocity and plasma pressure. While manifesting their subdominant role for this study, in that they can be calculated when the density and flow potential are obtained, they can be straightforwardly included in the code self-consistently, which we plan to do in the next model iteration. These fields are essential in connecting the ionosphere with the magnetospheric driver.

The paper is organized as follows. The HW model is introduced in the next section, and relevant aspects of its derivations are briefly discussed. In Sec.3, the linear stability analysis using a simple WKB-type approach is carried out. Sec.4 describes the simulation setup. The simulation results are presented in Sec.5 in two separate categories. The first set of simulations was executed using a symmetric density trough as an initial condition. While these results are limited to narrow troughs, the second category deals with the relaxation of a single-density wall. The results of these studies shed extra light on the relaxation of arbitrarily wide troughs and other ionospheric structures with sharp unidirectional density gradients. The paper closes with a summary and discussion of the results.

## 2 Model Equations

We describe the plasma dynamics using gyro-fluid type equations with resistivity in the plasma motion parallel to the magnetic field. We start by extending the HW (Hasegawa & Wakatani, 1987) two-field system of equations for the plasma density and electrostatic potential. The extended system also evolves the plasma pressure and the ion velocity component along the magnetic field. Depending on the spectrum extension to large wave numbers studied in this paper, we will also include finite Larmor radius (FLR) effects using the gyro-fluid methodology in the next version of this model.

The underlying idea behind the derivation of the HW model is to use the quasineutrality condition and to connect the density and potential perturbations, in the form  $\nabla \cdot j = 0$ . They also utilize the so-called “diamagnetic cancellations” between the FLR momentum heat flows and the diamagnetic velocity fluxes (e.g., Hinton & Horton, 1971), after obtaining the flow velocity from a drift approximation. In terms of the ordering,

188 the fields are decomposed in a usual manner, as, e.g., for the density  $n = n_0 + n'$ , with  
 189  $n' \ll n_0$  but  $\nabla n' \sim \nabla n_0$ . Keeping the above in mind, we adhere to a more formal but  
 190 systematic approach, starting from the Braginsky system (Braginskii, 1965).

191 The primary particle motion occurs in the forms of the ExB and diamagnetic drifts

$$\mathbf{V}_E = \frac{c}{B^2} \mathbf{B} \times \nabla \Phi, \quad \mathbf{V}_{\text{di}} = \frac{c}{enB^2} \mathbf{B} \times \nabla P_i. \quad (1)$$

192 Here  $\Phi$  is the full electrostatic potential, and  $P_i$  is the ion pressure. *In intense SAID*  
 193 *channels, the electron temperature typically exceeds ion temperature by a factor of 3-5*  
 194 *at heights above 300 km, while the ion temperature dominate at lower altitudes* (e.g., Liang  
 195 et al., 2021, 2022; E. V. Mishin & Streltsov, 2024). In this paper, we will evolve only the  
 196 density profile and keep the temperature profiles of both species fixed.

197 The velocity components in eq.(1) are largely incompressible and do not provide  
 198 evolution when inserted in the quasineutrality,  $\nabla \cdot \mathbf{j} = 0$ , equation. We, therefore, need  
 199 to add the ion polarization drift to  $V_{\text{di}}$ :

$$\mathbf{V}_p = -\frac{c^2 M}{eB^2} \left[ \frac{\partial}{\partial t} + (\mathbf{V}_E + \mathbf{V}_{\text{di}}) \cdot \nabla \right] \nabla \Phi$$

200 Using these expressions, including the diamagnetic cancellations mentioned above, the  
 201 continuity equation for ions, their parallel velocity  $V_{\parallel}$ , and thermal balance equations  
 202 can be written as follows (e.g., Horton, 1990):

$$\frac{\partial n_i}{\partial t} + \mathbf{V}_E \cdot \nabla n_i + \nabla_{\perp} \cdot (n_i \mathbf{V}_p) + \frac{\partial}{\partial z} n_i V_{\parallel} = 0 \quad (2)$$

$$\frac{\partial V_{\parallel}}{\partial t} + \mathbf{V}_E \cdot \nabla V_{\parallel} + V_{\parallel} \frac{\partial V_{\parallel}}{\partial z} = -\frac{e}{M} \frac{\partial \Phi}{\partial z} - \frac{1}{Mn} \frac{\partial P_i}{\partial z} - \nu_{\text{in}} V_{\parallel} \quad (3)$$

$$\frac{\partial P_i}{\partial t} + \mathbf{V}_E \cdot \nabla P_i + V_{\parallel} \frac{\partial P_i}{\partial z} + \gamma P_i \frac{\partial V_{\parallel}}{\partial z} = 0 \quad (4)$$

203 where  $\gamma$  is the ion adiabatic index, which can be fixed, e.g., at  $\gamma = 5/3$ ,  $P_i = nT_i$ ,  $\nu_{\text{in}}$   
 204 is the ion-neutral collision frequency, with other notation being standard. We have dropped  
 205 small particle and heat diffusivity terms, along with the molecular viscosity term. These  
 206 are expected to be predominantly turbulent, thus implicitly present in the above system.  
 207 However, we will use small dissipative terms for regularization purposes in computations.  
 208 After dropping some nonessential nonlinear terms discussed, e.g., in (Horton, 1990), eq.(2)  
 209 rewrites

$$\frac{\partial n_i}{\partial t} + \frac{c}{B} [\Phi, n_i] + \frac{c^2 M}{eB^2} \left( -n_0 \frac{\partial}{\partial t} \Delta_{\perp} \Phi - \frac{c}{B} n_0 [\Phi, \Delta_{\perp} \Phi] - \frac{c}{eB} [P_i, \Delta_{\perp} \Phi] \right) + \frac{\partial}{\partial z} n_i V_{\parallel} = 0$$

210 We have introduced the Poisson bracket as

$$[f, g] = \frac{\partial f}{\partial x} \frac{\partial g}{\partial y} - \frac{\partial f}{\partial y} \frac{\partial g}{\partial x}$$

211 Next, using the quasineutrality condition,  $n_e = n_i = n$ , and the electron continuity  
 212 equation within the guiding center approximation in eqs.(1),

$$\frac{\partial n_e}{\partial t} + \frac{c}{B} [\Phi, n_e] + \nabla_{\parallel} n_e V_{e\parallel} = 0, \quad (5)$$

213 where  $V_{e\parallel}$  is the parallel electron velocity, we relate the density and potential perturba-  
 214 tion as follows

$$\frac{c^2 M}{B^2} \left( n_0 \frac{\partial}{\partial t} \Delta_{\perp} \Phi + \frac{c}{B} n_0 [\Phi, \Delta_{\perp} \Phi] + \frac{c}{eB} [P_i, \Delta_{\perp} \Phi] \right) - \nabla_{\parallel} \cdot \mathbf{j} = 0 \quad (6)$$

215 Here,  $j_{\parallel}$  can be straightforwardly related to  $P$  and  $\Phi$  from the electron momentum bal-  
 216 ance:

$$j_{\parallel} = \frac{e}{m\nu_{ei}} (\nabla_{\parallel} P_e - en_e \nabla_{\parallel} \Phi)$$

217 Using the last three relations, we rewrite eqs.(5) and (6) as follows:

$$\frac{\partial n}{\partial t} + \frac{c}{B} [\Phi, n] - \frac{1}{m\nu_{ei}} (\nabla_{\parallel}^2 P_e - en_0 \nabla_{\parallel}^2 \Phi) = 0$$

$$\frac{\partial}{\partial t} \Delta_{\perp} \Phi + \frac{c}{B} [\Phi, \Delta_{\perp} \Phi] + \frac{c}{eBn_0} [P_i, \Delta_{\perp} \Phi] - \frac{M\omega_{ci}^2}{m\nu_{ei}en_0} (\nabla_{\parallel}^2 P_e - en_0 \nabla_{\parallel}^2 \Phi) = 0, \quad (7)$$

218 where  $\omega_{ci}$  is the ion Larmor frequency and  $\nu_{ei}$  is the electron-ion collision frequency. We  
 219 have not included collisions with neutrals in the friction force acting on electrons, assum-  
 220 ing that the  $\nu_{en}/\omega_{ce}$  is very small, *which is a valid assumption in the top ionosphere*. The  
 221 ionospheric flows of our interest often have a stronger density inhomogeneity than that  
 222 of the temperature by almost an order of magnitude. Under these conditions, the above  
 223 two equations form a closed system if we approximate the pressure as  $P_{e,i} = T_{e,i}n$ , as-  
 224 suming the base temperature profile  $T = T_0(y)$  (but not the density) being fixed, where  
 225  $y$  is the latitude. Note, however, that we call  $y$  latitude and  $x$  longitude here and below  
 226 just for simplicity. A more accurate definition of these coordinates is that  $y$  is aligned  
 227 with the density gradient,  $x$  is aligned with the constant density surface, and  $z$  is in the  
 228 magnetic field direction. In real geometry, especially in equatorial regions, this coordi-  
 229 nate system may be turned around the magnetic field and misaligned with longitude and  
 230 altitude.

