

First-Principles Explanation of the Drift Configuration Dependence of the Radial Electric Field and High-Confinement Access in Tokamaks

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The origin of the difference in the high-confinement (H -mode) power threshold between favorable and unfavorable drift configurations in tokamaks—experimentally linked to a deeper radial electric field (E_r) well in the former—remains unresolved. Using first-principles gyrokinetic simulations of edge and scrape-off-layer turbulence in the ASDEX Upgrade tokamak, we show that turbulence-driven poloidal flows generate this deeper E_r well in the favorable configuration through enhanced nonlinear turbulence-mean flow energy transfer. This transfer is significantly weaker in the unfavorable case, yielding a shallower E_r well, while turbulence intensity is simultaneously higher. Within the turbulence-flow shear suppression paradigm, the combination of stronger shear and reduced turbulence facilitates H -mode access in the favorable configuration. These results provide the first validated, self-consistent full- f gyrokinetic explanation of how drift configuration controls the nonlinear dynamics of profiles, E_r , flows, and turbulence, thereby setting the H -mode power threshold.

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In tokamaks, the direction of the ion ∇B drift (relative to the active X point) defines favorable (fav) and unfavorable (unfav) configurations. This change of direction significantly influences the H -mode power threshold, with unfav discharges generally requiring more than twice the threshold power than fav ones [1–5]. This difference observed in multiple devices [5–9] correlates with a deeper radial electric field (E_r) well near the separatrix in fav before the transition from low-confinement (L mode) to H mode. A threshold in E_r at the L - H transition has been identified in fav at the ASDEX Upgrade (AUG) [10], while it remains unclear in unfav. Understanding the mechanism underlying this difference in E_r is crucial for ITER and future fusion reactors, as it determines the heating power required and limits operational flexibility.

Recent experiments at AUG comparing fav and unfav discharges at matched parameters [5] show that neoclassical (NC) effects alone cannot explain the observed E_r differences, highlighting a significant role for non-NC physics. Proposed mechanisms include ion orbit losses [11–13], scrape-off-layer (SOL) flows [14], and nonlinear turbulence-mean flow interactions driven by the Reynolds

stress (RS) [4,15–20]. In particular, turbulence-mean flow energy transfer is highly sensitive to poloidal asymmetries in the RS, which can be driven by X -point geometry, flow shear, and radial transport [19,21]. Reduced models incorporating such asymmetries capture qualitative experimental trends [20], but a predictive, self-consistent, and nonlinear description of the interplay between equilibrium profiles, E_r , flows, and turbulence remains to be established. Existing studies—based on linear, simplified, or quasilinear approaches in simplified geometries—are not able to capture correctly the nonlocal coupling and nonlinear feedback between turbulence, mean flows, and E_r in realistic diverted configurations. In addition, edge turbulence requires a gyrokinetic (GK) treatment. Developing such a self-consistent and GK description is fundamental to determining whether these coupled effects fully account for the differences between fav and unfav and to quantifying their respective contributions.

In this Letter, we present the first GK study of fav and unfav discharges at AUG [5] using validated first-principles simulations of edge-SOL turbulence. We show that stronger *nonlinear* generation of poloidal flows near the separatrix explains the deeper E_r well in fav. An evolution equation for the *flux-surface-averaged* (FSA) poloidal kinetic energy reveals that these poloidal flows are driven by radially localized asymmetries in the total velocity stress Π , a full- f generalization of RS. These asymmetries arise from the orientation of $\mathbf{E} \times \mathbf{B}$ flow relative to the X point on the low-field side (LFS), producing a radial gradient of Π that enhances turbulence-flow energy transfer in fav. Meanwhile, turbulence intensity is stronger in unfav.

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Overall, the higher H -mode power threshold in unfav might result from both weaker turbulence-mean flow energy transfer and stronger turbulence levels.

Experiments at AUG—Dedicated deuterium experiments at AUG [5] compared a pair of lower single-null discharges, no. 36983 (fav) and no. 37375 (unfav), to investigate the effect of the ion ∇B drift direction on E_r . Both discharges were performed under matched conditions: line-averaged electron density $n_e \simeq 2.7 \times 10^{19} \text{ m}^{-3}$ (lower density branch), and 600 kW of electron cyclotron resonance heating in L mode. The on axis toroidal magnetic field and plasma current were $|B_\phi| = 2.5 \text{ T}$ and $|I_p| = 0.8 \text{ MA}$. Measurements of E_r were obtained from Doppler backscattering (DBS), while the electron temperature (T_e) and n_e were reconstructed via integrated data analysis [22]. The main ion temperature (T_i) was assumed to be equal to that of the impurity species measured by charge-exchange recombination spectroscopy (CXRS). The deuterium toroidal ($U_{\phi i}$) and poloidal ($U_{\theta i}$) velocities were inferred from impurity velocity measurements using the NC code NEOART [23].

Turbulence gyrokinetic modeling—We model these two discharges in the L -mode phase using the spectral version of the GENE-X code [24,25], which solves the full- f , electromagnetic, long-wavelength, and collisional GK Vlasov equation with X points. The details of the numerical setup are reported in the End Matter. The magnetic equilibrium at $t = 2.7 \text{ s}$ is considered for the fav discharge (no. 36983) in the L -mode phase (prior to the L - H transition triggered at 2.78 s with a power threshold of 1.1 MW). To minimize differences between the two configurations in the simulations, the unfav equilibrium (no. 37375) is generated by reversing the direction of B_ϕ from no. 36983, changing the direction of the ion ∇B drift and the magnetic helicity. A density source is introduced near the separatrix to mimic neutral ionization. To reach a quasisteady state, both simulations were run until $t \sim 2.5 \text{ ms}$. In the quasisteady state, the total power crossing the separatrix is approximately 0.5 MW in both cases. A time window of 0.1 ms is employed to perform the statistics.

