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Residual parallel Reynolds stress due to turbulence intensity gradient in tokamak plasmas

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A novel mechanism for driving residual stress in tokamak plasmas based on $k_{\parallel}$ symmetry breaking by the turbulence intensity gradient is proposed. The physics of this mechanism is explained and its connection to the wave kinetic equation and the wave-momentum flux is described. Applications to the H-mode pedestal in particular to internal transport barriers, are discussed. Also, the effect of heat transport on the momentum flux is discussed. © 2010 American Institute of Physics. [doi:10.1063/1.3503624]

I. INTRODUCTION

The anomalous angular momentum transport in magnetic confinement devices has become a popular subject in recent years. The renewed interest in this decades old problem is due to recent observations of rotation with no "external" torque input in a number of machines, in the high confinement mode (H-mode),[1–4] the low confinement mode (L-mode),[5–7] as well as internal transport barriers (ITBs).[8]

At the simplest level, the toroidal rotation with no torque input scales with ion pressure, temperature, or its gradient, whether the confinement is enhanced or not. In terms of its direction and scaling trends, L-mode rotation seems to be a complex phenomenon with various competing effects, which may frequently involve sign flips.[7,9,10] In contrast, H mode rotation is usually in the co-current direction, and the on-axis plasma velocity appears to increase with stored energy W_p and decrease with total plasma current I_p.[11] This empirical scaling, which has come to be called the Rice scaling, particularly manifests itself strongly in Alcator C-Mod and D-IIID,[1,2] and if taken at face value gives rather optimistic predictions for ITER's rotation speed.[11]

There are indeed different mechanisms to drive rotation without any external torque.[12] For instance, one particular mechanism that drives rotation, is the ripple in the toroidal magnetic field.[4,13] In principle, the mechanism is independent of the mode of tokamak operation. In Tore-Supra, where the magnetic field ripple is large, it consistently drives a countercurrent rotation.[13] Similarly in JT-60U, the ripple seems to drive a counter-current rotation that correlate well with ∇P_i,[14] and as the ripple is decreased, the H-mode rotation becomes more co-current.[4] More generally, the presence of any non-axisymmetry in the magnetic field configuration, be it due to MHD modes, magnetic islands, or external error fields can produce a neoclassical toroidal viscous force that set the offset value of toroidal rotation.[15–18] This was observed in various tokamaks including NSTX[19,20] and D-IIID.[21,22]

While there are different mechanisms, which give rise to "intrinsic" rotation (i.e., rotation with no applied torque) that may differently scale with different parameters dominant in different machines and in different modes of operation, there seems to be a universal part of this rotation that is associated with the self-organization process that leads to the formation of the H-mode. We argue that this distinction between intrinsic rotation (such as that driven by ripple or indirect wave momentum from the heating method) and L-H spin-up (i.e., the gain in rotation, associated with the plasma turbulence and its self-organization) is useful for an understanding of the differences between different observations, although the two might not always be additive. The jargon in this seems a bit confusing as some call this L-H spin-up component of plasma rotation as the "intrinsic rotation." It is obvious that whatever the name is, it is this component to which the Rice scaling applies. A residual of this should, in principle, be present even in the L-mode. However, if this is the case, it seems to be subdominant, and could have different scaling characteristics, which we do not know unless we know the physics of the dominant effect, which is possibly different in different plasma conditions and somehow subtract it. It is also interesting to note here that a similar spin-up is apparently associated with the transition to the I-mode, an improved L-mode confinement regime with a barrier in temperature but not in density,[23] and similarly to the L-H spin-up, L-I spin-up seems to also follow the Rice scaling in the Alcator C-Mod tokamak.[24]

On the theoretical front, there has been a multitude of attempts to explain this phenomenon. Note that an inward convective flux of momentum[25–27] (i.e., a "momentum

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pinch") that transports scrape-off layer (SOL) flows into the core, cannot explain the L-H spin-up, since the direction of the SOL flows can be changed by changing the location of the x-point (or where the plasma touches the wall), yet the direction of the H-mode rotation in the core remains unaltered.[28]

A theory that has some possibility of explaining this phenomenon is based on "residual" Reynolds stress[29,30] (or more generally, a residual component of the toroidal stress tensor). The residual stress corresponds to the part of the stress tensor Π, which persists even when the transported quantity and its derivatives vanish, so for instance, it can be proportional to the gradients of other fields, such as ∇P or ∇n. Such a term can exist for a vector quantity, such as momentum or scalar quantity, which can change its sign (since, for instance, if a positive, definite scalar quantity such as density is zero, its flux would also vanish). The residual stress can be proportional to the pressure gradient, or E_r shear (as well as other things), and may dominate the momentum transport within the H-mode pedestal. However the idea of a residual stress and that it plays a critical role in momentum transport is very general and is not exclusively linked to E_r shear.

