



HAL
open science

Full- f gyrokinetic simulations of Ohmic L-mode plasmas in linear and saturated Ohmic confinement regimes

Y. Idomura, G. Dif-Pradalier, X. Garbet, Y. Sarazin

► To cite this version:

Y. Idomura, G. Dif-Pradalier, X. Garbet, Y. Sarazin. Full- f gyrokinetic simulations of Ohmic L-mode plasmas in linear and saturated Ohmic confinement regimes. *Physics of Plasmas*, 2023, 30 (4), pp.042508. 10.1063/5.0127346 . cea-04264780

HAL Id: cea-04264780

<https://cea.hal.science/cea-04264780v1>

Submitted on 30 Oct 2023

HAL is a multi-disciplinary open access archive for the deposit and dissemination of scientific research documents, whether they are published or not. The documents may come from teaching and research institutions in France or abroad, or from public or private research centers.

L'archive ouverte pluridisciplinaire **HAL**, est destinée au dépôt et à la diffusion de documents scientifiques de niveau recherche, publiés ou non, émanant des établissements d'enseignement et de recherche français ou étrangers, des laboratoires publics ou privés.

full- f gyrokinetic simulations of ohmic L-mode plasmas in linear and saturated ohmic confinement regimes

Y. Idomura,¹ G. Dif-Pradalier,² X. Garbet,² Y. Sarazin,² and the Tore Supra Team^{2, a)}

¹⁾JAEA, CCSE, Kashiwa, Chiba 277-0871, Japan

²⁾CEA, IRFM, F-13108 St. Paul-lez-Durance cedex, France

(*Electronic mail: idomura.yasuhiro@jaea.go.jp)

(Dated: 7 March 2023)

Two time slices in linear and saturated ohmic confinement (LOC and SOC) regimes in a Tore Supra ohmic L-mode discharge are analyzed using nonlinear global full- f gyrokinetic simulations, and qualitative features of the LOC-SOC transition are reproduced. The exhaust of carbon impurity ions is caused by ion mixing, which is driven by the toroidal field stress. The intrinsic rotation develops in the opposite direction between the LOC and SOC phases, which is characterized by the different features of the mode asymmetry between trapped electron modes (TEMs) in the LOC phase and ion temperature gradient driven modes in the SOC phase, leading to the change of the profile shear stress. Here, the mode asymmetry or the ballooning angle depends not only on the profile shear and the E_r shear but also on the radial electric field E_r itself. The energy fluxes of electrons and deuterium ions are dominant in the LOC and SOC phases, respectively, and the ratio of the energy confinement time between two phases agree with the experimental value. Turbulent frequency spectra are characterized by quasi-coherent modes (QCMs) and broad-band spectra in the LOC and SOC phases, respectively. The QCMs are produced by a split of the toroidal mode number spectra of TEMs induced by the Doppler shift of poloidal $E \times B$ rotation due to E_r , which is sustained in the electron diamagnetic direction by the ripple induced counter-current rotation, and enhanced by higher temperature in the LOC phase.

I. INTRODUCTION

Ohmic L-mode discharges universally show a transition of the global energy confinement from the linear Ohmic confinement (LOC) regime, where the energy confinement time τ_E increases linearly with a density ramp, to the saturated Ohmic confinement (SOC) regime, where τ_E is saturated beyond certain critical density. This phenomenon is called as the LOC-SOC transition, and has been one of the long-standing issues in turbulent transport in magnetic confinement fusion devices^{1,2}. The LOC-SOC transition is important not only for the energy confinement but also for particle and momentum transport. During the LOC-SOC transition, the exhaust of impurity ions was observed in several devices^{3,4}, and the reversal or change of intrinsic rotation was universally observed in many devices². Since the control of impurity ions and intrinsic rotation is one of the critical issues in ITER, understanding of the LOC-SOC transition is important. The LOC-SOC transition is an interesting phenomenon also from the viewpoint of transport physics. Detailed turbulence measurements revealed the change of turbulent spectra, which are characterized by the so-called quasi-coherent modes (QCMs) and broad-band spectra in the LOC and SOC regimes, respectively⁵⁻⁷. Modulation experiments with edge cold pulses showed nonlocal transport in the LOC regime, while this feature disappeared in the SOC regime⁴.

The LOC-SOC transition has been analyzed based on theory of micro-instabilities^{3,8}, linear gyrokinetic calculations^{6,9-18}, nonlinear gyrokinetic simulations¹⁹⁻²⁶, and transport models^{21,22,27}. Many of these theoretical analyses showed transition of linearly most unstable modes from

trapped electron modes (TEMs) to ion temperature gradient driven modes (ITGs), while the importance of a mixed mode state including subdominant modes was pointed out in Ref.18. Here, the density ramp increases a collisional stabilization effect of TEMs, while the exhaust of impurity ions reduces a dilution stabilization effect of ITGs. Gyrokinetic simulations of ITGs and TEMs were validated against fluctuation measurements. In Ref.21, the LOC-SOC transition on Alcator C-Mod was analyzed using nonlinear global δf gyrokinetic simulations, and turbulent spectra obtained from the phase contrast imaging in the experiment and from its synthetic diagnostics in the simulation showed qualitative agreements. In Refs.23 and 28, the LOC-SOC transition on Tore Supra was analyzed using nonlinear local δf gyrokinetic simulations, and turbulent spectra obtained from the reflectometry measurement in the experiment and its synthetic diagnostics in the simulation showed QCMs and broad-band spectra in the LOC and SOC regimes, respectively. Regarding turbulent transport, many works analyzed heat transport, while understanding of particle and momentum transport has not been matured. On Tore Supra, carbon impurity content in the LOC regime is at a non-tracer level, and the mechanism of the exhaust of carbon impurity ions, which is linked with transport of electrons and deuterium ions through the ambipolar condition, is still an open issue. A study on momentum transport leading to the change of intrinsic rotation is very limited, because non-diffusive momentum transport is induced by various symmetry breaking effects, which are not determined within local gyrokinetic models. The Coriolis pinch effect and the residual stress effect, in which the profile shear was prescribed, were examined against the ASDEX Upgrade (AUG) experiment using local gyrokinetic calculations, and the residual stress due to the shear of radial electric field E_r , which is determined by neoclassical poloidal rotation, and the profile shear was shown to be important²⁰. On the other hand, in Ref. 24-26,

^{a)}Y. Peysson and Tore Supra Team, Nuclear Fusion 41, 1703 (2001).

1 intrinsic rotation profiles in the DIII-D experiment were repro-
 2 duced based on the balance between the diffusive momentum
 3 flux and the residual stress computed using nonlinear global
 4 δf gyrokinetic simulations with zero rotation. Here, the mo-
 5 mentum diffusivity is estimated using the heat diffusivity and
 6 a given Prandtl number, and the turbulent residual stress was
 7 correlated with the shear of turbulent zonal flows and the tur-
 8 bulent intensity gradient. In Ref. 29, intrinsic rotation profiles
 9 in the AUG experiment were analyzed using nonlinear global
 10 δf gyrokinetic simulations, where rotation and density pro-
 11 files are evolved, and it was shown that intrinsic flow gradi-
 12 ents at experimental levels were formed mainly by the profile
 13 shear stress. Therefore, theoretical understanding of momen-
 14 tum transport is not converged, and further global simulations
 15 are needed. Finally, in spite of the development of gyroki-
 16 netic modeling and theory, any single model cannot capture
 17 all aspects of the above complicated physics in the LOC-SOC
 18 transition in a self-consistent manner.

19 To address this challenging issue, we need global full- f gy-
 20 rokinetic models, which compute both plasma turbulence and
 21 plasma profiles in a self-consistent manner including multi-
 22 ple transport channels induced by collisional and turbulent
 23 transport. This capability enables us to study the following ef-
 24 fects, which are important for understanding transport physics
 25 in the LOC-SOC transition. Firstly, the formation of intrinsic
 26 rotation is directly computed under the strict toroidal angu-
 27 lar momentum conservation³⁰⁻³², in which in addition to
 28 the turbulent Reynolds stress, other mechanisms such as the
 29 neoclassical Reynolds stress, the toroidal field stress, and the
 30 $J \times B$ torque make significant contributions. Also, various
 31 symmetry breaking mechanisms such as the shear of the ra-
 32 dial electric field E_r and zonal flows, the profile shear, and
 33 the turbulent intensity gradient are involved consistently, and
 34 their relative importance can be examined. Secondly, E_r is
 35 determined self-consistently, and the resulting effects such as
 36 the Doppler shift of turbulent spectra, turbulence suppression
 37 by the E_r shear, and the residual stress induced by E_r and E_r
 38 shear are analyzed in a straightforward manner. Thirdly, self-
 39 consistent interaction of multiple transport channels gives the
 40 following effects. The collisional ion-electron coupling deter-
 41 mine the ion heating condition in Ohmic L-mode plasmas.
 42 Particle fluxes of electrons and multi-species ions satisfy the
 43 ambipolar condition, which is important in analyzing impu-
 44 rity transport at a non-tracer level. Particle and momentum
 45 transport channels are coupled through the toroidal angular
 46 momentum conservation. It is noted that in general, global δf
 47 gyrokinetic models or gradient driven simulations involve arti-
 48 ficial source/sink terms everywhere in a plasma to fix plasma
 49 profiles, and neoclassical physics is excluded. Therefore, im-
 50 portant conservation properties such as the toroidal angular
 51 momentum conservation and the ambipolar condition are not
 52 guaranteed or at least modified, which make study of the
 53 above effects difficult. In addition, former works on compar-
 54 isons of flux driven global full- f gyrokinetic simulations and
 55 gradient driven global δf gyrokinetic simulations showed sig-
 56 nificantly different properties in avalanche like nonlocal trans-
 57 port, zonal flows, and staircase structures^{33,34}. These motivate
 58 us to address the LOC-SOC transition via flux driven global

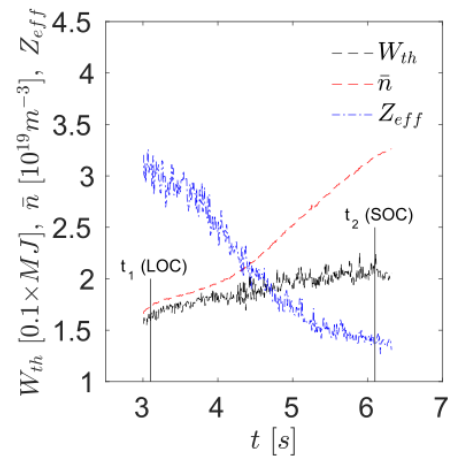


FIG. 1. Evolution of stored thermal energy W_{th} , line averaged density \bar{n} , and effective charge number Z_{eff} in Tore Supra ohmic L-mode discharge 48102. $t_1 \sim 3.1$ s is in a LOC phase, and $t_2 \sim 6.1$ s is in a SOC phase. Reproduced with permission from Plasma Phys. Control. Fusion 59, 064010 (2017)²³. Copyright 2017 IOP Publishing.

full- f gyrokinetic simulations.

In this study, we address the LOC-SOC transition on Tore Supra using the Gyrokinetic Toroidal 5D full- f Eulerian code GT5D³⁵⁻³⁷, which computes an electrostatic global full- f model in multi-species plasmas with kinetic electrons. The conditions of numerical experiments are based on Tore Supra ohmic L-mode discharge 48102, which were already analyzed in detail based on nonlinear local δf gyrokinetic simulations²³. The latter work showed validation studies on heat transport and turbulent spectra, while momentum transport and impurity transport have not been analyzed yet. In this work, we address all transport channels and the corresponding turbulent spectra with a single global full- f gyrokinetic model in a self-consistent manner.

The remainder of the paper is organized as follows. In Sec.II, the experimental data, calculation models, and simulation parameters are explained. In Sec.III, the validity of simulation parameters with scaled plasma sizes is discussed. In Sec.IV, numerical experiments in the LOC and SOC phases are presented, in which transport properties of particle, momentum, and energy transport are discussed, and turbulent spectra are analyzed. Finally, a summary is given in Sec.V.

II. CALCULATION MODEL

In this work, we analyze Tore Supra ohmic L-mode discharge 48102, which was discussed in detail in Refs.6, 14, 23, and 28. As shown in Figs.1 and 2, the discharge is characterized by a density ramp, leading to transition from the LOC regime to the SOC regime. In the LOC-SOC transition, the trend of the energy confinement time is changed from the linear increase proportional to the electron density to the saturation, the effective charge number Z_{eff} is decreased by ex-

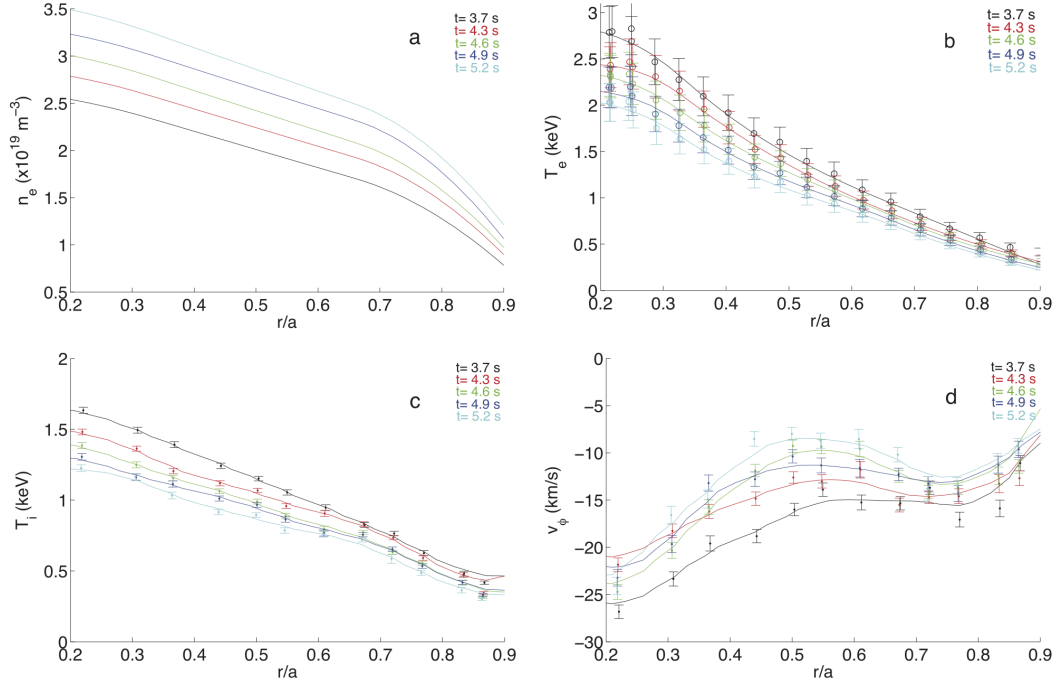


FIG. 2. (a) electron density, (b) electron temperature, (c) ion temperature, and (d) toroidal rotation profiles of carbon impurity ions at five time slices in Tore Supra ohmic L-mode discharge 48102. Reproduced with permission from Plasma Phys. Control. Fusion 57, 035002 (2015)¹⁴. Copyright 2015 IOP Publishing.