231 One advantage of neglecting the temperature variations is the similarity of this case  
 232 with a well-studied HW (Hasegawa & Wakatani, 1987) model of the DW turbulence. This  
 233 model assumes cold ions, so the third term in eq.(7) is absent. We keep it for assessing  
 234 its impact on the linear instability but will drop it for simplicity when performing the  
 235 simulations. Introducing the following dimensionless variables

$$\omega_{ci} t \rightarrow t; \quad x, y, z/\rho_s \rightarrow x, y, z;$$

236 where  $\rho_s = \sqrt{T_{00}/M}\omega_{ci}^{-1}$ , with  $T_{00}$  being the fiducial electron temperature, as is the  
 237 density  $n_{00}$  introduced below, we relate them to one of the  $y$ - ends of the  $(x, y)$  integra-  
 238 tion box. Depending on the particular simulation setup, these quantities may be the same  
 239 or different at the endpoints. Here, we adopt a geophysical convention of the equilibrium  
 240 gradients pointing to  $y$  rather than  $x$  (as in plasma literature) direction. Using the above  
 241 notation, we arrive at the HW equations, supplemented by the ion pressure term. We  
 242 will also normalize the dependent variables  $e\Phi/T_{00} \rightarrow \Phi$ ,  $n/n_{00} \rightarrow n$ . The term stem-  
 243 ming from the ion pressure  $P = P_i/n_{00}T_{00}$  will be used only for the linear instability.  
 244 The dimensionless equations then read:

$$\frac{\partial}{\partial t} \Delta_{\perp} \Phi + [\Phi + P, \Delta_{\perp} \Phi] + \alpha (\tilde{n} - \tilde{\Phi}) = 0 \quad (8)$$

$$\frac{\partial n}{\partial t} + [\Phi, n] + \alpha (\tilde{n} - \tilde{\Phi}) = 0 \quad (9)$$

245 We have also subtracted the zonal averaged,  $\langle \cdot \rangle$ , quantities in the terms under the  $\nabla_{\parallel}$   
 246 -operator by introducing

$$\tilde{n} = n - \langle n \rangle, \quad \tilde{\Phi} = \Phi - \langle \Phi \rangle$$

247 and replacing

$$\alpha = -\frac{V_{Te}^2}{\nu_{ei}\omega_{ci}}\nabla_{\parallel}^2 \rightarrow \frac{k_{\parallel}^2 V_{Te}^2}{\nu_{ei}\omega_{ci}}.$$

248 Subtracting  $\langle n \rangle$  and  $\langle \Phi \rangle$  from  $n$  and  $\Phi$  in the last terms of the above equations is jus-  
 249 tified in the magnetic plasma confinement contexts by the magnetic field shear that mixes  
 250 the zonally averaged variations of the respective quantities. In the SAPS applications,  
 251 this modification can be justified by a strong shear in the vertical velocity  $V_{\parallel}(y)$  (Huang  
 252 et al., 2007). As the magnetic shear in Tokamaks, this effect does not enter the HW sys-  
 253 tem explicitly, but it should also be taken into account in calculating the electron response.

254 Under zonally-averaged quantities, we understand averaging along the homogeneous  
 255 direction of the problem (i.e., longitude,  $x$ ) while the quasi-stationary variations of  $n$  and  
 256  $\Phi$  are in  $y$ -direction. It is assumed that  $\nabla_{\parallel} \langle \cdot \rangle \equiv 0$ . The rationale behind this assump-  
 257 tion is that the zonally-averaged perturbations vary much slower in time than  $\tilde{n}$  and  $\tilde{\Phi}$   
 258 and, therefore, they spread along the field during the time so that their variation in this  
 259 direction is much smoother:  $\nabla_{\parallel} \langle n \rangle, \nabla_{\parallel} \langle \Phi \rangle \ll \nabla_{\parallel} \tilde{n}, \nabla_{\parallel} \tilde{\Phi}$ . The simplification associ-  
 260 ated with the replacement of the  $\nabla_{\parallel}$ - operator by a constant can be justified, assuming  
 261 that the perturbations  $\tilde{n}$  and  $\tilde{\Phi}$  vary in  $z$ - direction as  $\exp(ik_{\parallel}z)$  and there exists a max-  
 262 imum growth rate in  $k_{\parallel}$ , thus defining  $\alpha$ , which we confirm in the next section. While  
 263 making this extension to the third dimension in the ambient magnetic field direction,  
 264 we simplify the dynamics across it by neglecting temperature and magnetic field gradi-  
 265 ents. Their contribution to the dynamics, including a gravity effect of magnetic field cur-  
 266 vature has been investigated in conjunction with the flute mode generation in the iono-  
 267 sphere (Sotnikov et al., 2008; V. Sotnikov et al., 2014). For similar reasons, the verti-  
 268 cal velocity  $V_{\parallel}$  is considered subdominant relative to the main  $n$  and  $\Phi$  variables, and  
 269 we do not evolve  $V_{\parallel}$  in this simple version of the HW model. It can be obtained *a pos-*  
 270 *teriori* from eq.(3) using  $\Phi$  obtained from the main model eqs.(8-9). Next, we examine  
 271 the linear stability of our model.

### 272 3 Linear Instability

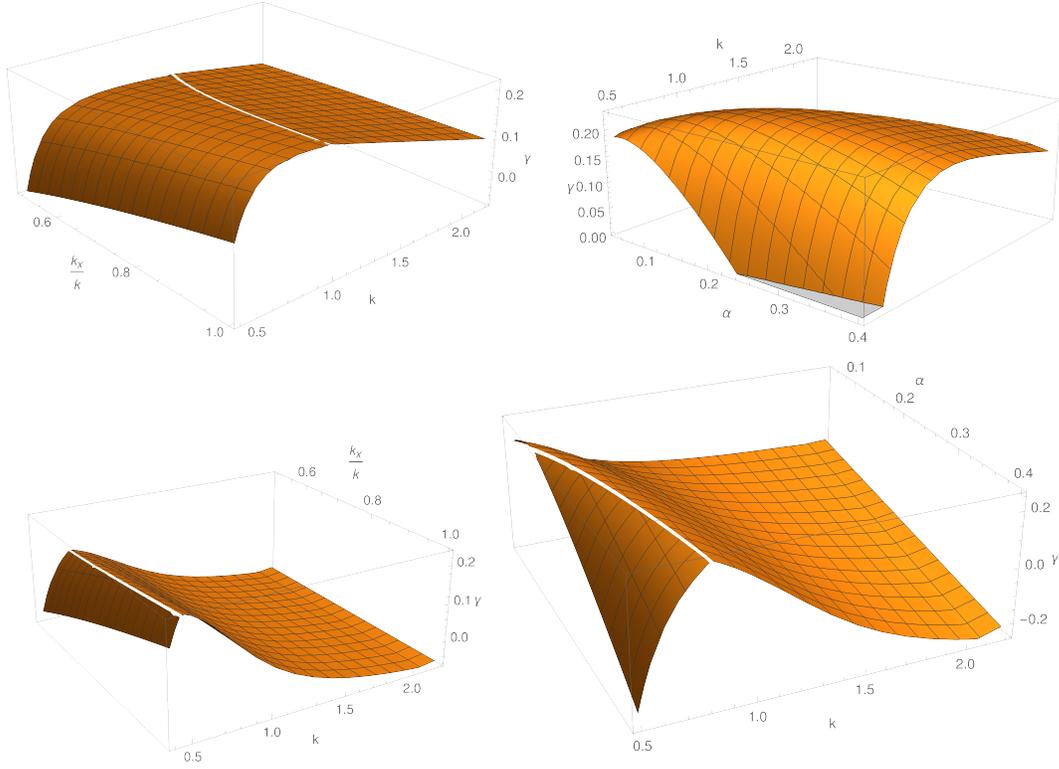
273 In this section, we perform a linear stability analysis of eqs.(8-9) using a simple short-  
 274 wave approximation. Namely, we consider the following instability driving parameters,  
 275 associated with the density and pressure gradients as constants, along with the adiabatic-  
 276 ity parameter  $\alpha$ , introduced earlier:

$$\kappa = \frac{1}{n_0 T_{0e}} \frac{\partial}{\partial y} n_0 T_{0i} \approx const, \quad \beta = \frac{1}{n_0} \frac{\partial}{\partial y} n_0 \approx const,$$

277 Here the zonal averaged profiles  $T_0(y)$  and  $n_0(y)$  are also normalized to  $T_{00}$  and  $n_{00}$ .  
 278 Linearizing then eqs.(8-9) with respect to  $\tilde{n}$  and  $\tilde{\Phi}$ , we obtain

$$\frac{\partial}{\partial t} \Delta_{\perp} \tilde{\Phi} - \kappa \Delta_{\perp} \tilde{\Phi}_x + \alpha (\tilde{n} - \tilde{\Phi}) = 0$$

$$\frac{\partial \tilde{n}}{\partial t} + \beta \tilde{\Phi}_x + \alpha (\tilde{n} - \tilde{\Phi}) = 0$$



**Figure 2.** **Left upper panel:** Instability growth rate as a function of wave number  $k$  and  $k_x$  - its components across plasma gradients and magnetic field, computed for  $\alpha = 0.1$ . **Right upper panel:** Growth rate shown as a function of  $k$  and  $\alpha$ , for  $k_x = k$ . The remaining parameters in eq.(10) are fixed at  $\beta = 2$ ,  $\kappa = 0.2$ . **Lower panels** The same as upper panels but for  $\kappa = 1.2$ . White lines on the surfaces mark a change of the square root branch in eq.(11) where the denominator of the expression under  $\tan^{-1}$  crosses zero.

