Stabilisation of short-wavelength instabilities by parallel-to-the-field shear in long-wavelength $E \times B$ flows

M. R. Hardman$^{1,2}$, M. Barnes$^{1,2}$, and C. M. Roach$^2$

$^1$Rudolf Peierls Centre for Theoretical Physics, University of Oxford, Oxford OX1 3PU, UK
$^2$Culham Centre for Fusion Energy, UKAEA, Abingdon OX14 3DB, UK

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Magnetised plasma turbulence can have a multiscale character: instabilities driven by mean temperature gradients drive turbulence at the disparate scales of the ion and the electron gyroradii. Simulations of multiscale turbulence, using equations valid in the limit of infinite scale separation, reveal novel cross-scale interaction mechanisms in these plasmas. In the case that both long-wavelength (ion-gyroradius-scale) and short-wavelength (electron-gyroradius-scale) linear instabilities are driven far from marginal stability, we show that the short-wavelength instabilities are suppressed by interactions with long-wavelength turbulence. Two novel effects contributed to the suppression: parallel-to-the-field-line shearing by the long-wavelength $E \times B$ flows, and the modification of the background density gradient by the piece of the long-wavelength electron adiabatic response with parallel-to-the-field-line variation. In contrast, simulations of multiscale turbulence where instabilities at both scales are driven near marginal stability demonstrate that when the long-wavelength turbulence is sufficiently collisional and zonally dominated the effect of cross-scale interaction can be parameterised solely in terms of the local modifications to the mean density and temperature gradients. We discuss physical arguments that qualitatively explain how a change in equilibrium drive leads to the observed transition in the impact of the cross-scale interactions.

1. Introduction

In a magnetised plasma, gradients in the mean temperature of the component particle species act as sources of free energy that drive instability. Instabilities saturate through nonlinear interactions to form turbulence; the nature of the turbulence is determined by the character of the underlying instabilities. In this paper we consider the effect of cross-scale nonlinear interactions in turbulence driven at the well-separated space-time scales associated with ion and electron dynamics, respectively. This represents a fundamentally different system to the one usually employed in turbulence studies, which consists of a single injection range, a single inertial range, and a single dissipation range.

Instabilities in magnetised plasmas typically have a structure that is elongated along field lines; this is a result of the Lorentz force which causes particles to perform gyro orbits in the plane perpendicular to the magnetic field line, whilst allowing unimpeded motion in the parallel-to-the-field direction. By definition, in a magnetised plasma the typical particle gyroradius is much smaller than the device scale $a$ that determines the parallel-to-the-field scale of instabilities. In the core of magnetic confinement fusion devices two of the dominant instabilities which drive turbulence are the ion temperature gradient (ITG)
and the electron temperature gradient (ETG) instabilities (Cowley et al. 1991; Romanelli 1989; Horton et al. 1988; Lee et al. 1987). The ITG instability drives turbulence at the scale of the ion thermal gyroradius $\rho_{\text{th,i}}$, at frequencies of order the ion transit frequency $v_{\text{th,i}}/a$, with $v_{\text{th,i}}$ the ion thermal speed. The ETG instability drives turbulence at the scale of the electron thermal gyroradius $\rho_{\text{th,e}}$, at frequencies of order the electron transit frequency $v_{\text{th,e}}/a$, with $v_{\text{th,e}}$ the electron thermal speed.

When the ion and electron temperatures are approximately equal the space and time scales associated with the ETG and ITG instabilities can be well separated: the separation of scales is determined by the square root of the electron-to-ion mass ratio $(m_e/m_i)^{1/2} \approx \rho_{\text{th,e}}/\rho_{\text{th,i}} \sim v_{\text{th,i}}/v_{\text{th,e}}$. For the deuterium ions commonly used in magnetic confinement fusion, the small value of $(m_e/m_i)^{1/2} \approx 1/60$ allows for the possibility that two distinct types of turbulence co-exist at disparate space-time scales. This possibility has sparked considerable interest in multiscale turbulence, see e.g. Maeyama et al. (2017a, 2015); Howard et al. (2016a, 2015, 2014); Bonanomi et al. (2018); Görler & Jenko (2008); Candy et al. (2007); Waltz et al. (2007); Staebler et al. (2016, 2017); Creely et al. (2019); Itoh & Itoh (2001).

Direct numerical simulations (DNS) of multiscale plasma turbulence show that the interactions between long-wavelength and short-wavelength scales can be important for determining the character of the turbulence and the flux of heat exhausted from the fusion plasma (Howard et al. 2016a, b; Maeyama et al. 2015, 2017a). The small value of $(m_e/m_i)^{1/2}$ makes DNS of multiscale plasma turbulence extremely challenging; the increased cost of multiscale DNS compared to conventional simulations scales with $(m_i/m_e)^{3/2}$. Recent multiscale DNS have been performed using the electron-to-deuterium mass ratio (e.g. Howard et al. (2016a, b); Bonanomi et al. (2018)) and electron-to-hydrogen mass ratio (e.g. Maeyama et al. (2015, 2017a)). The multiscale DNS provide evidence showing that the presence of short-wavelength turbulence can sometimes modestly enhance long-wavelength, ITG-driven fluxes (Howard et al. 2016a, b; Maeyama et al. 2015, 2017a). In contrast, multiscale DNS of microtearing mode (MTM) driven turbulence show that the presence of short-wavelength ETG modes can suppress the long-wavelength MTM (Maeyama et al. 2017a). Most relevant to the results presented in this paper, the multiscale DNS give clear evidence that shows that short-wavelength turbulence can be suppressed in the presence of long-wavelength turbulence as a result of cross-scale interaction (Maeyama et al. 2015, 2017a; Howard et al. 2016b).

In this paper we use numerical simulations of scale-separated turbulence to demonstrate the effect of long-wavelength turbulence on short-wavelength fluctuations. We show that strongly driven, long-wavelength turbulence can stabilise the short-wavelength ETG instability through cross-scale interaction. We show that this stabilisation is due to two effects: the parallel-to-the-field shearing of short-wavelength fluctuations by long-wavelength $E \times B$ drifts; and the modification of the background drives of instability by gradients of long-wavelength fluctuations. We show that the modification to the background drives can be characterised as having two main components: a perturbed Maxwellian piece that shifts the mean density and temperature gradients by a constant value; and a density gradient due to the long-wavelength electron adiabatic response that varies in the parallel-to-the-field direction. These mechanisms may explain the suppression of the short-wavelength modes observed in DNS of strongly driven turbulence. We show that parallel-to-the-field shearing can be a significant cross-scale interaction mechanism far above marginal stability, but that it can be less important in turbulence driven near marginal stability. We examine one example of near-marginal turbulence where we find that the effect of cross-scale interaction on the ETG instability can be...
described with the local modification of the density and temperature gradient length scales by gradients of long-wavelength fluctuations.

2. Scale-separated model

The local, $\delta f$ gyrokinetic equations describe the evolution of magnetised plasma turbulence driven by mean gradients in the limit that $\rho_{th,i}/a \to 0$ and $\rho_{th,e}/a \to 0$ (Catto 1978; Frieman & Chen 1982; Brizard & Hahm 2007); this model is the starting point in the derivation of scale-separated equations for multiscale, electrostatic turbulence (Hardman et al. 2019). We now briefly review this derivation, first presented in Hardman et al. (2019). Ultimately, we find that the local gradients of slowly evolving, long-wavelength turbulence act to modify the mean gradients and flows which drive (or suppress) rapidly evolving, short-wavelength fluctuations. In analogy to the coupling between turbulence and large-scale transport (Sugama & Horton 1997; Abel et al. 2013), we might expect short-wavelength turbulence to generate fluxes that diffuse long-wavelength turbulence. However, this effect is small by $(m_e/m_i)^{1/2}$ in the ordering, and hence too small to appear at leading order in the $(m_e/m_i)^{1/2}$ expansion that we employ.