1 hausting carbon impurity ions, and the relative intrinsic rota- 16
 2 tion direction with respect to the edge rotation velocity is re- 17
 3 versed (see Fig.2(d)). We focus on two time slices in the LOC 18
 4 and SOC regimes, $t_1 \sim 3.1$ s and $t_2 \sim 6.1$ s, which were ana- 19
 5 lyzed in Ref.23. In the latter work, it was shown that the LOC 20
 6 and SOC phases are respectively dominated by TEM and ITG, 21
 7 and the electron temperature gradient driven mode is unstable 22
 8 only in the SOC phase. 23

9 Numerical experiments are conducted using GT5D, which 24
 10 computes electrostatic ion scale turbulence driven by ITG and 25
 11 TEM. In the gyro-center coordinates, $\mathbf{Z} = (t; \mathbf{R}, v_{\parallel}, \mu, \alpha)$, a 26
 12 conservative form of the electrostatic gyrokinetic equation is 27
 13 commonly used for all particle species including electron, 28

$$\begin{aligned} & \frac{\partial \mathcal{J}_s f_s}{\partial t} + \nabla \cdot (\mathcal{J}_s \mathbf{R} f_s) + \frac{\partial}{\partial v_{\parallel}} (\mathcal{J}_s v_{\parallel} f_s) \\ & = \mathcal{J}_s \sum_s C(f_s, f_{s'}) + \mathcal{J}_s S_{src,s} + \mathcal{J}_s S_{snk,s}, \end{aligned} \quad (1)$$

$$\mathbf{R} = v_{\parallel} \mathbf{b} + \mathbf{v}_E + \mathbf{v}_D, \quad (2)$$

$$\mathbf{v}_E = \frac{c}{q_s B_{\parallel}^*} \mathbf{b} \times (q_s \nabla \langle \phi \rangle_{\alpha}), \quad (3)$$

$$\mathbf{v}_D = \frac{c}{q_s B_{\parallel}^*} \mathbf{b} \times (m_s v_{\parallel}^2 \mathbf{b} \cdot \nabla \mathbf{b} + \mu \nabla B), \quad (4)$$

$$v_{\parallel} = -\frac{\mathbf{B}^*}{m_s B_{\parallel}^*} \cdot (q_s \nabla \langle \phi \rangle_{\alpha} + \mu \nabla B) \quad (5)$$

14 where f_s denotes the guiding-center distribution function,
 15 $\mathcal{J}_s = m_s^2 B_{\parallel}^*$ is the Jacobian of the gyro-center coordinates,

$C(f_s, f_{s'})$ is a multi-species Coulomb collision operator³⁸,
 $S_{src,s}$ and $S_{snk,s}$ are respectively the source and sink terms, \mathbf{R}
 is the position of the guiding center, \mathbf{v} is the velocity of the guid-
 ing center, $v_{\parallel} = \mathbf{b} \cdot \mathbf{v}$ and $v_{\perp} = |\mathbf{b} \times \mathbf{v}|$ are respectively the ve-
 locities in the parallel and perpendicular directions to the mag-
 netic field, $\mu = m_s v_{\perp}^2 / 2B$ is the magnetic moment, α is the
 gyro-phase angle, $\mathbf{B} = B \mathbf{b}$ is the magnetic field, \mathbf{b} is the unit
 vector in the parallel direction, m_s and q_s are respectively the
 mass and charge of the particle species s , c is the velocity of
 light, $\Omega_s = q_s B / m_s c$ is the cyclotron frequency, $B_{\parallel}^* = \mathbf{b} \cdot \mathbf{B}^*$ is
 a parallel component of $\mathbf{B}^* = \mathbf{B} + (B v_{\parallel} / \Omega_s) \nabla \times \mathbf{b}$, ϕ is the elec-
 trostatic potential of turbulent fields, and the gyro-averaging
 operator is defined as $\langle \cdot \rangle_{\alpha} \equiv \oint d\alpha / 2\pi$.

The electrostatic potential is determined using the hybrid
 kinetic electron model³⁷, where the gyrokinetic Poisson equa-
 tion is modified as

$$\begin{aligned} & -\sum_s \nabla_{\perp} \cdot \frac{\rho_{ts}^2}{\lambda_{Ds}^2} \nabla_{\perp} \phi_{n \neq 0} + \frac{\alpha_p}{\lambda_{De}^2} \phi_{n \neq 0} \\ & = 4\pi \left[\sum_{s \neq e} q_s \int f_{s,n \neq 0} \delta([\mathbf{R} + \rho] - \mathbf{x}) d^6 Z \right. \\ & \quad \left. + q_e \int f_{e,t,n \neq 0} \delta([\mathbf{R} + \rho] - \mathbf{x}) d^6 Z \right], \end{aligned} \quad (6)$$

$$\begin{aligned} & -\sum_s \nabla_{\perp} \cdot \frac{\rho_{ts}^2}{\lambda_{Ds}^2} \nabla_{\perp} \phi_{n=0} \\ & = 4\pi \sum_s q_s \int f_{s,n=0} \delta([\mathbf{R} + \rho] - \mathbf{x}) d^6 Z. \end{aligned} \quad (7)$$

1 Here, $\mathbf{R} + \rho$ is the particle position, $\rho = \mathbf{b} \times \mathbf{v}/\Omega_s$ is the Lar-
 2 mor radius, $d^6Z = m_s^2 B^* d\mathbf{R} dv_{\parallel} d\mu d\alpha$ is the phase space vol-
 3 ume of the gyro-center coordinates, $\rho_{ts} = v_{ts}/\Omega_s$ is the Larmor
 4 radius evaluated with the thermal velocity $v_{ts} = (T_s/m_s)^{1/2}$,
 5 $\lambda_{Ds} = (T_s/4\pi n_s q_s^2)^{1/2}$ is the Debye length, n_s is the density,
 6 T_s is the temperature, n is the toroidal mode number, $f_{e,t}$ is
 7 a trapped part of the electron distribution function, and α_p is
 8 the flux-surface averaged fraction of passing electrons. The
 9 gyrokinetic Poisson equation is decomposed into $n \neq 0$ and
 10 $n = 0$ parts. The former is computed using adiabatic passing
 11 electrons and kinetic trapped electrons to avoid the so-called
 12 Ω_H mode³⁹. Here, a full kinetic electron distribution f_e is
 13 computed by Eq. (1), and in Eq. (6), its trapped part $f_{e,t}$ is
 14 extracted following a trapped-passing boundary at each position.
 15 The latter is solved using full kinetic electrons to satisfy
 16 the ambipolar condition, while $\phi_{n=0, m \neq 0}$ convective cells are
 17 filtered out from the solution to avoid the Ω_H mode, where m
 18 is the poloidal mode number. In the l.h.s., a linear polariza-
 19 tion density with a long wavelength approximation, $k_{\perp}^2 \rho_{ts}^2 \ll 1$
 20 is considered, which is valid for ion scale turbulence. The
 21 gyrokinetic Poisson operators including the ion polarization
 22 density and the adiabatic passing electron density are defined
 23 using the initial density and temperature.

24 The above full- f gyrokinetic model yields the following
 25 two balance relations. By taking the time derivative and the
 26 flux-surface average of Eq. (7) and substituting Eq. (1), the
 27 ambipolar condition is derived as

$$-\sum_s \frac{\rho_{ts}^2}{\lambda_{Ds}^2} \frac{\partial E_r}{\partial t} = 4\pi \sum_s q_s \langle f_s(\mathbf{R} \cdot \nabla r) - S_{src,s} - S_{snk,s} \rangle_{gf}, \quad (8)$$

28 where E_r is the radial electric field and the gyro/flux-surface
 29 average operator is defined as,

$$\langle A \rangle_{gf} = \left\langle \int A(\mathbf{Z}) \delta(\mathbf{R} + \rho - \mathbf{x}) d^6Z \right\rangle_f, \quad (9)$$

30 and $\langle \cdot \rangle_f$ is the flux-surface average operator. From the gyroki-
 31 netic equation and the Hamilton's equation for the canonical
 32 toroidal angular momentum, the toroidal angular momentum
 33 balance³² is given as

$$\begin{aligned} & \left\langle \frac{\partial m_s v_{\parallel} b_{\phi} f_s}{\partial t} \right\rangle_{gf} + \left\langle \frac{1}{\mathcal{J}} \frac{\partial}{\partial \mathbf{R}} \cdot (\mathcal{J} \mathbf{R} m_s v_{\parallel} b_{\phi} f_s) \right\rangle_{gf} \\ & - \left\langle \frac{q_s}{c} f_s \mathbf{R} \cdot \nabla \psi \right\rangle_{gf} + \left\langle f_s \frac{\partial \langle \phi \rangle_{\alpha}}{\partial \phi} \right\rangle_{gf} \\ & - \left\langle m_s v_{\parallel} b_{\phi} \sum_{s'} C(f_s, f_{s'}) \right\rangle_{gf} \\ & - \langle m_s v_{\parallel} b_{\phi} (S_{src,s} + S_{snk,s}) \rangle_{gf} = 0, \end{aligned} \quad (10)$$

34 where ψ is the poloidal flux, ϕ is the toroidal angle, and b_{ϕ} is
 35 the covariant toroidal component of \mathbf{b} . Here, the first term is
 36 the inertial term, the second term is the stress term, the third
 37 term is the radial current term, the fourth term is the toroidal
 38 field stress term, the fifth term is the collision term, and the
 39 sixth term is the source term. In the stress term, \mathbf{R} involve

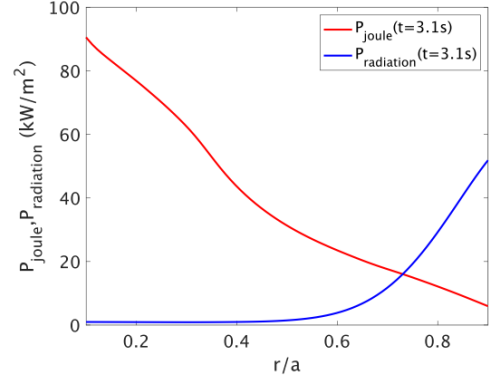


FIG. 3. Radial profiles of the ohmic heating and the radiation loss at $t = 3.1$.

both the magnetic drift \mathbf{v}_D and the $E \times B$ drift \mathbf{v}_E , which induce the neoclassical and turbulent Reynolds stress, respectively. Although the neoclassical Reynolds stress is negligible in the axisymmetric limit, it is greatly enhanced in the presence of non-axisymmetric turbulent fluctuations. The radial current term gives the $J \times B$ torque, when summed over all species. This term becomes important in the presence of finite particle transport. The toroidal field stress term is determined by the phase shift between the perturbed density and fluctuating toroidal electric fields, and is also interpreted as the polarization stress⁴⁰ or the generalized Maxwell stress³¹. The code has been verified through cross-code comparisons^{35,41,42}, neoclassical transport^{43,44}, convergence of E_r up to third order gyrokinetics³², and multi-species plasma transport³⁸.

In the numerical experiment, we consider deuterium (D) ions, carbon (C) impurity ions, and kinetic electrons with $m_D/m_e = 3,672$ in a circular concentric tokamak configuration, which is a good approximation for low β plasmas on Tore Supra. The device parameters are $B = 3.7\text{T}$, $R_0 = 2.38\text{m}$, $a = 0.7\text{m}$, $q = 0.76 \sim 4.6$, and $P_{joule} \sim 600\text{kW}$. The nominal plasma parameters at $r/a = 0.5$ in the LOC phase are $n_e \sim 1.88 \times 10^{19}\text{m}^{-3}$, $T_e = 1.54\text{keV}$, $T_D = 1.04\text{keV}$, $Z_{eff} \sim 3.06$, $v_e^* \sim 0.022$, and $\rho^{*-1} \sim 563$. On the other hand, those in the SOC phase are $n_e \sim 3.49 \times 10^{19}\text{m}^{-3}$, $T_e = 0.93\text{keV}$, $T_D = 0.73\text{keV}$, $Z_{eff} \sim 1.36$, $v_e^* \sim 0.164$, and $\rho^{*-1} \sim 672$. Here, B is the toroidal magnetic field, R_0 and a are the major and minor radii, r is the radial coordinate, q is the safety factor, P_{joule} is the ohmic heating power, $v_e^* \equiv qR_0/(\varepsilon^{3/2} v_{te} \tau_{e,D})$ is the normalized electron collisionality, $\varepsilon = r/R_0$, $\tau_{s,s'}$ is the collision time between the species s and s' , and $\rho^* = \rho_{tD}/a$.

Numerical experiments are conducted for $0 \leq r \leq a_0$ to avoid extremely low temperature below 100 eV, where $a_0 = 0.9a$. The boundary conditions of density, rotation, and temperature are given based on the experimental values at $r = a_0$, where rotation and temperature are assumed to be the same between D and C. In the initial condition, density and temperature profiles are given by the experimental data, while rotation profiles are set as rigid rotation with the edge rotation velocity, which is determined mainly by the neoclassical

1 toroidal viscosity induced by the toroidal ripple, and is almost
 2 unchanged through the LOC-SOC transition (see Fig.2(d)).
 3 Therefore, the edge rotation velocity is given by an approx-
 4 imate average value over the LOC-SOC transition. By using
 5 the same initial rotation profiles and boundary conditions, we
 6 compare the formation of intrinsic rotation between the LOC
 7 and SOC phases.

8 According to transport analysis, the radial deposition pro-
 9 files of the ohmic heating and the radiation loss are given as
 10 Fig.3, where the radiation loss makes a dominant contribution
 11 to the energy loss ($\sim 75\%$ and $\sim 50\%$ in the LOC and SOC
 12 phases, respectively). In this work, the time scale of numeri-
 13 cal experiment $\sim 10\text{ms}$ is significantly shorter than the time
 14 scales of the current diffusion and the particle confinement.
 15 Therefore, we do not compute the inductive toroidal electric
 16 field, and the resulting ohmic heating and Ware pinch. The
 17 ohmic heating is simply modelled as an on-axis electron heat-
 18 ing without particle and momentum inputs using a fixed heat
 19 source model, $S_{src,s} = v_{src,s} A_{src,s}(r)(f_{M1} - f_{M2})$, where $A_{src,s}$
 20 is the deposition profile, f_{M1} , f_{M2} are local Maxwellian distri-
 21 butions with different temperatures, and the heating rate $v_{src,s}$
 22 is determined to satisfy the target power input P_{in} . As the elec-
 23 tron collision time is sufficiently small $\tau_{ee} \sim 10R_0/v_0$, veloc-
 24 ity space perturbations due to the source term is expected to be
 25 quickly thermalized, leading to velocity distribution functions
 26 relevant for ohmic heating plasmas.

27 The sink model is somewhat complicated. The heat sink is
 28 given by the radiation loss for $r/a > 0.7$ or $r/a_0 > 0.8$. The
 29 edge rotation velocity is fixed by the momentum sink due to
 30 the neoclassical toroidal viscosity induced by the toroidal rip-
 31 ple. However, the momentum and particle sink due to neutral
 32 particles is uncertain, because the detailed distribution of neu-
 33 tral particles was not obtained. In order to model these prop-
 34 erties, we implement two sink models. One is a Krook type
 35 sink operator $S_{snk1,s} = v_{snk1,s} A_{snk1,s}(r)(f_s - f_{0s})$, which gives
 36 an effective boundary condition by keeping plasma paramet-
 37 ers at the boundary on average. Here, f_{0s} is the initial dis-
 38 tribution function, the sink rate is given as $v_{snk1,s} = 0.1v_0/a$,
 39 the reference velocity v_0 is given by the thermal velocity of
 40 D ions at 1keV, and the deposition profile $A_{snk1,s}$ is localized
 41 for $r/a_0 > 0.95$. It is noted that the above sink rate is chosen
 42 based on the sensitivity study in Ref. 45. Another is an ax-
 43 isymmetric variant of the Krook type sink operator $S_{snk2,s} =$
 44 $v_{snk2,s} A_{snk2,s}(r)(f_{s,n=0} - f_{0s})$, where $v_{snk2,s} = 0.01v_0/a$ and
 45 $A_{snk1,s}$ is distributed over $r/a_0 > 0.8$ following the radiation
 46 loss profile in Fig.3. The latter model modifies only an ax-
 47 isymmetric part, and a turbulent part is not affected. In addi-
 48 tion, the sink parameter is chosen so that $v_{snk2,s}$ is sufficiently
 49 smaller than the linear growth rates of TEM and ITG, while
 50 the sink effect is large enough to avoid deviations of density
 51 and temperature profiles from the experimental ones by accu-
 52 mulation of particles and energy induced by plasma transport.
 53 These features are important as a sink model distributed over
 54 turbulent regions. It is noted that this model also works as a
 55 particle source in the outer radii. However, $v_{snk2,s}$ is too weak
 56 to fix rotation profiles to the initial condition, and rotation pro-
 57 files freely evolve in this region.

58 In the numerical experiment, in order to save the com-

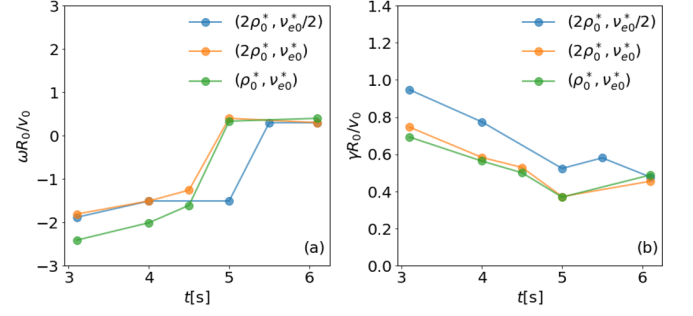


FIG. 4. (a) frequency ω and (b) growth rate γ of TEM and ITG at each time slice in the cases with $(n, \rho^*, v_e^*) = (50, \rho_0^*, v_{e0}^*)$, $(25, 2\rho_0^*, v_{e0}^*)$, and $(25, 2\rho_0^*, v_{e0}^*/2)$. $\omega < 0$ is the direction of the electron diamagnetic rotation.