279 Using the last two equations, we derive the following dispersion equation for the pertur-  
 280 bations of the form  $\exp(-i\omega t + ik_x x + ik_y y)$

$$\omega^2 + [\kappa k_x + i\alpha(k^{-2} + 1)]\omega + i\alpha k_x(\kappa - \beta k^{-2}) = 0 \quad (10)$$

281 The two roots for the frequency  $\omega$  can be written as

$$\omega = -i\alpha(1 + k^{-2}) - \kappa k_x \pm \sqrt{\kappa^2 k_x^2 - \alpha^2(1 + k^{-2})^2 + \frac{2i\alpha k_x(2\beta + \kappa - \kappa k^2)}{k^2}}, \quad (11)$$

282 while the growth rate

$$\gamma = \Im\omega = -\alpha(1 + k^{-2}) + \sqrt{r} \left| \sin \frac{\vartheta}{2} \right|,$$

where

$$r = \sqrt{\left(\kappa^2 k_x^2 - \frac{\alpha^2(k^2 + 1)^2}{k^4}\right)^2 + \left(\frac{2\alpha k_x(2\beta + \kappa - \kappa k^2)}{k^2}\right)^2},$$

$$\theta = \tan^{-1} \left[ \frac{2\alpha k^2 k_x (2\beta + \kappa - \kappa k^2)}{\kappa^2 k^4 k_x^2 - \alpha^2 (k^2 + 1)^2} \right]$$

283 This solution is shown in Fig.2 in two different representations for fixed values of param-  
 284 eters  $\beta = 2$  and two values of  $\kappa = 0.2$  and  $\kappa = 1.2$ . In the latter case, the excitation  
 285 zone is narrower and at lower values of  $k$  than in the former case. Note that smaller val-  
 286 ues of  $\kappa \ll \beta$  correspond to a temperature gradient being as strong but opposite to the  
 287 density gradient. This equilibrium corresponds to a pressure-balanced structure across  
 288 the magnetic field,  $nT \approx \text{const}$ , whereas the case  $\kappa \approx \beta$  corresponds to the case of con-  
 289 stant temperature. It is also worth mentioning that the growth rate maximizes at  $k_x \approx$   
 290  $k$ , and it also has a broad maximum in  $\alpha \sim 1$ , which selects a range of most unstable  
 291 waves in  $k_z$ . This is another justification of a single adiabaticity parameter  $\alpha$ , introduced  
 292 in the previous section.

293 The above instability analysis is oversimplified where the  $y$ - scale of perturbations  
 294 is not much shorter than the domain size as it does not adequately describe the eigen-  
 295 functions of the density and potential perturbations. Nevertheless, it is useful as a guid-  
 296 ance tool to select appropriate parameter values for simulations that we describe in the  
 297 next section.

## 298 4 Simulation Setup

299 We perform our simulations in a channel box that is periodic in  $x$ - direction (“lon-  
 300 gitude”) and has rigid boundaries in  $y$  direction (“latitude”, see text below eq.[7]). This  
 301 choice of simulation geometry reflects a density trough morphology, discussed in the In-  
 302 troduction. Given the computational limitations on the spatial scales, we have split the  
 303 simulations into two complementary sets. The first set addresses the sustainability of an  
 304 entire density trough against turbulent transport generated by gradient-driven instabil-  
 305 ities, developing primarily at its walls. Within this setting, the edge density is maintained  
 306 at the same level on both sides of the trough. The rate at which the influx of the outer  
 307 plasma refills the trough will quantify the efficiency of transport regulation by generated  
 308 ZFs.

309 Due to the vast difference in scales between a typical density trough latitudinal pro-  
 310 file and gradient-driven turbulence, the above-described simulation setup cannot ade-  
 311 quately address the formation and sustainability of a well-pronounced density trough.  
 312 In other words, the above full-trough simulation setup can capture only narrow troughs  
 313 and, thereby, relatively shallow ones with a density contrast, say,  $\Delta n \lesssim 10$ . Although  
 314 these studies are of significant interest, many troughs are at least a few degrees in lat-  
 315 itude. The question of our results’ scalability for the narrow troughs arises. Moreover,  
 316 our goal is to link the macroscopic trough sustainability to microturbulence excited at  
 317 its walls. If not self-regulated or otherwise suppressed, the turbulence would lead to a  
 318 rapid refill of the trough. However, to investigate how microturbulence generates mesoscale  
 319 ZFs in full-size density troughs would be computationally very expensive.

320 To alleviate this problem, the second set of simulations will be employed. It zooms  
 321 in the area of a sharp density gradient at one of the two walls of the trough. This setup  
 322 is motivated by observations of troughs having step-like walls that, as it seems, do not  
 323 interact with each other. We simulate an isolated wall by using an adjacent flat bottom  
 324 and top of the density profile as boundary conditions (BC). Such structures are often  
 325 observed (cf. Fig.1). In this setting, much fewer computational resources must be invested,  
 326 as it focuses on a steep part of the density profile that does not spread significantly in  
 327 the flat areas.

328 Based on the above considerations, the boundary condition in the second simula-  
 329 tion is set to maintain a fixed density contrast between the  $y$ - boundaries of the box, start-

330 ing from a step-like density distribution. In this case, particle flux initially drives the tur-  
 331 bulance where the gradient is sufficiently large. This setup is suitable for demonstrat-  
 332 ing a mesoscale character of zonal flow generation and flux suppression. It provides a  
 333 necessary link between microturbulence and the full-scale density trough phenomenon.

334 Returning to the full-trough simulation, we will start it from an inverted bell-shaped  
 335 density profile. Morphologically, it corresponds to a trough event, shown, e.g., in Fig.1  
 336 of (Foster & Burke, 2002) between  $50\text{--}54^\circ$  longitude, but much narrower in longitude.  
 337 Indeed, other data show local density variation in the troughs over scales less than  $0.5^\circ$   
 338 (Huang et al., 2007), which are often unresolved, e.g., Fig.1, at latitudes  $-41^\circ$  and  $-59^\circ$ .  
 339 We set periodic BC in  $x$ - direction,  $n(x + L_x, y, t) = n(x, y, t)$  and apply the same condi-  
 340 tion for  $\Phi$ . In two sets of simulations, we will use boundary conditions with and with-  
 341 out density contrast across the channel for the step-like and trough-like initial density  
 342 profiles. The stream function,  $\Phi = 0$  at  $y = 0, L_y$ . We start with a symmetric trough  
 343 with equal density values at the edges and a minimum in the middle. At  $t = 0$  a ran-  
 344 dom noise of weak fluctuations is added to this density profile. The initial potential,  $\Phi$ ,  
 345 contains only a weak random noise.

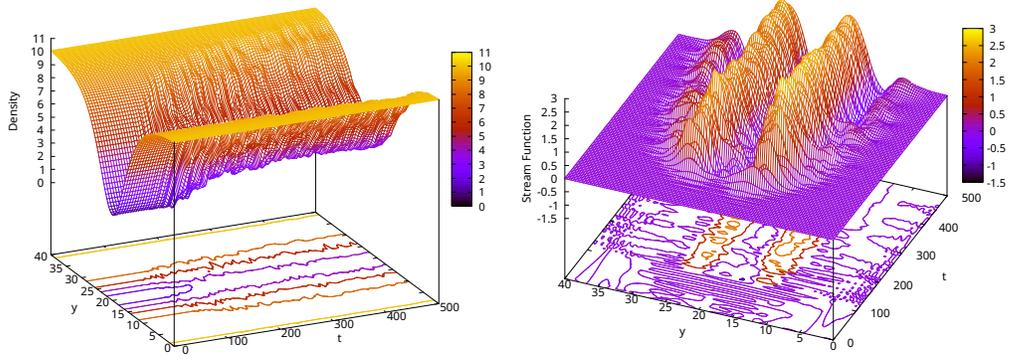
346 It should be noted, however, that most of the past numerical studies of systems sim-  
 347 ilar to eqs.(8-9) have been carried out in a flux-driven, doubly-periodic simulation box.  
 348 To model the relaxation of a density trough imposed as an initial condition, it is more  
 349 appropriate to apply the periodicity condition only across the plasma gradient. Apart  
 350 from a background density profile, the initial conditions are given by statistically isotropic,  
 351 random perturbations suppressed toward large wave numbers. The hyperviscosity/hyperdiffusivity  
 352 coefficients (standard for the spectral codes, not shown in eqs.[8-9] for clarity) are ap-  
 353 plied to suppress the “spill over” of higher harmonics.