Given its generality, the suggestion that an off-diagonal Reynolds stress persists is not a trivial one. In general, the off-diagonal term involves the average value of $k_\parallel$, which vanishes if $|\Phi_k|^2$ is symmetric with respect to $k_\parallel$. The standard formulation of drift-instabilities, which includes ion temperature gradient driven (ITG) mode, as well as trapped electron mode (TEM) does have this symmetry. Therefore one needs to consider additional processes to break this $k_\parallel$ symmetry and give a net wave momentum. Toroidal current generating an asymmetry in instability (e.g. Ref. 31), $E\times B$ shear,[29,32] Alfvénic turbulence,[33] charge separation induced by the polarization drift[34] and up-down asymmetry of flux surfaces[35] can be counted among the possible candidates. Parallel flow shear itself,[36] magnetic curvature (curvature from $B_\parallel^*$ in laboratory frame),[26] or the effect of Coriolis drift in rotating frame,[25] can also lead to $k_\parallel$ symmetry breaking but give diffusive and pinchlike contributions to the Reynolds stress. Here we only consider the truly off-diagonal (i.e., residual) terms, which do not contain the transported field itself, since only these can explain the formation of a nonvanishing field from an initial value of zero, or the anomalous residual "torque" that acts on the plasma when the field and its gradients are set to vanish.

In this paper, we suggest a simple mechanism for the generation of such an off-diagonal term from *turbulence intensity gradient*. We argue that the existence of a turbulence intensity gradient will locally break $k_\parallel$ symmetry and lead to a net residual stress. Since intensity gradients can be measured, and thus the residual stress separately estimated, the idea could, in principle, be tested. The idea presented here is consistent with (and inspired by) the fact that the turbulent momentum transport can be linked to turbulent wave-momentum transport and the wave kinetic equation, and computed by considering the wave-momentum flux.[30,37] In this formulation, two residual terms appear, one proportional to the $E\times B$ shear (which thus corresponds to Ref. 29), and another one proportional to the intensity gradient. Here we explain the direct mechanism for this second term. Note that the mechanism that is discussed here is based on symmetry breaking (i.e., the off-diagonal component of the cross-phase between $\tilde{v}_r$ and $\tilde{v}_\parallel$) due to intensity gradient, and is different from simply a $\partial_r\langle\tilde{v}_r\tilde{v}_\parallel\rangle\rightarrow\partial_r\langle|\tilde{\Phi}|^2\rangle\langle\cos\phi_{v_r,v_\parallel}\rangle$ dependence on turbulence intensity, which is also inevitably there. There are strong indications that such an effect is playing a role in ITG/TEM in momentum transport in gyrokinetic simulations.[38,39]

Note that the connection of this exercise to neoclassical theory is nontrivial. Here we explicitly compute only the parallel component of the Reynolds stress driven by intensity gradient, with a decomposition consisting of parallel and perpendicular directions (i.e., $\hat{\mathbf{r}}\times\hat{\mathbf{b}}$). In contrast, in neoclassical theory, it is common to use a decomposition of vector and tensor quantities in terms of "parallel" and toroidal directions. That is an oblique coordinate system well suited to global tokamak geometry, but not very useful to describe fluctuation dynamics. It is usually argued that, in an axisymmetric torus, the neoclassical viscous stress due to ion-ion collisions is dominant in the parallel direction (i.e., parallel in an oblique decomposition). This damps the ion poloidal flow to its neoclassical value and allows only the toroidal velocity to be anomalous. As long as the perpendicular Reynolds stress is not larger than the parallel Reynolds stress multiplied by the ratio of poloidal to toroidal magnetic fields, the parallel Reynolds stress will be the dominant drive term for this anomalous toroidal rotation.

In general, however, the toroidal projections of both the parallel and the perpendicular Reynolds stresses act to drive a toroidal flow, and while the poloidal projections compete against the parallel neoclassical viscosity and possibly get damped; the toroidal projections may drive anomalous toroidal rotation. While we use this basic picture to justify using the parallel Reynolds stress to represent the toroidal one, it should be noted that there are recent indications that poloidal rotation may also be anomalous[40,41] or at least radially varying deviations from an average neoclassical poloidal rotation should be expected.[42]

The remainder of this paper is organized as follows. In Sec. II, we derive the residual stress from intensity gradient, assuming mode rational surfaces are tightly packed. Then in subsection (a) of Sec. II we consider the opposite limit with little overlap between mode rational surfaces, which corresponds to weak or reversed shear case. In Sec. III, we discuss the effect of heat flux on momentum transport. Section IV provides a detailed discussion of the connection between wave kinetic formulation and the intensity gradient driven symmetry breaking. Section V contains results and conclusion.

II. RESIDUAL STRESS FROM TURBULENCE INTENSITY GRADIENT

The flux of angular momentum is related to the parallel component of the Reynolds stress

