

Fast Turbulence Phase Transition in a Flux-Driven Global Edge-SOL Simulation of a Tokamak Plasma

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A power ramp from L- to H-mode-like conditions is carried out in a 50 ms long global full- f turbulence simulation of ASDEX Upgrade with the GRILLIX code. After 32 ms of slow profile evolution, a sudden change in transport triggers a pedestal buildup within ~ 100 μ s. This transition results from an abrupt shift from drift-wave (DW) to kinetic-ballooning-mode (KBM) turbulence. The self-consistently evolved $E \times B$ flow shear first increases slowly and then sharply. Geodesic acoustic mode and Alfvénic flow oscillations intensify prior to the transition and subsequently die out. This is followed by slow oscillations of the transport and profiles, reminiscent of a dithering I-phase or edge-localized-mode (ELM) cycles observed experimentally. This study paves the way for a better understanding of confinement regime transitions.

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Introduction—Suppressed turbulent transport is of major interest for the operation of magnetic confinement fusion reactors. The H-mode [1,2] is routinely accessed on all major tokamaks today and is assumed to be accessible in future devices such as ITER. Many theories have been proposed to explain the L-H transition [3–7]. Nonetheless, 40 years after its discovery, reactor design still mostly relies on extrapolations from today's experiments [6,8,9], while first-principles-based predictive simulations of the phenomenon remain elusive [10–15].

We approach this challenge with the help of the GRILLIX code [16]. This tool has recently been validated against the ASDEX Upgrade (AUG) tokamak under stationary ITER baseline H-mode conditions [17], while L-mode simulations are carried out routinely [18–20]. Here, we report on a power ramp between the two regimes in a single confinement-time-long simulation. After slow transport-scale profile evolution, a fast transition between different turbulence regimes leads to a pedestal buildup, which to our knowledge is reproduced for the first time in a realistic global 3D simulation. Figure 1 provides a first glimpse of the turbulence morphology, which differs markedly between the two regimes, as discussed in more detail below.

Short code description—GRILLIX leverages the flux-coordinate independent (FCI) approach [21,22] for efficient

turbulence simulations across a domain from the pedestal top to the divertor [23]. The employed drift-fluid plasma model [17] is global and full- f , evolving fluctuations of arbitrary amplitude together with the plasma background. It is electromagnetic, includes electron and ion temperature dynamics, and features transcollisional extensions for the parallel heat conduction and viscosity [17]. Both neo-classical and turbulence driven $E \times B$ flow components develop self-consistently [17,24]. For a realistic plasma

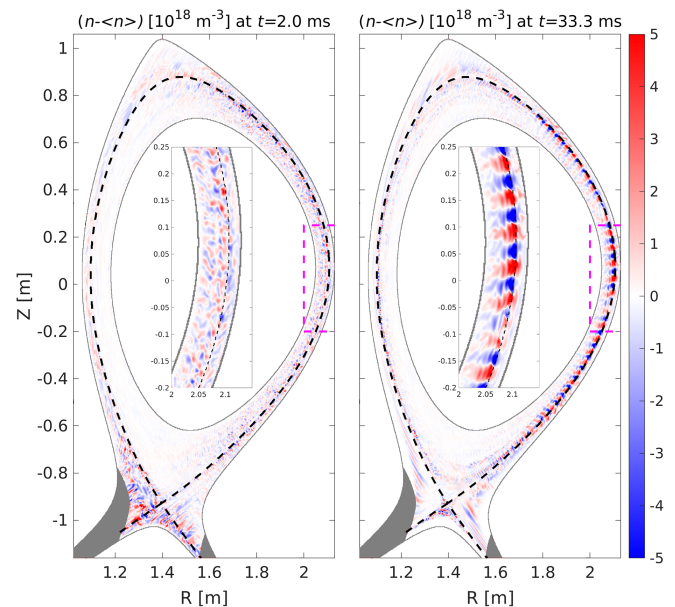


FIG. 1. Snapshots of plasma density fluctuations in a poloidal plane at 2 ms and 33.3 ms simulation time. Divertor and domain boundaries are marked in gray.