354 The most runs presented in this paper were performed on a grid with  $N_x \times N_y = 512 \times 512$   
 355 spatial points. However, they have been verified using limited runs on  $512 \times 1024$  grids,  
 356 which we also used for studies of fluctuation Fourier spectra. Given that the typical spec-  
 357 tra obtained from simulations, e.g., those presented in Sec.5.6 are very steep, the code  
 358 resolution is sufficient for the phenomena studied. Mode de-aliasing is achieved by trun-  
 359 cating the number of modes at approximately  $N_{x,y}/3$ . For example, we truncate the spec-  
 360 trum at  $160 \times 320$  Fourier harmonics for the finest resolution choice indicated above.

361 We have solved the system in eq.(8-9) using appropriate Fourier decompositions.  
 362 According to the Dirichlet boundary condition across the channel between  $y = 0, L_y$ ,  
 363 a mere sine expansion suffices for the  $y$ -direction. The  $y$ -dependent background den-  
 364 sity profile enters the equation only as a product with the plasma potential that vanishes  
 365 at both  $y$  boundaries. Therefore, an additional cosine expansion is not needed. We uti-  
 366 lize the spectral library ISPACK, documented in, e.g., [http://www.gfd-dennou.org/  
 367 library/spmodel/#label-2](http://www.gfd-dennou.org/library/spmodel/#label-2). The time-stepping is based on the second-order explicit  
 368 Adams–Bashforth algorithm.

## 369 5 Simulation Results

370 Before turning to the simulation results, discussing their scalability concerning the  
 371 box size and the density contrast imposed as a boundary and initial condition is worth-  
 372 while. Because of computational constraints, we cannot change these parameters at will  
 373 as a sizable rescaling exerts much stress on the code, thus increasing the computation  
 374 costs. In addition, we will precede our presentation of simulation results by a brief sub-  
 375 section of their outcome.



**Figure 3.** Time evolution of the density and potential averaged in zonal direction ( $x$ , periodic coordinate). The initial density profile is set to a V-shaped plasma depletion with a minimum at the middle of the trough, where the density  $n_0 = 1$ . The density is maintained at  $n_1 = 10$  at both edges of the trough. The adiabaticity parameter  $\alpha = 1.5$ .

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### 5.1 Scalability Issues

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Even the simplest version of eqs.(8-9) with  $P = 0$  does not allow us to rescale a given simulation result to larger boxes without rerunning the code. Indeed, by rescaling coordinates, time, and dependent variables as  $x, y, t \rightarrow \lambda x, \lambda y, \tau t$  and  $\Phi \rightarrow (\lambda^2/\tau) \Phi$ ,  $n \rightarrow (\lambda^2/\tau) n$ , respectively, we arrive at the following system

$$\frac{\partial}{\partial t} \Delta_{\perp} \Phi + [\Phi, \Delta_{\perp} \Phi] + \lambda^2 \tau \alpha (\tilde{n} - \tilde{\Phi}) = 0 \quad (12)$$

$$\frac{\partial n}{\partial t} + [\Phi, n] + \tau \alpha (\tilde{n} - \tilde{\Phi}) = 0 \quad (13)$$

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We have chosen the  $\lambda^2/\tau$  scaling factor to keep the convective parts of equations invariant. We also limit our discussion to the case of an equal box transformation in both directions by the same factor  $\lambda$ .

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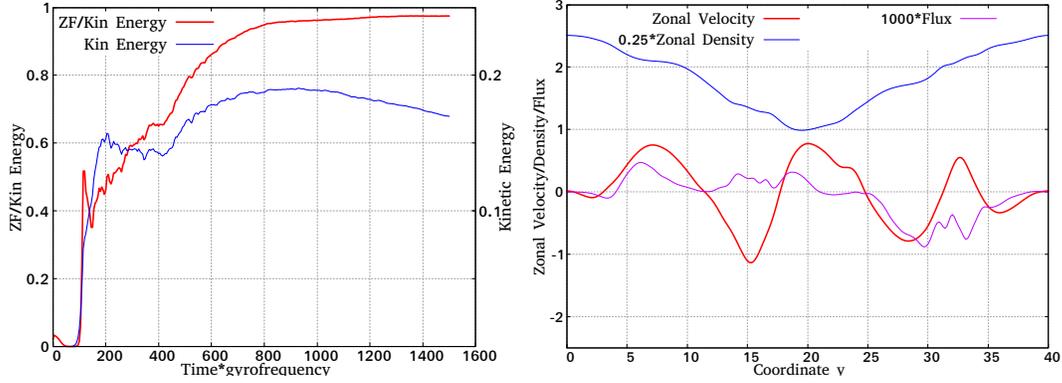
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Assume that we have performed a system scan in  $\alpha$  using our code for  $\lambda = \tau = 1$ . However, mapping  $\tau \alpha \rightarrow \alpha$ , for  $\tau \neq 1$  does not make the system invariant to the above scale transformations. Since we originally normalized the spatial scale to the ion sound Larmor radius,  $\rho_s$ , which is relatively small for the processes at hand, our interest is to increase the box size, that is, to get a handle on the case  $\lambda \ll 1$ . Therefore, any result obtained before rescaling, that is, for  $\lambda = \tau = 1$  cannot be mapped to the system behavior in significantly larger boxes, even if we ignore the additional problems associated with the insufficient short-scale resolution that arise from the box stretching without mesh refinement. As seen from the above equations, for  $\lambda \ll 1$  the density perturbations are subdominant to the potential and convected with the flow almost as a passive scalar. Note that stretching of the spatial variable generally reduces the growth rate, but increasing  $\tau$  does not alleviate the problem of the system's bias in favor of  $\Phi$ . These considerations will help us choose the simulation parameters and plan for the next steps. A simple way to restore the role of  $\Phi$  in the dynamical feedback loop of the system is to reinstate the ion pressure,  $P$ , in the equation for  $\Phi$  that we dropped in eq.(12). We defer these results to a future report.



**Figure 4.** Left panel: Time evolution of the ratio of the box-averaged ZF kinetic energy to the full kinetic energy (red) and kinetic energy (blue). Right panel: Zonally-averaged (over  $x$ ) profiles of density, velocity, and turbulent flux in  $y$ -direction at  $t = 1500$ , when most of the plasma kinetic energy is in the ZF component. Other simulation parameters, IC and BC, are the same as in Fig.3.

## 5.2 Expected Simulation Outcome

Notwithstanding the above limitations, the two-field model given by eqs.(8-9) and discussed in Sec.2 captures distinctively different turbulence and transport scenarios, including those described in the Introduction. Extensive studies of similar systems carried out primarily in the areas of magnetic fusion energy research and geostrophic flows can be distilled to two different transport regimes. An initially unstable plasma profile can relax to either of them by a turbulent transport, unstably driven by the plasma density or temperature gradients. However, the depth of the plasma relaxation will be markedly different in these regimes.

The first regime is characterized by strong-turbulence transport, driven by the DW instability. It leads to flattening the initially unstable density or temperature profile to a marginally stable one with strong residual turbulence and associated transport. The second regime evolves through a phase of self-organization of excited turbulence. Most of its energy is channeled into zonal flows that suppress the turbulence level and particle transport while the initially strong density gradient persists or suffers only a moderate flattening.

In distinguishing between these two regimes, there is no significant difference in whether the instability is driven by temperature and density gradient or just one of them. We, therefore, limit our study to a straightforward case in which the density gradient drives the instability, whereas the ion pressure is neglected compared to the electron pressure. The latter enters the equations through the normalization of  $\Phi$  and the adiabaticity parameter  $\alpha$ . Since only subdominant plasma temperature variations are typically observed in density troughs, we impose the condition  $P = 0$  in eq.(8). This reduces the number of instability control parameters from three to two, as we set  $\kappa = 0$  in the instability analysis presented in Sec.3. According to this analysis, the ion pressure gradient has no profound effect on the instability growth rate.