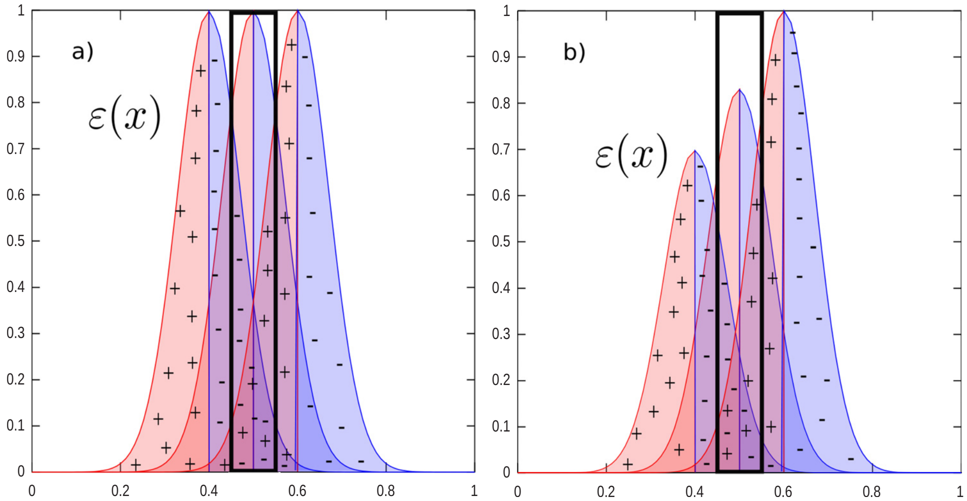

FIG. 1. (Color online) A cartoon of the symmetry breaking from intensity gradient. The flat intensity case (a) gives zero net $k_\parallel$ inside the box drawn in the center because the negative $k_\parallel$ contributed from the left eigenmode is canceled by the positive $k_\parallel$ contributed by the right eigenmode. In contrast, when an intensity gradient exists as in case (b), the positive $k_\parallel$ contributed by the right eigenmode exceeds the negative $k_\parallel$ contributed from the left eigenmode, and a net positive $k_\parallel$ results. Note that, in either case, the contribution of the central eigenmode to the "average $k_\parallel$ inside the box" vanishes. See the discussion on weak or reversed shear for the local $k_\parallel$ as opposed to average over the box.

$$\Gamma_\parallel = \langle \tilde{v}_r \tilde{v}_\parallel \rangle,$$

the gradient of which is the primary turbulent drive of toroidal rotation, or mean parallel velocity. A simple evolution equation for $\tilde{v}_\parallel$ would have the basic form

$$D_t \tilde{v}_\parallel - v_{ti} \frac{d\bar{v}_\parallel}{dx} \rho_i \partial_y \left(\frac{e}{T} \tilde{\Phi} \right) = - v_{ti}^2 \nabla_\parallel \left(\frac{e\tilde{\Phi}}{T} + \frac{\tilde{P}}{P_i} \right),$$

which allows the computation of the Reynolds stress as

$$\Gamma_\parallel = \langle \tilde{v}_r \tilde{v}_\parallel \rangle = - \mathrm{Re} \sum_k i v_{ti}^2 \rho_i k_\theta \left[\frac{v_{ti} k_\theta}{\omega_k} \frac{\rho_i}{v_{ti}} \frac{\partial U_\parallel}{\partial r} - \frac{v_{ti} k_\parallel}{\omega_k} \right] \times \left| \frac{e\tilde{\Phi}_k}{T_i} \right|^2 + \mathrm{Re} \sum_k v_{ti}^2 \rho_i \left[\frac{v_{ti} k_\parallel}{\omega_k} \tilde{v}_{r,k}^* \tilde{P}_k \right]. \quad (1)$$

Leaving the discussion of the last term to Sec. III, we focus on the first term. In toroidal geometry, we use $k_\theta = m/r$ and $k_\parallel = (m - nq)/qR$, so that

$$k_\parallel = - k_\theta \hat{s} \frac{(r - r_{n,m})}{qR},$$

the Reynolds stress in this notation becomes

$$R(r) = \langle \tilde{v}_r \tilde{v}_\parallel \rangle = - \mathrm{Re} \sum_{n,m} i v_{ti}^2 \left(\rho_i^2 k_\theta^2 \frac{\Omega_i}{\omega_k} \right) \times \left[\frac{\rho_i}{v_{ti}} \frac{\partial U_\parallel}{\partial r} + \frac{\hat{s}(r - r_{n,m})}{qR} \right] \times \left| \frac{e\tilde{\Phi}_{n,m}(r)}{T_i} \right|^2 .$$

This is the Reynolds stress at a given point r, however, the way it is computed involves a sum over all modes, each of which are localized at a different resonant surface. In order to compute it, we frequently assume that the mode rational surfaces are "tightly packed." This leads to

$$\sum_{m,n} \to \int dn \int dm,$$

following Ref. 43, we write $dn = (r/q) dk_\theta$ and $dm = |k_\theta| \hat{s} dr_{m,n}$.

A. The case of strong overlap among modes

In other words, by summing over different m's for a given n, we include the effect of modes localized at different spatial locations. However, of each of these modes, we only look at the contribution to a given radial location r.

The off-diagonal terms are proportional to $(r_{n,m} - r)$, and normally the $|e\tilde{\Phi}_{n,m}(r)/T_i|^2$ are centered around $r_{n,m}$ [i.e., they are even functions of $(r_{n,m} - r)$]. One can make a change of variables to $x = r_{n,m} - r$. Note that we choose the sign convention $dr_{nm} = dx$ so that it guarantees that a positive gradient in r_{nm} (i.e., the envelope increases as we move to rational surfaces localized at larger minor radii) implies a positive gradient in x also. Note that r is fixed and it corresponds to the point at which we want to compute the flux

$$\Gamma_\parallel(r) = \langle \tilde{v}_r \tilde{v}_\parallel \rangle = - \mathrm{Re}\, i v_{ti}^2 \int dk_\theta \int_{-\epsilon}^{+\epsilon} dx \left(\rho_i^2 k_\theta^2 \frac{\Omega_i}{\omega_k} \right) \frac{\hat{s} r}{q} |k_\theta| \times \left[\frac{\rho_i}{v_{ti}} \frac{\partial U_\parallel}{\partial r} - \frac{\hat{s} x}{qR} \right] \times \left| \frac{e\tilde{\Phi}(x)}{T_i} \right|^2 .$$