Validation of simulation results—The left column of Fig. 1 compares the outboard midplane (OMP) profiles (time and toroidally averaged) of n_e and T_i obtained from GENE-X with experimental measurements [5]. Overall, good agreement with experiments is observed in both configurations, albeit with slightly lower T_i gradients near the separatrix. Most importantly, the simulations reproduce the experimental finding that the profiles are nearly insensitive to the ion ∇B drift direction [5], although n_e is slightly higher in fav near $\rho_{\text{pol}} \sim 0.99$. Similar observations (but not shown) can be made for T_e , with $T_i \sim 1.3T_e$ at the separatrix.

The right column of Fig. 1 displays the OMP and FSA profiles of the deuterium toroidal (self-generated) and poloidal velocities. In GENE-X, the total (NC and

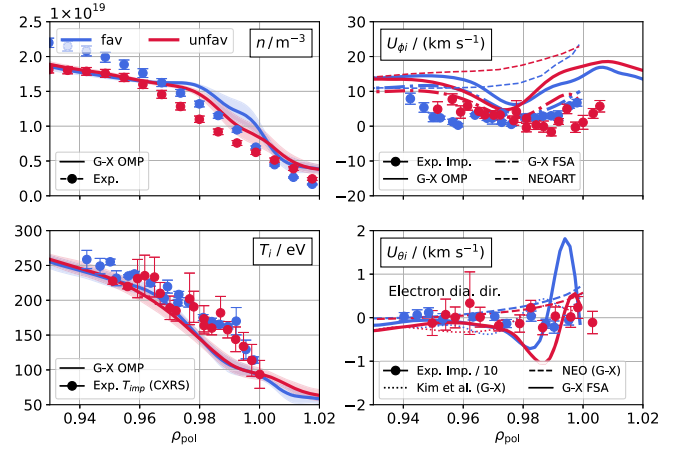


FIG. 1. Left: OMP n_e (top) and T_i (bottom) from GENE-X (G-X, solid) in fav (blue) and unfav (red). Shaded areas denote standard deviations (time and toroidal averages). Experimental data [5] are shown as circles. Right: OMP and FSA $U_{\phi i}$ (top) and $U_{\theta i}$ (bottom). CXRS impurity ($U_{\phi i}$ and $U_{\theta i}$) and NEOART ($U_{\phi i}$) main ion velocities are included [5]. NC predictions for $U_{\theta i}$ from Kim *et al.* [26] and NEO [27] are shown (dotted, dashed) using GENE-X profiles. The sign of unfav $U_{\theta i}$ is reversed for comparison.

turbulence-driven) ion velocity (U_i) is calculated from the zeroth-order velocity moment of the guiding-center drift corrected by a magnetization term [28,29]: $U_i = \int dW \mathbf{R} f_i / n_i + c \nabla \times \mathbf{m}_i / (q_i n_i)$, with $\mathbf{R}$ the guiding-center equation of motion, and $\mathbf{m}_i = -P_{\perp i} \mathbf{b} / B$ the classical magnetization. Then, $U_{\phi i} = U_i \cdot \hat{\mathbf{e}}_\phi \simeq U_i \cdot \mathbf{b}$ (with $\mathbf{b} = \mathbf{B} / B$) and $U_{\theta i} = U_i \cdot \hat{\mathbf{e}}_\theta$. As with n_e , T_i , T_e (not shown) $U_{\phi i}$ shows little dependence on the ∇B drift direction and is close to experiments [5]. Remarkably, the poloidal velocity $U_{\theta i}$ exhibits a clear difference: its maximal magnitude is much larger in fav than in unfav inside the E_r well located near $\rho_{\text{pol}} \sim 0.995$. In both cases, $U_{\theta i}$ is directed in the ion diamagnetic direction (in agreement with NC predictions) inside the edge ($\rho_{\text{pol}} \lesssim 0.98$) and reverses toward the electron diamagnetic direction near the separatrix.

To assess the NC level of $U_{\theta i}$, we compare the GENE-X results with the NC prediction of Kim [26] and NEO [27] calculations using the GENE-X profiles as inputs. Figure 1 demonstrates that $U_{\theta i}$ is consistent with NC close to the separatrix only in unfav. On the other hand, the amplitude of $U_{\theta i}$ is significantly larger in the vicinity of the separatrix in fav, indicating the presence of additional, non-NC mechanisms enhancing $U_{\theta i}$ in fav.

Radial electric field—In GENE-X, E_r is computed directly from the negative radial gradient of the electrostatic potential, ϕ_1 , obtained from the GK quasineutrality (QN) equation [30]. Figure 2 compares the simulated OMP E_r profiles with experimental DBS measurements [5] and NC NEO results (using GENE-X profiles). The simulations reproduce the overall experimental E_r behavior in both

the edge and SOL, with good agreement with experiments (within statistical and experimental uncertainties [5]). Enhanced geodesic acoustic mode (GAM) [31] activity is also found around the E_r well in unfav, being much smaller in fav (see Fig. 7). Consistent with experiments, a pronounced difference arises near the separatrix: the E_r well in fav is significantly deeper than in unfav. This deviation occurs despite nearly identical n_e , T_i , and $U_{\phi i}$ profiles in both cases (Fig. 1).