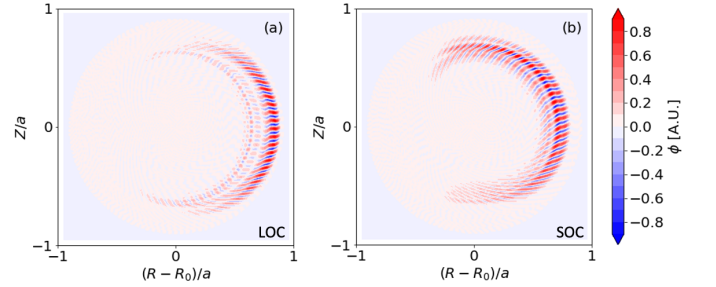


FIG. 5. Linear eigenfunctions $n = 25$ modes at (a) LOC and (b) SOC phases in the case with $\rho^* = 2\rho_0^*$ and $v_e^* = v_{e0}^*$.

putational cost, the plasma size is scaled by a half, and a $1/6$ wedge of the torus is computed using 5D grids $(N_R, N_\zeta, N_Z, N_{v\parallel}, N_{v\perp}) = (240, 48, 240, 96, 24)$. The $1/6$ wedge torus model is chosen based on the convergence studies in Ref. 46 and in Appendix A. Here, the computational domain in the velocity space is chosen as $-4v_{ts,r=0} \leq v_{\parallel} \leq 4v_{ts,r=0}$ and $0 \leq \sqrt{2\mu B/m_s} \leq 3v_{ts,r=0}$, where $v_{ts,r=0}$ is the thermal velocity at $r = 0$. It is noted that the contrast of temperature is $T_e(r=0)/T_e(r=a_0) \sim 10$ and $T_D(r=0)/T_D(r=a_0) \sim 5$, and the velocity space normalized by the thermal velocity at $r = a_0$ is expanded by ~ 3 times and ~ 2 times, respectively. The velocity grids are chosen to keep sufficient resolution from the core to the edge. GT5D was highly optimized on Fugaku, where a new mixed-precision communication-avoiding Krylov solver enabled efficient computation of kinetic electrons with the real mass ratio⁴⁷. Even with the reduced simulation parameters and the optimization techniques, numerical experiments require costly computation, and a single numerical experiment typically for $\sim 1,000R/v_{tD}$ required $\sim 1.5 \times 10^5$ CPU hours on Fugaku⁴⁸.

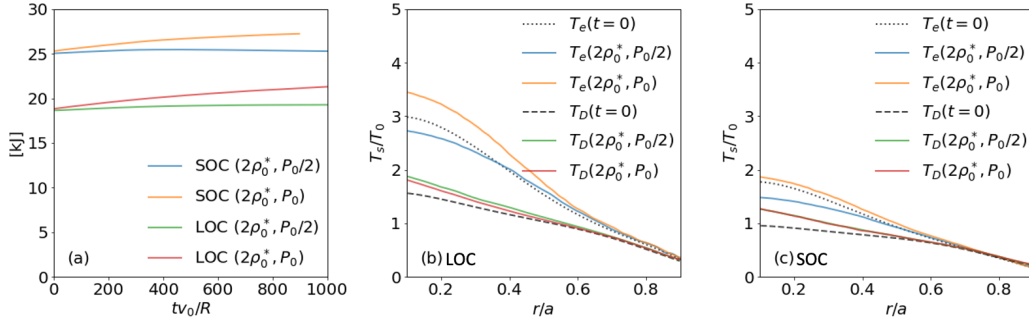


FIG. 6. (a) time history of the stored energy W_t and temperature profiles at $tv_0/R = 800$ in (b) LOC and (c) SOC phases in the cases with $(\rho^*, P_{in}) = (2\rho_0^*, P_0/2)$ and $(2\rho_0^*, P_0)$, respectively. In (b) and (c), the temperature is normalized by $T_0 = 1\text{keV}$ and dashed and dotted black curves show the initial condition given by the experimental data.

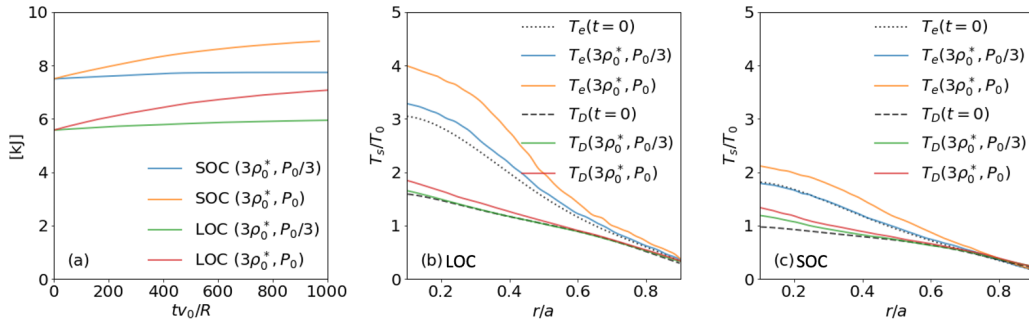


FIG. 7. (a) time history of the stored energy W_t and temperature profiles at $tv_0/R = 800$ in (b) LOC and (c) SOC phases in the cases with $(\rho^*, P_{in}) = (3\rho_0^*, P_0/3)$ and $(3\rho_0^*, P_0)$, respectively. In (b) and (c), dashed and dotted black curves show the initial condition given by the experimental data.

1 III. VALIDITY OF SCALED PLASMA SIZE

The scaling of the plasma size affects two important non-dimensional parameters, ρ^* and v_e^* . In this section, we discuss their influences on numerical experiments. Figure 4 shows the frequency and growth rate of $n = 50$ and $n = 25$ modes at each time slice in the cases with $\rho^* = \rho_0^*$ and $\rho^* = 2\rho_0^*$, respectively, where ρ_0^* is the experimental value. As shown in Ref.23, the most unstable modes in the LOC and SOC phases are given by TEM and ITG, respectively. The latter work also showed that in the LOC phase, ITG is stabilized by a dilution effect due to high Z_{eff} , while in the SOC phase, TEM is stabilized by high v_e^* due to the density ramp. It is noted that in the latter work, local gyrokinetic analyses were conducted at $r/a = 0.37$. However, in Fig.5, global eigenfunctions of both TEM and ITG are peaked in the outer radii. In the experimental parameters, $\rho^* = \rho_0^*$ and $v_e^* = v_{e0}^*$, the transition from TEM to ITG due to the density ramp occurs at $t = 4.5 \sim 5.0\text{s}$. In the experiment, the transition timing was shown to be at $t \sim 4.85^{28}$. However, in the 1/2 scale parameters, $\rho^* = 2\rho_0^*$ and $v_e^* = v_{e0}^*/2$, the transition timing is shifted to $t = 5.0 \sim 5.5\text{s}$ and the growth rate of TEM is increased by

20 $\sim 30\%$. This is attributed to the change of $v_e^* \propto R_0$, which characterizes the collisional stabilization effect on TEM. It is noted that the collisional stabilization affects only TEM, and the stability of ITG is similar in all the cases. On the other hand, by doubling the collisionality, the stability of TEM and the transition timing in the experimental parameters are recovered even with the 1/2 scale parameters. Therefore, in this work, we use the modified 1/2 scale parameters, in which the collisionality is adjusted to recover the linear stability in the original experimental parameters.

ρ^* scaling is essentially nonlinear and rather complicated. In the experiment, the confinement scaling was shown to follow the neo-Alcator scaling and the L-mode scaling in the LOC and SOC phases, respectively^{1,2}. The latter scaling has Bohm-like ρ^* dependency. On the other hand, according to the derivation of non-dimensional scaling in Ref.49, the former scaling gives gyro-Bohm like ρ^* dependency, while it does not satisfy the so-called Kadomtsev constraint. These controversial ρ^* dependencies between the LOC and SOC phases make the choice of input power difficult in scaled numerical experiments. In terms of the power balance, temperature profiles in a $1/N$ scale plasma become similar to the original ones by imposing the input power of $P_{in} = P_0/N$

1 and $P_{in} = P_0$, when the energy confinement follows the Bohm 55
 2 scaling and the gyro-Bohm scaling, respectively. Here, P_0 is 56
 3 the input power in the experiment. In order to determine a 57
 4 physically sound heating condition, we conduct power scan 58
 5 numerical experiments with the modified $1/N$ scale param- 59
 6 eters, in which the Bohm like heating condition $P_{in} = P_0/N$ 60
 7 and the gyro-Bohm like heating condition $P_{in} = P_0$ are com- 61
 8 pared. Figures 6 and 7 show the numerical experiments with 62
 9 the modified $1/2$ and $1/3$ scale parameters, respectively. In 63
 10 Fig.6(a), the stored energy in the Bohm like heating condition 64
 11 is almost unchanged both in the LOC and SOC phases. In 65
 12 Figs.6(b) and 6(c), the electron temperature is higher in the 66
 13 gyro-Bohm heating condition, while the stiff ion temperature 67
 14 is similar in both heating conditions. In Fig.7, these features 68
 15 are more pronounced in the modified $1/3$ scale parameters, 69
 16 and the increases of the stored energy and the electron tem- 70
 17 perature in the gyro-Bohm heating condition are clearly seen. 71
 18 It is noted that in Fig.6 (c), the electron temperature in the 72
 19 Bohm heating condition is lower than the experiment. This is 73
 20 attributed to the modified collisionality, which is increased to 74
 21 satisfy $\nu_e^* = \nu_0$. The SOC phase is characterized by high colli- 75
 22 sionality, and the increased collisionality leads to stronger ion- 76
 23 electron coupling. The Bohm like scaling in the SOC phase is 77
 24 consistent with the L-mode scaling¹ and our previous works 78
 25 on ρ^* scaling in L-mode plasmas^{33,46,50,51}. On the other hand, 79
 26 the Bohm like scaling in the LOC phase may not be consis- 80
 27 tent with the neo-Alcator scaling. This may be attributed to 81
 28 the lack of Kadomtsev constraint in it. This issue will be ad- 82
 29 dressed in future work. Following the results of the above 83
 30 power scan numerical experiments, in this work, we use the 84
 31 Bohm-like heating condition. 85

32 The use of scaled plasma sizes may also affect momen- 86
 33 tum transport. According to the higher order ballooning the- 87
 34 ory^{52,53}, the linear mode asymmetry characterized by the bal- 88
 35 looning angle θ_0 is derived as $\theta_0 \propto \rho^{*1/3}$, and the residual 89
 36 stress C_s given by the profile shear stress theory⁴¹ becomes 90
 37 $C_s \propto \rho^{*1/3}$, which is rather weak ρ^* dependency. However, 91
 38 nonlinear gyrokinetic simulations showed much stronger ρ^* 92
 39 dependencies of the residual stress and/or the resulting intrinsic 93
 40 rotation profiles^{29,31,33,54,55}. Therefore, the residual stress 94
 41 in scaled numerical experiments may be larger than that in the 95
 42 experiment. However, ρ^* scan numerical experiments typi- 96
 43 cally show similar shapes of intrinsic rotation profiles with 97
 44 different flow velocities^{33,55}, and thus, at least, the sign rela- 98
 45 tion of rotation changes from the initial rigid rotation is ex- 99
 46 pected to be preserved. 100

47 IV. NUMERICAL EXPERIMENT 101

48 A. Particle transport 102

49 We conduct two numerical experiments in the LOC and 103
 50 SOC phases up to $t\nu_0/R \sim 1,000$ using the modified $1/2$ scale 104
 51 parameters with $P_{in} = P_0/2$. Figure 8 shows the radial pro- 105
 52 files of the density, the parallel flow, and the temperature at 106
 53 the end of the numerical experiments. The time scale of the 107
 54 numerical experiment is significantly short compared to the 108

particle confinement time, and thus, the electron density is
 almost unchanged from the initial condition. However, in the
 LOC phase, where the charge densities of D and C are compa-
 rable, the density profiles of D and C respectively show peak-
 ing and relaxation, suggesting ion mixing (see Fig.8(a)). As
 the exhaust of C begins around $t \sim 3$ s in Fig.1(a), this result
 is qualitatively consistent with the experiment. On the other
 hand, in the SOC phase, where the charge density of C is at
 the tracer level, this kind of ion mixing does not occur. Fig-
 ure 9 shows the radial profiles of the radial current. **Both in
 the LOC and SOC phases, the steady ambipolar condition,
 $\sum_s q_s \langle f_s(\mathbf{R} \cdot \nabla r) \rangle_{gf} = 0$, is approximately satisfied.** However,
 the balance of the radial current among three species is quite
 different between the LOC and SOC phases. In the LOC
 phase, the D current and the C current greatly exceed the elec-
 tron current, and they balance with each other, leading to the
 ion mixing and the exhaust of C. On the other hand, the SOC
 phase shows a typical ambipolar condition, where the D cur-
 rent is comparable to the electron current and the balance is
 established mainly by the bulk species.

In the SOC phase, the particle confinement time τ_e and the
 impurity confinement time τ_C estimated using the volume av-
 eraged particle numbers and the particle fluxes at $r/a \sim 0.6$
 are respectively $\tau_e \sim 21.4\tau_E$ and $\tau_C \sim 3.4\tau_E$, where τ_E is esti-
 mated using the input power and the stored energy at $t\nu_0/R \sim$
 1,000 (see Fig.6(a)). The impurity confinement time in the
 SOC phase is in a similar range as that observed in the tracer
 impurity injection experiments on Tore Supra⁵⁶. In the LOC
 phase, the particle and impurity confinement times are esti-
 mated as $\tau_e \sim 29.4\tau_E$ and $\tau_C \sim 1.7\tau_E$, respectively. Although
 this result looks similar to the SOC phase, the impurity con-
 tent in the LOC phase is at a non-tracer level, $n_C q_C^2 > n_D q_D^2$,
 and the time scale of **inward D flux** is also comparable to the
 impurity confinement time, $\tau_D \sim 2.9\tau_E$, where τ_D is esti-
 mated using the absolute value of **D flux**. This asymmetric electron
 and D transport with an order of magnitude different particle
 confinement times suggest fast mixing of ions, leading to the
 exhaust of C impurity ions.

103 B. Momentum transport 104

105 Figures 8(b) and 8(e) show rotation profiles observed in
 106 the numerical experiments. Although the initial condition is
 107 given by rigid rotation with the same velocity for both D and
 108 C ions, their rotation profiles deviate from the initial condi-
 109 tion before the excitation of micro-instabilities. The rotation
 110 difference between D and C ions is determined by the neo-
 111 classical parallel momentum balance, and is produced by the
 112 pressure gradient, leading to the large rotation difference in
 113 the outer radii characterized by steep density profiles. This
 neoclassical effect generates the bootstrap current as shown
 in the benchmark of the parallel momentum balance between
 GT5D and the neoclassical theory in Appendix B. Because of
 this neoclassical effect, an effective boundary condition of the
 C rotation is deviated from the initial value of $U_C/\nu_0 \sim -0.05$
 to $U_C/\nu_0 \sim -0.08$. It is noted that the rotation difference is
 sustained by the collisional parallel momentum exchange be-

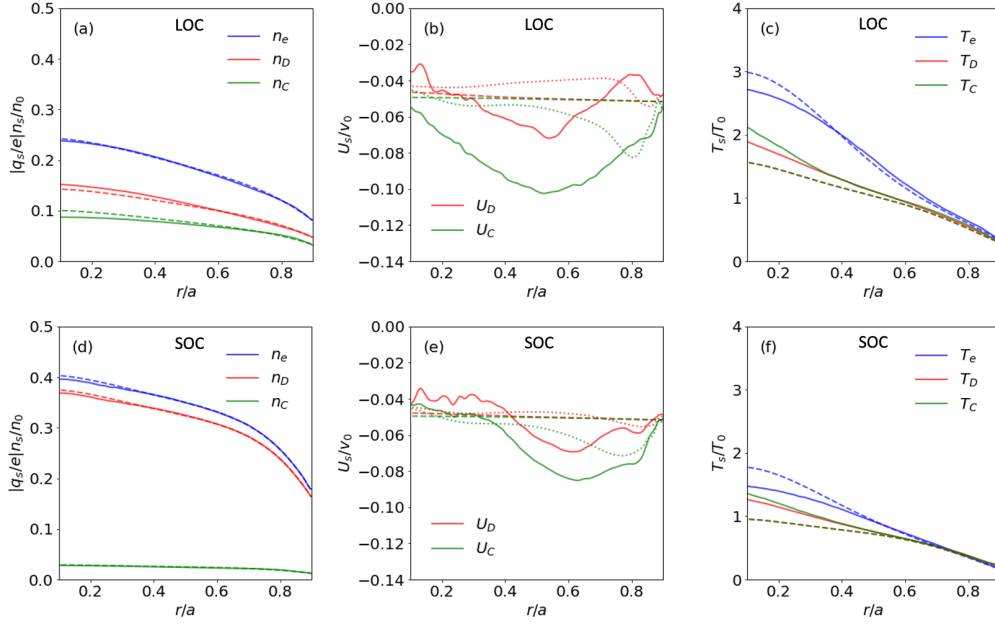


FIG. 8. Radial profiles of (a),(d) the density, (b),(e) the parallel flow, and (c),(f) the temperature averaged over $tv_0/R = 900 \sim 1000$. (a)-(c) and (d)-(f) respectively show the numerical experiments in the LOC and SOC phases. In (a) and (d), the density is normalized by $n_0 = 10^{20} \text{m}^{-3}$. In (a), (c), (d), and (f), dashed curves show the initial profiles given by the experimental data. In (b) and (e), dashed red and green curves show the initial profiles given by rigid rotation, and broken curves show rotation profiles in the neoclassical equilibrium states before the excitation of micro-instabilities ($tv_0/R = 50$).