In the continuum limit, one can replace the limits of integration with $-\infty$ to $+\infty$ instead of $-\epsilon$ to $+\epsilon$ (which is the smallness parameter as opposed to $\varepsilon(x)$, the turbulence intensity). Here, we write it this way to emphasize the connection to the discrete limit depicted, for instance, in Fig. 1. Now assume that we have an intensity gradient, so

$$\left| \frac{e\tilde{\Phi}(x)}{T_i} \right|^2 \to \varepsilon(x) = \varepsilon(0) + x \left. \frac{d\varepsilon(x)}{dx} \right|_{x=0} .$$

Using this, we can finally write

$$\Gamma_\parallel(r) = - v_{ti}^2 \int dk_\theta \int_{r-\epsilon}^{r+\epsilon} dr_{n,m} \left(\rho_i^2 k_\theta^2 \frac{\Omega_i}{\tau_k} \right) \frac{\hat{s} r}{q} |k_\theta| \times \left[\frac{\rho_i}{v_{ti}} \frac{\partial U_\parallel}{\partial r} \varepsilon(x) - \frac{\hat{s} x^2}{qR} \frac{d\varepsilon}{dx} \right], \quad (2)$$

or

$$\Gamma_\parallel(r) = -\chi_{\phi,0}\left(\varepsilon\frac{\partial U_\parallel}{\partial r} - v_{ti}\frac{\hat{s}}{qR}\frac{\langle x^2\rangle}{\rho_i^2}\rho_i\frac{\partial\varepsilon}{\partial r}\right), \quad (3)$$

where $\langle x^2\rangle \to \Delta^2$, where Δ is the mode width. Apparently the sign is such that (for $\hat{s}>0$), a positive gradient (such as the usual L-mode case) results in an outward flux of momentum (typically leading to counter-current rotation), while a negative gradient would result in an inward one (or co-current rotation).

Since fluctuation intensity is inevitably linked to profile gradients (e.g., ∇T), and most intrinsic mechanisms also lead to similar dependence, it is not easy to experimentally measure this effect. One way to distinguish the intensity gradient effects from that of profile gradients, may be to note that intensity gradient is in fact related to the "curvature" of the profile (e.g., $\partial^2 T/\partial r^2$) rather than its gradient. Physically, this is due to the fact that turbulence intensity is tied to available "free energy," which is ∇T. Thus its gradient should be linked to the radial derivative of available free energy or to profile curvature. One way to see that is to differentiate the Fick's law for heat flux

$$Q \propto -\chi_0\varepsilon\frac{\partial T}{\partial r},$$

assuming Q is independent of radius in steady state, we can write

$$\frac{\partial\varepsilon}{\partial r} \propto -\chi_0\frac{\varepsilon^2}{Q}\frac{\partial^2 T}{\partial r^2}.$$

This suggests that, if rotation is observed to be linked to profile curvatures more than its gradients (at least in some cases), which may be taken as an experimental indication that intensity gradients play an important role in intrinsic rotation. This effect would be particularly important at the top of the pedestal or an ITB where the profile curvature is most pronounced.

A similar argument can be raised about the dependence of intensity gradient on $E\times B$ shear. Note that while we neglect the direct effect of $E\times B$ shear on momentum flux, there is also the indirect effect via turbulence reduction. A simple expression for the turbulence reduction by $E\times B$ shear having the basic form

$$\varepsilon \sim \frac{\varepsilon_0}{1+\alpha\left(\frac{dv_E}{dr}\right)^2},$$

would suggest that

$$\frac{\partial\varepsilon}{\partial r} \sim -2\alpha\frac{\varepsilon^2}{\varepsilon_0}\frac{dv_E}{dr}\frac{d^2v_E}{dr^2},$$

which is interesting because it would suggest that the flux from Eq. (3) is toward an E_r "well." We can see this by noting that to the left of a well, $dv_E/dr<0$, $d^2v_E/dr^2>0 \Rightarrow \partial\varepsilon/\partial r>0$, and $\Gamma_\parallel$ is outward, while to the right of the well $dv_E/dr>0$, $d^2v_E/dr^2>0 \Rightarrow \partial\varepsilon/\partial r<0$, and $\Gamma_\parallel$ is inward. This means that an E_r well (for instance, as in a pedestal), pulls parallel momentum toward itself.

B. The case of little overlap among modes (weak or reversed shear)

It is possible that in some cases, the assumption of tightly packed modes or surfaces is not satisfied, and that their overlap is limited. Therefore we also consider the case of "weak overlap" among stationary modes, and only consider the sum over a few neighboring modes in order to compute the flux. Here, by stationary, we mean that while the waves radially propagate and get damped at their corresponding Landau resonance points defining an eigenmode. The eigenmodes themselves do not radially propagate. Let us define

$$\Gamma_\parallel^{(\mathrm{OD})} = \langle\tilde{v}_r\tilde{v}_\parallel\rangle^{\mathrm{OD}} = \mathrm{Re}\left[\sum_k iv_{ti}^2\rho_i k_\theta\left(\frac{v_{ti}k_\parallel}{\omega_k}\right)\right]\left|\frac{e\tilde{\Phi}_k}{T_i}\right|^2, \quad (4)$$

where the superscript OD indicates "off-diagonal," so that we can write

$$\Gamma_\parallel^{\mathrm{OD}}(r) = -\mathrm{Re}\sum_{m'}\sum_{n'} iv_{ti}^2\left(\rho_i^2\frac{m'^2}{r^2}\frac{\Omega_i}{\omega_k}\right)\frac{s(r-r_{n',m'})}{qR} \times\left|\frac{e\tilde{\Phi}_{n',m'}(r)}{T_i}\right|^2. \quad (5)$$