The OMP profile and velocity (Fig. 1) structures influence E_r through the radial force balance constraint. Thus, to identify the origin of the difference between E_r observed in Fig. 2, we consider the ion radial force balance, which holds in the full- f GK formalism [28,30]. Assuming isotropic ion pressure, the radial force balance reads as follows [29]: $E_r = \nabla_r P_i / (q_i n_i) + U_{i\phi} B_\theta - U_{i\theta} B_\phi$, with $E_r = -\nabla \phi_1 \cdot \hat{e}_\psi$. As shown in Fig. 1, the diamagnetic ($\nabla_r P_i / [q_i n_i]$) and toroidal rotation ($U_{i\phi} B_\theta$) terms are nearly identical in both configurations and dominate E_r for $\rho_{\text{pol}} \lesssim 0.98$, implying a small contribution from the poloidal flow ($-U_{i\theta} B_\phi$). However, within the E_r well ($\rho_{\text{pol}} \sim 0.995$), they nearly cancel out, leaving $E_r \approx -U_{i\theta} B_\phi$. Hence, the depth and the minimum value of the E_r well ($E_{r,\text{min}}$) are set by the amplitude and shape of the poloidal ion flow $U_{i\theta}$, which is much larger in amplitude in fav. With the same magnetic field strength,

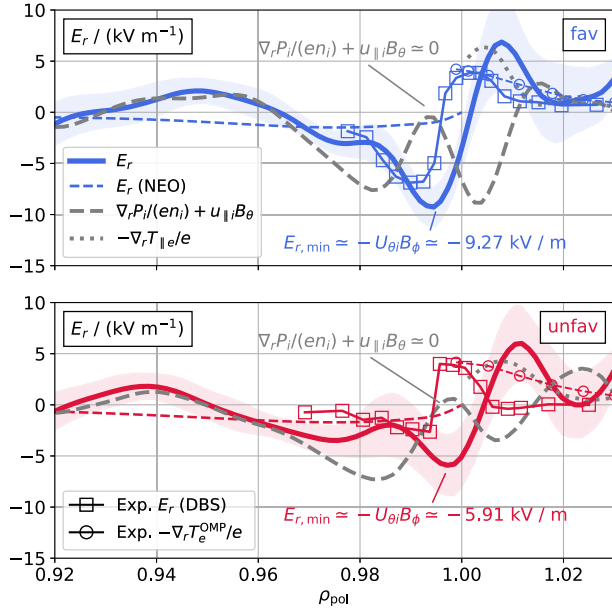


FIG. 2. OMP E_r from GENE-X (solid) in fav (top) and unfav (bottom). Shading indicates standard deviation. DBS measurements [5] are shown as squares with radial uncertainty $\Delta \rho_{\text{pol}} \approx 0.014$. NC E_r from NEO (using GENE-X profiles) is shown (dashed). In both cases $E_{r,\text{min}} \approx -U_{i\theta} B_\phi$. SOL estimates from experiment ($-\nabla T_e^{\text{OMP}}/e$, circles) and GK theory ($-\nabla T_{\parallel e}/e$ [30], dotted) are also shown.

the resulting inner $\mathbf{E} \times \mathbf{B}$ shear is approximately 1.6 times stronger in fav than in unfav [20].

The SOL E_r agree well with experiments [5] and conduction-limited GK [30], with only minor differences between fav and unfav. As a consequence, the outer $\mathbf{E} \times \mathbf{B}$ shear layer is 1.6 times larger in fav and exceeds the inner shear in amplitude in both fav and unfav. Altogether, Figs. 1 and 2 therefore indicate that the observed difference in the depth of the E_r well (and its shear) cannot originate from NC effects, but instead arises from nonlinear interactions between turbulence and poloidal flow, as we demonstrate below.

Edge transport and turbulence—We characterize here the edge transport and turbulence obtained from GENE-X. The top panel of Fig. 3 compares the simulated poloidal distribution of the (toroidally and time-averaged turbulent) radial ion particle flux, $\Gamma_i = n_i u_\psi$ (with $u_\psi = \mathbf{u} \cdot \hat{e}_\psi$ the radial component of the $\mathbf{E} \times \mathbf{B}$ velocity, $\mathbf{u} = c\mathbf{E} \times \mathbf{B}/B^2$) on the $\rho_{\text{pol}} = 0.995$ flux surface. In both configurations, Γ_i exhibits fine structures predominantly around the OMP position on the LFS and decays on the high-field side. While Γ_i features only moderate poloidal asymmetry, its magnitude is larger between the OMP and X-point position on the LFS in fav. Γ_i is also directed inward (outward) near the X point on the LFS in unfav (fav), reflecting the contribution from J^B (see the End Matter). This poloidal profile of Γ_i influences the poloidal distribution of the velocity stress Π (proportional to Γ_i) and hence its FSA and gradient, as we show below.