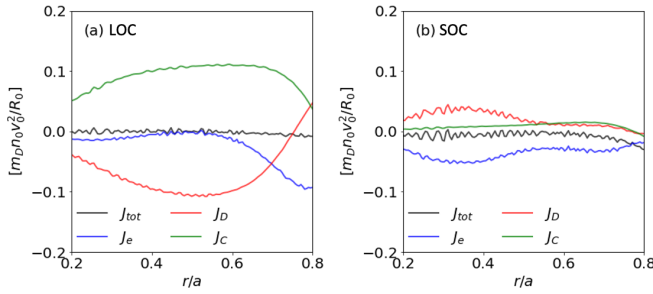


FIG. 9. Radial profiles of the radial current $J_s = \langle \frac{q_s}{c} J_s \mathbf{R} \cdot \nabla \psi \rangle_{gf}$ (the third term in Eq. (10)) averaged over $tv_0/R = 0 \sim 1000$. (a) and (b) respectively show the numerical experiments in the LOC and SOC phases.

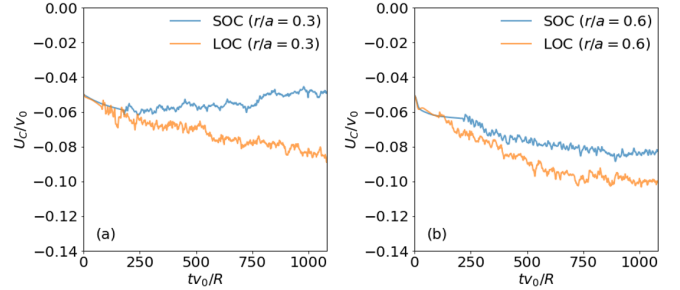


FIG. 10. Time histories of the carbon rotation at (a) $r/a = 0.3$ and (b) $r/a = 0.6$.

1 tween D and C ions, which is a stronger effect than the radial 13
 2 momentum transport, and thus, the axisymmetric sink $S_{snk2,s}$ 14
 3 with $v_{snk2,s} = 0.01v_0/a$ is too weak to compensate the rotation 15
 4 difference. 16

5 The intrinsic rotation shows qualitatively different profiles 17
 6 in the LOC and SOC phases. In the LOC phase, the change 18
 7 of C rotation in the counter-current direction develops glob- 19
 8 ally. Here, the negative direction corresponds to the counter- 20
 9 current direction. On the other hand, in the SOC phase, the 21
 10 change of C rotation in the plasma core is in the co-current 22

11 direction over significant radii $r/a < 0.5$ with respect to the
 12 effective boundary condition. In Fig. 10, it is shown that in the
 13 plasma core, the intrinsic rotation in the LOC and SOC phases
 14 develop in the opposite direction after the initial changes due
 15 to the neoclassical effect. In Fig. 11, the change of C rotation
 16 between the LOC and SOC phases shows a similar radial pro-
 17 file as the experiment.

18 Although the rotation change is consistent with the exper-
 19 iment, detailed profiles are different from the experimental
 20 observation. Figure 12 shows the time evolution of C rota-
 21 tion profiles and a comparison against the experiment. Here,
 22 the experimental profiles are shown by interpolated curves

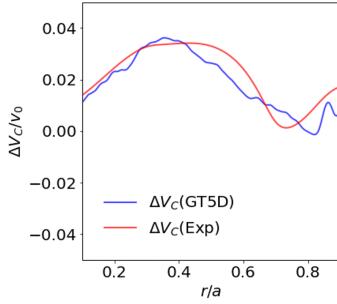


FIG. 11. The change of the carbon rotation between the LOC and SOC phases.

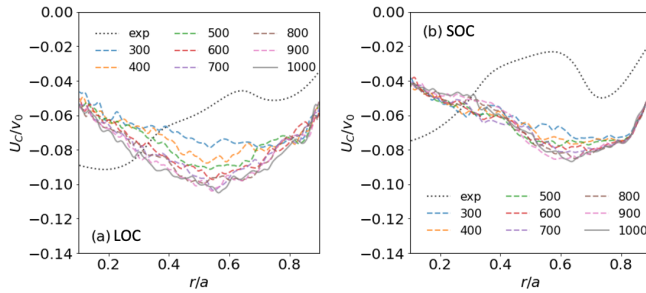


FIG. 12. The carbon rotation profiles at $tv_0/R = 300 \sim 1000$. Black broken curves show the carbon rotation profiles in the experiment. (a) and (b) respectively show the numerical experiments in the LOC and SOC phases.

1 for simplicity. Although the rotation profiles are still slowly
 2 evolving in time, qualitative features of profile shapes are un-
 3 changed. Both in the LOC and SOC phases, the C rotation
 4 profiles are shifted in the counter-current direction, because of
 5 the deviation of effective boundary conditions due to the neo-
 6 classical effect. In the LOC phase, the positive C rotation gra-
 7 dient U'_C in the outer radii (U'_C averaged over $r/a = 0.5 \sim 0.8$
 8 is $U'_C a/v_0 \sim 0.085$) is in a similar range as the experiment
 9 (U'_C averaged over $r/a = 0.3 \sim 0.9$ is $U'_C a/v_0 \sim 0.081$). In
 10 the SOC phase, the negative rotation gradient around the mid
 11 radius (U'_C averaged over $r/a = 0.4 \sim 0.6$ is $U'_C a/v_0 \sim -0.13$)
 12 is still weaker than the experiment (U'_C averaged over $r/a =$
 13 $0.6 \sim 0.7$ is $U'_C a/v_0 \sim -0.23$), and the position of negative ro-
 14 tation gradient is shifted from $r/a \sim 0.65$ to $r/a \sim 0.5$. Both
 15 in the LOC and SOC phases, the rotation gradients in the
 16 plasma core are opposite from the experiment, which may be
 17 attributed to the sawtooth activity inside the sawtooth inver-
 18 sion radius ($r/a \sim 0.25$)¹⁴.

19 In order to understand these intrinsic rotation profiles, the
 20 toroidal angular momentum balance (10) is analyzed. In 25
 21 Fig.13, the toroidal angular momentum balance is dominated 26
 22 by the radial current term and the toroidal field stress term, 27
 23 and the origin of the inertial term is not clearly seen. It is 28
 24 noted that the balance between the radial current term and the 29

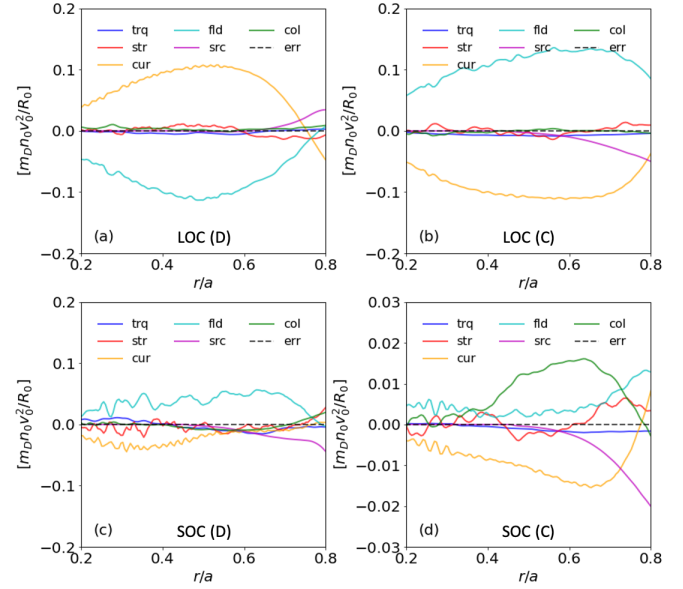


FIG. 13. The toroidal angular momentum balance, Eq. (10), is plotted for (a), (c) D and (b), (d) C. (a), (b) and (c), (d) respectively show the numerical experiments in the LOC and SOC phases. The radial profiles of the inertial term (trq, the first term), the stress term (str, the second term), the radial current term (cur, the third term), the toroidal field stress term (fld, the fourth term), the collision term (col, the fifth term), the source term (src, the sixth term), and the remaining error (err) are averaged over $tv_0/R = 0 \sim 1000$.

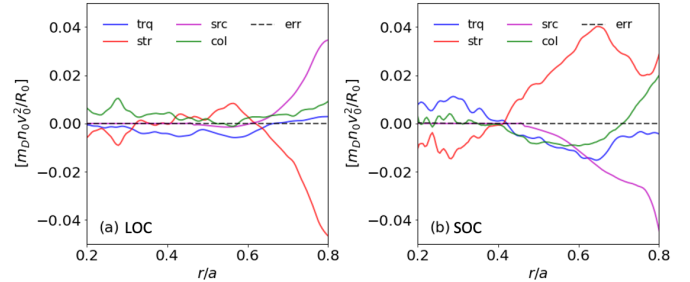


FIG. 14. The gyrokinetic equation part of the toroidal angular momentum balance, Eq. (11), is plotted for D. (a) and (b) respectively show the numerical experiments in the LOC and SOC phases. The radial profiles of the inertial term (trq, the first term), the partial stress term (str, the second term), the collision term (col, the third term), the source term (src, the fourth term), and the remaining error (err) are averaged over $tv_0/R = 0 \sim 1000$.

toroidal field stress term indicates that the ion mixing is pro-
 duced by the toroidal field stress term. In order to investigate
 the origin of the inertial term in detail, we remove the contri-
 butions from the radial current term and the toroidal field term
 by decomposing the toroidal angular momentum balance (10)

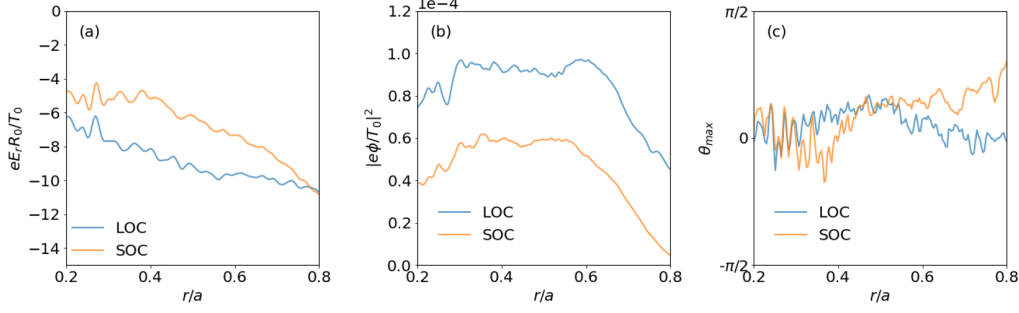


FIG. 15. Comparisons of radial profiles of (a) radial electric field E_r , (b) turbulent intensity $|e\phi/T_0|^2$, and (c) poloidal angle at the maximum turbulent intensity θ_{max} between the LOC and SOC phases. The profiles are averaged over $tv_0/R = 0 \sim 1000$.

1 into the gyrokinetic equation part,

$$\begin{aligned} & \left\langle \frac{\partial m_s v_{\parallel} b_{\phi} f_s}{\partial t} \right\rangle_{gf} + \left\langle \frac{m_s v_{\parallel} b_{\phi}}{\mathcal{J}} \frac{\partial}{\partial \mathbf{Z}} \cdot (\mathcal{J} \mathbf{Z} f_s) \right\rangle_{gf} \\ & - \left\langle m_s v_{\parallel} b_{\phi} \sum_{s'} C(f_s, f_{s'}) \right\rangle_{gf} \\ & - \left\langle m_s v_{\parallel} b_{\phi} (S_{src,s} + S_{snk,s}) \right\rangle_{gf} = 0, \end{aligned} \quad (11)$$

2 and the Hamilton's equation part,

$$\begin{aligned} & \left\langle \frac{f_s}{\mathcal{J}} \frac{\partial}{\partial \mathbf{Z}} \cdot (\mathcal{J} \mathbf{Z} m_s v_{\parallel} b_{\phi}) \right\rangle_{gf} \\ & - \left\langle \frac{q_s}{c} f_s \mathbf{R} \cdot \nabla \psi \right\rangle_{gf} + \left\langle f_s \frac{\partial \langle \phi \rangle / \alpha}{\partial \varphi} \right\rangle_{gf} = 0, \end{aligned} \quad (12)$$

3 following the derivation in Ref.32. Figure 14 shows the gyrokinetic equation part of the toroidal angular momentum balance, in which the inertial term is determined by the partial stress term, the collision term, and the source term. Except for the outer radii, where the source term becomes large, the inertial term is mainly correlated with the partial stress term, which is determined by plasma turbulence, while the collisional momentum exchange between D and C is non-negligible in the LOC phase. Therefore, we further investigate properties of turbulent momentum transport based on the mechanisms reviewed in Refs.57 and 58. It is noted that the source term develops in the opposite direction in the outer radii. In the numerical experiment, the momentum source is given by the Krook type sink models, $S_{snk,s1}$ and $S_{snk,s2}$, in which the sign of momentum source is determined in a passive manner by turbulent momentum transport. Therefore, turbulent momentum transport leads to the opposite net torque input between the LOC and SOC phases.

10 In the framework of local turbulence theory, the momentum flux Π_s may be decomposed as

$$\Pi_s = n_s m_s R_0 [-\chi_{\phi,s} R_0 \Omega'_{\phi,s} + V_{\phi,s} R_0 \Omega_{\phi,s} + C_s], \quad (13)$$

23 where $\Omega_{\phi,s}$ is the toroidal angular frequency, $\Omega'_{\phi,s} = \partial \Omega_{\phi,s} / \partial r$, $\chi_{\phi,s}$ is the momentum diffusion coefficient, $V_{\phi,s}$ is the pinch velocity, and C_s is the residual stress including

26 various symmetry breaking effects. Here, the residual stress comes mainly from the plasma shaping effect C_{PS} ⁵⁹, the profile shear C_{ρ^*} ⁴¹, the E_r shear C_E ^{60,61}, and the turbulence intensity gradient C_I ⁶². The symmetry breaking due to plasma shaping does not exist in the current circular concentric tokamak configuration. In Fig.15(a), the radial electric field E_r has negative shear on average, and the momentum transport due to the E_r shear stress $C_E \propto -dE_r/dr$ is expected to be outward both in the LOC and SOC phases. In Fig.15(b), the turbulent intensity $I = |e\phi/T_0|^2$ shows similar profiles, which have flat profiles for $r/a = 0.3 \sim 0.6$ and negative gradients for $r/a > 0.6$, and the momentum transport induced by the turbulent intensity shear stress $C_I \propto dI/dr$ seems to be similar between the LOC and SOC phases. It is noted that the turbulent intensity normalized by T_0 becomes smaller in the edge, where temperature becomes lower. However, $|e\phi/T_e|$ and $|\delta n_e/n_e|$ increases towards the edge. Finally, in Fig.15(c), an approximate ballooning angle θ_{max} , which is given by the poloidal angle at the maximum turbulent intensity within each magnetic surface, shows qualitative difference between the LOC and SOC phases. In the outer radii, θ_{max} in the LOC phase is slightly negative, while θ_{max} in the SOC phase is positive. The sign relation of the profile shear stress $C_{\rho^*} \propto -\theta_0 \sim \theta_{max}$ is consistent with the change of intrinsic rotation between the LOC and SOC phases.