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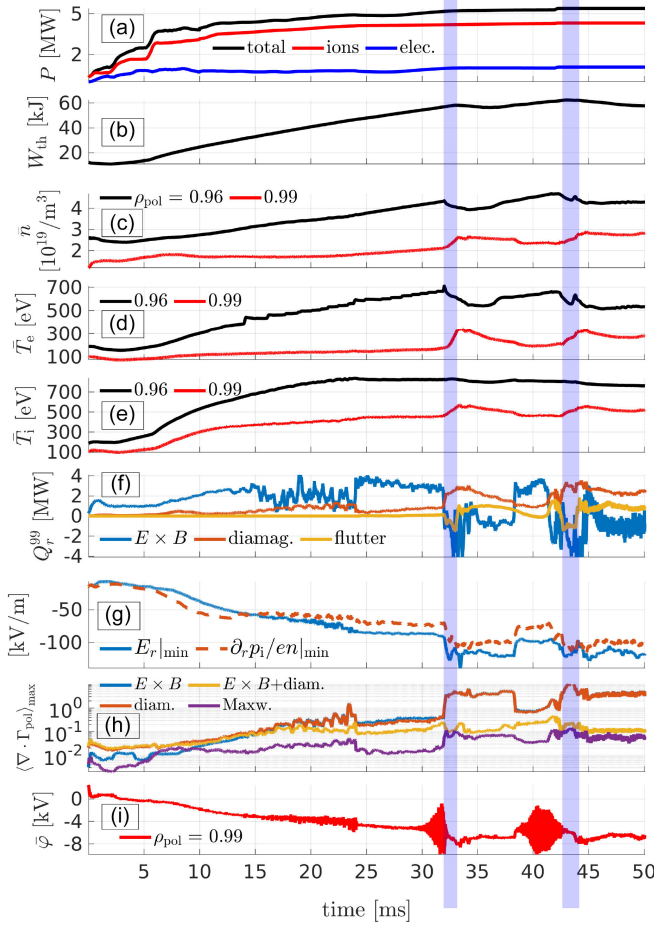


FIG. 2. Time traces of (a) total heating power, split between electrons and ions; (b) stored thermal energy $W_{th} = \int p dV$; (c) flux-surface averaged density at $\rho_{pol} = 0.96$ and 0.99 ; (d) same for electron and (e) ion temperature; (f) radial heat fluxes at $\rho_{pol} = 0.99$ separated into $E \times B$, diamagnetic, and magnetic flutter channels; (g) minimum E_r and ion pressure gradient; (h) polarization particle flux separated into $E \times B$ and diamagnetic Reynolds stress components, as well as Maxwell stress; and (i) flux-surface averaged electrostatic potential at $\rho_{pol} = 0.99$.

density source and divertor cooling, the code is coupled to a diffusive neutral gas model [18,20].

Simulation setup and caveats—With all of the above ingredients, the main trigger for the L-H transition is expected to be an increase in applied heating power, as displayed in Fig. 2(a). This is also the main difference with respect to the simulation setup detailed in our previous publication on the validation against an ASDEX Upgrade H-mode discharge (No. 40411) [17]. In fact, all physical and numerical parameters are the same as in Ref. [17], including the magnetic equilibrium. The domain is restricted to $\rho_{pol} \in (0.91, 1.035)$, as is depicted in Fig. 1. However, the initial profiles are chosen close to our L-mode validation [18,19], and the heating power ramp starts slowly from 700 kW. The outboard mid-plane (OMP) profiles at $t = 2$ ms are shown in Fig. 3 in blue.

The focus was put on increasing the ion heating since this is thought to be primarily responsible for the L-H

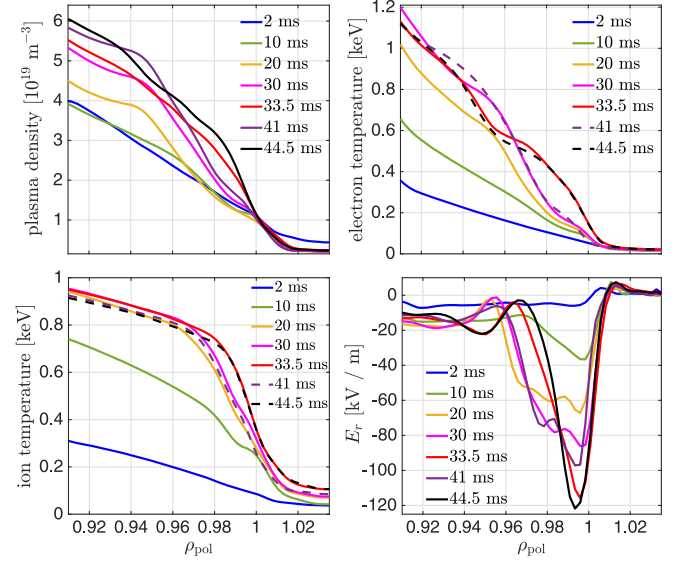


FIG. 3. Evolution of mean radial profiles of density, T_e , T_i , and radial electric field in the simulation, averaged over 300 μs .

transition [25]. We first tried to increase the heating in steps and then, in an attempt to speed up things, applied a large amount of heat at $\rho_{pol} = 0.91$ reaching up to 5.4 MW. However, the profiles evolved only slowly on the edge confinement time scale $\tau_{edge} = W_{th}/P \approx 10$ ms, with the heating power P and the thermal energy W_{th} stored in the simulation volume [shown in Fig. 2(b)]. Without the plasma core, W_{th} and τ_{edge} are much smaller than experimentally measured full-tokamak global quantities, but the simulation still turned out rather long: it took us a year to run it for 50 ms. A detailed discussion of simulation costs can be found in [17].