The bottom plots of Fig. 3 show the Fourier power spectra of ϕ_1 and n_e at $\rho_{\text{pol}} = 0.995$. The spectra in unfav exhibit higher amplitudes than in fav. The temperature spectra show analogous trends [32]. Note that the phase

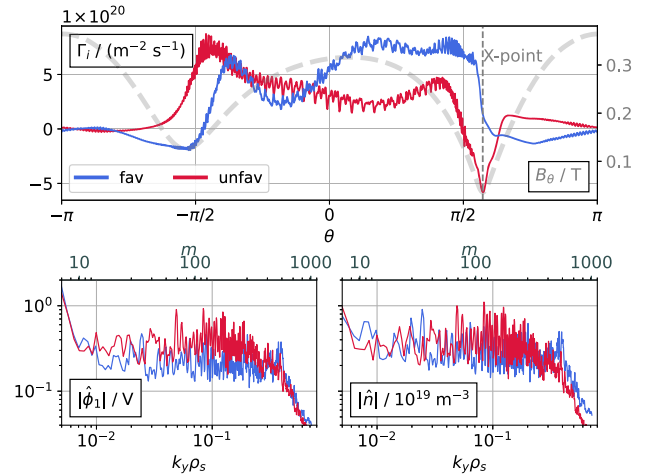


FIG. 3. Top: poloidal distribution of Γ_i (time and toroidal averages) obtained from GENE-X vs poloidal geometrical angle θ at $\rho_{\text{pol}} = 0.995$. Bottom: Fourier spectra of ϕ_1 (left) and n_e (right) vs $k_y \rho_s$ (ρ_s is the averaged ion sound Larmor radius) and poloidal mode number m .

shift characteristics remain qualitatively similar. Linear signatures of both electron drift waves and trapped electron modes are found in the edge [33], indicating comparable frequency-wave number relations in fav and unfav. Particle and heat transport crossing the separatrix are also of similar amplitudes. The present results demonstrate that the ion ∇B drift direction primarily influences the turbulence intensity being larger in unfav.

Nonlinear poloidal flow acceleration—From Fig. 2, we observe that $U_{i\theta}$ is dominated by the $\mathbf{E} \times \mathbf{B}$ velocity around the E_r well, such that $U_{i\theta} \simeq u_\theta$ ($u_\theta = \mathbf{u} \cdot \hat{\mathbf{e}}_\theta$). Thus, an equation for the *acceleration* of the FSA (zonal) poloidal flow momentum density, Mu_θ (with $M = \sum_\alpha n_\alpha m_\alpha$ the total mass density), can be obtained from the GK vorticity equation, Eq. (C1), and is given in Eq. (1):

$$\frac{\partial}{\partial t} (\langle Mu_\theta \rangle) = -\frac{d}{dr} \Pi + \Lambda. \quad (1)$$

From Eq. (1), the poloidal flow acceleration is driven by the negative radial gradient of the total (electrostatic) velocity stress, $\Pi = \langle M \rangle \langle u_\theta u_\psi \rangle$, with $\langle \langle f \rangle \rangle = \langle Mf \rangle / \langle M \rangle$ the Favre average operator [34]. The roles of diamagnetic currents and GAMs are discussed in the End Matter. We focus here on the drive due to Π . From Eq. (1), Π drives poloidal flow acceleration (damping) if its gradient is negative (positive): Fig. 5 shows a clear larger negative radial gradient in fav around the E_r well. Note that Π is related to Γ_i through $\Pi \simeq \langle m_i \Gamma_i u_\theta \rangle$. In particular, it constitutes the full- f generalization of the RS $\mathcal{R}$ [35]: Π can drive nonlinear poloidal flow acceleration, not only through $\partial_r \mathcal{R}$, but also through radial turbulent transport of poloidal momentum, poloidal density asymmetries, and triple fluctuation correlations. These effects, absent in previous formulations [18,20,21], remain finite for radially uniform $\mathcal{R}$.

For Π to provide a finite contribution in Eq. (1), radially dependent poloidal asymmetries in Π must develop [20,21]. These asymmetries can arise from the nonlinear interplay between magnetic-shear-induced (MSI) eddy tilting, X-point (or limiter) location, flow shear, and radial turbulent transport [17,18,21]. While the magnetic geometry fixes the MSI stress and X-point positions (identical in fav and unfav), the presence of a flow shear (inner or outer) layer can either nonlinearly amplify (fav) or reduce (unfav) the intrinsic MSI eddy tilting depending on the direction (sign of B_ϕ) of the $\mathbf{E} \times \mathbf{B}$ flow [17]. The top panels of Fig. 4 show snapshots of ϕ_1 fluctuations near the X point. In unfav, the $\mathbf{E} \times \mathbf{B}$ flow just inside the separatrix (in the outer shear) on the LFS is directed toward the X point and remains weak. Thus, turbulent structures are radially stretched mainly due to the MSI [21]. In contrast, in fav, the $\mathbf{E} \times \mathbf{B}$ flow is stronger near the separatrix and directed away from the X point, shearing the turbulent structures and in the opposite direction (toward the OMP) in the outer shear layer, corresponding to a situation where the flow