At the same time, omitting the ion pressure greatly simplifies the energetics of the system at hand. Indeed, we can define the total energy as follows:

$$E_{\text{tot}} \equiv \frac{1}{2} \int \int \left[ |\nabla\Phi|^2 + (n - \eta)^2 \right] dx dy, \quad (14)$$

428 where  $\eta$  is an arbitrary constant. This constant can be introduced because eqs.(8-9) are  
 429 invariant against the replacement  $n \rightarrow n + \text{const}$ . Meanwhile,  $E_{\text{tot}}$  strictly decreases  
 430 with time:

$$\frac{dE_{\text{tot}}}{dt} = -\alpha \int \int (\tilde{\Phi} - \tilde{n})^2 dx dy$$

431 As indicated earlier, we start the simulation with a weak noise of  $\tilde{\Phi}$  and  $\tilde{n}$ , and a den-  
 432 sity profile  $n(0, y) = n_0(y)$ , with a typically high density contrast across the channel,  
 433  $n_{\text{max}} \gg n_{\text{min}}$ . The flow potential can thus grow significantly only if the instability de-  
 434 creases the box-averaged value of  $(n - \eta)^2$ . Since at  $t = 0$ ,  $\Phi \approx 0$ , the most informa-  
 435 tive choice of  $\eta$  with regard to the grows of  $|\nabla\Phi|^2$ , is to minimize the second term un-  
 436 der the integral in eq.(14). This requirement leads to  $\eta = \bar{n}_0$ , where the bar denotes  
 437 an averaged value over the interval  $0 < y < L_y$ .

438 Another integral quantity that constrains the energy pathways of the system is based  
 439 on the conservation of any arbitrary function,  $F$ , of the potential vorticity,  $w = \Delta\Phi -$   
 440  $n$ :

$$\frac{d}{dt} \int \int F(w) dx dy = 0$$

441 For example, we can use the enstrophy

$$\mathcal{E} = \frac{1}{2} \int \int (\Delta\Phi - n)^2 dx dy = \text{const},$$

442 to control the code accuracy.

### 443 5.3 Transport Barriers in Density Troughs

444 To understand the density profile evolution in a trough, which undergoes unsta-  
 445 ble growth of DWs presumably driven at its walls, we set the following initial conditions  
 446 in a box  $(0, 0, L_x, L_y)$ . To mimic a simple symmetric trough profile, the simulation starts  
 447 with a V-shaped latitudinal density distribution of the form

$$n(y) = n_1 - \frac{n_2}{\cosh\left(\frac{y-L_y/2}{L_y\delta}\right)},$$

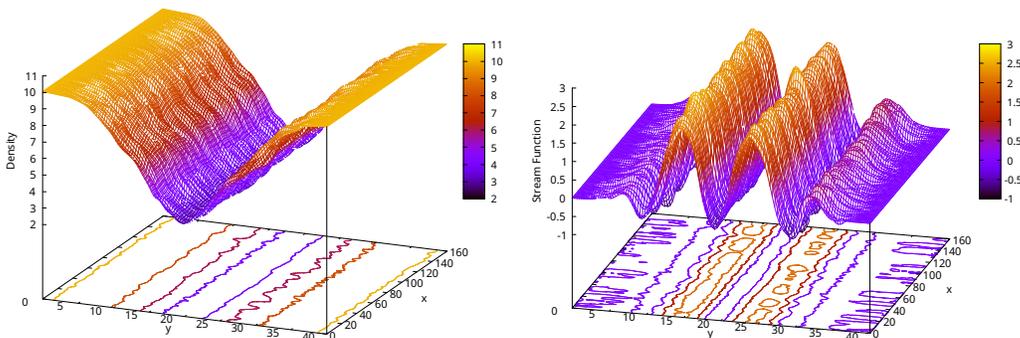
448 and zero mean potential. Both quantities are, however, superposed by small-amplitude  
 449 random fluctuations added as Fourier harmonics with random phases and amplitudes.  
 450 To accelerate the initial phase of the system evolution, we suppress the random harmon-  
 451 ics exponentially toward the highest wave numbers. The above density profile typically  
 452 stays frozen, and the plasma potential fluctuations do not grow above the initial noise  
 453 level for the first tens of time units ( $\omega_{ci}^{-1}$ ). After the most unstable modes grow to am-  
 454 plitudes sufficient to produce significant anomalous diffusion, the density profile starts  
 455 rapidly flattening. However, in the case of  $\alpha \sim 1$ , the ZF also grows so that the den-  
 456 sity flattening stops because of the backreaction from the ZF. If the established ZF is  
 457 strong, the zonally-averaged (in  $x$  direction) density profile freezes again, apart from its  
 458 turbulent fluctuations around the mean value.

459 An example of the density trough relaxation outlined above is shown in Fig.3. Dur-  
 460 ing the first 50-100 time units, the gradient-driven instability grows sufficiently for the  
 461 turbulent diffusion to refill the bottom of the density trough. The refill material is sup-  
 462 plied from the steep parts of the initial density profile where the instability is the strongest.  
 463 The areas adjacent to the edges of the trough remain undisturbed since the density gra-  
 464 dient there is subcritical regarding the instability. This behavior is consistent with the

465 instability analysis in Sec.3. The surface plots show that the turbulence, time-averaged  
 466 zonal density, and potential profiles rapidly spread toward the trough edges. Had they  
 467 proceeded at the same pace, the relaxed part of the density profile would have reached  
 468 the trough edges in a few tens of time units. Instead, the system comes to a quasi-stationary  
 469 state after  $t \approx 200$ , so the relaxation front does not reach the edges. The density gra-  
 470 dient relaxation proceeds further for smaller  $\alpha \ll 1$ , which we will consider in the next  
 471 section by focusing on one wall of a trough.

472 Time-asymptotic profiles of zonally-averaged density, velocity, and a strongly sup-  
 473 pressed particle flux are shown in the right panel of Fig.4 for the case considered above.  
 474 Comparing the left panel with Fig.3, we see that the fast density relaxation ceases when  
 475 the ZF absorbs a significant part of the system’s kinetic energy. As the wave amplitudes  
 476 are not growing and even a slow decay of the plasma kinetic energy is evident from Fig.4,  
 477 the density profile remains steep enough to produce linear instability. However, its growth  
 478 is suppressed by the excited ZF. Apart from arresting the growth of wave amplitudes,  
 479 it slows down turbulent transport. The area of its suppression extends toward the trough  
 480 edges, where the plasma is locally stable. This configuration fully preserves the trough  
 481 from edge erosion, thus maintaining the integral plasma depletion indefinitely within the  
 482 integration time. Again, such a relaxation outcome is impossible for significantly smaller  
 483 adiabaticity parameter  $\alpha$  values.

484 Now, we briefly consider the role of plasma parameters in selecting the trough re-  
 485 laxation scenario. We reduced the number of control parameters to two, of which  $\beta$  con-  
 486 trols the initial instability growth, but the evolving density profile then takes over its role.  
 487 As we emphasized, the adiabaticity parameter,  $\alpha$ , is more critical for the regime selec-  
 488 tion. The observed turbulent transport barriers are formed for  $\alpha \approx 1$ . For  $\alpha \ll 1$ , the  
 489 transport barriers do not emerge or are weak, so the wall erosion proceeds to the trough  
 490 edges. Projecting this situation on a real trough, we may predict that it will be broad-  
 491 ened and refilled with the external plasma if  $\alpha \ll 1$ . Again, we do not include the field-  
 492 aligned transport here, which may be ineffective in filling the trough regardless of our  
 493 model limitations (e.g., Huang et al., 2007). To investigate a trough relaxation in more  
 494 detail, in the next Section we focus on one particular wall of the trough, as the overall  
 495 picture of the full trough is mainly symmetric.



**Figure 5.** Time-asymptotic distribution of the density (left panel) and plasma potential (right panel) at  $t = 500$ . The simulation parameters, IC and BC, are the same as in Fig.3.

#### 496 5.4 Sustainability of Density Walls

497 In this section, we study the density gradient relaxation in the DW-ZF turbulence  
 498 by focusing on one of the walls of a density trough or any other structure where an iso-

499 lated density jump is present. This approach allows us to gain more insight into the for-  
 500 mation of transport barriers pertinent to broader and deeper troughs. The DW insta-  
 501 bility is expected to originate at the maximum density gradient. Numerous observations  
 502 show that density depletions exceed two orders of magnitude (e.g., Huang et al., 2007),  
 503 occurring over a  $\sim 1^\circ$  latitudes. While the results of the previous section are formally  
 504 applicable to narrow density troughs, by focusing on one of the walls, we can measure  
 505 the transport barrier penetrability in *broader density structures*. The bottom of such troughs  
 506 is often flat, but the walls are steep. Therefore, a step-like initial density profile should  
 507 adequately describe the trough relaxation at either of its two walls. The imposition of  
 508 the channel boundary conditions at  $y = 0, L_y$  allows us to infer how far the relaxing  
 509 part of the initial density profile may spread as the outer plasma percolates through the  
 510 barrier to the bottom of the trough.