The sum over n' involves at least $\{n-1,n,n+1\}$. Thus the flux in Eq. (5) is nonzero mainly because not all $\Phi_{n',m'}$'s have exactly the same shape. In particular, we assume that their amplitudes differ. Let us take the distance between two neighboring mode rational surfaces as $\Delta = r_{n+1,m} - r_{n,m}$ and define further $r = r_{n,m}+\delta r$ (i.e., $\delta r<\Delta$)

$$\Gamma_\parallel^{\mathrm{OD}}(r) = -\mathrm{Re}\sum_m iv_{ti}^2\left(\rho_i^2\frac{m^2}{r^2}\frac{\Omega_i}{\omega_k}\right)\frac{\hat{s}}{qR}\left[(\delta r+\Delta)\times\left|\frac{e\tilde{\Phi}_{n-1,m}(r_{n-1,m}+\Delta+\delta r)}{T_i}\right|^2 + \delta r\left|\frac{e\tilde{\Phi}_{n,m}(r_{n,m}+\delta r)}{T_i}\right|^2 + (\delta r-\Delta)\times\left|\frac{e\tilde{\Phi}_{n+1,m}(r_{n+1,m}-\Delta+\delta r)}{T_i}\right|^2\right]. \quad (6)$$

At a lowest order approximation, we will use a simple Gaussian instead of the actual eigenfunction. For a Gaussian centered around $r_{n,m}$

$$f(r_{n,m}+\Delta) = f_0 e^{-\Delta^2/2\sigma^2},$$

is the attenuation at a distance Δ from the central value. This allows us to write

$$\left|\frac{e\tilde{\Phi}_{n\pm1,m}(r_{n\pm1,m}+\Delta+\delta r)}{T_i}\right|^2 \to \left|\frac{e\tilde{\Phi}_{n\pm1,m}(r_{n\pm1,m}+\delta r)}{T_i}\right|^2 e^{-\Delta^2/2\sigma^2} \to \varepsilon_{n\pm1,m}e^{-\Delta^2/2\sigma^2},$$

where σ corresponds to the mode width, which is expected to be slightly larger than Δ so that there is some overlap (on the other hand in the limit $\sigma \gg \Delta$, one can safely assume the mode rational surfaces are tightly packed), and $\varepsilon_{n,m}$ is the intensity of the mode n,m near the mode rational surface n,m (within a distance $\delta r \ll \Delta$ of it). Within this local approximation (6) becomes

$$\Gamma_{\parallel}^{\mathrm{OD}}(r) = -\,\mathrm{Re}\sum_m i v_{ti}^2\left(\rho_i^2\frac{m^2}{r^2}\frac{\Omega_i}{\omega_k}\right)\frac{\hat{s}}{qR}\left[[\varepsilon_{n,m} + (\varepsilon_{n-1,m} + \varepsilon_{n+1,m})e^{-\Delta^2/2\sigma^2}]\delta r + \Delta^2\frac{(\varepsilon_{n-1,m} - \varepsilon_{n+1,m})}{\Delta}e^{-\Delta^2/2\sigma^2}\right], \quad (7)$$

note that in the limit of small Δ

$$\frac{(\varepsilon_{n-1,m} - \varepsilon_{n+1,m})}{\Delta} \to -\frac{d\varepsilon}{dr}.$$

What is rather clear in Eq. (7) is that, if $\Delta^2/2\sigma^2 \ll 1$, one recovers a similar expression as before [e.g., Eq. (2)] with the sum over m corresponding to the k_θ integration. However on the other hand if $\Delta^2/2\sigma^2 > 1$, only the term proportional to $\delta r = r - r_{n,m}$ survives. This roughly corresponds to an internal barrier with distant flux surfaces, as would correspond to a weak or reversed shear ITB (However note that the representation employed here, is not formally suitable to describe a reversed shear ITB case).

This gives a locally nonzero flux "toward" the mode rational surface in a weak shear case and "away from" it in the reversed shear case. In normal shear this term would cause a radial modulation of the flux that averages out if the mode rational surfaces are close together. However if we have a barrier around a single dominant mode rational surface, this term may be the primary effect since the effects due to mode rational surface interactions cease to be important.

In general, the effect mentioned here would be largest if there is a *sharp decrease or increase* in the turbulence intensity (such as in a barrier or a pedestal). For example for the case of sudden decrease (as a function of r), there would exist a "last active mode rational surface" whose negative $k_\parallel$ part could not be balanced (since there would be no modes on the right (e.g., see Fig. 1 for a cartoon) giving rise as a result to a net parallel (and residual) Reynolds stress.

III. THE EFFECT OF HEAT FLUX

The simple picture that we described in this paper up to this point does not have a dependence on the direction of mode rotation. This is mainly due to the fact that up to now, we have neglected the last term in Eq. (1), which represents the effect of heat flux on momentum flux. It is common to use, for instance, a particular model for the evolution of P such as

$$D_t\tilde{P} - v_{ti}\frac{dP}{dx}\rho_i\partial_Y\tilde{\Phi} = 0,$$

in order to solve for $\tilde{P}$ and substitute into the expression for flux.