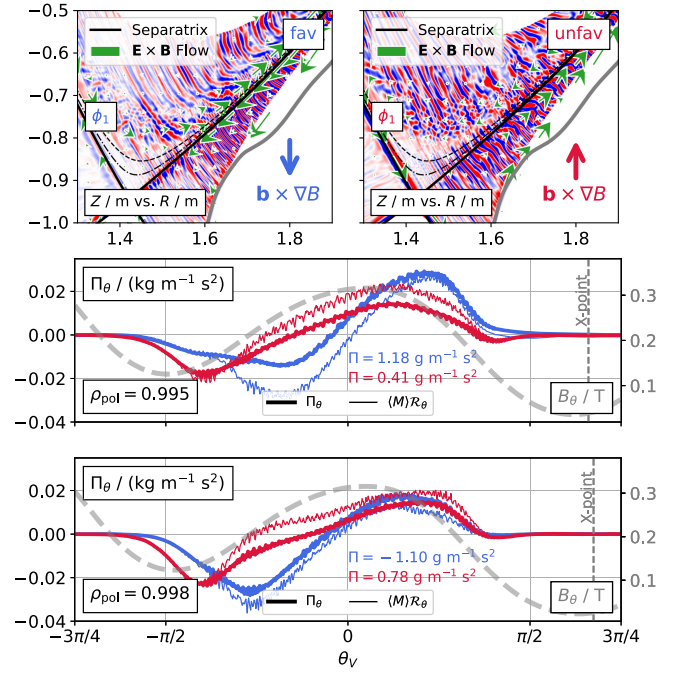


FIG. 4. Top: snapshots of ϕ_1 and mean $\mathbf{E} \times \mathbf{B}$ flow (arrows) near the X point in fav (left) and unfav (right). Center and bottom: poloidal distributions of Π_θ (thick) and $\mathcal{R}_\theta$ (thin) at $\rho_{pol} = 0.995$ and 0.998 vs θ_v [20]. The X point is indicated by the vertical dashed line.

shear ($\alpha_u^{(u)}$) and MSI tilt ($\alpha_u^{(s)}$) angles are in the same direction, i.e., $\alpha_u^{(u)} \alpha_u^{(s)} > 0$ in fav [17,18,20]. This localized $\mathbf{E} \times \mathbf{B}$ flow shear in the outer shear radially modulates the tilt of the turbulent eddies and the poloidal distribution of Π , thereby introducing a finite Π gradient and yielding self-amplification.

The middle and bottom panels of Fig. 4 show the time-averaged poloidal distribution of Π_θ ($\Pi = \int d\theta_v \Pi_\theta$) as a function of θ_v on two flux surfaces: inside the E_r well and closer to the separatrix (outer shear layer). We also display the corresponding poloidal distribution $\mathcal{R}_\theta$ of the RS $\mathcal{R}$ to assess the importance of the full- f approach. In unfav, Π_θ exhibits weak radial variation and is poloidally asymmetric around the OMP ($\theta_v = 0$), such that its FSA remains constant on both flux surfaces (small radial gradient). Π_θ closely follows $\mathcal{R}_\theta$, indicating that full- f corrections are small in this case. Note that the poloidal asymmetries in $\mathcal{R}_\theta$ arise from the competition between flow shear and MSI eddy tilting [20]. In contrast, in fav, Π_θ also displays a poloidal asymmetry around the OMP, but its amplitude varies radially: inside the E_r well, Π_θ becomes negative with reduced amplitude above the OMP and positive with larger amplitude below, leading to a net positive value of Π . This asymmetry is strongly correlated with the poloidal distribution of Γ_i (see Fig. 3). Near the separatrix, Π_θ remains asymmetric, but its amplitude becomes comparable on both sides of the OMP, yielding $\Pi < 0$. The radial

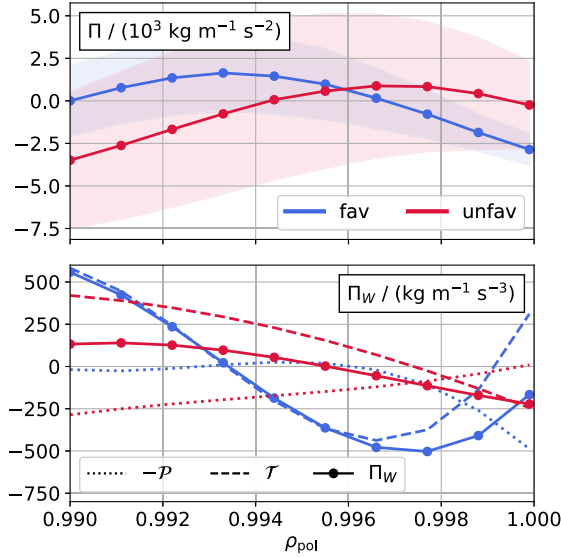


FIG. 5. Top: time-averaged radial profiles (with standard deviations) of Π . Bottom: time-averaged radial profiles of the stress work Π_W (solid), transport $\mathcal{T}$ (dashed), and negative production $-\mathcal{P}$ (dotted) contributions in fav (blue) and unfav (red).

variation of Π induces a finite negative radial gradient of about 7 larger in fav than in unfav. Noticeable deviations between Π_θ and $\mathcal{R}_\theta$ are observed in fav, particularly inside the E_r well, indicating non-negligible full- f contributions in Π (triple fluctuation correlation). The radial dependence of Π_θ in fav modifies the radial profile (and gradient) of Π shown in Fig. 5, thereby influencing the efficiency of the turbulence-mean flow interaction.

