Dynamics of zonal-flow-like structures in the edge of the TJ-II stellarator

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Abstract
The dynamics of fluctuating electric field structures in the edge of the TJ-II stellarator, which display zonal-flow-like traits, is studied. These structures have been shown to be global and affect particle transport dynamically (Alonso J et al 2012 Nucl. Fusion 52 063010). In this paper we discuss the possible drive (Reynolds stress) and damping (neoclassical viscosity, geodesic transfer) mechanisms for the associated $E \times B$ velocity. We show that (a) while the observed turbulence-driven forces can provide the necessary perpendicular acceleration, a causal relation could not be firmly established, possibly because of the locality of the Reynolds stress measurements, (b) the calculated neoclassical viscosity and damping times are comparable to the observed zonal-flow relaxation times and (c) although an accompanying density modulation is observed to be associated with the zonal flow, it is not consistent with the excitation of pressure sidebands, as those present in geodesic acoustic oscillations, caused by the compression of the $E \times B$ flow field.

(Some figures may appear in colour only in the online journal)

1. Introduction

Mass flows are an active research topic in magnetic confinement fusion due to their importance for plasma stability and confinement. In recent years, both experimental and theoretical efforts have been made to improve the understanding of the momentum transport mechanisms that determine the observed plasma rotation profiles in tokamaks (see, e.g. [2]). In stellarators without a direction of symmetry in the magnetic field strength, the neoclassical non-ambipolar fluxes are formally dominant in determining the equilibrium radial electric field and plasma rotation. Mass flows in such a system experience a magnetic viscous damping in all directions, which tends to keep their amplitude low [3].

Fluctuating, flux surface-constant zonal flows have a specific interest, for they are thought to be instrumental in turbulent transport control and to be a natural state of the drift-wave-type turbulence. The question as to what extent different magnetic configurations can host zonal-flow structures has attracted some attention [4–6]. For non-symmetric systems, Sugama and Watanabe [4] found the suggestive result that neoclassical optimization could lead to reduced ZF damping and a turbulent transport optimization as a ‘welcome side-effect’. Previously, this was found experimentally in LHD’s inward-shifted configuration [7].

Experimentally, zonal flows manifest themselves as time correlations in two electric potential measurements taken on the same flux surface. The parallel connection length between the probes must be longer than the typical parallel length scale of the field-aligned, turbulent eddies. Evidence of long-range correlations (LRCs) in distant electric potential signals was first found in the CHS stellarator [8] with a double heavy ion beam probe system and later in the TJ-II stellarator with a double Langmuir probe system [9]. Shortly afterwards other devices reported similar findings (see [10] and references therein).

In this paper we continue the analysis of the spatio-temporal characteristics of zonal-flow-like floating potential structures initiated in [1]. In that reference, a method to extract the ZF component from a set of long-range correlated signals was introduced and tested. The reconstructed structure showed actual traits of a zonal flow as a collective, fluctuating...
and transport-regulating structure (see the next section). Our concern in this work is the dynamics of these structures: their drive and damping mechanisms. We investigate the involvement of the Reynolds stress in the generation of the zonal-flow structures and the neoclassical collisional damping and geodesic transfer as possible damping mechanisms.

2. Experimental set-up description

Experiments were carried out in the four-period, flexible, low-shear stellarator TJ-II using electron cyclotron heated plasmas. The typical magnetic field strength is \( \sim 1 \text{T} \) and the rotational transform \( i/2\pi(a) = 1.65 \) at the last closed flux surface (LCFS) (typical minor radius \( a \) is 20 cm), dropping slowly to a central value of \( i/2\pi(0) = 1.55 \). Electron temperatures range from 1 to 0.8 keV at the centre to an edge value of 100–50 eV. The floating potential measurements presented in this work were obtained with a 2D array of \( 5 \times 4 \) tungsten pins, with radial and poloidal separations of 6 mm and 3 mm, respectively (see figure 1). Simultaneous measurements were made with another pin in a different sector of the machine, having a long parallel connection length with the 2D probe (\( \gtrsim 100 \text{m} \)). The sampling rate of the floating potential signals was 2 MHz.

This double-probe set-up, or a similar one, has been used in TJ-II [9] and other devices [10] to monitor LRCs in floating potential signals. In TJ-II distant probe correlation values of 0.5 and larger are consistently observed close below a critical line-averaged density \( n_{cr} \approx 0.6 \times 10^{19} \text{m}^{-3} \). Above this density, the edge radial electric field changes sign from positive (electron root) to negative (ion root). This is due to the increase in electron collisionality, which modifies the electron radial fluxes and the ambipolar electric field [11].

In the next section we analyse several plasma discharges with line-averaged densities close to the critical density that display LRCs. So as to investigate the dynamics of the collective, zonal-flow-like, potential oscillations, we make use of an analysis technique introduced in our previous work [1] that we briefly explain below.

3. Analysis of the dynamics of the zonal-flow like events

In a previous work [1] we introduced a method to extract the spatial and temporal characteristics of the zonal-flow-like component of the long-distance correlated floating potential signals. This extraction is based on the singular value decomposition. Here we give a brief explanation of the results of that analysis and refer the interested reader to the original reference for details.