However, this leads to a double propagator ($\mathrm{Re}[-i/\omega_k^2]$ in a simple quasilinear calculation) to which the usual recipe for causality [i.e., using $\omega_k \to \omega_k^{(r)} + i|\gamma_k|$] is not immediately applicable. This is due to the fact that, while $\mathrm{Re}[-i/\omega_k^2]$ depends on the sign of $\omega_k^{(r)}$, the decorrelation rate ($\tau_c^{Q.L.} \propto |\gamma_k|/\omega_k^2$) does not, and it is not immediately clear if one should use τ_c^2 or $\mathrm{Re}[-i/\omega_k^2]$ as the propagator. In order to avoid this confusion, here we present a derivation that simply expresses this term in terms of heat flux.

Now let us consider the last term in Eq. (1)

$$\Gamma_{\parallel}^{(H)} = +\,\mathrm{Re}\sum_k v_{ti}^2\rho_i\left[\frac{v_{ti}k_\parallel}{\omega_k}\tilde{v}_{r,k}^*\tilde{P}_k\right],$$

we use the superscript H to indicate the connection to heat flux. In general

$$\Gamma_{\parallel}^{(H)} = \sum_k v_{ti}^2\rho_i\left[\mathrm{Re}\left(\frac{v_{ti}k_\parallel}{\omega_k}\right)\overbrace{\mathrm{Re}(\tilde{v}_{r,k}^*\tilde{P}_k)}^{Q} - \mathrm{Im}\left(\frac{v_{ti}k_\parallel}{\omega_k}\right)\mathrm{Im}(\tilde{v}_{r,k}^*\tilde{P}_k)\right], \quad (8)$$

where Q is the heat flux. Notice that the second term,

$$-\sum_k v_{ti}^2\rho_i\,\mathrm{Im}\left(\frac{v_{ti}k_\parallel}{\omega_k}\right)\mathrm{Im}(\tilde{v}_{r,k}^*\tilde{P}_k) = -\,\mathrm{Re}\left[\sum_k i v_{ti}^2\rho_i k_\theta\left(\frac{v_{ti}k_\parallel}{\omega_k}\right)\right]\mathrm{Re}(\tilde{\Phi}_k^*\tilde{P}_k),$$

has essentially the exact same form with the other off-diagonal term (4), which we have computed in Sec. II. This term would be driven by the "in-phase" oscillations of $\tilde{P}_k$ and $\tilde{\Phi}_k$. The sign of this term depends on the phase between the two fields. If the cross-phase $\phi = \phi_p - \phi_\phi$ is between $-\pi/2 < \phi < \pi/2$, this term is in the opposite direction to the primary off-diagonal term (i.e., away from a maximum of turbulence intensity), while if $\pi/2 < \phi < 3\pi/2$, it is in the same direction (i.e., toward a maximum of intensity gradient). The phase between fluctuating pressure and electrostatic field is a property of the type of microturbulence (i.e., ITG versus TEM) that exists in the plasma. However it is well-known that the phase might change in the nonlinear regime especially in the presence of sheared flows.

Finally, the first term in Eq. (8) is also interesting because it links the momentum flux to heat flux

$$\Gamma_{\parallel}^{(H_1)} = \sum_k v_{ti}^2\rho_i\left[\frac{\omega_k^{(r)}v_{ti}k_\parallel}{\omega_k^2 + \gamma_k^2}\right]Q_k.$$

In order to formulate it in the same way as before we can use $\omega_k^{(r)} \equiv (v_{\mathrm{ph}}/v_{ti})\rho_i k_\theta$, where v_{ph} is the phase velocity of the microturbulence. In simple quasilinear theory, the phase velocity determines the sign. However it is not uncommon to observe, ITG modes rotating in the electron direction in fully developed case, due to nonlinear frequency shift (Thus, in fact the sign of this term is strictly speaking undetermined in this simple picture since causality argument does not by itself set a direction of rotation). Assuming that the heat flux is proportional to turbulence intensity (i.e., $Q_k=(Q/\varepsilon)\varepsilon_k$), we write

$$\Gamma_{\parallel}^{(H_1)} \approx \frac{1}{v_{ti}}\left\langle \frac{v_{\rm ph}}{|\gamma_k|}\right\rangle \frac{Q}{\varepsilon}\left(\sum_k v_{ti}^2\rho_i \frac{|\gamma_k|}{\omega_k^2+\gamma_k^2}\rho_i k_\theta v_{ti}k_{\parallel}\left|\frac{e\Phi_k}{T_i}\right|^2\right),$$

where the part in parenthesis is again nothing but $\Gamma_{\parallel}^{(\rm OD)}$ as defined in Eq. (4). With this, we can write

$$\Gamma_{\parallel}^{(H_1)} \approx \frac{1}{v_{ti}}\left\langle \frac{v_{\rm ph}}{\gamma_k}\right\rangle (\Gamma_{\parallel}^{(\rm OD)}/\varepsilon)Q,$$

or explicitly in terms of the intensity gradient as

$$\Gamma_{\parallel}^{(H_1)} \approx \frac{1}{v_{ti}}\left\langle \frac{v_{\rm ph}}{|\gamma_k|}\right\rangle \chi_{\phi,0}\left(v_{ti}\frac{s}{qR}\frac{\Delta^2}{\rho_i^2}\rho_i\frac{1}{\varepsilon}\frac{\partial\varepsilon}{\partial r}\right)Q.$$

Note that this is a nonlinear term since it is proportional to both the intensity and the pressure gradient (i.e., via heat flux). Its sign is determined by the sign of $v_{\rm ph}\times\varepsilon^{-1}\partial\varepsilon/\partial r$. For linear ITG this has the opposite trend to that of $\Gamma_{\parallel}^{(\rm OD)}$.