Nonlinear energy transfer—To assess the efficiency of energy transfer from turbulence to the mean poloidal flow, we consider the FSA poloidal kinetic energy, K_θ , whose evolution equation is given by Eq. (F1): K_θ is driven by the negative of the work done by the stress Π , Π_W . Π_W includes the contributions [see Eq. (F2)] from turbulent transport of mean kinetic energy ($\mathcal{T}$), and the energy production by turbulence $\mathcal{P}$ (see, such that $\Pi_W = \mathcal{T} - \mathcal{P}$ [36]. Figure 5 shows the profiles of Π_W , and of $\mathcal{T}$ and $\mathcal{P}$ obtained from the simulations. Π_W exhibits a substantially larger and negative amplitude around the E_r well in fav (see Fig. 2). On the other hand, Π_W decreases almost linearly toward the separatrix unfav and remains much smaller than in fav. These differences in Π_W indicate a stronger nonlinear energy transfer from turbulence into K_θ in fav. Figure 5 shows $\mathcal{T}$ dominates within the E_r well. In contrast, $-\mathcal{P}$ is dominant in the outer shear layer near the separatrix [21]. In unfav, $\mathcal{T}$ and $\mathcal{P}$ have opposite signs and comparable amplitudes, but remain small in magnitude. Figure 5 demonstrates that the energy transfer from turbulence to mean flow (mediated by Π_W) provides the poloidal flow acceleration near the separatrix in fav.

Conclusions—First-principles GK simulations of edge and SOL turbulence in realistic diverted geometry

reproduce the experimentally observed difference in E_r well depth between the fav and unfav configurations at AUG [5]. The simulations show that this difference arises from nonlinear turbulence-driven poloidal flows, which set both the depth and shear of the E_r well. We demonstrate that the efficiency of this nonlinear energy transfer is enhanced (fav) or reduced (unfav) depending on the direction of localized $\mathbf{E} \times \mathbf{B}$ shear near the LFS separatrix. From the evolution equation of the mean poloidal kinetic energy, the strength of the turbulence-mean flow energy transfer is identified to be driven by the radial gradient of the full- f velocity stress Π . Finally, turbulence activity is found to be larger in unfav. These results provide the first self-consistent GK explanation of the drift configuration dependence of the E_r well depth, identifying turbulence-mean flow energy transfer (mediated by the gradient of Π) as the key nonlinear mechanism yielding the deeper E_r well in fav. With the stronger turbulence in unfav, this nonlinear mechanism driven by the $\mathbf{E} \times \mathbf{B}$ drift direction between the OMP and X point on the LFS may underlie the observed configuration dependence of the H -mode (and I mode) power threshold [16,37]. Finally, flux-driven (fixed input power) simulations are required to further test this mechanism against experimental power thresholds. Finally, we note that similar trends are also observed in simulations of other devices, such as TCV [38], indicating the universality of the present findings.

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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

Magnetic field orientation—The magnetic field $\mathbf{B}$ can be written as $\mathbf{B} = B_\phi \hat{\mathbf{e}}_\phi + B_\theta \hat{\mathbf{e}}_\theta$, with $\hat{\mathbf{e}}_\phi$ (B_ϕ) and $\hat{\mathbf{e}}_\theta$ (B_θ) the toroidal and poloidal unit vectors (components of $\mathbf{B}$). We choose $\hat{\mathbf{e}}_\phi$ to be such that $B_\phi > 0$ in both configurations. We invert *only* the direction of $\hat{\mathbf{e}}_\phi$ between fav and unfav, while the poloidal magnetic field has the same direction. As a consequence, $\hat{\mathbf{e}}_\theta$ points away from the active X point on the LFS in *both configurations*. This implies that the direction of the $\mathbf{E} \times \mathbf{B}$ inside the E_r well (see Fig. 4) is positive (negative) in fav (unfav). The unit radial vector is defined as $\hat{\mathbf{e}}_\psi = \nabla\psi/|\nabla\psi|$ (ψ is the poloidal flux function), such that $\hat{\mathbf{e}}_\theta = \sigma_B \hat{\mathbf{e}}_\phi \times \hat{\mathbf{e}}_\psi$ (with $\sigma_B = 1$ in fav and -1 in unfav). We note that, to facilitate comparison between the two configurations, the signs of the velocities in unfav are adjusted to match that of fav.

Simulation setup—We summarize here the numerical setup employed in GENE-X. While a comprehensive description of the numerical methods implemented in GENE-X can be found in [24,25], only the most relevant aspects are highlighted here. At the boundaries of the numerical domain (extending from $\rho_{\text{pol}} = 0.9$ to 1.03), the distribution functions are assumed to be Maxwellian [25], with density and temperatures fixed at values close to experiments. The electromagnetic fields are set to zero. To avoid artificial flow shear near the edge

boundary, Neumann boundary conditions on flux variables are applied. For the numerical resolution, we use six parallel (Hermite) and four perpendicular (Laguerre) spectral coefficients in velocity space [30], 32 poloidal planes, and a constant grid spacing of ~ 1.4 mm (about 1.5 ion thermal gyroradius ρ_i near the separatrix) within the poloidal plane.

To mimic the effects of neutral ionization in the edge, we introduce a poloidally homogeneous, axisymmetric, radially localized Gaussian energy-conserving density source centered at $\rho_{\text{pol}} = 0.99$, with a constant amplitude of $1.5 \times 10^{22} \text{ s}^{-1}$ for both fav and unfav. In the present setup, particle and energy sources and sinks arise from a combination of Dirichlet boundary conditions at the edge and SOL boundaries—where the profiles are fixed to values close to experiments—and perpendicular hyperdiffusion applied within a narrow buffer region between ghost points and the numerical domain. Upon discretization, this approach is equivalent to a Krook-like operator that relaxes the distribution function toward a Maxwellian with fixed density and temperature near the boundary. Acting as a particle and energy reservoir, this numerical strategy injects or removes particles and energy to enforce the boundary conditions. As a consequence, the present simulations can be considered effectively flux driven, but with adaptive sources, while the profiles between the edge and SOL boundaries are free to evolve.