The normalized profile shear stress is given as⁴¹

$$\frac{R_0 C_{\rho^*,s}}{v_{ts} \chi_{\phi,s}} = -\frac{\hat{s} \theta_0}{2|q|k_{\theta} \rho_{ts}} \left(\frac{R_0}{L_{n,s}} + 4 - \frac{1}{(\sqrt{2}qk_{\theta} \rho_{ts})^2} \right), \quad (14)$$

where θ_0 is the ballooning angle, \hat{s} is the magnetic shear, k_{θ} is the poloidal wavenumber, and $L_{n,s} = |n_s / \nabla n_s|$ is the density gradient scale length. Here, the mode asymmetry θ_0 and the corresponding linear growth rate γ are given by the higher order ballooning theory^{52,53} as,

$$\theta_0 = -\text{sign}(\hat{s} \omega'_r) \left| \frac{\omega'_r}{2k_{\theta} \gamma_0 \hat{s}} \right|^{1/3}, \quad (15)$$

$$\gamma = \gamma_0 \cos(\theta_0), \quad (16)$$

where γ_0 is the linear growth rate without mode asymmetry, ω_r is the mode frequency, and $\omega'_r = \partial \omega_r / \partial r$. Although Ref.41

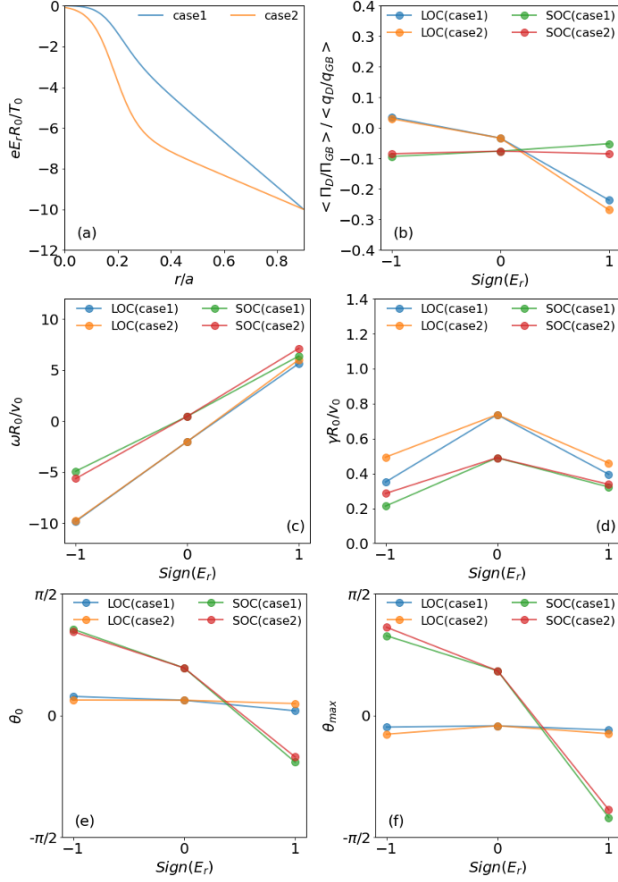


FIG. 16. Linear calculations of $n = 25$ modes of TEM in the LOC phase and ITG in the SOC phase with externally imposed positive and negative E_r . (a) shows the radial profiles of E_r used in case 1 and case 2. (b) shows the volume averaged quasilinear flux ratio between turbulent heat and momentum transport. The heat and momentum fluxes are respectively normalized by $q_{GB} = \chi_{GB} n_0 T_0 / a$ and $\Pi_{GB} = \chi_{GB} n_0 v_0 / a$, where $\chi_{GB} = (v_0^3 / \Omega_D^2) / a$. (c) and (d) show the frequency and the growth rate, respectively. (e) and (f) shows the ballooning angle θ_0 , which is defined at the minimum radial wavenumber k_r , and an approximate ballooning angle, which is given at the maximum amplitude of ϕ , respectively.

discussed the profile shear stress without E_r , the higher order ballooning theory can be naturally extended including the E_r shear^{63,64}. Therefore, the profile shear stress involves the E_r shear stress, when the mode asymmetry is affect also by E_r . In Ref.45, the profile shear stress including E_r was analyzed for ITG with adiabatic electrons using linear global calculations with externally imposed E_r , and the relation between the E_r shear and the profile shear stress was clarified.

In this work, we apply this approach to TEM in the LOC phase and ITG in the SOC phase to understand the change of θ_0 between the LOC and SOC phases. Based on the E_r profiles in Fig.15(a), model E_r profiles are chosen as $eE_r R_0 / T_0 = sign(E_r)(c_0 + c_1 r / a_0)(\tanh[(r/a_0 - 0.2)/0.1] + 1)/2$ with $(c_0, c_1) = (0, 10)$ (case 1) and $(5, 5)$ (case 2), in 34

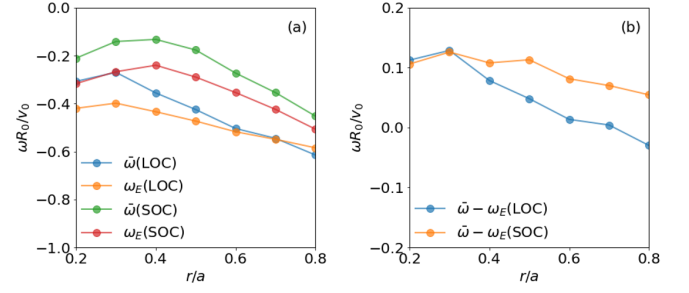


FIG. 17. (a) shows the weighted turbulence frequency $\tilde{\omega}$ and the Doppler shift due to $E \times B$ rotation ω_E of $n = 24$ modes at each radius. (b) shows of the frequency of $n = 24$ modes corrected by subtracting the Doppler shift, $\tilde{\omega} - \omega_E$. The frequency spectra are obtained from the time series data for $t\nu_0/R = 600 \sim 1000$.

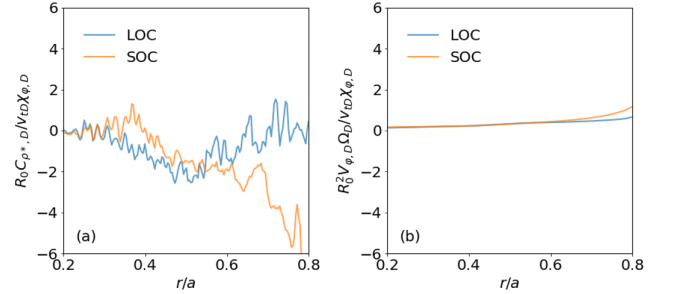


FIG. 18. Radial profiles of (a) the normalized profile shear residual stress $R_0 C_{\rho^*,D} / v_{iD} \chi_{\phi,D}$ and (b) the normalized curvature pinch $R_0^2 V_{\phi,D} \Omega_D / v_{iD} \chi_{\phi,D}$. The plasma profiles are given by the simulation data in Fig.8. In estimating the profile shear stress, $k_{\theta} \rho_D = 0.35$ is assumed, and θ_0 is given by θ_{max} in Fig.15(c).

which E_r at the outer radii becomes similar but the E_r shear differ by two times (see Fig.16(a)). In Fig.16, (b) the volume averaged quasi-linear flux ratio between turbulent heat and momentum transport, (c) the frequency, (d) the growth rate, (e) the ballooning angle θ_0 defined by the poloidal angle at the minimum radial wavenumber k_r , and (f) the approximate ballooning angle θ_{max} defined by the poloidal angle at the maximum amplitude of ϕ are shown for case 1 and case 2 with negative and positive E_r . It is noted that linear calculations with E_r show oscillatory linear growth, which is a typical feature of the so-called Floquet ballooning mode⁶⁵, and the above results are time averaged. In Fig.16(c), the frequencies of both TEM and ITG are affected by the Doppler shift, and become the same sign as in the numerical experiment. Here, case 1 and case 2 give almost the same frequencies, because they have similar E_r . In Fig.16(e), the changes of θ_0 are similar between case 1 and case 2, while the theory predicts $\theta_0 \propto \omega_r^{1/3}$. This result suggests that θ_0 is determined not only by the E_r shear but also by E_r itself. This feature was observed also in Refs.45 and 64. Accordingly, in Fig.16(d), the

1 stabilization effect due to E_r becomes similar between case 1
 2 and case 2. Another important finding is that TEM in the LOC
 3 phase shows much smaller change of θ_0 than ITG in the SOC
 4 phase. This may be attributed to the higher linear growth rate
 5 and the narrower mode width of TEM in the LOC phase. In
 6 Fig.16(f), TEM in the LOC phase and ITG in the SOC phase
 7 respectively show slightly negative θ_{max} and positive θ_{max} at
 8 the negative E_r , which is consistent with the numerical exper-
 9 iment (see Fig.15(c)). In Fig.16(b), the resulting momentum
 10 transport at the negative E_r becomes outward and inward in
 11 the LOC and SOC phases, respectively.

12 In Fig.17(a), the weighted turbulence frequency $\bar{\omega}(r, n) =$
 13 $\sum_{\omega} \omega |\hat{\phi}_{n, \omega}|^2 / \sum_{\omega} |\hat{\phi}_{n, \omega}|^2$ of $n = 24$ mode at each radius is in
 14 the electron diamagnetic direction over the whole radii, where
 15 $\hat{\phi}_{n, \omega}$ is the Fourier component of ϕ at $\theta = 0$ with respect
 16 to the toroidal mode number n and the frequency ω . How-
 17 ever, if one corrects the frequency by subtracting the Doppler
 18 shift of poloidal $E \times B$ rotation due to the radial electric field
 19 $\omega_E = k_{\theta} v_{E\theta} = \langle \mathbf{v}_E \cdot \nabla \theta / |\nabla \theta| \rangle_f$. The SOC
 20 phase is characterized by the positive frequency or ITG over
 21 the whole radii. On the other hand, the LOC phase shows ITG
 22 only in the plasma core, and the outer radii are character-
 23 ized by the negative frequency or TEM. These results show that
 24 the different features of the mode asymmetry between TEM
 25 in the LOC phase and ITG in the SOC phase, which depend
 26 not only on the profile shear and the E_r shear but also on E_r
 27 itself, are key physics in understanding the change of intrinsic
 28 rotation during the LOC-SOC transition.

29 Finally, we compare the relative magnitude of the profile
 30 shear stress against the Coriolis pinch given as⁵⁷,
 31

$$\frac{R_0^2 V_{\phi, s} \Omega_s}{v_{ts} \chi_{\phi, s}} = \left(-\frac{4T_s}{T_e} - \frac{R_0}{L_n} \right) \frac{R_0 \Omega_s}{v_{ts}}. \quad (17)$$

32 Figure 18 shows (a) the normalized profile shear stress and (b)
 33 the normalized Coriolis pinch, where the plasma profiles are
 34 given by the simulation data in Fig.8. In estimating the profile
 35 shear stress, $k_{\theta} \rho_{iD} = 0.35$ is assumed and the mode asymme-
 36 try is given by θ_{max} in Fig.15(c). The Coriolis pinch is out-
 37 ward both in the LOC and SOC phases, and is negligible for
 38 the intrinsic rotation with the low Mach number. The negligi-
 39 ble impact of the Coriolis pinch was also shown in Ref. 20.

40 C. Energy transport

41 In Figs.8(c) and 8(f), T_e in the plasma core is slightly lower
 42 than the experiment, while T_D and T_C are slightly higher than
 43 the experiment. As already discussed, the increased collision-
 44 ality leads to stronger ion-electron coupling, and the differ-
 45 ence between T_e and T_D in the plasma core becomes smaller
 46 than the experiment. Figure 19 shows the surface integrated
 47 energy flux $Q_s = \hat{Q}_s S$ observed in the quasi-steady phase of
 48 the numerical experiments, where $\hat{Q}_s \equiv \langle \int [m_s v_{\parallel}^2 + \mu B] f_s (\mathbf{v}_D +$
 49 $\mathbf{v}_E) \cdot \nabla r d^3 v \rangle_f$ is the energy flux density and $S = 2\pi r R_0$. In the
 50 current numerical experiments, the deposition profiles of the
 51 ohmic heating and the radiation loss overlap with each other

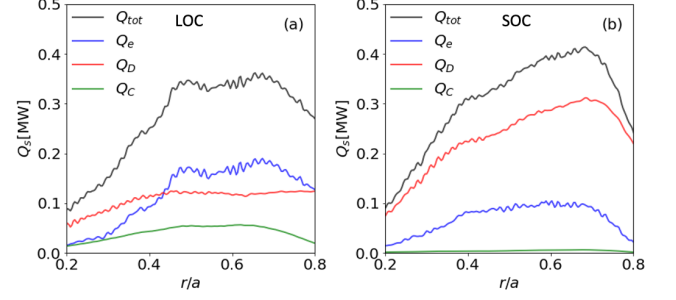


FIG. 19. Radial profiles of the surface integrated energy flux Q_s averaged over $t\nu_0/R = 900 \sim 1000$. (a) and (b) respectively show the numerical experiments in the LOC and SOC phases.

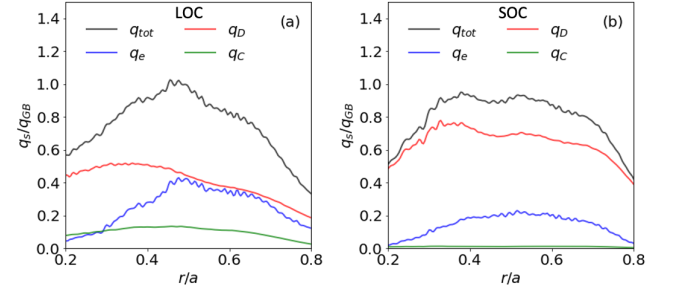


FIG. 20. Radial profiles of the heat flux density q_s averaged over $t\nu_0/R = 900 \sim 1000$. (a) and (b) respectively show the numerical experiments in the LOC and SOC phases.

(see Fig.3), and a source free region does not exist. Therefore,
 both in the LOC and SOC phases, the total energy flux slightly
 exceeds the input power of $P_{in} = P_0/2 = 300\text{kW}$, while the
 stored energy is in the quasi-steady state (see Fig.7(a)). In the
 LOC phase, the electron energy flux is dominant, while both
 the D and C energy fluxes make significant contributions. On
 the other hand, in the SOC phase, the D energy flux is dom-
 inant, and the C energy flux is negligible. These results are
 qualitatively consistent with the experiment and the local gy-
 rokinetic analyses in Ref.23, while several differences are not
 captured in the current numerical experiments.

Figure 20 shows the heat flux density $q_s = \hat{Q}_s - 5/2T_s\Gamma_s$,
 where $\Gamma_s \equiv \langle \int f_s (\mathbf{v}_D + \mathbf{v}_E) \cdot \nabla r d^3 v \rangle_f$ is the particle flux
 density. In Ref.23, the power balance analysis at $r/a =$
 0.37 gave the ion and electron heat fluxes of $(q_i, q_e) =$
 $(4.5 \pm 1.0\text{kWm}^{-2}, 6.7 \pm 1.0\text{kWm}^{-2})$ and $(q_i, q_e) =$
 $(14.0 \pm 3.0\text{kWm}^{-2}, -1.0 \pm 3.0\text{kWm}^{-2})$ in the LOC and SOC
 phases, respectively. However, in the LOC phase shown in
 Fig.20(a), the ion heat fluxes $q_D + q_C$ exceeds q_e at
 $r/a = 0.37$. In the SOC phase shown in Fig.20(b), we do
 not observe the electron heat pinch over the whole radii. The
 detailed ratios between the ion and electron energy fluxes
 may be affected either by the enhanced ion-electron coupling
 and by the modelling of sources and sinks. On the other hand,
 the mechanism of the

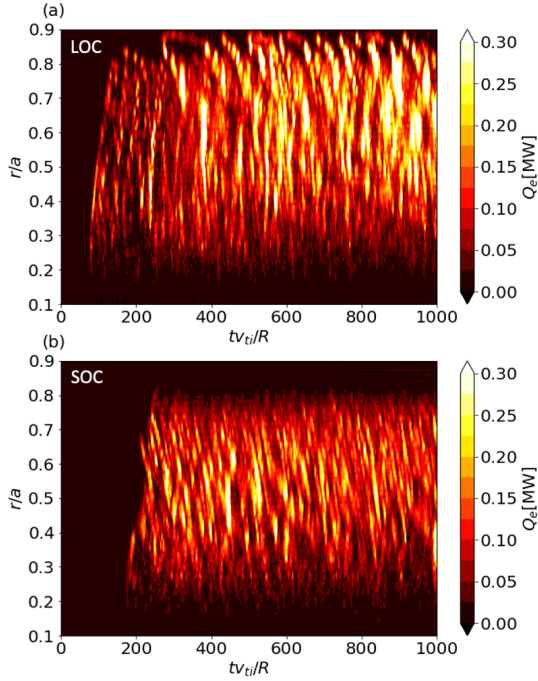


FIG. 21. Spatio-temporal evolution of the electron energy flux in (a) LOC and (b) SOC phases.