A large matrix \( M \) if formed defining its entries as

\[
M_{ij} = \phi_i(t_j),
\]

where \( \phi_i(t_j) \) is the \( i \)th floating potential evaluated at the \( j \)th time sample. Here \( t_j = (j - 1)\Delta t \) and \( \Delta t^{-1} = 2 \times 10^6 \text{Hz} \). The size of \( M \) is thus \( m \times n \), where \( m \) is the number of floating potential signals (20 + 1 in our case) and \( n \) is the number of samples in the time window under analysis (6 \( \times 10^4 \) in our case).

The SVD of \( M \) can be written as the sum

\[
M = \sum_a \sigma_a u^a(v^a)^\dagger.
\]

where \( \{u^a\}_{a=1..m} \) and \( \{v^a\}_{a=1..n} \) are the orthonormal vector sets and are called topos and chronos, respectively. The spatial information is coded in the topos set whereas the temporal evolution is contained in the chronos set. The singular values \( \sigma_a \) are non-negative and are sorted in decreasing order, i.e. \( \sigma_1 \geq \sigma_2 \geq \cdots \geq \sigma_n \geq 0 \). Their squared values are the contributions of each of the modes \( a \) to the total signal energy, i.e. \( E = \sum_{ij} |\phi_i(t_j)|^2 \equiv \sum_{a=1}^{\text{rank}(M)} \sigma_a^2 \).

In [1] we tentatively identified the zonal-flow component as

\[
\Phi_{ZF} = \sigma_1 u^1(v^1)^\dagger,
\]

This identification was motivated by the empirical observations that this perfectly correlated spatio-temporal structure was (a) low-\( \kappa_0 \), (b) broadband spectrum but dominated by low (\( \lesssim 10 \text{kHz} \)) frequencies (no clear peak indicative of GAMs was observable), (c) explained most of the observed LRC. Consistently, we then showed that (d) it was collective, i.e. fluctuations occurred in the 2D probe and the distant probe simultaneously, and (e) the zonal-flow amplitude defined as \( A_{ZF}(t_j) = (v^1)^\dagger(t_j) \) dynamically modulated the outward particle transport, as shown by all H-alpha monitors around the device. We refer the reader to the cited reference for a more detailed explanation of these points.

3.1. Conditionally averaged \( E_r \) excursion and dynamical parameters

In this section, we make use of the extraction technique presented above to study the temporal and spatial characteristics of the ZF structure. Regarding the spatial structure we are interested to know what is the radial profile of the potential modulation and also to verify our definition (equation (2)) showing that the modulations are flux-surface collective. This is carried out in figures 2 and 3.
Figure 2. Long-range correlated, zonal-flow-like floating potential modulations for three TJ-II shots. Each profile is conditionally average to $A_{ZF}$ values in the different regions of its pdf, as shown in the inset. The conditionally averaged distant probe potential is shown using squares. The bars in the left figure correspond to 1 standard deviation.

Figure 3. Typical evolution of the floating potential excursions shown in figure 2. Each trace corresponds to the poloidal average of the floating potential signals at each radius. This spatial average is then de-offset and conditionally averaged to $A_{ZF}$ values in different regions of its probability density function (PDF, shown in the inset). The different colours code different deviations from the mean, i.e. most likely, profile shown in black. The modulations in the distant probe are plotted using square symbols located around the radial position of the peak LRC. In these shots the distant probe was held fixed. It displays modulations of the same sign though of somewhat lower amplitude as those in the 2D probe.

In figure 3 we show how the floating potential excursions typically (conditionally averaged to $A_{ZF}/\sigma > 1$) proceed in time. Again, the collective nature of the excursions is observed in the similar waveforms of the 2D and distant probe evolutions. A temporal asymmetry is noticeable in these structures. In the following sections we obtain the dynamical parameters associated with the time evolution of $E \times B$ velocity and compare the result with the turbulent Reynolds stress and neoclassical viscosity as possible drive and damping mechanisms.

3.2. Relation to turbulent Reynolds stress forces

A typical $E \times B$ velocity excursion is shown in figure 4(a). This is a derived quantity from the floating potential evolution shown in figure 3 for the shot TJII#26240 by subtracting the traces at $\rho = 0.86$ and $\rho = 0.90$ to get an estimate of the radial electric field at $\rho = 0.88$. We model the evolution as the competition of a drive term $R(t)$ and a magnetic viscous damping term, i.e.

$$\frac{d}{dt}V_E = R(t) - \nu V_E.$$  

We can get an estimate of the damping time $\nu^{-1}$ by fitting an exponential to the $V_E$ decay in figure 4(a). The clear peak in the typical floating potential evolution (figure 3) is indicative of a sudden change in the dynamics (e.g. the forcing) of the structure. We therefore assume that flow freely decays and the forcing $R$ vanishes on average (recall this waveform is conditionally averaged) for $t \gtrsim 30 \mu s$. The so obtained decay time is $\nu^{-1} = 25 \mu s$, similar to the one obtained in [1] and in biasing relaxation experiments [12]. This value is then used to integrate equation (3) to infer the acceleration $R(t)$.
assessing whether the viscosity is constant in time (figure 4(b)). A time-dependent viscosity has been derived in [13], which shows that the effective collision frequency involved in the viscosity is augmented by the addition of the growth rate of the ZF. Using this simplifying hypothesis we ignore this correction, and note that in general this leads to an underestimation of the inferred acceleration.