GK vorticity equation—We provide a streamline derivation, the GK vorticity equation, which we use to obtain the FSA poloidal flow acceleration, given in Eq. (1). Using the QN equation, $\omega = \sum_{\alpha} q_{\alpha} n_{\alpha}$, with the electrostatic vorticity variable $\omega = \nabla \cdot \mathbf{\Omega}$ ($\mathbf{\Omega} = c^2 \mathbf{ME}/B^2$, $\mathbf{E} = -\nabla \phi_1$) and applying the FSA operator, $\langle f \rangle = (V')^{-1} \oint dS f / |\nabla \psi| := \int d\theta_V \langle f \rangle_{\phi} / (2\pi)$ [with $V(\psi)$ the volume enclosed by the ψ flux surface with $V' = dV/d\psi$, dS the surface element of the flux surface ψ , θ_V the flux-surface average angle defined in [20], and $\langle \cdot \rangle_{\phi}$ the toroidal average], we derive, by taking the zeroth-order moment of the GK Vlasov equation,

$$\frac{\partial}{\partial t} \langle \omega \rangle + \frac{d}{dr} \langle J_{\perp} \rangle = 0, \quad (\text{C1})$$

with d/dr the radial derivative along $\hat{\mathbf{e}}_{\psi}$. Here, $J_{\perp} = \mathbf{J}_{\perp} \cdot \hat{\mathbf{e}}_{\psi}$ is the radial component of the *total* and perpendicular *gyrocenter* current density, $\mathbf{J}_{\perp} = \sum_{\alpha} \mathbf{J}_{\perp \alpha} = \sum_{\alpha} q_{\alpha} \int dW \mathbf{R}_{\perp} f_{\alpha}$. Here, $\mathbf{J}_{\perp \alpha}$ can be separated into different contributions: $\mathbf{J}_{\perp \alpha} = \mathbf{J}_{\alpha}^{\text{es}} + \mathbf{J}_{\alpha}^{\text{em}} + \mathbf{J}_{\alpha}^{\text{pol}}$, with $\mathbf{J}_{\alpha}^{\text{es}} = q_{\alpha} n_{\alpha} \mathbf{u}$ (electrostatic), $\mathbf{J}_{\alpha}^{\text{em}} = J_{\parallel \alpha} \mathbf{B}_1/B$ (flutter with $J_{\parallel \alpha} = q_{\alpha} n_{\alpha} U_{\parallel \alpha}$ the parallel current), $\mathbf{J}_{\alpha}^{\text{B}} = c P_{\parallel \alpha} \nabla \times \mathbf{b}/B + c P_{\perp \alpha} \mathbf{b} \times \nabla B/B^2$ (diamagnetic), $\mathbf{J}_{\alpha}^{\text{pol}} = -cm_{\alpha} n_{\alpha} \mathbf{b} \times \nabla u^2/(2B)$ (polarization). Hereafter, we neglect $\mathbf{J}^{\text{em}}$, which is associated with the Maxwell stress, due to the low values of electron plasma beta in the present cases. Note that in Eq. (C1), the contribution from the density source vanishes since no vorticity is injected.

Radial current contributions—We briefly discuss the contributions from the radial electrostatic ($J^{\text{es}} = \mathbf{J}^{\text{es}} \cdot \hat{\mathbf{e}}_{\psi}$) and diamagnetic ($J^{\text{B}} = \mathbf{J}^{\text{B}} \cdot \hat{\mathbf{e}}_{\psi}$) currents. These currents can either drive or damp the mean vorticity ω (mean poloidal flow) depending on their sign (relative to the poloidal flow velocity). J^{es} typically acts as a driving term through nonlinear interactions with turbulence, whereas J^{B} generally damps poloidal flows. In particular, in the absence of turbulence, J^{B} damps poloidal flows via parallel ion viscosity (NC damping [23,26]). In Fig. 6, the FSA and time-averaged radial profiles of J^{B} and J^{es} are shown in fav and unfav. Here, the signs are defined relative to u_{θ} such that positive (negative) values indicate drive (damping). The contribution from J^{B} is negative in both cases, but its amplitude is larger in unfav around the E_r well, indicating stronger damping. In contrast, the electrostatic radial current J^{es} is positive and stronger in the fav case near the E_r well, providing a stronger drive in fav. The combination of J^{B} and J^{es} therefore suggests a stronger net drive of the poloidal flow in fav. However, we note that the present analysis is qualitative since a rigorous balance analysis of Eq. (C1) is hindered by the separation of characteristic timescales, with J^{B} evolving