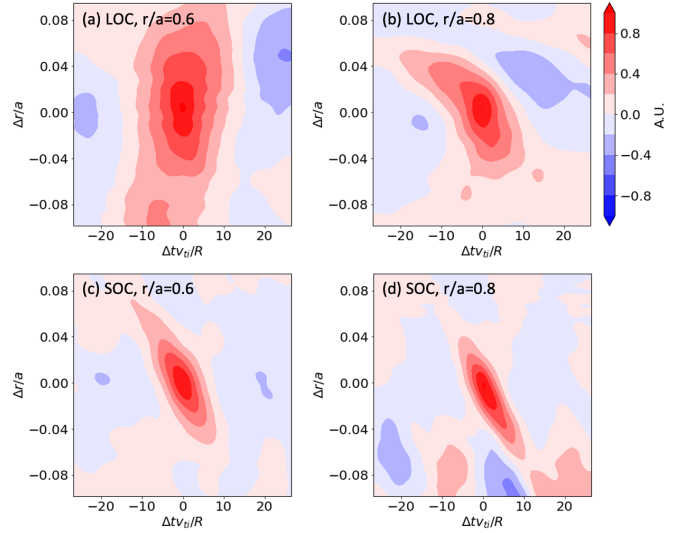


FIG. 22. Spatio-temporal auto-correlation functions of the electron energy flux evaluated for (a) LOC phase, $r/a = 0.6$, (b) LOC phase, $r/a = 0.8$, (c) SOC phase, $r/a = 0.6$, and (d) SOC phase, $r/a = 0.8$. Δt and Δr respectively show delays in time and radius. The electron energy flux data in Fig.21 is processed for $tv_{0}/R = 400 \sim 1000$.

1 electron heat pinch is still an open issue. It is noted that in
 2 Ref.66, off diagonal terms in the transport matrix was derived
 3 for the curvature pinches, where an Onsager symmet-
 4 rical contribution to the heat flux gives heat pinch when off-
 5 diagonal particle pinch occurs.

6 Figures 21 and 22 show the spatio-temporal evolution of the
 7 electron energy flux and the corresponding spatio-temporal
 8 auto-correlation functions. Both the LOC and SOC phases are
 9 characterized by bursts of avalanche like nonlocal transport,
 10 while their dynamics is somewhat different. The SOC phase
 11 is characterized by ballistic inward propagation of global
 12 avalanches from the edge to the core. This feature is clearly
 13 seen by the spatio-temporal autocorrelation functions in Figs.
 14 22(c) and 22(d). In Ref.36, it was shown that the direc-
 15 tion of the avalanche propagation in ITG turbulence is de-
 16 termined by the asymmetry induced by the E_r shear, which
 17 results in the inward and outward propagation with the neg-
 18 ative and positive E_r shear, respectively. The inward propa-
 19 gation of avalanches is consistent with the negative E_r shear
 20 in Fig.15(a). On the other hand in the LOC phase, the prop-
 21 agation width is limited, and avalanches show discontinuities
 22 around $r/a \sim 0.7$. In Fig.22 (b), the spatio-temporal au-
 23 tocorrelation function at $r/a \sim 0.8$ shows that inward propa-
 24 gation of avalanches is limited and the correlation length be-
 25 comes shorter inside $r/a \sim 0.8$, while in Fig.22(a), the spatio-
 26 temporal autocorrelation function at $r/a \sim 0.6$ suggest global
 27 mode structures with small time delay. This may be attributed
 28 to the change of dominant turbulence drive from ITG to TEM,
 29 which was shown in Fig.17(b).

D. Quasi-coherent modes

Figures 23(a) and 23(c) show the frequency spectra of the electrostatic potential measured at $\theta = 0, \phi = 0$, and $r/a = 0.7$. Since the spectrum is symmetric with respect to the sign of frequency, only the negative frequency part is plotted. Here, $r/a = 0.7$ is characterized by TEM and ITG in the LOC and SOC phases, respectively (see Fig.17(b)). The spectrum in the LOC phase shows a few quasi-coherent peaks, which were called as the quasi-coherent modes (QCMs). On the other hand, the SOC phase is characterized by a broad-band spectrum without such coherent peaks. These spectra are qualitatively consistent with those in the experiment²⁸, while the peak positions of QCMs are different. In the experiment, the QCM peak appeared in the 50kHz range, which correspond to $\omega R/v_0 \sim 0.5$. Although the secondary peak with $n = 18$ is in this range, the primary peak appear in the lower frequency range. This difference may be attributed to the magnetic ripples, which can enhance E_r and the resulting Doppler shift⁶⁹.

Figures 23(b) and 23(d) show the frequency spectrum of each n component measured at $\theta = 0$ and $r/a = 0.7$, which is non-symmetric with respect to the sign of frequency because of the resonance condition depending on n . It is noted that $\bar{\omega}$ in Fig.17 is estimated by the weighted average using the $n = 24$ spectrum. The n spectra clearly show the structures of the QCMs and the broad-band fluctuations, which have the following features. Firstly, the width of each n spectrum is narrower in the LOC phase. The half height width of the $n = 12$ spectrum is estimated as $\sim 0.035v_0/R$ and $\sim 0.065v_0/R$ in the LOC and SOC phases, respectively. This may be attributed to the difference of nonlinear spectral broad-

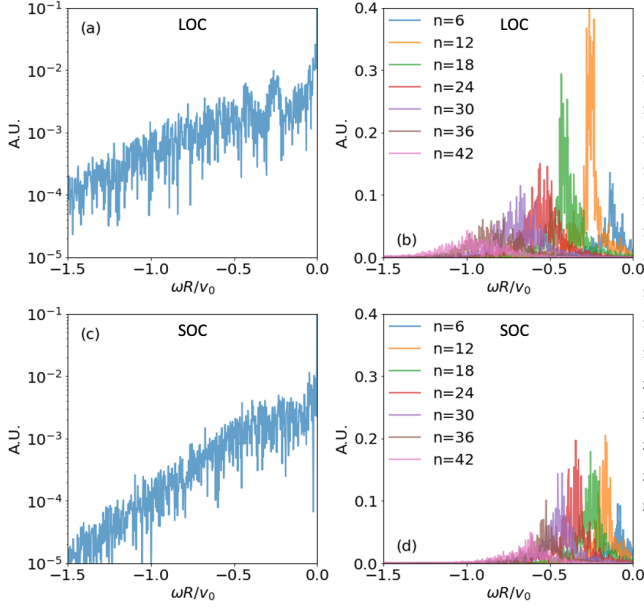


FIG. 23. Frequency spectra of (a),(c) the electrostatic potential measured at $(r, \theta, \varphi) = (0.7a, 0, 0)$, $|\hat{\phi}_\omega|$, and (b),(d) the toroidal mode number component of the electrostatic potential measured at $(r, \theta) = (0.7a, 0)$, $|\hat{\phi}_{\omega,n}|$. (a),(b) and (c),(d) respectively show the numerical experiments in the LOC and SOC phases. The frequency spectra are obtained from the time series data for $t\nu_0/R = 600 \sim 1000$.

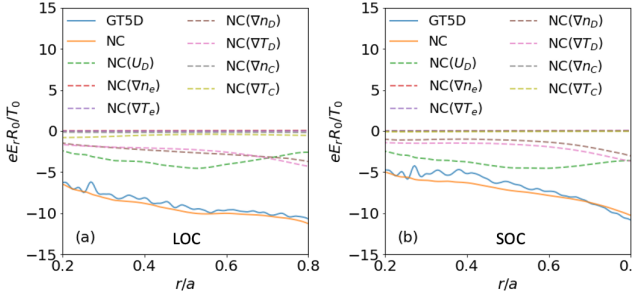


FIG. 24. Radial profiles of E_r observed in the numerical experiment (GT5D) and E_r estimated by the parallel momentum balance based on the Hirshman-Sigmar's moment approach (NC)^{67,68}. The D component of the parallel momentum balance is computed using the D parallel flow U_D , the density n_s , and the temperature T_s from GT5D, and terms proportional to U_D and the gradients of n_s and T_s are also shown. Both E_r and the plasma profiles are averaged over $t\nu_0/R = 0 \sim 1000$. (a) and (b) respectively show the numerical experiments in the LOC and SOC phases.

ening between TEM and ITG. Secondly, the frequency gap of peaks in the n spectra is larger in the LOC phase. The mode frequency with the Doppler shift is given by $\omega_{shift} = \omega + k_\theta v_{E\theta} \sim (nq/r)(v_{ph} + v_{E\theta})$, and the frequency gap between adjacent n modes becomes $\Delta\omega_{shift} = (q/r)(v_{ph} + v_{E\theta})$, where $v_{ph} = \omega/k_\theta$ is the phase velocity. As shown in Fig.17(a),

both in the LOC and SOC phases, the Doppler shift is in the electron diamagnetic direction, $v_{E\theta} < 0$, and the frequencies of TEM and ITG are shifted in the electron diamagnetic direction. On the other hand, the phase velocity of TEM and ITG are in the electron diamagnetic direction, $v_{ph} < 0$, and in the ion diamagnetic direction, $v_{ph} > 0$, respectively. In addition, the magnitude of $v_{E\theta}$ becomes larger in the LOC phase. Therefore, the frequency gap becomes larger in the LOC phase. Because of these two features, the n spectra in the LOC phase split to form the QCMs. It is noted that in Ref.23, the former effect was shown, while the latter effect was not taken into account, because the simulation was conducted using a local δf gyrokinetic model without E_r .

In Fig.24, E_r observed in the numerical experiment is compared against the neoclassical theory. By substituting the D parallel flow U_D , the density n_s , and the temperature T_s from GT5D for the D component of the parallel momentum balance (B7), E_r is written as

$$E_r = -\frac{1}{A_{21} + A_{22} + A_{23}} \left[\frac{|\nabla\psi|}{F} \langle BU_D \rangle_f - A_{21} \frac{T_e}{q_e} \frac{\partial \ln p_e}{\partial r} - A_{22} \frac{T_D}{q_D} \frac{\partial \ln p_D}{\partial r} - A_{23} \frac{T_C}{q_C} \frac{\partial \ln p_C}{\partial r} + A_{24} \frac{T_e}{q_e} \frac{\partial \ln T_e}{\partial r} + A_{25} \frac{T_D}{q_D} \frac{\partial \ln T_D}{\partial r} + A_{26} \frac{T_C}{q_C} \frac{\partial \ln T_C}{\partial r} \right], \quad (18)$$

where F is the toroidal flux, $p_s = n_s T_s$ and A_{ij} is given by elements of the matrix $\mathbf{A} = (\mathbf{L} - \mathbf{M})^{-1} \cdot \mathbf{M}$ (see Appendix B for detailed definitions). In Fig.24, the dominant contributions come from the parallel flow, the density gradient of D, and the temperature gradient of D. It is noted that because of the charge dependency, the contribution from C is negligible even in the LOC phase. The contributions from the density and temperature gradients are larger in the LOC phase, reflecting the higher temperature. On the other hand, the largest contribution comes from the counter-current rotation, and its magnitude is comparable between the LOC and SOC phases, because the boundary condition or the baseline rotation is determined by the toroidal ripple. The large negative E_r induced by the ripple induced counter-current rotation is a unique feature on Tore Supra^{69,70}.

V. SUMMARY

In this work, we analyzed Tore Supra ohmic L-mode discharge 48102 focusing on two time slices in the LOC and SOC phases using the global full- f gyrokinetic Eulerian code GT5D. Because of limited computational resources, the numerical experiments were conducted using the modified 1/2 scale parameters, in which the plasma size is reduced by a half, while the collisionality is enhanced to keep the experimental value of v_e^* . The modified 1/2 scale parameters keep the stability of TEM in the original experimental parameters. However, collisional ion-electron coupling becomes stronger, leading to smaller temperature difference between ions and electrons. The heating condition was chosen based on power scans using the modified 1/2 and 1/3 scale parameters. The

1 power scan numerical experiments showed the Bohm like 53
 2 scaling both in the LOC and SOC phases, and the Bohm like
 3 heating condition, in which the input power is scaled as pro- 54
 4 portional to the plasma size, was adopted in the numerical ex- 55
 5 periments. The Bohm like scaling is consistent with the L-
 6 mode scaling observed in the SOC phase, while its relevance
 7 in the LOC phase, which is characterized by the neo-Alcator
 8 scaling, is still an open issue. 56

9 The numerical experiments captured qualitative features 57
 10 of transport phenomena in the LOC-SOC transition on Tore 58
 11 Supra. Firstly, the exhaust of C is induced by the ion mix- 59
 12 ing, which is driven by the toroidal field stress. Secondly, the 60
 13 intrinsic rotation develops in the opposite directions between
 14 the LOC and SOC phases, which is characterized by the dif- 61
 15 ferent features of TEM in the LOC phase and ITG in the SOC
 16 phase, leading to the change of the profile shear stress. Here, 62
 17 the mode asymmetry or the ballooning angle depends not only
 18 on the profile shear and the E_r shear but also on E_r itself. 63
 19 It is noted that although the rotation change is reproduced, 64
 20 detailed intrinsic rotation profiles are still different from the 65
 21 experiment, and further investigations on momentum source 66
 22 and sink models are needed. 67
 23 Thirdly, the ratio of the stored energy between two time slices and the temperature profiles 68
 24 in the experimental data are approximately recovered in the 69
 25 quasi-steady state, where the energy fluxes of electrons and D 70
 26 are dominant in the LOC and SOC phases, respectively. 71
 27 Finally, turbulent spectra in the LOC and SOC phases are char- 72
 28 acterized by QCMs and broad-band spectra, which are pro- 73
 29 duced by the difference of Doppler shift in the LOC and SOC 74
 30 phases. The ripple induced counter-current rotation produce 75
 31 the Doppler shift in the electron diamagnetic direction, which 76
 32 enhances the frequency gap of each n spectrum, leading to 77
 33 QCMs. 78
 34

35 As already discussed, further studies are required to under- 82
 36 stand ρ^* scaling in the LOC phase. Although Tore Supra 83
 37 showed the change of intrinsic rotation in the co-current di-
 38 rection during the LOC-SOC transition, many other devices 84
 39 showed the change of intrinsic rotation in the opposite direc-
 40 tion². Another important issue is the change of nonlocal trans- 85
 41 port between the LOC and SOC phases, which was experi- 86
 42 mentally observed by applying edge cold pulses². In future 87
 43 work, these important issues will be addressed using global
 full- f gyrokinetic simulations. 88

44 ACKNOWLEDGMENTS

45 This work was supported by the MEXT Japan (Program for 93
 46 Promoting Researches on the Supercomputer Fugaku "Explora- 94
 47 tion of Burning Plasma Confinement Physics", and Grant 95
 48 No.22K03584) and the NIFS Collaborative Research Program 96
 49 (NIFS16KNST103). This research used computational re-
 50 sources of the supercomputer Fugaku provided by the RIKEN
 51 Center for Computational Sciences (hp210178,hp220165), 95
 52 and the ICEX provided by the JAEA. 96

DATA AVAILABILITY STATEMENT

The data that support the findings of this study are available from the corresponding author upon reasonable request.