In the limit that $(m_e/m_i)^{1/2} \to 0$, multiscale turbulence can be meaningfully decomposed into components: the long-wavelength “ion scale” (IS) and the short-wavelength “electron scale” (ES). We assume that IS turbulence varies only on $\rho_{th,i}$ perpendicular-to-the-field scales in space and on $a/v_{th,i}$ scales in time, whereas ES turbulence has much finer $\rho_{th,e}$ perpendicular-to-the-field structures, and evolves on the rapid $a/v_{th,e}$ time scale. Both IS and ES turbulence have parallel-to-the-field scales of order $a$. Here, $\rho_{th,p} = v_{th,p}/\Omega_p$, with $p$ the particle species index, $v_{th,p} = \sqrt{2T_p/m_p}$, $\Omega_p = Z_peB/m_pc$ the species cyclotron frequency, $T_p$ the species temperature, $Z_p$ the species charge number, $e$ the unit charge, $B$ the magnetic field strength, and $c$ the speed of light. The distribution function $\delta f$ and electrostatic potential $\phi$ of the turbulence are sums of the IS and ES components;

$$\delta f = \bar{\delta f} + \tilde{\delta f}, \quad \phi = \bar{\phi} + \tilde{\phi}, \quad (2.1)$$

where $\bar{\chi}$ and $\tilde{\chi}$ are the IS and ES pieces of any quantity $\chi$, respectively. To carry out this decomposition, we introduce the ES average $\langle \cdot \rangle^{ES}$, an average over $\rho_{th,e}$ perpendicular-to-the-field scales and $a/v_{th,e}$ times. We assume that ES turbulence is statistically periodic on $\rho_{th,e}$ scales, i.e., $\langle \tilde{\delta f} \rangle^{ES} = 0$. This allows us to use $\langle \cdot \rangle^{ES}$ to extract the scale-separated, coupled gyrokinetic equations for $\bar{\delta f}$ and $\tilde{\delta f}$ from the gyrokinetic equations for $\delta f$.

The IS component of the turbulence is evolved with the usual gyrokinetic equation for the ion species, and a parallel-orbit-averaged equation for the electron species. The IS turbulence evolves independently of the ES turbulence to leading order in the $(m_e/m_i)^{1/2}$ expansion. At the ES the ion species have a Boltzmann response. The evolution of the ES turbulence is governed by the short-wavelength electron gyrokinetic equation (Hardman et al. 2019)

$$\left( \frac{\partial}{\partial t} + v_E \cdot \nabla \right) \tilde{g} + v_{\parallel} \mathbf{b} \cdot \nabla \theta \frac{\partial \tilde{g}}{\partial \theta} + \left( v_M + \tilde{v}_E \right) \cdot \nabla \tilde{g} + \tilde{v}_E \cdot \left( \nabla F_0 + \nabla \tilde{g} \right) = \frac{eF_0}{T_e} \left( v_{\parallel} \mathbf{b} \cdot \nabla \theta \frac{\partial \tilde{\varphi}}{\partial \theta} + v_M \cdot \nabla \tilde{\varphi} \right) + \tilde{C}, \quad (2.2)$$

where we express the ES gyrokinetic equation in terms of the particle guiding centre $\mathbf{R} = \mathbf{r} + \mathbf{b} \times v/\Omega_e$, with $\mathbf{r}$ the particle position, $\mathbf{b}$ the magnetic field direction, $v$ the
particle velocity; $\varepsilon$ the particle kinetic energy; $\lambda = \mu/\varepsilon$ the pitch angle, with $\mu$ the magnetic moment; $\sigma$ the sign of the particle velocity in the magnetic field direction $v_\parallel = \mathbf{v} \cdot \mathbf{b} = \sigma(2\varepsilon/m_p)^{1/2}(1 - \lambda B)^{1/2}$; the gyroaveraged ES potential $\tilde{\varphi} = \langle \hat{\varphi} \rangle^0$, with $\vartheta$ the angle of gyration about $\mathbf{b}$ and $\langle \cdot \rangle^\vartheta$ the electron gyroaverage at fixed $\mathbf{R}$; the ES gyroaveraged electron distribution function $\tilde{g} = \langle \delta f_e \rangle^\vartheta = \tilde{h} + e\tilde{\varphi}F_0/T_e$, with $\delta f_e(r) = \tilde{h}(\mathbf{R}) + e\hat{\varphi}(r)F_0/T_e$, $\tilde{h}$ the non-Boltzmann part of the ES electron distribution function, and $F_0$ the Maxwellian electron mean distribution function. The poloidal angle $\theta$ is the parallel-to-the-field coordinate; $\mathbf{v}_M$ is the electron magnetic drift; and $\mathcal{C}$ accounts for the effects of collisions on $\tilde{g}$. The long-wavelength and short-wavelength $\mathbf{E} \times \mathbf{B}$ drifts are $\mathbf{v}_E = (c/B)\mathbf{b} \times \nabla \hat{\varphi}$ and $\tilde{\mathbf{v}}_E = (c/B)\mathbf{b} \times \nabla \tilde{\varphi}$, respectively; and $\tilde{g} = \delta f_e$ is the long-wavelength, IS, electron distribution function. Gyroaverages do not appear in the definitions of $\mathbf{v}_E$ or $\nabla \tilde{g}$ as $\rho_{th,e}$ is much smaller than the spatial scale of the IS structures. Equation (2.2) is closed by the ES quasineutrality relation $Z_i\tilde{n}_i = \tilde{n}_e$, with $\tilde{n}_p$ the ES fluctuating density of species $p$. The ES quasineutrality relation can be written
\[
\int d^3\mathbf{v} |_r \left( \tilde{g}(\mathbf{R}) + \frac{eF_0}{T_e} \left( \tilde{\varphi}(r) - \tilde{\varphi}(\mathbf{R}) \right) \right) = -\frac{Z_i e\tilde{\varphi}(r)}{T_i} n_{0e},
\]
where $n_{0e}$ is the mean electron density, and the subscript $r$ on the volume element indicates that the integral is to be taken at fixed particle position.

Cross-scale interaction is mediated by the terms $\mathbf{v}_E \cdot \nabla \tilde{g}$ and $\tilde{\mathbf{v}}_E \cdot \nabla \tilde{g}$ in the electron gyrokinetic equation, equation (2.2). The gradient $\nabla \tilde{g}$ and the drift $\tilde{\mathbf{v}}_E$ are constant on $a/v_{th,e}$ time scales, and so appear as additional “equilibrium” terms in the short-wavelength electron gyrokinetic equation. By inspecting equation (2.2) we see that we are able to ascribe simple physical interpretations to the cross-scale terms: the term $\mathbf{v}_E \cdot \nabla \tilde{g}$ represents an advection term that will introduce velocity shear; and the gradient $\nabla \tilde{g}$ appears in the term $\tilde{\mathbf{v}}_E \cdot (\nabla F_0 + \nabla \tilde{g})$, i.e., $\nabla \tilde{g}$ modifies the usual equilibrium drive of instability $\nabla F_0$. Equations (2.2) and (2.3) are solved in a thin ES flux tube embedded within a larger IS flux tube, see figure 1 for a cartoon. In each ES flux tube $\nabla \tilde{g}$ and $\mathbf{v}_E$ take values which are constant in the perpendicular-to-the-field plane; the effects of the perpendicular derivatives of $\nabla \tilde{g}$ and $\mathbf{v}_E$ are small by $(m_e/m_i)^{1/2}$ and so only appear at higher order. However, both $\nabla \tilde{g}$ and $\mathbf{v}_E$ can vary along the field line: this results in a background drive of instability that is no longer uniform in $\theta$, and in suppression of short-wavelength instabilities due to parallel-to-the-field $\mathbf{E} \times \mathbf{B}$ shearing by long-wavelength flows. The component of $\mathbf{v}_E$ due to the piece of the electrostatic potential $\tilde{\varphi}$ that is constant on the flux surface may be removed from equation (2.2) by boosting to a toroidally rotating frame [Hardman 2019]; this component of $\mathbf{v}_E$ can only change the instability frequency by a Doppler shift. We will discuss the physical interpretation of these cross-scale interaction mechanisms further in section 4.