Particularly under H-mode conditions, magnetic fluctuations in a full- f turbulence simulation must be carefully filtered to avoid double counting the Shafranov shift and deterioration of the field-aligned numerics. To this end, we have recently developed a dynamical high-pass filter [26]: the magnetic shift is defined for the parallel vector potential via $\partial_r A_s = (A_{||} - A_s)/\lambda$, with $\lambda = 2$ ms, and the magnetic fluctuations are computed as $A_1 = A_{||} - \langle A_s \rangle_\phi$. We started the long simulation presented here with an earlier model [17] and switched to the new Zhang filter [26] at $t = 24$ ms. A discontinuity from the transition can be seen in the results, but we have carefully verified that this does not affect any key conclusions for $t > 24$ ms (see End Matter Appendix A). Note that the consistency of magnetic fluctuations with an evolving equilibrium remains an active area of research [26,27].

Macroscopic results—In the first 30 ms, the profiles and stored energy evolve slowly and continuously due to the applied heating. Notably, the dynamics differs between various observables and radii: at $\rho_{pol} = 0.99$, plasma density n and electron temperature T_e saturate within 5 ms, ion temperature T_i saturates only after 12 ms, and

radial electric field E_r reaches a plateau only at 25 ms; at $\rho_{\text{pol}} = 0.96$, T_i saturates after 25 ms, but n and T_e do not. We discuss this phase further in the *flow analysis* section. For now, let us focus on the discontinuous jump which occurs at 32 ms: particularly T_e , but also n and T_i , suddenly increase at $\rho_{\text{pol}} > 0.97$. At the same time, there is a drop at $\rho_{\text{pol}} < 0.97$. The morphology of density fluctuations in Fig. 1 changes drastically: broadband turbulent chaos is replaced with large but very coherent structures. In Fig. 3, we see that a pedestal with steep gradients develops in the outer part of the profiles (near the separatrix), while the inner part flattens. Also the E_r well develops slowly with the applied heating power, but then strongly deepens in the outer edge at 32 ms. This is at least qualitatively reminiscent of an L-H transition [2,6]. The new state is not yet stationary: the system relaxes back after a couple ms, and the process repeats at 42 ms. This is similar to either edge localized mode (ELM) cycles [28,29] or dithering I-phase oscillations [30–32], which could also be related [33].

Quantitatively, the E_r is significantly larger than typically measured in ASDEX Upgrade. As we see in Fig. 2(g), it mostly follows the ion pressure gradient, so the pedestal that we obtain is particularly steep. There are a number of possible explanations, such as the absence of gyrokinetic effects (e.g., trapped-electron, microtearing, or ETG modes [12,34–36]) or peeling modes (ELMs) [28,29] in our model (further discussed in End Matter Appendix B). Simulating the latter together with turbulence is challenging since large magnetic perturbations destroy field alignment [26]. Still, it is interesting that a fast pedestal buildup arises in our simulation, and the reasons for this are worth exploring.

In Fig. 2(f), we show flux-surface integrated heat fluxes at $\rho_{\text{pol}} = 0.99$, $Q_{r,j}^{99} = \oint Q_j r R d\phi d\theta_s$, divided into three channels j :

$$Q_{E \times B} = \frac{5}{2} n (T_e + T_i) \mathbf{v}_{E \times B} \cdot \mathbf{e}_r, \quad (1)$$

$$Q_{\text{diamag}} = \frac{5}{2} n (T_i^2 - T_e^2) \left(\nabla \times \frac{\mathbf{B}}{e B^2} \right) \cdot \mathbf{e}_r, \quad (2)$$

$$Q_{\text{flutter}} = \left(\frac{5}{2} n (T_e v_{\parallel} + T_i u_{\parallel}) + q_{\parallel e} + q_{\parallel i} \right) \mathbf{b}_1 \cdot \mathbf{e}_r. \quad (3)$$

The ion contribution is generally dominant due to the injected ion heating. $Q_{E \times B}$ is electrostatic turbulent transport; Q_{diamag} is essentially neoclassical transport due to up-down temperature asymmetries [37]; and Q_{flutter} is electromagnetic turbulent transport due to magnetic fluctuations $\mathbf{b}_1 = \mathbf{B}_1/B_0$, which can have a radial component. In the early L-mode phase, the transport is dominated by $Q_{E \times B}$. Then, with the heating power ramp and profile steepening, Q_{diamag} increases. Q_{flutter} remains negligible until 32 ms. But then, both $Q_{E \times B}$ and Q_{flutter} become negative: turbulent transport becomes electromagnetic and actually changes direction. This is balanced by the increased neoclassical

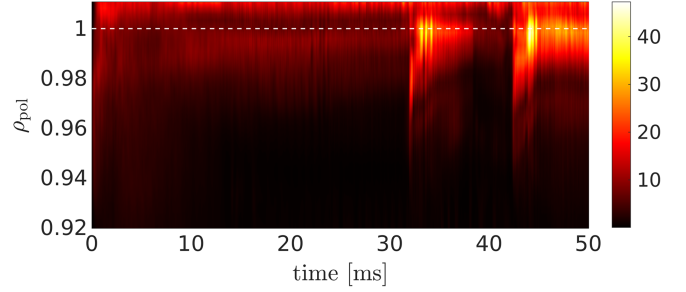


FIG. 4. Evolution of OMP density fluctuation amplitudes $\delta n = \sqrt{\langle n^2 \rangle_\phi / \langle n \rangle_\phi^2 - 1}$, computed via toroidal averages.

transport outward, the net transport becoming close to zero. Still, thermodynamically, heat is expected to flow down the gradient in all channels [38], neoclassical and turbulent. Hence, we are observing transient nonequilibrium behavior, only possible in a full- f flux-driven simulation. This is what leads to the abrupt profile steepening. The profiles at $\rho_{\text{pol}} < 0.97$ relax on the other hand: also there electromagnetic transport kicks in, but is directed radially outward. Next, we want to understand the reasons for this transition.