IV. CONNECTION TO WAVE KINETIC FORMULATION AND TO SPREADING

If we put aside the coupling to heat transport, the basic idea that the intensity gradient leads to symmetry breaking can also be directly obtained from wave kinetics. It was recently shown that, for electrostatic turbulence, the nonresonant particle momentum flux is equal to the wave-momentum flux[30,44] (up to neoclassical effects). This allows one to formulate the plasma rotation problem in terms of wave-momentum transport by taking the $k_{\parallel}$ moment of the wave kinetic equation. In other words, if $N_k=\sigma_k\varepsilon_k$ is the wave action density, we can symbolically write

$$\frac{\partial}{\partial t}\langle\int \sigma_k\varepsilon_k k_{\parallel}dk\rangle = -\frac{\partial}{\partial r}\langle\int k_{\parallel}v_{gr}\sigma_k\varepsilon_k dk\rangle = -\frac{\partial}{\partial r}\Gamma_{\parallel}^{(W)}.$$

In this formulation, it is obvious that a spatial gradient of turbulence intensity would lead to a transport of wave momentum. Moreover, since using the conservation of wave-plasma momentum, we can write

$$\Gamma_{\parallel}^{(\rm NR)} = \Gamma_{\parallel}^{(W)},$$

where NR corresponds to nonresonant, and W corresponds to the wave-momentum flux. This wave momentum flux driven by the intensity gradient would also drive a flux of nonresonant plasma momentum.

In order to show the link between wave-momentum flux and turbulence spreading, we use an expression for the diagonal part of turbulence spreading (i.e., flux of turbulent fluctuation intensity)

$$\Gamma_{\varepsilon} \propto -\chi_{\varepsilon}\frac{\partial\varepsilon}{\partial r}, \quad (9)$$

which is nothing but a Fick's law for turbulence intensity and can be obtained from mixing length arguments[45] or rigorous two-scale closure calculations.[46] Multiplying both sides of Eq. (9) by $\langle k_{\parallel}\rangle$, we get the part of the wave-momentum flux that is driven by turbulence intensity flux,

$$\Gamma_{\parallel}^{(w)} = -\chi_{\phi}^{(w)}\frac{\partial\varepsilon}{\partial r}, \quad (10)$$

where

$$\chi_{\phi}^{(w)} \equiv \frac{\Sigma_k\chi_{\varepsilon,k}k_{\parallel}\sigma_k\varepsilon_k}{\Sigma_k\sigma_k\varepsilon_k},$$

which has the same sign as the average wave-momentum density. This is of course only a "part" of the total wave-momentum flux, since there is also the transport of $\langle k_{\parallel}\rangle$ itself, as well as an inevitable off diagonal wave-momentum flux term (i.e., $\Gamma_{\parallel}^{(w)}$ proportional to things other than $\partial\langle k_{\parallel}\varepsilon\rangle/\partial r$).

Note that, in the case of electrostatic turbulence, it is the total wave-momentum flux that should equal the total nonresonant particle momentum flux. This means that the off-diagonal component of the wave-momentum flux should correspond to the diagonal component of the plasma momentum flux (that is, it should be proportional to $\partial v_{\phi}/\partial r$!) in order that the two are actually equal. In other words, an expression of the form,

$$\Gamma_{\parallel} = \Gamma_{\parallel}^{(w)} = -\chi_{\phi}^{(w)}\frac{\partial\varepsilon_k}{\partial r} - \chi_{\phi}\frac{\partial v_{\phi}}{\partial r}, \quad (11)$$

can be used to describe the momentum flux of waves and nonresonant particles, where what is diagonal and what is off-diagonal changes depending on which quantity we consider. Notice that in the above expression, $\chi_{\phi}^{(w)}$ is proportional to the wave momentum so that it can, in principle, have either signs.

Of course, in general this is not the whole flux of plasma momentum either. There are terms, as we have shown, related to temperature and density gradients (for instance, through the connection to heat flux and particle transport), the effects of resonant particles, fast particle losses, electromagnetic fluctuations and other things. In particular, the introduction of a sheared $E\times B$ flow would modify the above picture, which we know to be important in the pedestal region. The only reason we drop these terms is to keep the focus on the physics of the intensity gradient.

The weak point in the above argument is the sudden transition from the discrete "localized eigenmodes" picture to the fully developed turbulence picture intrinsic in the wave kinetic description (at least the way we use it). The connection can be made clearer if we note that the wave kinetic flux is by definition proportional to $v_{\rm gr}\propto 2k_y\langle x\rangle$, at which point the intensity gradient can be employed to obtain a nonzero $\langle x\rangle$, which gives $v_{\rm gr}\propto\partial_r\varepsilon$ as a nonlinear group velocity. Since the wave kinetic flux is $\Gamma^{(w)}\propto v_{\rm gr}\varepsilon$, we again recover the form (10) with $\chi_{\phi}^{(w)}$ proportional to ε.