To enable the study of cross-scale interaction in this multiscale framework, the cross-scale terms $\mathbf{v}_E \cdot \nabla \tilde{g}$ and $\tilde{\mathbf{v}}_E \cdot \nabla \tilde{g}$ in equation (2.2) were implemented in the $\delta f$ gyrokinetic code GS2 ([Kotschenreuther et al. 1995, Hardman 2019]).

3. Numerical results.

In this paper we address the key question of how the linear stability of the ETG mode is affected by cross-scale interaction. We focus on a plasma in which single-scale microturbulence is well understood: we use a simple magnetic geometry consisting of concentric circular flux surfaces, with parameters largely corresponding to the widely-used Cyclone Base Case benchmark (CBC) ([Dimits et al. 2000]). In the simulations we
Figure 1: A cartoon of a large-scale, IS flux tube with several narrow ES flux tubes embedded within it.

Figure 2: (a) The IS potential $\bar{\phi}$ in units of $\rho_{th,i}^* = \rho_{th,i}/a$, shown as a function of time (volume-averaged), and (b) at the outboard midplane at $t \approx 1560a/v_{th,i}$. Times at which the IS turbulence is sampled are indicated with vertical lines, and sampled $(x_s, y_s)$ positions are indicated by crosses.

take $a$ to be the half-diameter of the last closed flux surface; we take the normalised minor radius at the centre of the flux tube $\rho_0 = r_0/a = 0.54$, with $r_0$ the minor radius of the flux surface; the major radius $R = 3.0a$; the safety factor at the centre of the flux tube $q_0 = 1.4$; the magnetic shear $\dot{s} = \rho d\ln q/d\rho = 0.8$; equal ion and electron temperatures $T_i = T_e = T$; the normalised temperature gradient $a/L_T = -d\ln T/d\rho = 2.3$; the normalised density gradient $a/L_n = -d\ln n/d\rho = 0.733$; the normalised self collision frequencies $\nu_{ii}/v_{th,i} = \nu_{ee}/v_{th,e} = 10^{-2}$; and the electron-to-deuterium-ion mass ratio $m_e/m_i = 2.7 \times 10^{-4} \approx 1/3670$. In the IS simulations, instead of explicitly implementing the parallel-orbit-averaged equation for the electron species, we impose that the passing, nonzonal electron response is Boltzmann, whilst continuing to solve for the remainder of the electron distribution function with the usual GS2 algorithm. In the limit that $(m_e/m_i)^{1/2} \to 0$, this has the effect of solving the electron parallel-orbit-averaged equation numerically. The GS2 algorithm employs $(\varepsilon, \lambda)$ velocity coordinates, and retains the bounce points in $\theta$ for all particles – in this scheme particles with large parallel velocities have a response that is an orbit average over the fluctuations. To verify that $m_e/m_i = 2.7 \times 10^{-4}$ is sufficiently small for the IS turbulence here to be in the $(m_e/m_i)^{1/2} \to 0$ limit, we performed additional IS simulations with $m_e/m_i = 2.7 \times 10^{-12}$ and a Boltzmann passing, nonzonal electron response. Identical IS fluxes and cross-scale interactions with ETG instabilities were found at $m_e/m_i = 2.7 \times 10^{-12}$ and $m_e/m_i = 2.7 \times 10^{-4}$. Physically, this is because the electron bounce time is shorter than the correlation time of the IS turbulence considered here.

To assess the impact of cross-scale interaction on the ETG instability we carry out
the following numerical investigation. We perform simulations of long-wavelength ITG-driven turbulence, using the resolutions detailed in appendix A.1. We compute a sample of $\nabla \tilde{g}(\theta, \varepsilon, \lambda, \sigma)$ and $\tilde{v}_E(\theta)$ at 6 times $t_s$, on a $6 \times 5$ grid in radial and binormal position, $x_s$ and $y_s$, respectively. For each $t_s$ at every sampled $(x_s, y_s)$ we compute the ETG linear growth rate $\gamma(t_s, x_s, y_s)$, including the cross-scale interaction terms $\tilde{v}_E \cdot \nabla \tilde{g}$ and $\tilde{v}_E \cdot \nabla \tilde{g}$ in equation (2.2). The resolutions used for the linear calculations are detailed in appendix A.2. This statistical approach is necessary because $\tilde{v}_E$ and $\nabla \tilde{g}$ must be calculated from turbulent fields which may vary intermittently in time and space. Figure 2a shows the spatially-averaged IS potential $\langle \tilde{\phi} \rangle_{\psi, \psi_0}(L)$, with sampled times indicated with a vertical line; and figure 2b shows the IS potential $\tilde{\phi}(x, y)$ at the outboard midplane at time $t \approx 1560 a/v_{th,i}$ – the last sampled time in figure 2a, with sampled positions indicated with crosses. The $(x, y)$ coordinates are proportional to the flux-surface and field-line labels, $\psi$ and $\alpha$, respectively, that define the magnetic field $\mathbf{B} = \nabla \alpha \times \nabla \psi$: we use the perpendicular-to-the-field coordinates $(x, y)$ that are defined to have units of length; i.e., $x = (\psi - \psi_0)(d\psi/dx)^{-1}$ and $y = (\alpha - \alpha_0)(d\alpha/dy)^{-1}$, respectively, with $(\psi_0, \alpha_0)$ the coordinates of the centre of the IS flux tube, $d\psi/dx = B_{ref} \rho_0 q_{0}/q_{0}$ and $d\alpha/dy = B_{ref} \rho \partial \rho/\partial \psi |_{\rho_0}$, where $B_{ref}$ is a reference magnetic field. We take $B_{ref}$ to be the toroidal magnetic field strength on the flux surface $\psi_0$ at the major radius $R_{ref} = (R_{max} + R_{min})/2$, where $R_{max}$ and $R_{min}$ are the maximum and minimum major radial positions taken in the flux tube, respectively. We note that we choose to define the poloidal angle $\theta$ in the simulations to be such that $\mathbf{b} \cdot \nabla \theta$ is constant in $\theta$.

Figure 3 shows the linear growth rate $\gamma$ spectrum in the absence of cross-scale interaction. In figure 3a we show $\gamma(k_y)$ for modes with $\hat{\theta}_0 = 0$, where $\hat{\theta}_0$ is the poloidal angle at which the wave fronts of the mode align with the minor radial direction of the flux surface. Modes with $\hat{\theta}_0 = 0$ have radially aligned wave fronts at the outboard midplane. Figure 3a shows that for the simulation parameters used in this paper, see curve i), there is a natural stable gap in binormal wave number $k_y$ between the ITG and ETG modes: this serves to define the cut-off between the IS and ES. We note that this separation of scales can become larger when both the drives of instability $a/L_{T_i}$ and $a/L_{T_e}$ are reduced, see curve ii), whereas the separation can disappear for sufficiently strong drive.