Microscopic turbulence analysis—A transport reduction could be expected to result from reduced fluctuation amplitudes. However, Fig. 4 shows that the density fluctuation amplitudes at the OMP actually substantially increase at 32 ms, consistent with the 2D morphology in Fig. 1. The same is observed also for temperature and potential fluctuations. Fluctuation amplitudes alone also could not explain the changed direction of the transport. A deeper analysis of the microscopic turbulence characteristics is required.

Besides fluctuation amplitudes, turbulent transport is equally determined by correlations between fluctuating quantities. We focus on the heat convected by the particle flux, as the analysis of heat conduction due to temperature fluctuations yields the same results. The total radial particle flux can be written as [39,40]

$$\Gamma_r = \oint \oint n v_{E,r} r R d\phi d\theta_s = \frac{\langle R \rangle_{\theta_s}}{2\pi} \oint d\phi \sum_{k_y} \hat{\Gamma}(k_y), \quad (4)$$

$$\hat{\Gamma}(k_y) = \frac{r}{\langle B \rangle_{\theta_s}} k_y |\hat{n}(k_y)| |\hat{\phi}(k_y)| \sin[\alpha_{n,\phi}(k_y)]. \quad (5)$$

Here, $v_{E,r}$ is the radial $E \times B$ velocity, and the flux-surface integral goes over the toroidal angle ϕ and the poloidal ballooning angle θ_s . The binormal coordinate is $y = r\theta_s$. Transforming the density and electrostatic potential into binormal Fourier space, we can write the Fourier particle flux $\hat{\Gamma}$ in terms of fluctuation amplitudes and the phase factor $\sin[\alpha_{n,\phi}(k_y)]$, where the phase shift is $\alpha_{n,\phi}(k_y) = \text{Im} \log(\hat{n}^* \hat{\phi})$.

For exactly in- or out-of-phase fluctuations, $\alpha_{n,\phi} = 0$ or $\pm\pi$, the transport vanishes irrespectively of the fluctuation

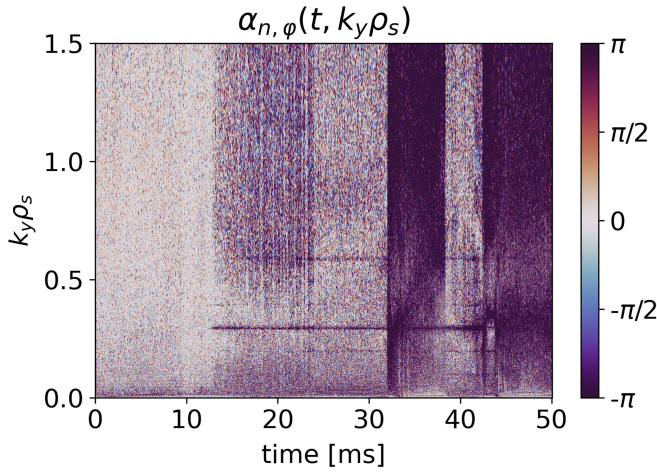


FIG. 5. Evolution of phase shifts between density and electrostatic potential fluctuations vs $k_y \rho_s$ at $\rho_{\text{pol}} = 0.99$.

amplitudes. Outward transport corresponds to $\alpha_{n,\phi} > 0$, required for microinstability growth, while $\alpha_{n,\phi} < 0$ leads to inward transport. The relevance of the phase shift analysis is that it characterizes very well the microinstabilities underlying the turbulence [39,41]. For instance, the adiabatic response in drift-wave (DW) turbulence leads to nearly in-phase fluctuations of density, temperature, and electrostatic potential (see End Matter Appendix C) [19]. On the other hand, the destabilization of (ideal) ballooning modes by magnetic curvature leads to phase shifts around $\pi/2$. In practice, many other instabilities exist with various flavors: the ion temperature gradient (ITG) mode, for example, can have ballooning character for T_i , but drift-wave character for n , but it is irrelevant here since we have $\eta_i = \partial_r \ln T_i / \partial_r \ln n < 1$ below the threshold.

Important for our analysis is the kinetic ballooning mode (KBM) [42,43]. Despite the name, the key difference from ideal MHD ballooning is captured by our two-fluid model [44]: the diamagnetic flux of ion temperature leads to an earlier but less explosive destabilization of the mode due to the magnetic resonance. More importantly, finite Larmor radius effects stabilize the KBM, leading to a phase shift of $\pm\pi$ [45]. The KBM has been predicted for H-mode pedestals by local gyrokinetic simulations [46]; it plays a key role in the EPED pedestal model [47] and has been observed experimentally [48].