511 The initial density profile is then set in the form of a step,

$$n_0(y) = n_1 + \frac{n_2 - n_1}{\tanh \frac{L_y - a}{L_y \delta} + \tanh \frac{a}{L_y \delta}} \left( \tanh \frac{y - a}{L_y \delta} + \tanh \frac{a}{L_y \delta} \right),$$

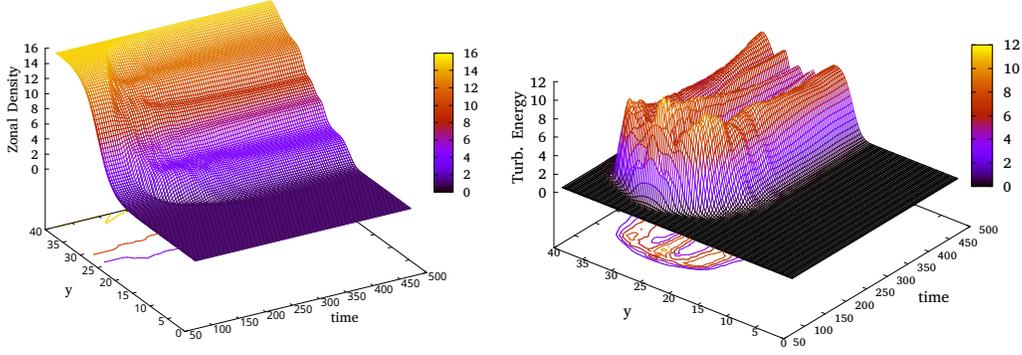
512 raising from  $n_1$  to  $n_2$ . Its raise, centered at  $y = a$ , is smoothed out by a *tanh* function  
 513 over the scale  $L_y \delta$ . Compared to the full-trough relaxation presented in the preceding  
 514 section, the wall relaxation more clearly reveals the fine structure of the relaxed density  
 515 step. The “staircase” profile seen in Fig.6 forms, following a transition period. It lasts  
 516 for several tenths of  $\omega_{ci}^{-1}$  after the onset of instability, during which the staircase fully  
 517 develops and saturates. Staircase formation is not unusual in the turbulent transport phe-  
 518 nomena in various settings (e.g., Balmforth et al., 1998; Dif-Pradalier et al., 2015; Dritschel  
 519 & McIntyre, 2008; Guo et al., 2019; Y. J. Kim et al., 2022; Malkov & Diamond, 2019;  
 520 Milovanov et al., 2021; Pružina et al., 2022) and references therein. The alternation of  
 521 steep and flat pieces of the density profile mirrors variations in the anomalous plasma  
 522 diffusivity. A noteworthy aspect of this profile is that the transport-driving turbulence  
 523 is not significantly diminished after the relaxation is completed and the density profile  
 524 freezes. The total plasma energy, eq.(14),

$$E_{\text{tot}} \equiv \frac{1}{2} \int \int [|\nabla \Phi|^2 + (n - \bar{n}_0)^2] dx dy > \frac{1}{2} \int \int (|\nabla \Phi|^2 + \tilde{n}^2) dx dy \equiv E_{\text{turb}} \quad (15)$$

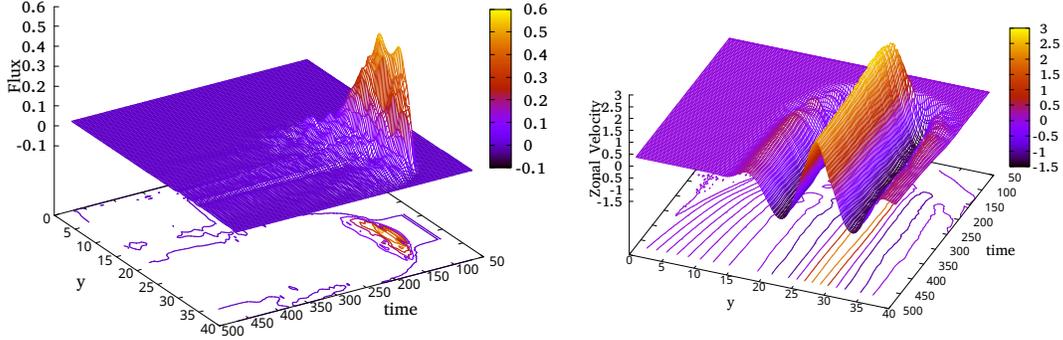
525 is not conserved by eqs.(8-9), cf. eq.(14). Since the simulation starts from  $\Phi \approx 0$ , the  
 526 free energy for the DW instability is in the  $(n - \bar{n}_0)^2$  term of the total energy. We in-  
 527 cluded here the ZF energy into a turbulent component on the r.h.s. of the inequality above  
 528 to show that the ZF taps energy from  $(n - \bar{n}_0)^2$ . The relaxation is similar to the one pre-  
 529 sented in the preceding section. It begins with an unstable growth of DWs that quickly  
 530 saturates when a strong ZF establishes and the particle flux through the transport bar-  
 531 rier drops precipitously.

## 532 5.5 Density Wall Thickening and Destruction

533 The simulation results reported in the preceding subsection have demonstrated a  
 534 minimal density wall thickening after its relaxation in the case of adiabaticity param-  
 535 eter  $\alpha \sim 1$  and the density contrast across the wall  $\Delta n \sim 10$ . Although the DW in-  
 536 stability proved strong, and the ensuing anomalous diffusion typically starts dismantling  
 537 the wall at a high pace, it slows down abruptly and seemingly comes to its end, e.g., Fig.6  
 538 (left panel), as soon as ZF streams arise. The resulting density profile acquires an ap-  
 539 pearance of a staircase, but the wall thickens by a mere factor of two or so. Given the  
 540 narrowness of the observed density depletions (Huang et al., 2007), it might not even be  
 541 detectable. The flux-stopping ZF becomes progressively more laminar toward the end  
 542 of the relaxation and the plasma diffusive flux drops sharply from its initial burst ob-  
 543 served at the outset of relaxation, Fig.7.



**Figure 6.** Left panel: Relaxation of a step-like initial density profile. Right panel: Time evolution of the turbulent energy profile, averaged only in  $x$ -direction, i.e., as a function of time and  $y$  (cf. eq.[15]). The simulation parameters are as follows:  $n_1 = 1$ ,  $n_2 = 15$ ,  $\alpha = 1$ .

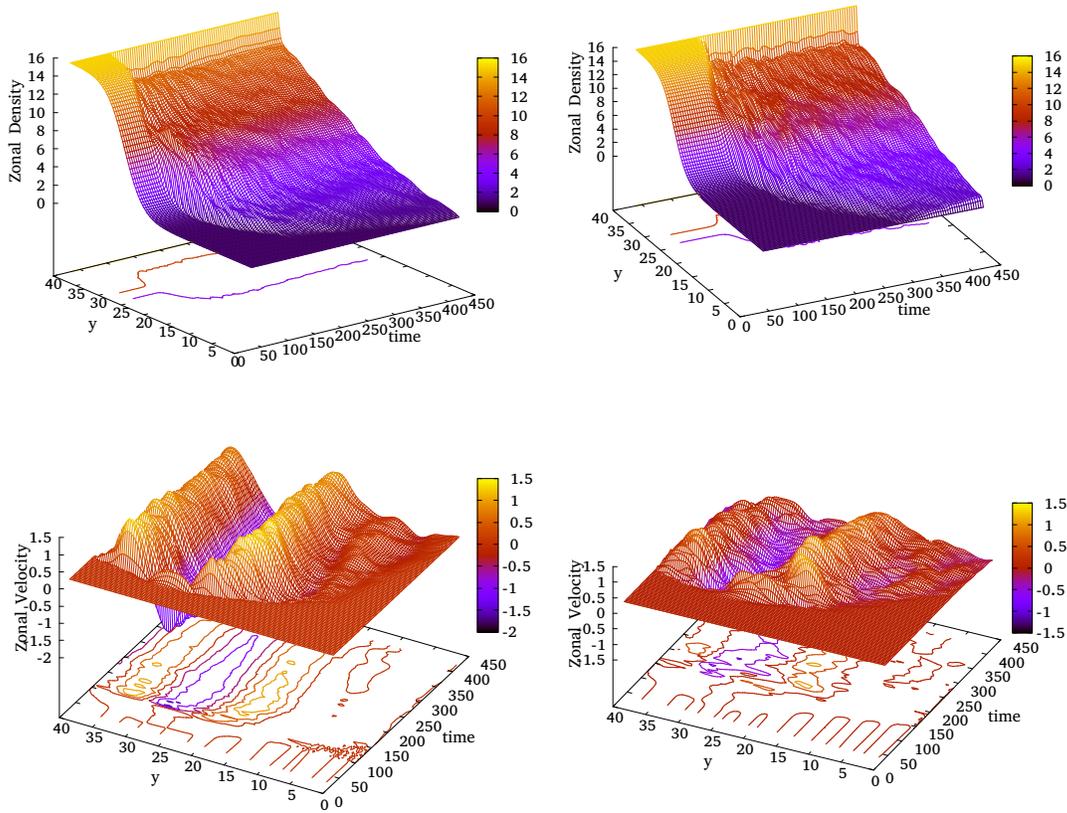


**Figure 7.** Same as in Fig.6 but for the diffusive particle flux  $\langle \tilde{n} \tilde{V}_y \rangle$  and zonal velocity  $\langle V_x \rangle$ .