Figure 3: (a) The linear growth rate without cross-scale effects for three cases: i) for the parameters given at the start of section 3 ($a/L_{T_i} = a/L_{T_e} = 2.3$ and $a v_{pp}/v_{th,p} = 10^{-2}$); ii) for the case of reduced drive, with $a/L_{T_i} = a/L_{T_e} = 1.38$, holding $a v_{pp}/v_{th,p} = 10^{-2}$; and iii) for the case of reduced collision frequency, with $a/L_{T_i} = a/L_{T_e} = 2.3$ and $a v_{pp}/v_{th,p} = 10^{-4}$. (b) The ETG growth rate in the absence of cross-scale interaction for the parameters given at the start of section 3 ($a/L_{T_e} = 2.3$ and $a v_{ee}/v_{th,e} = 10^{-2}$). Modes within the dashed curve are unstable.
Figure 4: The ETG growth rate in the presence of cross-scale interaction due to IS turbulence. Figures (a) and (b) show growth rates calculated at different \((x_s, y_s)\) and times \(t_s\). Note the suppression of the growth rate in (a) and (b) compared to Figure 3b.

or low collisionality, see, e.g., curve iii). Collisional dissipation creates a stable gap in \(k_y\) between the ITG and the ETG modes through the following effects. Firstly, collisional dissipation stabilises both the weakly growing ETG modes with \(k_y \rho_{\text{th},i} < 1\), and the ITG modes with \(k_y \rho_{\text{th},e} \gtrsim 1\). Secondly, collisions have the effect of detrapping trapped electrons with an enhanced collision frequency \(\nu_{\text{ee}} \sim \nu_{\text{ee}} / \epsilon\), where \(\epsilon = r_0 / R\) is the inverse aspect ratio. When the mode frequency \(\omega\) is sufficiently small, i.e., \(\omega / \nu_{\text{ee}} \lesssim 1\), the trapped electron drive of trapped electron modes (TEM) \cite{Adam1976} and ITG modes is reduced, and hence TEM and ITG modes that lie in the range \(k_y \rho_{\text{th},i} \gtrsim 1\) are stabilised, cf. \cite{Roach2009}. In figure 3b we plot the full ETG linear growth rate spectrum \(\gamma(k_y, \theta_0)\) in the absence of cross-scale interaction. Note that the most unstable modes occur at \(\theta_0 = 0\). The dashed curve in figure 3b indicates the stability boundary where \(\gamma = 0\), calculated with an interpolated spline fit. Figures 4a and 4b show the ETG growth rates in the presence of cross-scale interactions at two of the sampled IS positions and times. Comparing figure 3b with figures 4a and 4b shows that the effect of cross-scale interaction can be dramatic: in these examples the ETG mode is largely suppressed, and completely stabilised, respectively.

To determine the average, or typical effect of IS turbulence on the ETG instability we calculate the maximum ETG growth rate \(\gamma_{\text{max}}\) at every \((t_s, x_s, y_s)\), where we note that the maximum \(\gamma\) may occur at a \((k_y, \theta_0)\) that is different from the single-scale case. Figure 5a shows that the majority of IS drifts and gradients had the effect of stabilising the ETG mode: we find that \(\gamma_{\text{max}} < 0\) at 115 \((t_s, x_s, y_s)\) within the sample of 180, with only a single \((t_s, x_s, y_s)\) showing a \(\gamma_{\text{max}}\) larger than the maximum growth rate in the absence of cross-scale interaction. This is illustrated further in figure 5b which shows that the average of the sampled ETG growth rates \(\langle \gamma \rangle_{t_s, x_s, y_s}\) is negative, for all \((k_y, \theta_0)\). In figure 6 we show the distribution of \(k_y\) and \(\theta_0\) of the fastest growing ETG modes in the presence of IS turbulence, with only modes that have \(\gamma_{\text{max}} > 0\) included in the histograms. Note the significant spread in \(k_y\) and \(\theta_0\) that results from the cross-scale interactions with the IS turbulence.

These results indicate that strongly driven ITG turbulence can suppress, or stabilise, strongly driven ETG instabilities. This is in agreement with results from multiscale DNS which show reduced short-wavelength heat transport in the presence of strongly driven ITG instabilities, see, e.g. \cite{Howard2016b,Maeyama2015}. In the following sections we explore the physical mechanisms that result in the stabilisation of the ETG instability in the \((m_e/m_i)^{1/2} \rightarrow 0\) limit.
Figure 5: (a) The maximum ETG growth rate $\gamma_{\text{max}}$ computed including the cross-scale terms $\mathbf{v}_E \cdot \nabla \tilde{g}$ and $\tilde{v}_E \cdot \nabla \tilde{g}$ in equation (2.2). We see that the typical effect of cross-scale interaction for these parameters is to stabilise the ETG mode. The dashed vertical line indicates the value of $\gamma_{\text{max}}$ for the ETG mode in the absence of cross-scale interaction. (b) The average ETG growth rate $\langle \gamma \rangle_{t_s,x_s,y_s}$ as a function of $(k_y, \hat{\theta}_0)$.

Figure 6: (a) The $k_y$ of the fastest growing ETG mode. (b) The $\hat{\theta}_0$ of the fastest growing ETG mode. The dashed lines indicate the $k_y$ and $\hat{\theta}_0$ of the ETG mode in the absence of IS turbulence. Only modes with $\gamma_{\text{max}} > 0$ are included in the figures.

4. Physical Mechanisms and Interpretation

The cross-scale terms in equation (2.2), $\tilde{v}_E \cdot \nabla \tilde{g}$ and $\tilde{v}_E \cdot \nabla \tilde{g}$, have intuitive physical interpretations. The gradient $\nabla g$ modifies the local gradients of the background Maxwellian distribution function $\nabla F_0$, and thus modifies the instability drive. The drift $\tilde{v}_E$ is a cross-field flow that advects ES fluctuations and that can vary strongly in the parallel-to-the-field direction. This parallel-to-the-field variation introduces a new parallel length scale and means that $\tilde{v}_E$ cannot simply be removed from the equations by changing to a rotating or boosted frame. The most unstable toroidal instabilities typically have a parallel wave number $k_\parallel$ that is set by the connection length $qR$, i.e. $k_\parallel qR \sim 1$, cf. [Parisi et al. (2020)]. Parallel-to-the-field shearing of the ETG fluctuations acts to increase $k_\parallel$ by imposing a parallel length scale such that $k_\parallel qR \gtrsim 1$, and hence stabilises toroidal modes. We illustrate the physical picture for parallel-to-the-field shearing in figure 7. Figure 7a depicts the field-aligned structure of the toroidal ETG mode in the absence of parallel shearing. Figure 7b shows the result of imposing a cross-field flow that varies along the magnetic field line: the $k_\parallel$ of the ETG mode is now set by the parallel-to-the-field structure of $\mathbf{v}_E$, and not by the connection length $qR$. The drift $\mathbf{v}_E$ varies strongly in
Figure 7: (a) Cartoon of field-aligned contours of ETG-driven potential fluctuations in the absence of parallel-to-the-field shear. (b) The same ETG-driven mode in the presence of an $\mathbf{E} \times \mathbf{B}$ drift which varies in the direction of $\mathbf{b}$ on a scale shorter than the connection length by a factor of order unity: the parallel-to-the-field scale of the mode is shortened.

the direction of $\mathbf{b}$ because at a specific $x_s$ and $y_s$ the IS potential $\overline{\phi} = \overline{\phi}(\theta)$. We note that we may express $\nabla \overline{g} = \nabla \overline{h} + \nabla (e \overline{\phi}/T_e) F_0$, with $\overline{h}$ the nonadiabatic response of electrons at IS. In the orbit-averaged model $\overline{h}$ is a constant in $\theta$, at fixed $(\varepsilon, \lambda)$, and is zero in the passing piece of the velocity space for modes which oscillate in $y$ (nonzonal modes) [Hardman et al. 2019]. Despite this, $\overline{g}$ can vary strongly in the direction of $\mathbf{b}$ because of the contribution from the adiabatic response $e \overline{\phi} F_0/T_e$. Hence, $\nabla \overline{g}$ contains a contribution to the density gradient $\nabla (e \overline{\phi}/T_e) F_0$ that has the effect that the total modified drive of instability $\nabla \overline{g} + \nabla F_0$ drives the mode nonuniformly in $\theta$: this effect can also impose an increased $k_\parallel$ on the mode. We note that the contribution to $\nabla \overline{g}$ from the nonadiabatic electron response $\nabla \overline{h}$ contains modifications to the background density and temperature gradients, and gradients of higher-order velocity moments that may drive or suppress instabilities. It is worth noting that the velocity moments of $\overline{h}$ do in general depend on $\theta$: this is a consequence of the fact that trapped particles can only access a limited range of $\theta$, and the fact that the nonzonal, passing electron nonadiabatic response vanishes in the orbit-averaged model.