In our previous, validated simulations, we have found drift-wave turbulence in L-mode [19,23] (consistent also with local gyrokinetic simulations [49]) and KBM turbulence in H-mode conditions [17]. Figure 5 shows the phase shift distribution in binormal wave number space, normalized to the sound Larmor radius, as a function of time at $\rho_{\text{pol}} = 0.99$. Indeed, early in the L-mode phase we find $\alpha_{n,\phi} \approx 0$. With increasing power and gradient buildup, $\alpha_{n,\phi}$ becomes more nonadiabatic to carry the increased heat flux. A KBM precursor appears in the spectrum at $t \approx 13$ ms, when T_i at $\rho_{\text{pol}} = 0.99$ saturates. However, a

large part of the spectrum remains DW-like. At $t = 24$ ms, we see that the improved Shafranov shift model [26] makes the spectrum somewhat more adiabatic, i.e., details of the implementation of magnetic fluctuations affect the turbulence. Still, it is clearly less adiabatic than early in the L-mode phase.

The most interesting transition happens at $t = 32$ ms: suddenly, the whole spectrum shifts to $\alpha_{n,\phi} = \pm\pi$. This means that we have a very abrupt DW-KBM transition. This microscopic transition in the phase shifts is the reason for the macroscopic abrupt change in the transport: before the KBM fully settles in, it causes inward fluxes and a pedestal buildup. The duration of 1–2 ms for this transient inward transport is consistent with how long even stationary near-equilibrium GRILLIX simulations would need to saturate, i.e., the global turbulence saturation time, set primarily by the SOL transit time of ~ 1 ms for AUG [17,50]. Note how, in the fluctuations in Fig. 4, the burst begins simultaneously at $\rho_{\text{pol}} \in (0.96, 0.99)$ and then shifts outward. This means that the same mode causes the increased radial transport at $\rho_{\text{pol}} = 0.96$ as well as the inward transport at $\rho_{\text{pol}} = 0.99$ (see Fig. 2). However, the steeper gradients then lead to outward KBM transport, relaxing the pedestal and abruptly transitioning back into the DW regime. The cycle then repeats. The reason for these I-phase-like [51–53] transport oscillations appears to be the sharp increase of KBM growth rates beyond the threshold [45], which relaxes the gradients again after the nonlinearly caused pedestal overshoot. Ultimately, this is expected to settle on the marginal stability threshold of the KBM.

Flow analysis— $E \times B$ shear is crucial for the L-H transition [15,54–56], and indeed it becomes quite large in the presented simulation. However, we find that it suppresses DW turbulence gradually, while the experimentally observed fast change in plasma profiles [2] is due to a subsequent transition to a different kind of turbulence, KBMs in our case. As shown in Fig. 2(g), $\partial_r p_i / en|_{\text{min}}$ plateaus at $t \approx 12$ ms, while $E_r|_{\text{min}}$ keeps increasing. Thus, in this phase, we see a significant increase of plasma rotation (since the radial force balance $E_r = \partial_r p_i / en - v_\theta B_\phi + v_\phi B_\theta$ is necessarily fulfilled [24,57,58]). On the other hand, at $t > 25$ ms, $E_r|_{\text{min}}$ seems to mostly follow the ion pressure gradient, as is also found experimentally in H-mode on ASDEX Upgrade [31,59]. There is a constant offset due to contributions from poloidal and toroidal flows, but it is smaller in comparison and remains similar across the fast transition at $t = 32$ ms. To analyze the turbulent drive of poloidal rotation, we show the Reynolds stress [24,60] in Fig. 2(h). To avoid separating the non-Boussinesq [61] and compressibility contributions, we show the whole divergence of the nonlinear polarization particle flux,

$$\nabla \cdot (\Gamma_{\text{pol}}) = \nabla \cdot \left[-\frac{nM_i}{eB^2} (\mathbf{v}_E \cdot \nabla) \left(\nabla_\perp \varphi + \frac{\nabla_\perp p_i}{en} \right) \right], \quad (6)$$

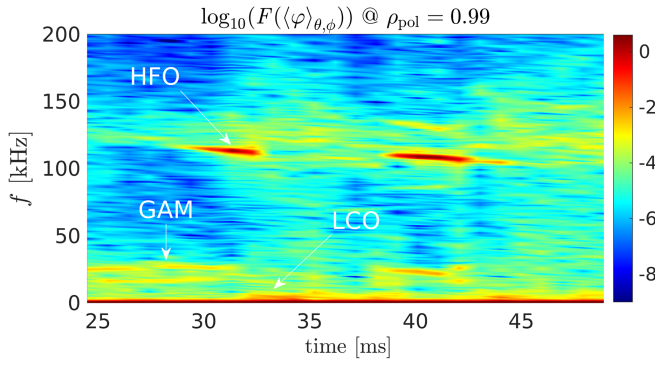


FIG. 6. Evolution of the frequency spectrum of the flux-surface averaged electrostatic potential, featuring geodesic acoustic mode (GAM), high frequency oscillations (HFO), and limit cycle oscillations (LCO).