A. Effects on heat and particle transport

We showed that a gradient in the fluctuation intensity leads to a transport of momentum. It is also well-known that such a gradient leads to transport of fluctuation intensity itself (i.e., spreading). It is then natural to ask, if it has a direct effect on particle and heat transport. One such effect might be associated with the higher order terms in an expansion in terms of $k_{\|}v_{\|}/\omega$ that arise in the quasilinear expression for flux that would normally vanish due to symmetry (note that since the plasma rotates, $\langle v_{\|}\rangle$ is not zero). However this effect would be of the form (i.e., $\Gamma_n^{(\text{ext.})}\equiv\Gamma_n-\Gamma_n|_{\partial\varepsilon/\partial x=0}$ is the additional term in particle flux driven by the intensity gradient effect)

$$\Gamma_n^{(\text{ext.})}\propto-\langle k_{\|}k_y\rangle\langle v_{\|}\rangle\frac{\partial n}{\partial r},$$

where the symmetry breaking is required so that $\langle k_{\|}k_y\rangle\neq 0$, which can be explained by an intensity gradient. This effect would be nonlinear in the gradients [i.e., $(\partial\varepsilon/\partial r)(\partial n/\partial r)$].

However there is also an indirect way by which the mechanism that is discussed here would effect particle and heat transport in tokamaks. This is via turbulence spreading. We usually tend to think of anomalous transport as a local process. This means that if the turbulence drive is absent in a certain region, we would take the flux in that region as being neoclassical. However the lack of turbulence in one region implies the existence of a turbulence intensity gradient (assuming there is turbulence somewhere in the machine). This would lead to turbulence spreading, and as we showed in this paper, to momentum transport. This would in turn lead to anomalous particle and heat transport even in regions where the local analysis would predict neoclassical transport.

V. RESULTS AND CONCLUSION

We have shown that a radial gradient of turbulence intensity, leads to parallel wave-number symmetry breaking and thus to a residual contribution to the parallel Reynolds stress. This is an important observation, since residual stress is essential for intrinsic rotation, or for anomalous momentum transport problem. Intensity gradients are of particular interest since they occur in regions of strong profile curvature, and thus may be found at the two ends of H-mode pedestals and ITBs. These are structures that can support large temperature gradients and thus can drive intrinsic rotation. Intensity gradients will also naturally result from L-mode profile structure (i.e., fluctuation intensity rises toward the L-mode edge), and so this effect may be a relevant symmetry breaking mechanism for the L-mode. The L-mode rotation may in turn affect the L-H power threshold. Thus, the mechanism described here, may be of great interest to various aspects of enhanced confinement.

We have likewise shown that a locally nonzero residual stress (with a flux toward or away from a mode rational surface) survives also in the limit of well separated mode rational surfaces, such as weak or reversed shear profiles. In addition, we pointed out the connection between heat flux and that of residual momentum flux, showing that an off-diagonal momentum flux term directly proportional to heat flux does indeed exist.

Using the mechanism of turbulence intensity gradient as the source of residual stress, we observe that momentum transport can be coupled to turbulence spreading at a phenomenological level. This may explain the intrinsic spin-up (of plasma, by turbulence intensity gradient) occurring at a faster time scale,[47] possibly corresponding to the spreading time scale, compared with the time scale at which the rotation profile naturally evolves (i.e., transport time scales). The spreading time scale is roughly the geometric mean[45] of the time scale associated with the microinstability (e.g., the linear growth) and the time scale associated with transport (e.g., roughly $\tau\propto a^2/D$, where a is the system size and D is the diffusion coefficient).

Note that we do not present these results as a complete model of anomalous momentum transport. This is rather obvious, since we do not have any profile evolution (neither of density nor of temperature), nor any detailed study of the structure of the flux (for example the E_r shear effect, which is known to be important, is dropped for convenience). We believe that the approach is justified however, because we introduce a new explanation of the residual stress (and thus intrinsic rotation), and think that the narrow focus on the simple idea helps understanding of the particular mechanism better than an all inclusive model.

The effect that is discussed here, could possibly be seen in L-mode, where there are large intensity gradients near the edge, and the effect of $E\times B$ shear is less pronounced compared to the H-mode. We expect the H-mode case still dominated by the $E\times B$ shear effects. Of course, it should also be noted that the two effects are linked since the $E\times B$ shear (and in particular the curvature of the radial electric field $\partial^2E_r/\partial r^2$) inevitably leads to an intensity gradient.

Finally, we also noted that one way to decouple the effect of an intensity gradient from that of temperature gradient, experimentally, is to look for dependence of toroidal rotation on profile curvature (i.e., $\partial^2T/\partial r^2$). Such dependence, if observed, would strongly suggest that the effect of intensity gradient on intrinsic rotation is indeed important. A similar study of distinguishing effects of E_r curvature as opposed to E_r shear would also shed light on the role of intensity gradient.

It is noteworthy that the edge intensity gradient may well change sign from L to H, which would suggest that the effect discussed here would also change sign. We think that this effect is robust and ubiquitous. As such, for instance, it should be present even in a case where neoclassical effects are known to be important. In the general case, of a neoclassical structure of the transport equations (say of v_ϕ and v_θ), the effect that is mentioned here will be purely for the parallel direction (hence would appear in both v_ϕ and v_θ but would not enter the force balance, since the contributions from v_ϕ and v_θ to the radial electric field would cancel each other).

A similar term exists for the direction perpendicular to the magnetic field (and radius). This term can be treated using an approach based on fluctuating radial current due to nonambipolarity of turbulent transport.[48] Thus, in its general form, the neoclassical momentum transport equations would be modified by adding turbulent stresses including the residual contributions in parallel and perpendicular directions. Note that perpendicular contribution would dominate the poloidal flow, as suggested in Ref. 48, and is important even for the toroidal direction (even though it is multiplied by

B_θ/B_ϕ). We leave the discussion of the effect of this perpendicular residual stress on toroidal rotation to a future publication.

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