Appendix A: Wedge size convergence

In this appendix, we verify the convergence of the numerical experiments with respect to the wedge size. We repeat the numerical experiments with the 1/3 wedge torus model, and compare plasma profiles and turbulent frequency spectra. In Fig.25, we compare radial profiles of the density, the parallel flow, and the temperature obtained from the numerical experiments with the 1/3 and 1/6 wedge torus models, which show negligible difference between these models. This indicates the corresponding particle, momentum, and energy transport transport is also converged between the 1/3 and 1/6 wedge torus models. Figure 26 shows the frequency spectra of the electrostatic potential measured in the numerical experiments with the 1/3 wedge torus model. Here, the definitions of plots are the same as those in Fig.23. Figures 26(a) and 26(c) respectively show the QCMs in the LOC phase and the broad-band spectrum in the SOC phase as in Fig.23. However, the structures of QCMs are slightly different from the 1/6 wedge torus model. In the 1/6 wedge torus model, the primary and secondary peaks are formed by adjacent toroidal modes, $n = 12$ and $n = 18$. In the 1/3 wedge torus model, the toroidal mode numbers of the primary and secondary peaks are unchanged, while there exists a subdominant mode with $n = 15$ in between them. Therefore, in addition to the enhancement of the frequency gap between adjacent toroidal modes due to the Doppler shift, some nonlinear coupling mechanism to selectively excite $n = 12$ and $n = 18$ modes exists. This issue will be addressed in future works.

Appendix B: Parallel flows of bulk and impurity ions

In this appendix, we verify the difference of parallel flows of bulk and impurity ions based on the neoclassical parallel momentum balance equations. In the neoclassical benchmark, we consider multi-species plasmas with kinetic electrons, D ions, and C impurity ions in a circular concentric tokamak configuration with $R_0/a = 5$, $a/\rho_{tD} = 150$, and $q(r) = 0.85 + 2.18(r/a)^2$, $m_D/m_e = 3,672$, and $Z_{eff} = 4$. The initial density and temperature profiles are given as

$$n_s(r) = n_{0s} \exp \left[-\frac{\Delta_{ns}}{L_{ns}} \tanh \left(\frac{r-r_0}{\Delta_{ns}} \right) \right], \quad (B1)$$

$$T_s(r) = T_{0s} \exp \left[-\frac{\Delta_{Ts}}{L_{Ts}} \tanh \left(\frac{r-r_0}{\Delta_{Ts}} \right) \right], \quad (B2)$$

where $R_0/L_{ns} = R_0/L_{Ts} = 2.22$, $\Delta_{ns} = \Delta_{Ts} = 0.3a$, and $r_0 = 0.5a$. The normalized collisionality $v_{s,s'}^* \equiv qR/(\epsilon^{3/2}v_{Ts}\tau_{s,s'})$ at the mid-radius is given as $v_{e,D}^* = 0.08$, $v_{D,D}^* = 0.08$, and $v_{C,D}^* = 2.49$, where $\tau_{s,s'}$ is the collision time. The benchmark

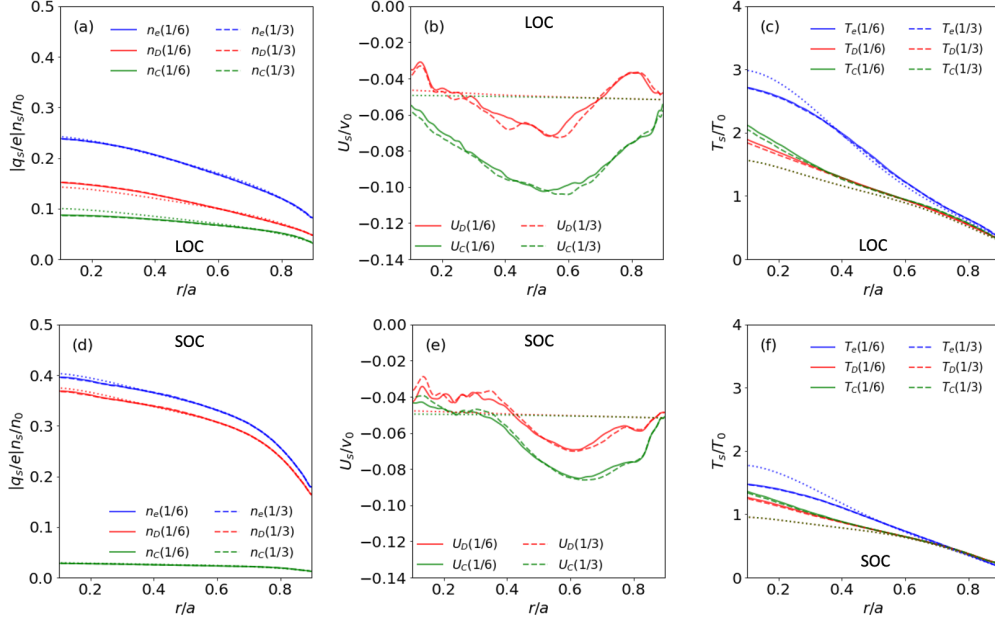


FIG. 25. Radial profiles of (a),(d) the density, (b),(e) the parallel flow, and (c),(f) the temperature averaged over $t\nu_0/R = 900 \sim 1000$. (a)-(c) and (d)-(f) respectively show the numerical experiments in the LOC and SOC phases. Solid and dashed curves respectively show the results with the 1/6 and 1/3 wedge torus models, and dotted curves show the initial profiles.

1 calculation is performed in the axisymmetric limit on 4D grids ¹⁴
² $(N_R, N_\zeta, N_Z, N_{v\parallel}, N_{v\perp}) = (160, 1, 160, 96, 24)$. ¹⁵

³ Figure 27 shows a comparison of parallel currents between ¹⁶
⁴ GT5D and the neoclassical theory^{67,68}. Here, the neoclassical ¹⁷
⁵ results are computed using the Matrix Inversion code⁶⁸, which ¹⁸
⁶ computes the parallel components of the momentum and heat ¹⁹
⁷ flow balance equations ²⁰

$$\sup>21
$$[\mathbf{M} - \mathbf{L}]\vec{X} = \mathbf{M}\vec{V}. \quad (\text{B3}) \sup>22$$$$

⁸ Here, \mathbf{L} and \mathbf{M} are matrices given by the parallel friction co-
⁹ efficients and the viscosity coefficients, respectively. The par- ²³
¹⁰ allel components of flow and heat flux \vec{X} and the thermody- ²⁴
¹¹ namic force \vec{V} are respectively given as ²⁵

$$\sup>26
$$\vec{X} = \begin{bmatrix} \langle \mathbf{B} \cdot \mathbf{U}_e \rangle \\ \langle \mathbf{B} \cdot \mathbf{U}_D \rangle \\ \langle \mathbf{B} \cdot \mathbf{U}_C \rangle \\ 2\langle \mathbf{B} \cdot \mathbf{q}_e \rangle / (5p_e) \\ 2\langle \mathbf{B} \cdot \mathbf{q}_D \rangle / (5p_D) \\ 2\langle \mathbf{B} \cdot \mathbf{q}_C \rangle / (5p_C) \end{bmatrix}, \quad \vec{V} = \begin{bmatrix} V_{1e} \\ V_{1D} \\ V_{1C} \\ V_{2e} \\ V_{2D} \\ V_{2C} \end{bmatrix}, \quad (\text{B4}) \sup>30$$$$

¹² where the thermodynamic force is given as

$$\sup>31
$$V_{1s} = -\frac{F}{B\nabla\psi} \left[\frac{d\phi}{dr} + \frac{1}{q_s n_s} \frac{dp_s}{dr} \right], \quad (\text{B5}) \sup>32$$$$

$$\sup>33
$$V_{2s} = -\frac{F}{B\nabla\psi} \frac{dT_s}{dr}. \quad (\text{B6}) \sup>34$$$$

¹³ By inverting Eq.(B3) as

$$\sup>35
$$\vec{X} = A\vec{V}, \quad (\text{B7}) \sup>36$$$$

and substituting ϕ , n_s , and T_s from GT5D, one can estimate
parallel flows for each species, where $A = [\mathbf{L} - \mathbf{M}]^{-1} \cdot \mathbf{M}$.
Also, the bootstrap current is estimated as $\langle \mathbf{B} \cdot \mathbf{J} \rangle = \sum_s q_s n_s \langle \mathbf{B} \cdot \mathbf{U}_s \rangle$.
The benchmark result shows good quantitative agree-
ments, indicating that GT5D can correctly compute the differ-
ence of ion parallel flows in multi-species plasmas. In the cur-
rent benchmark case involving a high fraction of C impurity
ions, the C rotation develops in the counter-current direction
with respect to D rotation.

¹ITER Physics Expert Groups on Confinement and Transport and Confinement Modelling and Database, I. P. B. Editors, and I. EDA, “Chapter 2: Plasma confinement and transport,” Nucl. Fusion **39**, 2175–2249 (1999).

²J. Rice, J. Citrin, N. Cao, P. Diamond, M. Greenwald, and B. Grierson, “Understanding LOC/SOC phenomenology in tokamaks,” Nucl. Fusion **60**, 105001 (2020).

³X. Garbet, J. Payan, C. Laviron, P. Devynck, S. K. Saha, H. Capes, X. P. Chen, J. P. Coulon, C. Gil, and G. R. Harris, “Turbulence and energy confinement in TORE SUPRA Ohmic discharges,” Nuclear fusion **32**, 2147 (1992).

⁴J. E. Rice, C. Gao, M. L. Reinke, P. H. Diamond, N. T. Howard, H. J. Sun, I. Cziegler, A. E. Hubbard, Y. A. Podpaly, and W. L. Rowan, “Non-local heat transport, rotation reversals and up/down impurity density asymmetries in Alcator C-Mod ohmic L-mode plasmas,” Nuclear Fusion **53**, 033004 (2013).

⁵A. Krämer-Flecken, V. Dreval, S. Soldatov, A. Rogister, and V. Vershkov, “Turbulence studies with means of reflectometry at TEXTOR,” Nuclear fusion **44**, 1143 (2004).

⁶H. Arnichand, R. Sabot, S. Hacquin, A. Krämer-Flecken, X. Garbet, J. Citrin, C. Bourdelle, G. Hornung, J. Bernardo, C. Bottereau, F. Clairet, G. Falchetto, and J. Giacalone, “Quasi-coherent modes and electron-driven turbulence,” Nucl. Fusion **54**, 123017 (2014).

⁷H. Arnichand, R. Sabot, S. Hacquin, A. Krämer-Flecken, C. Bourdelle, J. Citrin, X. Garbet, J. C. Giacalone, R. Guirlet, and J. C. Hillesheim, “Dis-

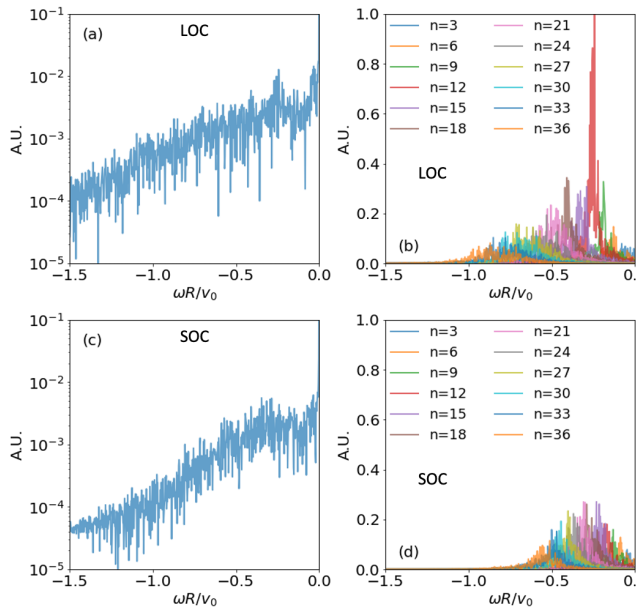


FIG. 26. Frequency spectra of the electrostatic potential in the numerical experiments with the 1/3 wedge torus model. The definitions of plots are the same as those in Fig.23.

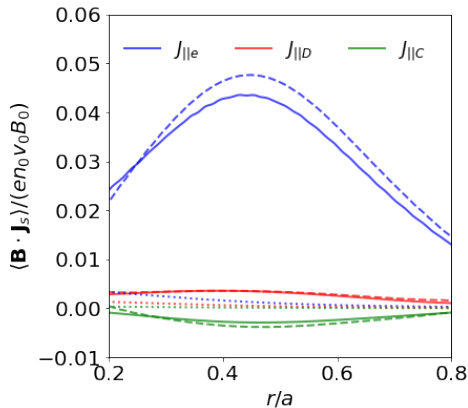


FIG. 27. Radial profiles of the parallel current of electrons, deuterium ions, and carbon impurity ions. Broken lines show the initial condition given by non-shifted Maxwellian distributions. Solid and dashed lines respectively show the results from GT5D and neoclassical theory (the matrix inversion code) at $t_{vTD}/R_0 = 250$.

- 1 discriminating the trapped electron modes contribution in density fluctuation
2 spectra,” Nuclear fusion **55**, 093021 (2015).
3 ⁸F. Romanelli, W. M. Tang, and R. B. White, “Anomalous thermal confinement in
4 ohmically heated tokamaks,” Nuclear fusion **26**, 1515 (1986).
5 ⁹C. L. Rettig, T. L. Rhodes, J. N. Leboeuf, W. A. Peebles, E. J. Doyle, G. M.
6 Staebler, K. H. Burrell, and R. A. Moyer, “Search for the ion temperature
7 gradient mode in a tokamak plasma and comparison with theoretical
8 predictions,” Physics of Plasmas **8**, 2232–2237 (2001).
9 ¹⁰J. E. Rice, M. J. Greenwald, Y. A. Podpaly, M. L. Reinke, P. H. Diamond,
10 J. W. Hughes, N. T. Howard, Y. Ma, I. Cziegler, B. P. Duval, P. C. Enn-