flux-surface averaged and maximized over the radius, which drives vorticity via $\partial_t \Omega = -\nabla \cdot (\Gamma_{\text{pol}})$ [24]. While the first term in the bracket constitutes the usual $E \times B$ Reynolds stress, which is responsible for the inverse cascade [53,62], the second diamagnetic term stems from finite Larmor radius corrections and breaks vorticity conservation [41,63]. In the early L-mode phase, consistent with [63], the diamagnetic part is actually the largest. With increasing heating power, the $E \times B$ part increases, driving the rotation. The two become comparable in magnitude at $t \approx 15$ ms, but with opposite signs. Therefore, they largely cancel, and the sum plateaus at $t \approx 20$ ms. At 32 ms, we see a large jump in magnitude of the $E \times B$ part of the Reynolds stress, but it is balanced by the diamagnetic part and the sum remains similar. On the other hand, in presence of magnetic fluctuations, it is interesting to also analyze the Maxwell stress contained in $\langle \nabla \cdot j_{\parallel} \mathbf{b}_1 \rangle \sim -\mu_0 \partial_r^2 \langle b_1^r b_1^\theta \rangle$. This term also increases at 32 ms and becomes comparable to the total Reynolds stress. In conclusion, the turbulence transition occurs in two phases: (1) a gradually increasing Reynolds stress, $E \times B$ shear, and DW turbulence suppression until $t = 32$ ms and (2) a fast transition to KBM turbulence.

In addition to mean flow acceleration, strong oscillations of the zonally averaged electrostatic potential in Fig. 2(i) indicate oscillating flows, driven by the geodesic coupling [64] and Stringer spin-up [51]. Figure 6 shows the frequency spectrum for these oscillations: early before the transition, at 25–30 ms, we clearly see geodesic acoustic modes (GAMs) in the spectrum [65,66]. GAMs are often found in L-mode and tend to disappear in H-mode [17,66]. They provide an energy sink for drift-wave turbulence, taking energy out of the zonal flow [64,66]. However, at 29–32 ms, the large amplitude oscillation actually has a significantly higher frequency of ~ 120 kHz, in the (kinetic) Alfvénic range [52]. Alfvénic oscillations indeed couple to GAMs as a part of the dissipation cycle in drift-wave turbulence [64], and they can also be driven by ballooned turbulent transport [51,52]. The amplitude of these Alfvénic oscillations reaches a level comparable to

the stationary zonal flow, before they disappear and a transition occurs. The frequency matches with the experimentally observed high frequency oscillations (HFO), which occur as precursors in I-phase bursts [32,33,52]. On the other hand, after 32 ms, the HFO and GAMs disappear. Instead, a low frequency oscillation emerges, likely related to I-phase limit cycle oscillations (LCO) [31,32,52]. Thus, our simulation is perhaps more comparable with a dithering L-I rather than an L-H transition, but the two are closely related.

Conclusions—We have shown that rapid phase transitions occur in global full- f turbulence simulations of the plasma edge subject to a power ramp. Macroscopically, the phase transition is characterized by an abrupt steepening of the edge pressure gradient due to a transient quench of the radial transport. Microscopically, the underlying reason is the quench of nonlinearly sustained drift-wave turbulence by mean and oscillating flows and a literal phase transition to kinetic ballooning mode turbulence. This opens up the opportunity to investigate in further studies how well such simulations compare to experimental L-H or L-I transitions in detail. Meanwhile, many key characteristics are indeed observed, such as abrupt profile steepening, a rapid transition in fluctuation amplitudes, dying out of GAMs, and I-phase-like transport oscillations and bursts. Finally, we remark that a precise determination of the L-H power threshold in general may require full- f gyrokinetic simulations [67] in conjunction with MHD peeling-ballooning modes [29], which is complicated due to the ELMs destroying the field alignment, necessary for affordable turbulence simulations [26]. Meanwhile, our drift-fluid simulations are likely useful for investigating regime transitions in collisional ELM-free scenarios such as the X-point radiator [68,69] and the quasicontinuous exhaust regime [70].

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Data availability—The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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End Matter

Appendix A: Impact of the magnetic fluctuation filter—As briefly discussed in the main text, and detailed in [26], our model employs a high-pass filter for magnetic fluctuations. The filtered fluctuating magnetic potential which enters parallel operators is defined as $A_1 = A_{\parallel} - \langle A_s \rangle_{\phi}$, with the full dynamic potential $A_{\parallel}$ computed from Ohm's law, the toroidal average $\langle \rangle_{\phi}$, and $\partial_t A_s = (A_{\parallel} - A_s)/\lambda$. This dynamical model for the magnetic shift A_s corresponds to an exponentially weighted average of $A_{\parallel}$ in time, with the additional toroidal average having only a minor effect. As discussed in Ref. [26], this approach has several advantages over simpler methods [17,71–73] (including our previous, more naive method [17], which was using

only a toroidal and a discrete time average over 20 μ s intervals), which primarily utilize spatial filtering: it provides a smoother approximation of A_s and in particular better preserves the 3D mode structure of A_1 .