To shed light on the physical mechanisms underlying the result shown in figure 5, we separately consider the effect of including $\nabla \overline{g}$ and $\overline{\mathbf{v}}_E$. In figure 8a we show the histogram of the maximum ETG growth rate $\gamma_{\text{max}}$ computed including only the cross-scale term $\overline{\mathbf{v}}_E \cdot \nabla \overline{g}$ (setting $\nabla \overline{g} = 0$). In figure 8b we show the histogram of $\gamma_{\text{max}}$ computed including only the cross-scale term $\overline{\mathbf{v}}_E \cdot \nabla \overline{g}$ (setting $\overline{\mathbf{v}}_E = 0$). We see by comparing figures 5a, 8a, and 8b that both the effect of parallel-to-the-field shearing and modifications to the background gradients are important in the ETG stabilisation shown in figure 5. From figure 8a we see that the effect of $\overline{\mathbf{v}}_E$ appears to be uniformly stabilising; i.e., there are no instances where $\gamma_{\text{max}}$ becomes larger than the single-scale ETG growth rate. However, in figure 8b we see that the effect of $\nabla \overline{g}$ almost always stabilises the ETG mode, but in rare instances can make the ETG mode more unstable. In the following sections we study these cross-scale physics mechanisms in more detail and provide qualitative explanations for these results.

We note that our decomposition of the cross-scale terms, while physically motivated, is not unique. In particular, we recall that we can further decompose $\nabla \overline{g} = \nabla \overline{h} + \nabla (e \overline{\phi}/T_e) F_0$. In figures 8c and 8d we show the cross-scale effects of $\nabla \overline{h}$ and $\nabla (e \overline{\phi}/T_e) F_0$ on the ETG mode separately. The piece due to the nonadiabatic response $\nabla \overline{h}$ can both drive and suppress the ETG mode, whereas the piece due to the adiabatic response...
\[ \nabla (e\bar{\phi}/T_e)F_0 \] appears to only suppress the ETG mode. We give a qualitative explanation for this observation in section 6.

5. Evidence for parallel-to-the-field shearing

In this section we show that the physical picture given in figure 7 is consistent with the simulation results presented in section 3: the IS turbulence indeed shortens the parallel-to-the-field scale of the ETG modes. To demonstrate this we first consider the structure of the cross-scale terms and the structure of the ETG eigenmode in the ballooning coordinate, labelled by \( \theta \). The fastest-growing ETG mode from a single-scale simulation with CBC parameters occurs at \( k_y \rho_{th,e} = 0.58 \) and \( \hat{\theta}_0 = 0 \). Figure 9a compares the potential eigenmode for the \( k_y \rho_{th,e} = 0.58 \) and \( \hat{\theta}_0 = 0 \) mode from a single-scale simulation, to the \( k_y \rho_{th,e} = 0.58 \) and \( \hat{\theta}_0 = 0 \) mode obtained at a specific example space-time location in the IS turbulence. We see that the cross-scale interaction introduces oscillatory features into the potential eigenmode. This is a result of the variation of \( v_E \) and \( \nabla \bar{g} = \nabla (e\bar{\phi}/T_e)F_0 + \nabla \bar{h} \) along the magnetic field line. To illustrate this, in figure 9b we show the imposed IS \( E \times B \) drift frequency

\[ \nabla_{E \times B} = c\kappa_\rho \frac{\rho_{th}}{B_{ref}} \left( k_x \frac{\partial \bar{\phi}}{\partial y} - k_y \frac{\partial \bar{\phi}}{\partial x} \right), \]  

(5.1)

with \( k_x \) the field-aligned radial wave number of the ETG mode, i.e., the wavenumber \( k = k_x \nabla x + k_y \nabla y \), and the geometrical factor \( \kappa_\rho = B_{ref}(dx/d\psi)(dy/d\alpha) \simeq 1.01 \) for CBC
Figure 9: (a) The real part of the normalised potential eigenmode for the fastest-growing ETG mode for i) a specific instance of IS gradients and NoIS) the eigenmode in the absence of cross-scale coupling. The normalised growth rates are $\gamma/(v_{th,e}/a) = 0.045$ and 0.094, respectively. For simplicity, we chose to compare a case where the fastest-growing mode in the presence of IS turbulence appears at $k_y \rho_{th,e} = 0.58$ and $\hat{\theta}_0 = 0.0$. (b) The thermal magnetic drift frequency $\omega_M$ compared to the IS $E \times B$ drift frequency $\omega_{E \times B}$. (c) The drive frequencies due to equilibrium density gradients $\omega^*_n$, and due to gradients in the adiabatic part of the IS, electron distribution function $\omega^*_\phi$. (d) The drive frequencies due to equilibrium temperature gradients $\omega^*_T$, and gradients of temperature and density in the nonadiabatic part of the IS, electron distribution function, $\omega^*_{\delta T_e}$ and $\omega^*_{\delta n,h}$, respectively.

parameters. The frequency $\omega_{E \times B}$ is compared to the thermal magnetic drift frequency

$$\omega_M = \frac{v^2_{th,e}}{\Omega_e} k \cdot b \times \left( \frac{\nabla B}{B} + b \cdot \nabla b \right). \quad (5.2)$$

In figure 9c we show the drive frequencies associated with the $\nabla (e\phi/T_e) F_0$ piece of $\nabla g$, $\omega^*_\phi = -\omega_{E \times B}$, compared to the background density gradient drive $\omega^*_n = -ck_y T_e/eB_{ref} L_n$; and in figure 9d we show the background temperature gradient drive frequency $\omega^*_T = -ck_y T_e/eB_{ref} L_{T_e}$, and the drive frequencies due to the density $\delta n_h$ and temperature $\delta T_e$ moments of $\nabla h$,

$$\omega^*_{\delta n,h} = \frac{c k_y T_e}{e B_{ref}} \left( \frac{k_y}{n} \frac{\partial \delta n_h}{\partial x} - \frac{k_x}{n} \frac{\partial \delta n_h}{\partial y} \right) \quad (5.3)$$
Figure 10: A comparison of the power spectrum of the electrostatic potential $\tilde{\phi}$ of the fastest growing ETG ballooning mode as a function of $k_\parallel$ for several cases. In case NoIS) we show the power spectrum of the fastest growing ballooning mode in the absence of cross-scale coupling. The remaining curves show the $(t_s, x_s, y_s)$-averaged power spectrum of the fastest growing ETG mode in the case that: i) both cross-scale terms $\mathbf{v}_E \cdot \nabla \tilde{g}$ and $\tilde{v}_E \cdot \nabla \tilde{g}$ are included; ii) only the shearing term $\mathbf{v}_E \cdot \nabla \tilde{g}$ is included; iii) only the modification to $\nabla F_0$ from the adiabatic response of electrons is included, i.e., $\tilde{v}_E \cdot \nabla (e\phi/T_e) F_0$; and iv) only the modification to $\nabla F_0$ from the nonadiabatic response of electrons is included, i.e., $\tilde{v}_E \cdot \nabla \tilde{h}$. To measure the spread of the distribution of $k_\parallel$ we use $\sigma_{k_\parallel}$, defined in equation (5.6). We find i) $\sigma_{k_\parallel} = 0.065$, ii) $\sigma_{k_\parallel} = 0.064$, iii) $\sigma_{k_\parallel} = 0.043$, iv) $\sigma_{k_\parallel} = 0.031$, and NoIS) $\sigma_{k_\parallel} = 0.027$.

and

$$\mathbf{\omega}^{*}_{\delta r_e} = \frac{c k_\rho T_e}{e B_{ref}} \left( k_y \frac{\partial \tilde{M}_e}{\partial x} - k_x \frac{\partial \tilde{M}_e}{\partial y} \right),$$

(5.4)

respectively. There are two key feature to note in figures 9b, 9c, and 9d. Firstly, the frequencies that mediate the effects of IS turbulence on ETG modes are comparable in amplitude to the equilibrium drive and thermal magnetic frequencies. Secondly, the frequencies vary in ballooning $\theta$ on the scale of the ETG eigenmode in 9a. In the strongly driven turbulence that we consider here, we have found that the $\theta$ dependence of $\mathbf{\omega}_{E \times B}$ cannot be described by a simple function with few parameters. In consequence, simplified quantitative modelling of the effects of parallel-to-the-field shearing by $\mathbf{\tilde{v}}_E$ is challenging. As we discuss in section 6, we find that it is possible to approximate $\mathbf{\omega}^{*}_{\delta n,h}$ and $\mathbf{\omega}^{*}_{\delta r_e}$ by a constant value.