- 11 ever, D. Ernst, C. L. Fiore, C. Gao, J. H. Irby, E. S. Marmor, M. Porkolab,
12 N. Tsujii, and S. M. Wolfe, “Ohmic energy confinement saturation and core
13 toroidal rotation reversal in Alcator C-Mod plasmas,” Physics of Plasmas
14 **19**, 056106 (2012).
15 ¹¹C. Sung, A. White, N. Howard, C. Oi, J. Rice, C. Gao, P. Ennever,
16 M. Porkolab, F. Parra, D. Mikkelsen, D. Ernst, J. Walk, J. Hughes, J. Irby,
17 C. Kasten, A. Hubbard, M. Greenwald, and the Alcator C-Mod Team,
18 “Changes in core electron temperature fluctuations across the ohmic en-
19 ergy confinement transition in Alcator C-Mod plasmas,” Nucl. Fusion **53**,
20 083010 (2013).
21 ¹²R. McDermott, C. Angioni, G. Conway, R. Dux, E. Fable, R. Fischer,
22 T. Pütterich, F. Ryter, E. Viezzer, and the ASDEX Upgrade Team, “Core
23 intrinsic rotation behaviour in ASDEX Upgrade ohmic L-mode plasmas,”
24 Nucl. Fusion **54**, 043009 (2014).
25 ¹³C. Gao, J. Rice, H. Sun, M. Reinke, N. Howard, D. Mikkelsen, A. Hub-
26 bard, M. Chilenski, J. Walk, J. Hughes, P. Ennever, M. Porkolab, A. White,
27 C. Sung, L. Delgado-Aparicio, S. Baek, W. Rowan, M. Brookman,
28 M. Greenwald, R. Granetz, S. Wolfe, E. Marmor, and The Alcator C-Mod
29 Team, “Non-local heat transport in Alcator C-Mod ohmic L-mode plas-
30 mas,” Nucl. Fusion **54**, 083025 (2014).
31 ¹⁴J. Bernardo, C. Fenzi, C. Bourdelle, Y. Camenen, H. Arnichand, J. P. S.
32 Bizarro, S. Cortes, X. Garbet, Z. O. Guimarães-Filho, T. Aniel, J.-F. Artaud,
33 F. Clairet, P. Cottier, J. Gunn, P. Lotte, and the Tore Supra Team, “Density
34 impact on toroidal rotation in Tore Supra: Experimental observations and
35 theoretical investigation,” Plasma Phys. Control. Fusion **57**, 035002 (2015).
36 ¹⁵Y. Shi, J. Kwon, P. Diamond, W. Ko, M. Choi, S. Ko, S. Hahn, D. Na,
37 J. Leem, J. Lee, S. Yang, K. Lee, M. Joung, J. Jeong, J. Yoo, W. Lee, J. Lee,
38 Y. Bae, S. Lee, S. Yoon, K. Ida, and Y.-S. Na, “Intrinsic rotation reversal,
39 non-local transport, and turbulence transition in KSTAR L-mode plasmas,”
40 Nucl. Fusion **57**, 066040 (2017).
41 ¹⁶D. Na, Y.-S. Na, C. Angioni, S. Yang, J. Kwon, H. Jhang, Y. Camenen,
42 S. Lee, Y. Shi, W. Ko, J. Lee, T. Hahn, and KSTAR Team, “A compre-
43 hensive study on rotation reversal in KSTAR: Experimental observations and
44 modelling,” Nucl. Fusion **57**, 126008 (2017).
45 ¹⁷J. A. Lee, W. Lee, J. M. Kwon, S. H. Ko, J. Leem, G. S. Yun, H. K. Park,
46 Y. S. Park, K. W. Kim, and N. C. Luhmann Jr, “Observation of electron
47 driven quasi-coherent modes and their connection with core intrinsic rota-
48 tion in KSTAR ECH and ohmic L-mode plasmas,” Physics of Plasmas **25**,
49 022513 (2018).
50 ¹⁸N. Cao, J. Rice, P. Diamond, A. White, S. Baek, M. Chilenski, J. Hughes,
51 J. Irby, M. Reinke, P. Rodriguez-Fernandez, and the Alcator C-Mod Team,
52 “Hysteresis as a probe of turbulent bifurcation in intrinsic rotation reversals
53 on Alcator C-Mod,” Nucl. Fusion **59**, 104001 (2019).
54 ¹⁹L. Lin, M. Porkolab, E. M. Edlund, J. C. Rost, M. Greenwald, N. Tsujii,
55 J. Candy, R. E. Waltz, and D. R. Mikkelsen, “Studies of turbulence and
56 transport in Alcator C-Mod ohmic plasmas with phase contrast imaging and
57 comparisons with gyrokinetic simulations,” Plasma Phys. Control. Fusion
58 **51**, 065006 (2009).
59 ²⁰C. Angioni, R. M. McDermott, F. J. Casson, E. Fable, A. Bottino, R. Dux,
60 R. Fischer, Y. Podoba, T. Pütterich, F. Ryter, and E. Viezzer, “Intrinsic
61 Toroidal Rotation, Density Peaking, and Turbulence Regimes in the Core
62 of Tokamak Plasmas,” Phys. Rev. Lett. **107**, 215003 (2011).
63 ²¹M. Porkolab, J. Dorris, P. Ennever, C. Fiore, M. Greenwald, A. Hubbard,
64 Y. Ma, E. Marmor, Y. Podpaly, M. L. Reinke, J. E. Rice, J. C. Rost, N. Tsujii,
65 D. Ernst, J. Candy, G. M. Staebler, and R. E. Waltz, “Transport and
66 turbulence studies in the linear ohmic confinement regime in Alcator C-
67 Mod,” Plasma Phys. Control. Fusion **54**, 124029 (2012).
68 ²²P. Ennever, M. Porkolab, J. Candy, G. Staebler, M. L. Reinke, J. E. Rice,
69 J. C. Rost, D. Ernst, C. Fiore, J. Hughes, J. Terry, and Alcator C-Mod
70 Team, “The effects of dilution on turbulence and transport in C-Mod ohmic
71 plasmas and comparisons with gyrokinetic simulations,” Phys. Plasmas **22**,
72 072507 (2015).
73 ²³J. Citrin, H. Arnichand, J. Bernardo, C. Bourdelle, X. Garbet, F. Jenko,
74 S. Hacquin, M. J. Pueschel, and R. Sabot, “Comparison between measured
75 and predicted turbulence frequency spectra in ITG and TEM regimes,”
76 Plasma Phys. Control. Fusion **59**, 064010 (2017).
77 ²⁴W. X. Wang, B. A. Grierson, S. Ethier, J. Chen, E. Startsev, and P. H.
78 Diamond, “Understanding and predicting profile structure and parametric
79 scaling of intrinsic rotation,” Physics of Plasmas **24**, 092501 (2017).

- ²⁵B. A. Grierson, W. X. Wang, S. Ethier, G. M. Staebler, D. J. Battaglia, J. A. Boedo, J. S. DeGrassie, and W. M. Solomon, "Main-ion intrinsic toroidal rotation profile driven by residual stress torque from ion temperature gradient turbulence in the DIII-D tokamak," *Physical Review Letters* **118**, 015002 (2017).
- ²⁶B. A. Grierson, C. Chrystal, S. R. Haskey, W. X. Wang, T. L. Rhodes, G. R. McKee, K. Barada, X. Yuan, M. F. F. Nave, A. Ashourvan, and C. Holland, "Main-ion intrinsic toroidal rotation across the ITG/TEM boundary in DIII-D discharges during ohmic and electron cyclotron heating," *Physics of Plasmas* **26**, 042304 (2019).
- ²⁷I. Erofeev, E. Fable, C. Angioni, R. McDermott, and The ASDEX Upgrade Team, "Theory-based modeling of LOC-SOC transitions in ASDEX Upgrade," *Nucl. Fusion* **57**, 126067 (2017).
- ²⁸H. Arnichand, J. Citrin, S. Hacquin, R. Sabot, A. Krämer-Flecken, X. Garbet, C. Bourdelle, C. Bottereau, F. Clairet, J. C. Giacalone, Z. O. Guimarães-Filho, R. Guirlet, G. Hornung, A. Lebschy, P. Lotte, P. Maget, A. Medvedeva, D. Molina, V. Nikolaeva, D. Prisiazhniuk, the Tore Supra and the ASDEX Upgrade teams, "Identification of trapped electron modes in frequency fluctuation spectra," *Plasma Phys. Control. Fusion* **58**, 014037 (2016).
- ²⁹W. Hornsby, C. Angioni, Z. Lu, E. Fable, I. Erofeev, R. McDermott, A. Medvedeva, A. Lebschy, A. Peeters, and The ASDEX Upgrade Team, "Global gyrokinetic simulations of intrinsic rotation in ASDEX Upgrade Ohmic L-mode plasmas," *Nucl. Fusion* **58**, 056008 (2018).
- ³⁰B. Scott and J. Smirnov, "Energetic consistency and momentum conservation in the gyrokinetic description of tokamak plasmas," *Physics of Plasmas* **17**, 112302 (2010).
- ³¹J. Abiteboul, X. Garbet, V. Grandgirard, S. J. Allfrey, P. Ghendrih, G. Latu, Y. Sarazin, and A. Strugarek, "Conservation equations and calculation of mean flows in gyrokinetics," *Physics of Plasmas* **18**, 082503 (2011).
- ³²Y. Idomura, "Accuracy of momentum transport calculations in full- f gyrokinetic simulations," *Comput. Sci. Disc.* **5**, 014018 (2012).
- ³³M. Nakata and Y. Idomura, "Plasma size and collisionality scaling of ion temperature-gradient-driven turbulence," *Nucl. Fusion* **53**, 113039 (2013).
- ³⁴G. Dif-Pradalier, G. Hornung, X. Garbet, P. Ghendrih, V. Grandgirard, G. Latu, and Y. Sarazin, "The $E \times B$ staircase of magnetised plasmas," *Nucl. Fusion* **57**, 066026 (2017).
- ³⁵Y. Idomura, M. Ida, T. Kano, N. Aiba, and S. Tokuda, "Conservative global gyrokinetic toroidal full- f five-dimensional Vlasov simulation," *Computer Physics Communications* **179**, 391–403 (2008).
- ³⁶Y. Idomura, H. Urano, N. Aiba, and S. Tokuda, "Study of ion turbulent transport and profile formations using global gyrokinetic full- f Vlasov simulation," *Nucl. Fusion* **49**, 065029 (2009).
- ³⁷Y. Idomura, "A new hybrid kinetic electron model for full- f gyrokinetic simulations," *Journal of Computational Physics* **313**, 511–531 (2016).
- ³⁸Y. Idomura, K. Obrejan, Y. Asahi, and M. Honda, "Dynamics of enhanced neoclassical particle transport of tracer impurity ions in ion temperature gradient driven turbulence," *Physics of Plasmas* **28**, 012501 (2021).
- ³⁹W. W. Lee, "Gyrokinetic particle simulation model," *Journal of Computational Physics* **72**, 243–269 (1987).
- ⁴⁰C. J. McDevitt, P. H. Diamond, Ö. D. Gürcan, and T. S. Hahm, "Toroidal Rotation Driven by the Polarization Drift," *Phys. Rev. Lett.* **103**, 205003 (2009).
- ⁴¹Y. Camenen, Y. Idomura, S. Jolliet, and A. Peeters, "Consequences of profile shearing on toroidal momentum transport," *Nucl. Fusion* **51**, 073039 (2011).
- ⁴²Y. Asahi, V. Grandgirard, Y. Idomura, X. Garbet, G. Latu, Y. Sarazin, G. Dif-Pradalier, P. Donnel, and C. Ehrlacher, "Benchmarking of flux driven full- F gyrokinetic simulations," *Phys. Plasmas* **24**, 102515 (2017).
- ⁴³S. Satake, Y. Idomura, H. Sugama, and T.-H. Watanabe, "Benchmark test of drift-kinetic and gyrokinetic codes through neoclassical transport simulations," *Computer Physics Communications* **181**, 1069–1076 (2010).
- ⁴⁴S. Matsuoka, Y. Idomura, and S. Satake, "Neoclassical transport benchmark of global full- f gyrokinetic simulation in stellarator configurations," *Phys. Plasmas* **25**, 022510 (2018).
- ⁴⁵Y. Idomura, "Full- f gyrokinetic simulation over a confinement time," *Physics of Plasmas* **21**, 022517 (2014).
- ⁴⁶S. Jolliet and Y. Idomura, "Plasma size scaling of avalanche-like heat transport in tokamaks," *Nucl. Fusion* **52**, 023026 (2012).
- ⁴⁷Y. Idomura, T. Ina, Y. Ali, and T. Imamura, "Acceleration of Fusion Plasma Turbulence Simulations using the Mixed-Precision Communication-Avoiding Krylov Method," in *SC20: International Conference for High Performance Computing, Networking, Storage and Analysis* (IEEE, Atlanta, GA, USA, 2020) pp. 1–13.
- ⁴⁸M. Sato, Y. Ishikawa, H. Tomita, Y. Kodama, T. Odajima, M. Tsuji, H. Yashiro, M. Aoki, N. Shida, I. Miyoshi, K. Hirai, A. Furuya, A. Asato, K. Morita, and T. Shimizu, "Co-Design for A64FX Manycore Processor and "Fugaku"," in *SC20: International Conference for High Performance Computing, Networking, Storage and Analysis* (IEEE, Atlanta, GA, USA, 2020) pp. 1–15.
- ⁴⁹Y. Sarazin, J. Hillairet, J.-L. Duchateau, K. Gaudimont, R. Varennes, X. Garbet, P. Ghendrih, R. Guirlet, B. Pégourié, and A. Torre, "Impact of scaling laws on tokamak reactor dimensioning," *Nucl. Fusion* **60**, 016010 (2020).
- ⁵⁰Y. Idomura and M. Nakata, "Plasma size and power scaling of ion temperature gradient driven turbulence," *Physics of Plasmas* **21**, 020706 (2014).
- ⁵¹Y. Idomura, "Isotope and plasma size scaling in ion temperature gradient driven turbulence," *Physics of Plasmas* **26**, 120703 (2019).
- ⁵²J. W. Connor, J. B. Taylor, and H. R. Wilson, "Shear damping of drift waves in toroidal plasmas," *Physical review letters* **70**, 1803 (1993).
- ⁵³J. Y. Kim and M. Wakatani, "Radial structure of high-mode-number toroidal modes in general equilibrium profiles," *Physical review letters* **73**, 2200 (1994).
- ⁵⁴R. E. Waltz, G. M. Staebler, and W. M. Solomon, "Gyrokinetic simulation of momentum transport with residual stress from diamagnetic level velocity shears," *Physics of Plasmas* **18**, 042504 (2011).
- ⁵⁵R. Buchholz, Y. Camenen, F. J. Casson, S. R. Grosshauser, W. A. Hornsby, P. Migliano, and A. G. Peeters, "Toroidal momentum transport in a tokamak due to profile shearing," *Physics of Plasmas* **21**, 062304 (2014).
- ⁵⁶M. Mattioli, R. Giannella, R. Myrnas, C. Demichelis, B. Denne-Hinnov, T. D. De Wit, and G. Magyar, "Laser blow-off injected impurity particle confinement times in JET and Tore Supra," *Nuclear Fusion* **35**, 1115 (1995).
- ⁵⁷A. Peeters, C. Angioni, A. Bortolon, Y. Camenen, F. Casson, B. Duval, L. Fiederspiel, W. Hornsby, Y. Idomura, T. Hein, N. Kluy, P. Mantica, F. Parra, A. Snodin, G. Szepesi, D. Strintzi, T. Tala, G. Tardini, P. de Vries, and J. Weiland, "Overview of toroidal momentum transport," *Nucl. Fusion* **51**, 094027 (2011).
- ⁵⁸P. Diamond, Y. Kosuga, Ö. Gürcan, C. McDevitt, T. Hahm, N. Fedorczak, J. Rice, W. Wang, S. Ku, J. Kwon, G. Dif-Pradalier, J. Abiteboul, L. Wang, W. Ko, Y. Shi, K. Ida, W. Solomon, H. Jhang, S. Kim, S. Yi, S. Ko, Y. Sarazin, R. Singh, and C. Chang, "An overview of intrinsic torque and momentum transport bifurcations in toroidal plasmas," *Nucl. Fusion* **53**, 104019 (2013).
- ⁵⁹Y. Camenen, A. G. Peeters, C. Angioni, F. J. Casson, W. A. Hornsby, A. P. Snodin, and D. Strintzi, "Transport of Parallel Momentum Induced by Current-Symmetry Breaking in Toroidal Plasmas," *Phys. Rev. Lett.* **102**, 125001 (2009).
- ⁶⁰R. R. Dominguez and G. M. Staebler, "Anomalous momentum transport from drift wave turbulence," *Physics of Fluids B: Plasma Physics* **5**, 3876–3886 (1993).
- ⁶¹Ö. Gürcan, P. H. Diamond, and T. S. Hahm, "Spatial and spectral evolution of turbulence," *Physics of plasmas* **14**, 055902 (2007).
- ⁶²Ö. D. Gürcan, P. H. Diamond, P. Hennequin, C. J. McDevitt, X. Garbet, and C. Bourdelle, "Residual parallel Reynolds stress due to turbulence intensity gradient in tokamak plasmas," *Physics of Plasmas* **17**, 112309 (2010).
- ⁶³J. Y. Kim, Y. Kishimoto, M. Wakatani, and T. Tajima, "Poloidal shear flow effect on toroidal ion temperature gradient mode: A theory and simulation," *Physics of Plasmas* **3**, 3689–3695 (1996).
- ⁶⁴Y. Kishimoto, J.-Y. Kim, W. Horton, T. Tajima, M. J. LeBrun, and H. Shirai, "Toroidal mode structure in weak and reversed magnetic shear plasmas and its role in the internal transport barrier," *Plasma Phys. Control. Fusion* **41**, A663–A678 (1999).
- ⁶⁵R. E. Waltz, R. L. Dewar, and X. Garbet, "Theory and simulation of rotational shear stabilization of turbulence," *Physics of Plasmas* **5**, 1784–1792 (1998).
- ⁶⁶X. Garbet, N. Dubuit, E. Asp, Y. Sarazin, C. Bourdelle, P. Ghendrih, and G. T. Hoang, "Turbulent fluxes and entropy production rate," *Physics of Plasmas* **12**, 082511 (2005).

- 1 ⁶⁷S. Hirshman and D. Sigmar, "Neoclassical transport of impurities in toka- 8
2 mak plasmas," Nucl. Fusion **21**, 1079–1201 (1981). 9
- 3 ⁶⁸M. Kikuchi, M. Azumi, S. Tsuji, K. Tani, and H. Kubo, "Bootstrap current 10
4 during perpendicular neutral injection in JT-60," Nucl. Fusion **30**, 343–355 11
5 (1990). 12
- 6 ⁶⁹R. Varennes, X. Garbet, L. Vermare, Y. Sarazin, G. Dif-Pradalier, V. Grand- 13
7 girard, P. Ghendrih, P. Donnel, M. Peret, K. Obrejan, and E. Bourne, "Syn-
ergy of Turbulent Momentum Drive and Magnetic Braking," Phys. Rev.
Lett. **128**, 255002 (2022).
- ⁷⁰C. Fenzi, X. Garbet, E. Trier, P. Hennequin, C. Bourdelle, T. Aniel,
G. Colledani, P. Devynck, C. Gil, and Ö. Gürçan, "On plasma rotation
with toroidal magnetic field ripple and no external momentum input," Nu-
clear Fusion **51**, 103038 (2011).