The new filtering technique was developed while the long and costly simulation presented in this letter was already running. As it becomes increasingly critical with increasing β , i.e., much more in H- than in L-mode conditions [17,26], we decided to simply switch to the improved method before the transition. Some artifacts from the switch can be seen in the results in the main text: Figs. 2(f) and 5 show that the improved filter leads to more adiabatic and coherent drift-wave turbulence. However, we verified and show below that this is simply due to the model

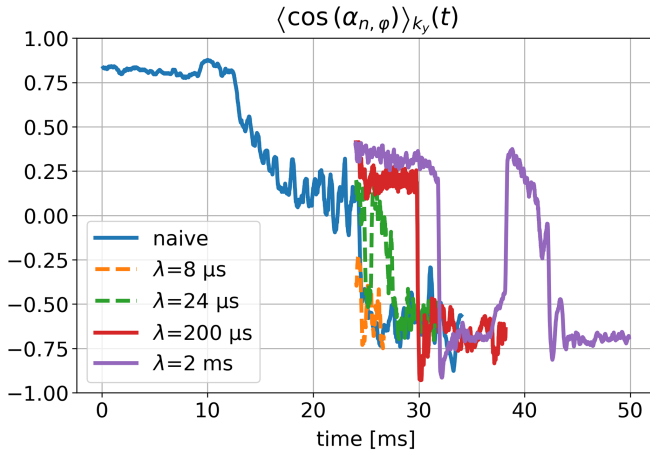


FIG. 7. Scan of the DW-KBM transition dynamics, characterized by the k_y -averaged $\cos(\alpha_{n,\varphi})$ of the phase shift shown in Fig. 5, as a function of the filter width λ in the magnetic fluctuation filter (smoothed over 0.2 ms intervals). See text for the definition of the filter.

improvement, which slightly decreases the transport and thus prolongs the L-mode phase, not the switch of the model *per se*.

The main result of this Letter lies in the fast DW-KBM turbulence transition, characterized by a jump of the phase-shifts from around zero to around $\pm\pi$ in Fig. 5. To turn this into a 1D observable, we take a k_y average. But first, we compute the cosine of $\alpha_{n,\varphi}$, which is 1 for ideal drift-waves and -1 for stable KBMs ($\langle \alpha_{n,\varphi} \rangle_{k_y}$ is the same for DW and KBM, as $\pm\pi$ averages to zero). In between, as the phase distribution becomes broader and more incoherent, the k_y -averaged $\cos(\alpha_{n,\varphi})$ approaches zero ($\alpha_{n,\varphi}$ does not become $\pi/2$). We calculate this quantity using our original, naive method (with which we ran to $t = 24$ ms) as well as the new Zhang filter with $\lambda \in \{8, 24, 200, 2000\} \mu\text{s}$ (the values correspond to $\{0.3, 1, 10, 100\} R_0/c_{s0}$, with $R_0 = 1.65$ m and $c_{s0} = \sqrt{100 \text{ eV}/2m_u}$), whereby our reference simulation used $\lambda = 2$ ms. The result is shown in Fig. 7.

The values of λ were chosen in accordance with the detailed analysis in Ref. [26]: they must be neither too small nor too large. With a too small λ , high-frequency turbulent magnetic fluctuations are removed together with the magnetic shift (as $\omega < 1/\lambda$ are filtered out from A_1), essentially inhibiting magnetic flutter and artificially destabilizing the turbulence. With a too large λ , the magnetic shift is not evolved quickly enough, and the background component B_s in the magnetic flutter may become too large, which may deteriorate the field-aligned numerics if $B_s > B_0 h_\perp/h_\parallel \approx 10^{-3} B_0$ is reached (where $h_\perp$ and $h_\parallel$ are the perpendicular and parallel grid resolutions, respectively, and B_0 the background magnetic field). The latter can be compensated by increasing parallel grid resolution, but this makes the simulations much more expensive and destroys the whole point of field alignment. Our analysis in [17,26]

suggests that, while for L-mode DW turbulence $\lambda \gtrsim 10 \mu\text{s}$ is sufficient, in typical ASDEX Upgrade conditions, for H-mode KBM turbulence a $\lambda > 200 \mu\text{s}$ should be used to preserve the mode structure. The maximum value of λ is determined by the maximum tolerable B_s , which depends on the actual time evolution of the plasma background. We have checked that a 10 times larger value than the minimum required for H-mode KBM turbulence, i.e., 2 ms, is still numerically sound. Hence, this is why the simulation with $\lambda = 2$ ms became the reference simulation.

Coming back to Fig. 7, first, we find that the DW-KBM transition occurs also with the naive filter at around 25 ms. The Zhang filter reproduces this transition with $\lambda = 8 \mu\text{s}$ —which is, however, too small. Moving to larger λ , we find somewhat increased adiabaticity and thus a prolonged DW L-mode phase. Nonetheless, it always collapses eventually into a KBM spectrum—proving the robustness of our results.