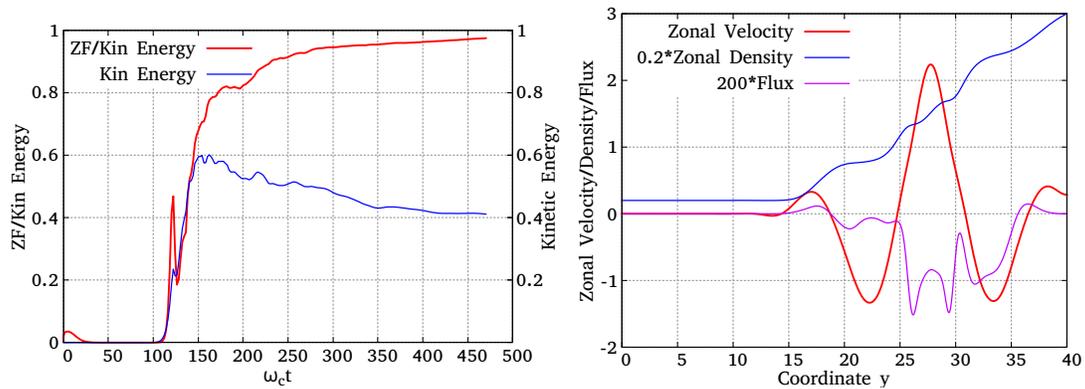
544 That said, the question to be posed is how far the system parameters can deviate  
 545 from those above, in fact, optimal parameter values. In particular, one may ask whether  
 546 the plasma diffuses far enough from the edge into the depleted density region, e.g., to  
 547 refill a broad density trough. A thorough parameter scan is beyond the scope of this paper,  
 548 mainly because such a scan should go deeper than scanning in  $\alpha$  and  $\Delta n$ . As we  
 549 discussed in Sec.5.1, a box resizing should make a significant impact on the result by ef-  
 550 fectively ascribing two different values to the parameter  $\alpha$ : One is in the density, and an-  
 551 other is in the vorticity equation. Here, we present the effect of changing  $\alpha$  on the den-  
 552 sity wall sustainability and destruction.

553 We show in Fig.8 another simulation of turbulent relaxation in time of the same  
 554 density wall as shown in Fig.6. However, this simulation was carried out for two smaller  
 555 values of  $\alpha$ , namely  $\alpha = 0.1$  (left panel) and  $\alpha = 0.03$  (right panel). The other pa-  
 556 rameters are the same as listed in Fig.6. It can be concluded from these three runs that  
 557 the decrease of  $\alpha$  leads to less efficient transport suppression. The same initial density  
 558 profile attached to the high-density side of the channel broadens faster and reaches farther  
 559 toward the opposite side for smaller  $\alpha$ .

560 We also show in Fig. 8 the zonal flow profile, averaged along the channel, as a func-  
 561 tion of time. A close relation between the flux suppression efficiency and the ZF strength  
 562 and regularity is seen from these plots. Note that for the smallest  $\alpha = 0.03$ , the ZF is

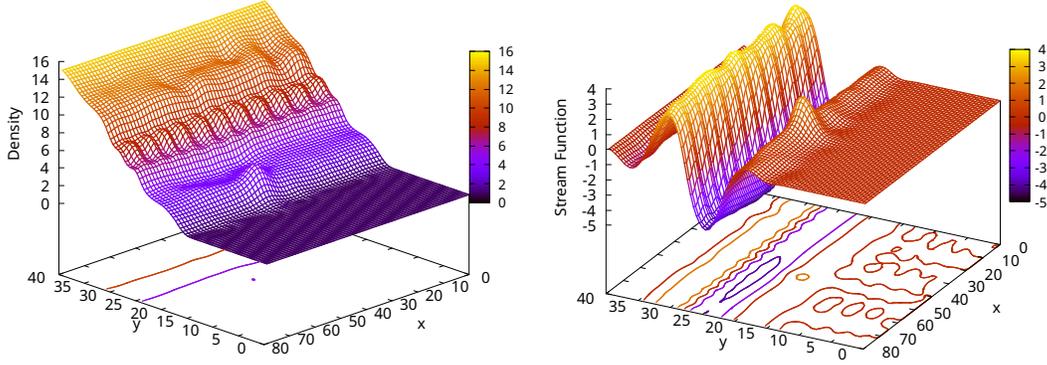


**Figure 8.** Density wall relaxation for two values of the adiabaticity parameter  $\alpha = 0.1$  (left panel) and  $\alpha = 0.03$  (right panel) to be compared with the density profile shown for the same values of other parameters in Fig.6 but for  $\alpha = 1$ .



**Figure 9.** Same as Fig.4 but for wall simulation shown in Fig.6.

563 highly turbulent and relatively slow compared to the fast and laminar flow for  $\alpha \approx 1$ ,  
 564 shown in Fig.7.



**Figure 10.** Same as in Fig.5 but for the step-like initial distribution of plasma density shown in Figs.6 and 7.

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## 5.6 Spectra of DW-ZF turbulence

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Power spectra of residual fluctuations that, as we will show, persist for a long time after an initial density gradient has relaxed, are of significant practical importance. These fluctuations must profoundly impact radio wave transmission used for communication, geopositioning, and ionosphere probing purposes, to name a few (e.g., Kelly, 2012). Specific aspects of electromagnetic wave scattering from turbulent density layers have been discussed in (Sotnikov et al., 2008; V. I. Sotnikov et al., 2010; V. Sotnikov et al., 2014). As established earlier in this section, the fluctuations are confined to steep density gradients and strong shear flows. No significant turbulence spreading effects beyond its generation domain are observed for  $\alpha \sim 1$ . This turbulence confinement to the steep density gradient is still pertinent to a sufficiently broad range of the adiabaticity parameter  $\alpha$  and other simulation parameters, such as the box size and aspect ratio. For significantly smaller values of  $\alpha$ , however, initially sharp density contrasts broaden considerably after the instability progresses toward saturation. At the same time, limits are imposed on our simulation in this case. Nevertheless, sustainable boundaries with steep density gradients change the characteristics of the waves that pass through or are reflected by these boundaries. Moreover, a density trough with sharp edges will likely work as a duct for certain waves.

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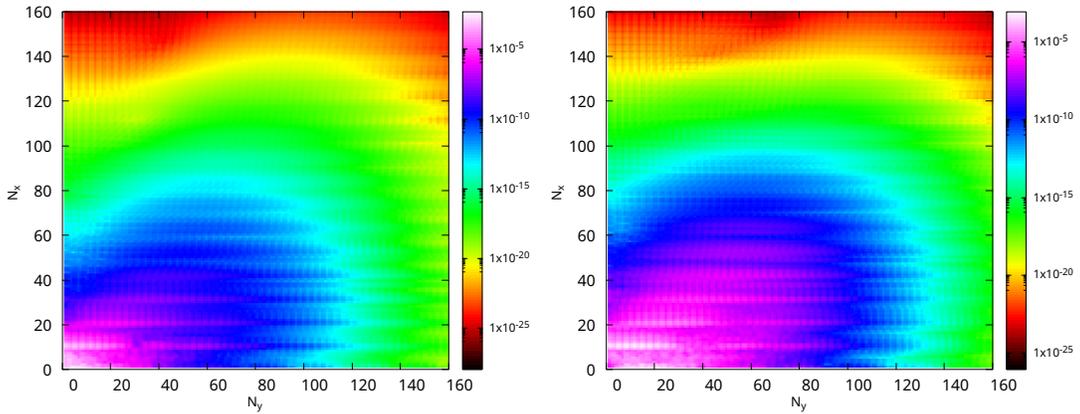
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More specific effects of the relaxed density gradients on passing or reflected waves can be established by focusing on particular waves. Although this task is out of the scope of the present paper, a brief characterization of fluctuations and regular plasma motion inside the gradient areas is in order. The simulation results provide helpful information for wave interaction analyses with the relaxed plasma gradient. The most salient aspect of the regular motion is a strong shear with alternating flow directions. Given that it is formed in the area where the density drops may reach two orders of magnitude or even more, these flows should significantly impact the wave reflection from the dense plasma and mode conversion. The structure of the flow and its intensity can be gleaned from, e.g., Fig.9 and 10. These plots show rather typical flow and density profiles emerging after the relaxation of the initial density gradient with a moderate density contrast  $\Delta n/n_0 \sim 10$ . The number of strong ZF streams, three in this case, has increased in broader boxes for higher density contrasts and more strongly perturbed initial conditions. However, with varying simulation parameters, the flow preserves its morphology of alternating shear layers.