To assess whether or not the physical picture in figure 7 is consistent with the simulation results, we calculate the $k_\parallel$ spectrum of the fastest growing ETG mode at each $(t_s, x_s, y_s)$ in the sample. For a single ballooning mode, we define the $k_\parallel$ spectrum to be $|\tilde{\phi}(k_\parallel)|^2$, where

$$\tilde{\phi}(k_\parallel) = \int_{-\infty}^{\infty} \exp[i k_\parallel z] \tilde{\phi}(z) dz,$$

(5.5)

cf. Parisi et al. (2020), with $z = \theta / \mathbf{b} \cdot \nabla \theta$, and $\tilde{\phi}(z)$ is normalised so that the maximum value of $\tilde{\phi}(z)$ is 1 — this maximum may occur away from $z = 0$. We note that we define the $\theta$ coordinate in the simulations such that $\mathbf{b} \cdot \nabla \theta$ is constant in $\theta$. In figure 10 we compare $|\tilde{\phi}(k_\parallel)|^2$ for the fastest growing ETG mode in the absence of cross-scale interaction with the average $k_\parallel$ spectrum $\langle |\tilde{\phi}(k_\parallel)|^2 \rangle$ of the fastest growing ETG modes in
several calculations including various combinations of the cross-scale interaction terms. Figure 10 indicates that cross-scale interaction introduces tails in the $k_\parallel$ spectra. These tails correspond to the oscillatory structure of the eigenmode shown in figure 9a. To quantify the size of these tails, we introduce a measure of the spread of the distribution of $k_\parallel$:

$$\sigma^2_{k_\parallel} = \frac{\int k_\parallel^2 |\tilde{\phi}(k_\parallel)|^2 dk_\parallel}{\int |\phi(k_\parallel)|^2 dk_\parallel}.$$  \hspace{1cm} (5.6)

For the fastest-growing ETG mode in the absence of cross-scale interaction $\sigma_{k_\parallel} = 0.027$. For the average $k_\parallel$ spectrum of the fastest-growing ETG modes in the presence of all cross-scale interaction terms we find that $\sigma_{k_\parallel} = 0.065$. This indicates that the effect of cross-scale interaction is indeed to introduce larger $k_\parallel$ into the ETG mode. Figure 10 indicates that the dominant component of the $k_\parallel$ tails is introduced by the parallel shearing term $v_E \cdot \nabla \tilde{g}$. We also note that the term $\tilde{v}_E \cdot \nabla \tilde{h}$ introduces almost no $k_\parallel$ component: cross-scale interaction via $\tilde{v}_E \cdot \nabla \tilde{h}$ leaves the form of the ETG eigenmode unchanged. This is consistent with the fact that $\tilde{h}$ is constant in $\theta$ at fixed $(\varepsilon, \lambda, \sigma, x, y)$ in the parallel-orbit-averaged model.

6. Modifications to the drives of instability

The physical interpretation of the impact of $\nabla \tilde{g}$ in equation (2.2) is that $\nabla \tilde{g}$ modifies the background gradients in $\nabla F_0$ which drive the instability. As discussed in section 4, $\nabla \tilde{g}$ is in general a complicated function of $(\theta, \varepsilon, \lambda)$. In this section we examine to what extent we can model $\nabla \tilde{g}$ by simple modifications to the values of $a/L_e$ and $a/L_n$ locally in the IS flux tube. For this we use the results of the simulations presented in figures 8b, 8c and 8d. This simplified model should be valid when the plasma is collisional and the fluctuations are zonally dominated: in this limit, $\tilde{h}$ can be approximated as a perturbed Maxwellian that has no variation within a flux surface. We realise this limiting scenario in section 7.

We present the maximum ETG growth rate for the CBC magnetic geometry as a function of the background temperature and density gradients, $\gamma_{\text{CBC max}}(a/L_e, a/L_n)$, in figure 11a. We define effective background temperature and density gradients,

$$\frac{a}{L^\text{eff}_{T_e}} = \frac{a}{L_{T_e}} - \left\langle \frac{a \kappa_\rho}{T_e} \frac{\partial \overline{\delta T_e}}{\partial x} \right\rangle_\theta,$$  \hspace{1cm} (6.1)

$$\frac{a}{L^\text{eff}_{n,h}} = \frac{a}{L_n} - \left\langle \frac{a \kappa_\rho}{n} \frac{\partial \overline{\delta n_h}}{\partial x} \right\rangle_\theta,$$  \hspace{1cm} (6.2)

and

$$\frac{a}{L^\text{eff}_{n}} = \frac{a}{L_n} - \left\langle \frac{a \kappa_\rho}{n} \frac{\partial \overline{\delta n_h}}{\partial x} \right\rangle_\theta,$$  \hspace{1cm} (6.3)

respectively, with $\left\langle \cdot \right\rangle_\theta = \int_{-\pi}^{\pi} \cdot d\theta / B \cdot \nabla \theta$ an average in $\theta$ between $(-\pi, \pi)$ at fixed $(x, y)$, $\overline{\delta n_h}$ and $\overline{\delta T_e}$ the density and temperature moments of $\tilde{h}$, respectively, and $\overline{\delta n} = \overline{\delta n_h} + ne\overline{\phi}/T_e$. The normalised density gradient scale $a/L^\text{eff}_{n,h}$ measures the effective change to $a/L_n$ by density gradients in the long-wavelength nonadiabatic response of electrons, whereas $a/L^\text{eff}_{n}$ measures the effective change to $a/L_n$ by gradients in the total long-wavelength density fluctuation.
In figure 11b we show that there is a strong correlation between the $\gamma_{\text{max}}$ calculated using equation (2.2) including only the cross-scale term due to the IS, nonadiabatic electron response $\hat{\mathbf{v}}_E \cdot \nabla h$ and the model growth rate $\gamma_{\text{max}}^{\text{CBC}}(a/L_{T_e}, a/L_n)$. In contrast, in figure 11c, we show that the cross-scale interaction due to $\nabla g$ cannot be modelled by $\gamma_{\text{max}}^{\text{CBC}}(a/L_{T_e}, a/L_{n})$. Instead, the model $\gamma_{\text{max}}^{\text{CBC}}$ almost always overestimates the growth rate of the ETG modes. We can understand these results by referring to figure 10. Cross-scale interaction due to the electron nonadiabatic response $\nabla h$ did not change the $k_{\parallel}$ spectrum of the ETG eigenmodes: $\nabla h$ has the effect of making the mode more or less unstable by changing the drives only. In contrast the difference between $\nabla g$ and $\nabla h$, $\nabla(\epsilon\phi/T_e)F_0$, had the effect of introducing larger $k_{\parallel}$ components into the mode, which cannot be captured by simple modifications to $a/L_n$ and $a/L_{T_e}$. The fact that $\gamma_{\text{max}}^{\text{CBC}}(a/L_{T_e}, a/L_{n})$ typically overestimates the ETG growth rate in figure 11c indicates that the effect of the parallel-to-the-field variation in $\nabla(\epsilon\phi/T_e)F_0$ is stabilising. These results suggest that in general we can characterise $\nabla g$ as having two distinct components: firstly, a perturbed Maxwellian that simply shifts the equilibrium density and temperature gradients by constant values; and secondly, a density gradient from the long-wavelength electron adiabatic response that varies in the parallel-to-the-field coordinate. Whilst the former component may in general stabilise or destabilise the ETG
mode, the latter component almost always appears to be stabilising. The effect of $\nabla g$ on an ETG mode depends on the relative importance of these two components.