Therefore, we conclude that the change to an improved magnetic fluctuation filter only improves the validity of the simulation results. Note that, since the filter corresponds to an exponentially weighted average of $A_\parallel$, the previous solution is damped out as $e^{-t/\lambda}$ and thus should not affect a transition several λ later. While low λ converges to our previous, naive filter, a larger $\lambda = 2$ ms is more correct, and the Zhang filter allows for a variation during a simulation as it preserves C^0 smoothness of A_s . Nonetheless, as the precise choice of λ does affect particularly the timing of the transition, it is desirable to further improve the treatment of magnetic fluctuations and the consistency with the magnetic equilibrium in full- f turbulence simulations, as discussed in Ref. [26].

Appendix B: Model limitations—ELMs are only triggered once the pedestal becomes steep enough, and ELM cycles are slow, taking many ms and thus more than typical turbulence saturation times [29]. Therefore, we expect the transition dynamics to be independent of their inclusion, but they would then limit the pedestal to a lower height. The L-H transition is known to depend primarily on the ion heat flux [9]. Therefore, we believe that it is unlikely that TEM, MTM, or ETG modes are critical for the L-H transition as they primarily transport electron heat. Nonetheless, they could also affect the profiles. Furthermore, as the two heat channels are coupled, any additional transport mechanisms could decrease the edge confinement time, i.e., how fast profiles react to changes in heating power, and thus speed-up the L-H transition. We also note that L-mode turbulence was found to be of drift-wave type in AUG and JET also in gyrokinetic simulations [49], while in H-mode TEMs are expected to be suppressed by the flow shear [35]. However, besides individual linear instabilities, also important might be the intricate nonlinear interactions between them, including multiscale effects, as well as neoclassical transport and equilibrium evolution.

Appendix C: Phase relations in DW turbulence—A detailed linear and nonlinear analysis of drift-waves in our two-fluid system has been carried out in Refs. [19,39,41]. Here, we summarize the main results for the phase relations. First, the adiabatic response refers to the (electrostatic) force balance in Ohm's law $e\nabla_{\parallel}\phi = (\nabla_{\parallel}p_e)/n$. In the isothermal limit, it leads to $\alpha_{n,\phi} = 0$. Finite transport requires a finite phase shift. In the DW system, analyzed in slab geometry without curvature, the phase shift results from additional terms in Ohm's law which inhibit the adiabatic response. The full linearized Ohm's law reads as

$$\beta_0 \frac{\partial \tilde{A}_{\parallel}}{\partial t} + \mu \frac{\partial \tilde{j}_{\parallel}}{\partial t} = \frac{\partial}{\partial z} \left(\tilde{n} + (1 + \theta) \tilde{T}_e - \tilde{\phi} \right) - \beta_0 (2 + \theta) \omega_* \frac{\partial \tilde{A}_{\parallel}}{\partial y} - \eta_{\parallel 0} \tilde{j}_{\parallel}, \quad (\text{C1})$$

where the two terms on the left side are magnetic induction and electron inertia, respectively; the first bracket on the right side is the force balance, including the thermal force with $\theta = 0.71$; the second term is magnetic flutter acting on the background radial pressure gradient; and the last term is resistivity. Solving the full linear system [19], the phase shift is

$$\alpha_{n,\phi} = \underbrace{[2\gamma\beta_0\omega_R - \gamma\beta_0(2+\theta)\omega_*K]}_{\text{induction}} + \underbrace{[2\gamma\mu\omega_R K^2]}_{\text{inertia}} + \underbrace{[\omega_R\eta_{\parallel 0}K^2]}_{\text{resistivity}} / k_z^2, \quad (\text{C2})$$

where ω_R and γ are the real frequency and the growth rate, respectively, such that $\omega = \omega_R + i\gamma$, β_0 is the electron beta, ω_* is the diamagnetic frequency, $K = k_y \rho_s$, μ is the electron to ion mass ratio, $\eta_{\parallel 0}$ is the resistivity, and k_z is the parallel wave number. DWs can thus be destabilized by magnetic induction, electron inertia, and resistivity, while magnetic flutter is stabilizing.

Importantly, we note that, in general, DWs are driven by the full electron pressure gradient, including both n and T_e . Since T_e directly enters Ohm's law through the electron pressure gradient and the thermal force, it also participates in the adiabatic response. T_i , on the other hand, does not. For $\eta_i = \partial_r \ln T_i / \partial_r \ln n \gtrsim 1$, this leads to the separate ITG instability. However, in our case $\eta_i < 1$, T_i oscillates nearly in-phase with the density. Specifically, our linear analysis [19] shows that

$$\frac{\alpha_{T_e,\phi}}{\alpha_{n,\phi}} = 1 + \frac{2\theta - 1}{3} \epsilon + O(\epsilon^2), \quad (\text{C3})$$

$$\frac{\alpha_{T_i,\phi}}{\alpha_{n,\phi}} = 1 - \frac{1}{3} \epsilon + O(\epsilon^2), \quad (\text{C4})$$

with $\epsilon = K|\omega|^2 / (\omega_* \omega_R) < 1$. The phase shifts for the temperatures thus tend to be similar to those of the density in DW turbulence.

Linear analysis is also useful for KBMs, but it requires ballooning coordinates (not slab geometry) and warm ions. This will be reported in a separate publication.