**Figure 11.** Density (left) and vorticity (right) power spectra, shown as functions of mode numbers in  $y$  and  $x$  directions.

598 These flows can be included in the overall power spectrum, manifesting several en-  
 599 hanced harmonics in  $k_y$  direction, or can be considered as a quasistationary background  
 600 flow with superposed fluctuations. The choice of representation, again, depends on the  
 601 wave scales that interact with the flow. For instance, the incoherent scatter radar tech-  
 602 nique that implies a radio frequency well above 100 MHz for some radar facilities with  
 603 submeter wavelengths would probably rely on the second choice. If we project our re-  
 604 sults to a larger box (see, however, Sec.5.1), the scale of the shear flow may reach hun-  
 605 dreds of meters or even more. However, depending on other control parameters, primar-  
 606 ily the adiabaticity  $\alpha$  and the density contrast  $\Delta n$ , the ZF profile may become much more  
 607 oscillatory, which would *de facto* decrease the scale of the ZF.

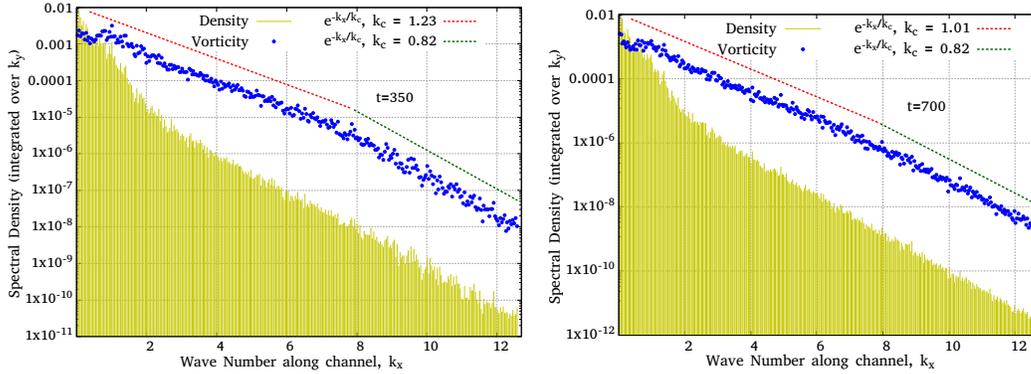
608 To demonstrate a typical structure of the spectra emerging after the density gra-  
 609 dient relaxation, we show in Fig.11 a 2D spectral density as a function of mode num-  
 610 bers in  $x$  and  $y$  direction. The spectrum is anisotropic at lower  $k$  and becomes more isotropic  
 611 toward shorter scales of perturbations (higher  $N_x$  and  $N_y$ ). One can also see visible mod-  
 612 ulations in the  $x$  direction likely produced by a nonlinear coupling of a few of the most  
 613 unstable modes. For example, the flow is modulated at  $N_x=11$  in the case shown in Fig.10.  
 614 In addition, an isolated vortex is seen in this figure, both on the flow stream function  
 615 and the density surface plot. The weak short-scale modulation in  $N_y$ , Fig.11, is due to  
 616 the limited mode number and can be removed using finer grids, but without significant  
 617 effects on the results.

618 More detailed information on the power spectra associated with the propagation  
 619 of disturbances in the flow along the channel, i.e., DW propagation, can be obtained by  
 620 integrating the spectrum over  $k_y$ . Such a spectrum is shown in Fig.12 for two times of  
 621 integration. In both cases, the spectrum has an exponential form, which is best seen in  
 622 a linear-log format. However, the spectrum shown at  $t = 350$  (as other earlier spectra)  
 623 has a distinct break at  $k_x \approx 8$  with a clear steepening at higher  $k_x$ , especially pronounced  
 624 in the flow vorticity. The exponential dependence of the spectrum below and above the  
 625 break can be fitted as  $\propto \exp(-k_x/k_c)$  with the “cut-off”  $k_c = 1.23$  and  $0.82$ , respec-  
 626 tively.

627 At later times (e.g.,  $t = 700$ , shown in Fig.12) the spectrum before the break steep-  
 628 ens, showing a trend to asymptotically form a single exponent at high  $k_x$  with a cut-off  
 629  $k_c = 0.82$ . This happens earlier for the density fluctuations, as may be seen from the

630 figure, but the overall spectra at higher  $k_x$  are very similar for density and vorticity per-  
 631 turbations. One may conclude that the coupling of the two equations ( $\alpha$ - terms ) are sub-  
 632 dominant at short scales, and both  $n$  and  $\Delta\Phi$  are convected with the same stream func-  
 633 tion  $\Phi$ . This conclusion can be drawn directly from the equations.

634 At lower  $k_x \gtrsim 1$ , the intensity of the density perturbations drops precipitously,  
 635 while the vorticity level decays at approximately the same rate as at higher  $k_x$  values  
 636 below the break. Both vorticity and density fluctuations have a spectral maximum at  
 637  $k_x \approx 1$ , which stems from the maximum growth rate (outer scale of turbulence). Apart  
 638 from the spectrum straightening at higher  $k_x$ , the scattering of modes around the respec-  
 tive exponential forms (straight lines in the linear-log format shown) diminishes.



**Figure 12.** Power spectrum of density and vorticity perturbations integrated over  $k_y$ , The box size is  $L_x \times L_y = 80 \times 40$ , the respective number of grid points 1024x512.

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## 640 6 Conclusions and Discussion

641 In this paper, we have investigated mesoscale dynamics in the Earth’s ionosphere  
 642 that we expect to accompany phenomena associated with *stormtime plasma density troughs*.  
 643 Our focus was on the sharp plasma interfaces between high and low-density regions. These  
 644 are thought to result from large-scale ionosphere instabilities, e.g., those of the Rayleigh-  
 645 Taylor variety. They can also be formed by a rapid plasma evacuation along the mag-  
 646 netic field related to the ionosphere connection with the magnetosphere. Irrespective of  
 647 their genesis, the density contrasts may undergo a secondary instability at much shorter  
 648 scales, not accounted for by the primary mechanisms of their generation.

649 The development of a micro-scale instability in a certain parameter range, primar-  
 650 ily where an adiabaticity parameter,  $\alpha = k_z^2 V_{Te}^2 / \omega_{ci} \nu_{ei}$ , does not strongly deviate from  
 651  $\alpha \sim 1$ , has profound consequences for mesoscopic and eventually macroscopic dynam-  
 652 ics of the adjacent environment. We have investigated them in two basic settings. The  
 653 first setting mimics a density trough, albeit limited in its transverse scale due to the code  
 654 restrictions. To ease this limitation, we have studied in more detail the case of a den-  
 655 sity wall, which may represent either a part of a wide trough or an isolated density de-  
 656 pletion in the ionosphere.

657 In both cases, we found that an unstable growth of plasma drift waves is accom-  
 658 panied by the generation of strong mesoscopic zonal flows (ZF) between the high and  
 659 low-density regions, directed across the magnetic field. The shearing effect of the gen-  
 660 erated ZF on the transport-driving turbulence eddies lead to a substantial reduction of  
 661 anomalous plasma diffusion. Consequently, we predict phenomena such as equatorial spread-

F plasma bubbles and density troughs (e.g., Huang et al., 2007), to be considerably more sustainable than mere drift-wave unstable density depletions. At the same time, a SAID phenomenon has been found (He et al., 2017) to have a large dispersion in its duration. One may assume that the underlying mechanisms behind the most protracted SAID events are sensitive to, and may even result from, specific ionospheric conditions at their locales. Whether these conditions correlate with the ZF generation requirements merits a separate observation analysis. The present study concludes that long-lifetime SAID are consistent with a strong flux suppression by the ZF generation. The required strong ZF streams are often seen near sharp density gradients in many observations (e.g., Huang et al., 2007).

We have also extracted detailed perturbation spectra from our simulation results. These data are necessary for probing the ionosphere by radar transmission using coherent and incoherent scatter and scintillation techniques (e.g., Kelly, 2012). In most cases with  $\alpha \sim 1$ , our code renders exponentially decaying fluctuation spectra. They can be straightforwardly fed into wave scattering models based on any observation techniques mentioned above. The choice of the density fluctuation spectrum is particularly crucial for the scintillation techniques.

## Data Availability Statement

We have used the ISPACK software library for FFT transforms, documented in and publically available from, e.g., <http://www.gfd-dennou.org/library/spmodel/label-2>. The output from FORTRAN subroutines has been ported to Gnuplot, graphic software, also publically available at the Open Source Software repository at <https://sourceforge.net/p/gnuplot/gnuplot-main/ci/master/tree/>.

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