Finally, we return to the result presented in figure 8b: the average effect of $\nabla g$ was to suppress ETG instability. We can now understand this result qualitatively. Firstly, the parallel-to-the-field variation in the adiabatic response $\nabla (e\bar{\phi}/T_e)F_0$ suppressed the instability by increasing the $k_\parallel$ in the ETG mode. Secondly, note that for particular the CBC equilibrium we have chosen, $a/L_n$ maximises the ETG growth rate at fixed $a/L_T$, as shown in figure 11a: any modification to $a/L_n$ will reduce the ETG growth rate $\gamma_{\text{max}}$. The histogram of density gradient length scales $a/L_{\text{eff}}$ shown in figure 11d indicates that the typical modification to $a/L_n$ is sufficient to reduce $\gamma_{\text{max}}$ by an amount of order unity. This second stabilising effect is likely to be special to this equilibrium.

7. Cross-scale interaction in near-marginal turbulence

The simulation results presented in the previous sections were obtained for parameters where there is a strong background drive of instability for ITG modes at scales comparable to $\rho_{\text{th,i}}$ and ETG modes at scales comparable to $\rho_{\text{th,e}}$. In this section we present results for the case where $a/L_T$ and $a/L_n$ are reduced so that both the ITG and ETG modes are driven near marginal stability. Weakly driven IS turbulence can reach a Dimits-shift regime where the fluxes are low, and a long-wavelength, long-lived zonal flow regulates the IS dynamics; see, e.g., Dimits et al. (2000); Rogers et al. (2000). This represents a very different state from strongly driven, critically balanced, IS turbulence, cf. Barnes et al. (2011). Consequently, the relative importance of parallel-to-the-field shearing, parallel-to-the-field variation in the IS density gradient, and constant modifications to the background drives may differ in turbulence driven near marginal stability.

To obtain a near-marginal case, we use the parameters described in section 3 but reduce the normalised background temperature gradient length scales to $a/L_T = a/L_n = 1.38$. We use the resolutions given in appendix A.1 and we take $m_e/m_i = 2.7 \times 10^{-12}$ in the IS simulation. The IS turbulence is evolved until a long-lived, zonally dominated state emerges. We sample the IS gradients $\nabla \bar{g}$ and drifts $\bar{\mathbf{v}}_E$ at radial coordinates in the IS flux tube $x_s$, for a single time $t_s$ and binormal position $y_s$. In figure 12a we show $\bar{\phi}(x, y)$.
at $t_s$ for the near-marginal IS turbulence, with the sampled radial points taken along the dashed line. At each of the 128 equally spaced $x_s$ positions we perform an ETG linear calculation, including the effects of cross-scale interaction via equation (2.2) and using resolutions detailed in appendix A.2. In contrast to the strongly driven case, in figure 12b we show that the maximum growth rate of the ETG mode $\gamma_{\text{max}}$ – calculated including both cross-scale interaction terms, $\bar{\nu}_E \cdot \nabla \bar{g}$ and $\bar{v}_E \cdot \nabla \bar{g}$ – is strongly correlated with the model growth rate $\gamma_{\text{CBC}}^{\text{max}}(a/L_{\text{Te}}^{\text{eff}}, a/L_n^{\text{eff}})$, with the effective density gradient scales $a/L_{\text{Te}}^{\text{eff}}$ and $a/L_n^{\text{eff}}$ defined in equations (6.1) and (6.3), respectively, and $\gamma_{\text{CBC}}^{\text{max}}$ shown in figure 11a. The imperfections in the correlation in figure 12b are primarily due to parallel-to-the-field shearing from the term $\bar{v}_E \cdot \nabla \bar{g}$ in equation (2.2) (cf. figure 10 which shows that $\bar{v}_E \cdot \nabla \bar{g}$ injects larger $k_\parallel$ components than $\bar{v}_E \cdot \nabla (e \phi/T_e) F_0$). We infer that the effect of parallel-to-the-field shearing is weak in this near-marginal turbulence. We verified that the average $k_\parallel$ spectrum of the fastest growing eigenmodes was identical in the cases with and without cross-scale interaction.

We can understand the result in figure 12b with the following intuition. The wavenumber spectrum of near-marginal turbulence is dominated by the zonal mode, a mode which is constant on each flux surface and only depends on the radial coordinate. In consequence, perpendicular gradients in near-marginal turbulence are dominantly in the radial direction, and the variation of the the electrostatic potential $\phi$ in the parallel-to-the-field direction in the turbulence is weak. In this regime, the $E \times B$ advection term $\bar{\nu}_E \cdot \nabla \bar{g}$ only modifies the ETG frequency by a Doppler shift (Hardman et al. 2019). Hence, the dominant cross-scale interaction term that alters the ETG growth rate is due to the radial gradients of the IS electron distribution function, i.e., $\nabla \bar{g} \simeq \nabla x \partial \bar{g} / \partial x$. In addition, if the plasma is sufficiently collisional then $\bar{g}$ is restricted to be a perturbed Maxwellian of only density and temperature gradients, and in consequence, the effect of cross-scale interaction in near-marginal turbulence may be parameterised with the effective background gradient length scales $a/L_{\text{Te}}^{\text{eff}}$ and $a/L_n^{\text{eff}}$. The results presented in this section demonstrate that this near-marginal cross-scale interaction regime can be realised.

8. Discussion

The results presented in this paper were obtained in the limit $(m_e/m_i)^{1/2} \rightarrow 0$ with a theory that assumes scale-separation between long-wavelength and short-wavelength structures. This theory may not describe multiscale turbulence accurately for realistic values of $(m_e/m_i)^{1/2}$ under a variety of circumstances; e.g., when radially elongated ES structures (Dorland et al. 2000; Jenko et al. 2000; Jenko & Dorland 2002) grow to be as large as a typical IS eddy; or when there is no significant dissipation at wavelengths between the scales of the IS and ES turbulence. Nonetheless, in this paper we have demonstrated that significant cross-scale interactions do persist in the limit $(m_e/m_i)^{1/2} \rightarrow 0$: strongly driven, long-wavelength turbulence can stabilise a strongly driven ETG instability; and long-wavelength turbulence driven near marginal stability can change the growth rate of a weakly driven ETG mode by an order unity factor. The physical mechanisms responsible for these cross-scale interactions were shown to be, firstly, the parallel-to-the-field shearing of short-wavelength fluctuations by long-wavelength $E \times B$ drifts, and secondly, the modification of the background drives of instability by gradients in long-wavelength fluctuations. We demonstrated that the modification to the background drives could be characterised by two distinct pieces – a perturbed Maxwellian that affects a constant modification to the mean density gradient
a/L_n and temperature gradient a/L_{T_e}, and a density gradient from the long-wavelength electron adiabatic response that varies in the parallel-to-the-field direction. In the case of strongly driven turbulence, the novel effects of both parallel-to-the-field E × B shearing and parallel-to-the-field variation in the long-wavelength density gradient were found to be stabilising. In the case of near-marginal turbulence, the effect of parallel-to-the-field E × B shearing was weak, parallel-to-the-field variation in the long-wavelength density gradient was unimportant, and the effect of cross-scale interaction could be parameterised solely with the effective background gradient length scales a/L_{n_{eff}} and a/L_{T_{eff}}.