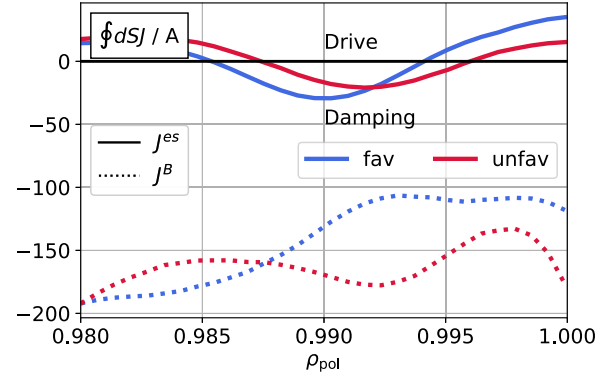


FIG. 6. Surface-integrated and time-averaged radial profiles of total $J^{\text{es}} = \mathbf{J}^{\text{es}} \cdot \hat{\mathbf{e}}_{\psi}$ (solid) and $J^{\text{B}} = \mathbf{J}^{\text{B}} \cdot \hat{\mathbf{e}}_{\psi}$ (dotted). Positive (negative) values (relative to the direction of u_{θ}) indicate drive (damping).

on longer collisional timescales and J^{es} on shorter turbulence timescales. Consequently, we focus here on the role of J^{es} in driving mean poloidal flows.

GAM oscillations—The coupling between mean poloidal flows and diamagnetic currents can transfer energy to pressure sidebands via the excitation of GAMs [39–41]. These pressure sidebands can in turn couple to parallel currents and dissipate energy through electron resistivity or can be nonlinearly depleted [39], thereby limiting the growth of mean (zonal) poloidal flows. In the present simulations, the GAM activity, which is inferred from the GAM amplitude A_{GAM} , is found to be stronger in unfav around the E_r well and reaches a minimum at the location of $E_{r,\text{min}}$, as shown in Fig. 7. In contrast, A_{GAM} is much smaller in fav. The difference in the GAM amplitude between fav and unfav, also evidenced in the standard deviations in Fig. 5, arises from nonlinear interactions with turbulence via the velocity stress Π . This is confirmed by a detailed Fourier analysis of Π (not shown here) revealing a clear peak at the linear GAM frequency in unfav. This

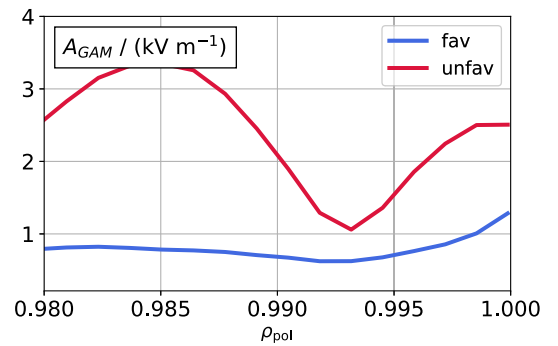


FIG. 7. Radial profiles of the GAM amplitude (A_{GAM}) measured around the E_r well in the frequency (f) band $10 \text{ kHz} \lesssim f \lesssim 30 \text{ kHz}$.

enhanced GAM activity suggests a larger energy transfer through geodesic coupling in unfav, which may further inhibit the buildup of mean poloidal flows and contribute to a shallower E_r well in unfav. A more quantitative analysis of GAM drive and interaction with turbulence [41] is deferred to future work.

Mean poloidal flow acceleration equation—From the vorticity equation in Eq. (C1), an evolution equation for the mean (FSA) poloidal flow acceleration driven by turbulence can be derived. Using the QN condition in the electrostatic current, such that $J^{\text{es}} = u_\psi \nabla \cdot \mathbf{\Omega}$, and integrating over ψ , the *acceleration* of the FSA poloidal flow momentum density (Mu_θ) is obtained and is given by Eq. (1). Note that the damping contribution of J^B in Eq. (1) and the constant of integration are absorbed in Λ . The gradient of the velocity stress Π appears due to the combination of the radial components of J^{es} and J^{pol} assuming that $\nabla \cdot \mathbf{u} \simeq 0$ and neglecting the variation of B . Finite Larmor radius corrections to Π [42,43] are absent here, but will be subject to future work. Note that

the total stress Π and the RS stress $\mathcal{R} = \langle \widetilde{u_\psi} \widetilde{u_\theta} \rangle$ ($\tilde{f} = f - \langle f \rangle$), are related by $\Pi = \langle M \rangle \mathcal{R} + \langle \tilde{M} \widetilde{u_\theta} \widetilde{u_\psi} \rangle + \langle M \rangle (\langle u_\theta \rangle \langle u_\psi \rangle - \langle \widetilde{u_\theta} \rangle \langle \widetilde{u_\psi} \rangle)$ [35].

From Eq. (1), the evolution equation for the FSA poloidal kinetic energy, $K_\theta = \langle Mu_\theta \rangle^2 / 2$, is

$$\frac{1}{\langle M \rangle} \frac{\partial K_\theta}{\partial t} = -\Pi_W + \langle u_\theta \rangle \langle \Lambda \rangle. \quad (\text{F1})$$

Here, $\Pi_W = \langle u_\theta \rangle d_r \Pi$ denotes the stress work associated with Π and can be written as [36]

$$\begin{aligned} \Pi_W &= \frac{d}{dr} (\langle u_\theta \rangle \Pi) - \Pi \frac{d}{dr} \langle u_\theta \rangle \\ &:= \mathcal{T} - \mathcal{P}, \end{aligned} \quad (\text{F2})$$

with $\mathcal{T}$ representing the turbulent transport (flux) of mean kinetic energy and $\mathcal{P}$ the energy production by turbulence [36]. Note that the sign of K_θ and Π_W is invariant between fav and unfav, being quadratic in u_θ .