We can qualitatively explain the difference in the nature of the cross-scale interactions that we find in the cases of near-marginal and strongly driven turbulence. Near-marginal ITG turbulence is dominated by zonal modes, see e.g., Dimits et al. (2000); Rogers et al. (2000). Zonal modes have a structure that depends only on radial position, and hence the long-wavelength E × B drift \( \vec{v}_E \) only has the effect of shifting the ETG fluctuation frequency by a Doppler shift – for zonal modes where \( \vec{\phi} = \bar{\phi}(x) \) the term \( \vec{v}_E \cdot \nabla \bar{g} \) in equation (2.2) can be removed by boosting to a toroidally rotating frame (Hardman 2019). In this situation where the effect of parallel shearing is negligible, the dominant cross-scale interaction arises from the radial gradient \( \nabla g \approx \nabla x \partial g/\partial x \). If \( g \) is sufficiently close to a perturbed Maxwellian – perhaps as a result of moderate collisionality – then we find that \( \partial g/\partial x \) has dominant contributions from a density gradient and a temperature gradient. In contrast, strongly driven ITG turbulence is not zonally dominated, but consists of critically balanced turbulent eddies (Barnes et al. 2011) that can have a parallel correlation length as large as the device scale \( qR \). Hence, \( \nabla g \) and \( \vec{v}_E \) can contain parallel-to-the-field variation that results in the nonuniform drive and the shearing apart of short-wavelength modes, respectively. In critically balanced turbulence, eddies with a greater \( k_y \) have a faster nonlinear turnover time, and hence a shorter parallel correlation length and a greater characteristic \( k_{||} \times k_y^{4/3} (R/L_{T_i})^{4/3} \) (Barnes et al. 2011). A range of \( k_y \) modes in the inertial range can contribute to the gradient \( \nabla g \) and the drift \( \vec{v}_E \), with the result that \( \nabla g / \partial x \) and \( \vec{v}_E \) can have parallel-to-the-field oscillations on scales shorter than \( qR \). This observation leads us to expect that the impact of parallel-to-the-field shearing increases with increasing \( R/L_{T_i} \).

A criterion for when to expect parallel-to-the-field E × B shear to suppress ETG modes can be obtained from a simple quasilinear argument. The effective shearing rate \( \hat{\omega}_E \) should scale with the parallel-to-field variation in the cross-scale advection operator \( \vec{v}_E \cdot \nabla \), i.e., \( \hat{\omega}_E \sim k_{ETG} \partial \vec{v}_E / \partial \theta \). We can expect parallel-to-the-field shear suppression when the effective shear rate

\[
\hat{\omega}_E \gtrsim \gamma_{ETG};
\]  

(8.1)
i.e., when the shear rate is faster than the typical ETG mode growth rate \( \gamma_{ETG} \), at a typical ETG mode wave number \( k_{ETG} \). A quasilinear estimate gives \( \vec{v}_E \sim \gamma_{ITG}/k_{ITG} \), with \( \gamma_{ITG} \) and \( k_{ITG} \) the typical ITG mode growth rate and wave number, respectively. If we assume that the variation of \( \vec{v}_E \) in the parallel-to-the-field direction \( \partial \vec{v}_E / \partial \theta \sim \vec{v}_E \), then we find that we can expect parallel-to-the-field shear suppression of ETG turbulence when

\[
\frac{\gamma_{ITG}}{k_{ITG}} \gtrsim \frac{\gamma_{ETG}}{k_{ETG}}.
\]  

(8.2)

This result is consistent with observations of the behaviour of multiscale turbulence made in light of some DNS (Staebler et al. 2016; 2017; Creely et al. 2019), though their interpretation differs from the one given here: Staebler et al. (2016; 2017); Creely et al. (2019) neglect the variation of E × B drifts along field lines, and so neglect the key
physical mechanism that is critical to the quasilinear argument: the effect of parallel-to-the-field $E \times B$ shear. We note that parallel-to-the-field flow shear has previously been considered as a part of a model of hyperviscous dissipation of ITG driven turbulence (Smith, 1997).

The suppression of short-wavelength turbulence seen in some DNS is often assumed to be the result of perpendicular-to-the-field $E \times B$ shearing by long-wavelength turbulence. In fact, that mechanism does not appear in the leading-order, scale-separated ES equation (2.2) (Hardman et al., 2019). In the limit that $(m_e/m_i)^{1/2} \to 0$, ES structures do not have a large enough spatial extent to be affected by perpendicular-to-the-field shear in IS $E \times B$ flows.

Finally, we note that the scale-separated model should be modified if factors that we have assumed to be of order unity become large enough to interfere with the $m_e/m_i^{1/2}$ expansion; possible examples of such parameters include the ratio $(\gamma_{ETG}/k_{ETG})/(\gamma_{ITG}/k_{ITG})$, the degree of spatial anisotropy in the turbulence, and the ratio of the zonal to nonzonal fluctuation amplitudes (Hardman et al., 2019).

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The simulations presented in this work were performed using the GS2 branch [https://bitbucket.org/gyrokinetics/gs2/branch/ms_pgelres], with the latest revision at the time of writing being commit ade5780. GS2 input files for the simulations presented in this work may be found at [https://doi.org/10.5287/bodleian:b7x0008d].

Appendix A. Simulation Resolutions

This appendix details the resolutions used for the simulations presented in this paper.

A.1. Nonlinear simulations

For the nonlinear simulations presented in section 3, we use a flux tube with a perpendicular cross-section of $40\pi \rho_{th,i} \times 40\pi \rho_{th,i}$, and a single poloidal turn in the parallel-to-the-field direction. We use $n_y = 21$ toroidal modes with $k_y > 0$ and $k_{y,\text{max}} \rho_{th,i} = 1.05$. We use $n_x = 255$ radial modes with $k_{x,\text{max}} \rho_{th,i} = 6.38$. We specify that each $(k_x, k_y)$ mode has a $[-\pi, \pi]$ extent in $\theta$, with $n_\theta = 33$ points. We use a $\lambda$ grid with $n_\lambda = 27$. To describe passing particles we take $n_{\lambda,p} = 11$ points, chosen using Gauss-Radau rules (Hildebrand, 1987). For the trapped particles we employ a nonspectral, unequally spaced grid in $\lambda$, with $n_{\lambda,t} = 16$ points at the outboard midplane. Following Barnes et al. (2010), we use an $\varepsilon$ grid derived from a spectral speed $v = \sqrt{2\varepsilon/m}$ grid with $n_\varepsilon = 12$ points.
For the near-marginal nonlinear simulation presented in section 7, we use the resolutions described above, with \( n_\varepsilon = 16 \).

### A.2. Linear simulations

For the linear simulations presented in sections 3 - 7, we use ballooning modes with an extent \([-3\pi, 3\pi]\) in the ballooning angle. For compatibility with the nonlinear simulations, each \([-\pi, \pi]\) segment has \( n_\theta = 33 \) points, and we take \( n_\lambda = 27 \). For the near-marginal linear simulations presented in section 7 and the linear simulations used to calculate the growth rate \( \gamma_{\text{CBC}}^{\text{max}} \) in figure 11a, we take \( n_\varepsilon = 16 \). For the linear simulations presented in sections 3 - 6 (excluding those used to calculate \( \gamma_{\text{CBC}}^{\text{max}} \)) we take \( n_\varepsilon = 12 \). For all ES linear simulations we consider 21 toroidal modes in the wave number range \( 0.04 \leq k_y \rho_{\text{th},e} \leq 0.86 \) (\( 2.5 \leq k_y \rho_{\text{th},i} \leq 52.5 \)), and we consider 21 values of \( \hat{\theta}_0 \) in the range \(-\pi < \hat{\theta}_0 < \pi \). The growth rate \( \gamma_{\text{CBC}}^{\text{max}} \) in figure 11a is computed for \( \hat{\theta}_0 = 0 \) and \( 0.04 \leq k_y \rho_{\text{th},e} \leq 0.86 \).

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