A natural question is what physical mechanism provides the required force or acceleration. The pin distribution of 2D probe allows us to compute the turbulent flux of poloidal (or more precisely perpendicular) momentum $\Gamma_{B0} = \bar{v}_{E,\theta} \bar{v}_{E,\theta} (i.e. the perpendicular Reynolds stress (RS)) at the location of the six central pins. We can then estimate the RS acceleration, given by the divergence of the flux, as

$$a_{RS} = -\frac{d}{dr} \langle \bar{v}_{E,\theta} \bar{v}_{E,\theta} \rangle,$$

(4)

where the radial derivative is approximated by finite differences and the poloidal average $\langle \cdot \rangle$ is the mean of the poloidally separated $\Gamma_{B0}$ measurements at each radial position. Figure 4(b) shows the conditional average of $a_{RS}(\rho = 0.88)$ over the same time windows used to get the $V_E$ evolution shown in figure 4(a). There is no observable causal relation between the locally measured $a_{RS}$ and the $V_E$ excursions. The inset in the figure shows the PDF of $a_{RS}$, which nevertheless shows that the inferred acceleration fall within the support of the PDF. In other words: we observe RS boosts capable of providing the required acceleration, were they global, i.e. representative of the flux surface averaged RS, but observe no temporal correlation between those boosts and the $V_E$ excursions. The analysis of the other two shots casts similar conclusions.

We emphasize that this conclusion is subject to several experimental approximations in the calculation of the Reynold stress acceleration. First, we implicitly take a poloidally and toroidally localized measurement to be representative of the flux surface average. Second, we approximate the fluctuations in the plasma potential by those of the probe floating potential, thus neglecting electron temperature fluctuations.

3.3. Comparison of relaxation time with the neoclassical collisional damping rate

We begin this section by introducing a basic model of flow and viscosity. We follow [14] for the notation and choice of Hamada coordinates $(V, \theta, \xi)$. In these coordinates the magnetic field is written as

$$B = 2\pi \chi'(V) e_\theta + 2\pi \psi'(V) e_\xi,$$

(5)

where $\chi$ and $\psi$ are the poloidal and toroidal magnetic fluxes, respectively, and prime ($'$) denotes derivation with respect to their argument (the volume enclosed by the flux surface $V$). The $s$-species flows are assumed to be dynamically incompressible and tangent to the flux surfaces. They are compactly written as

$$u_s = E_s(\psi) e_\theta + \Lambda_s(\psi) B,$$

(6)

in terms of two flux constants $E$ and $\Lambda$. The former can be written in terms of the $s$-pressure $p_s$ and electric potential $\phi$ as $E_s(\psi) = 2\pi ((p_s'(\psi)/Z_s n_s e) + \psi'(\psi))$. With this definition, and the property $e_\theta \times B = \nabla \psi/2\pi$ (valid for any magnetic coordinate system) the perpendicular $E \times B$ plus diamagnetic flow is recovered from $B \times (equation (6))$. The parallel part of $e_\theta$ is the Pfirsch–Schlueter flow, as the Hamada poloidal covariant base vector satisfies $\nabla \cdot e_\theta = 0$ and $\langle e_\theta \cdot B \rangle = (\mu_0/2\pi) I_T (= 0$ for a zero toroidal current stellarator). The $\Lambda_s$ term is then the so-called ‘bootstrap’ flow $\Lambda_s = (u_s \cdot B) / \langle B^2 \rangle$.

In our extraction of the zonal-flow dynamical parameters, namely acceleration and relaxation time, we assumed a freely decaying flow and fitted an exponential to get a damping rate. Consistently with this interpretation, we would expect the decay of $E_r \equiv -\phi' (r)$ to be determined by the return of neoclassical currents to an ambipolar equilibrium only delayed by the ion inertia through the ion polarization current. To make this more explicit we start from the momentum balance equation summed over species

$$m_i \frac{d}{dt} \langle u_i \rangle + \nabla \cdot \Pi_i + \Pi_e = j \times B,$$

(7)

where we keep only the ion’s inertia ($m_e \ll m_i$). Projecting this equation on $e_\theta$ and taking flux surface averages we obtain

$$m_i \frac{d}{dt} \langle e_\theta \cdot n u_i \rangle + \langle e_\theta \cdot \nabla \cdot \Pi_i \rangle + \langle e_\theta \cdot \nabla \cdot \Pi_e \rangle = \frac{(j \cdot \nabla \psi)}{2\pi} = 0,$$

(8)
where the quasi-neutrality condition $\nabla \cdot j = 0$ is used to obtain the last identity. The first term on the lhs of equation (8) is the ion polarization current, while the second and third terms are recognized as the ion and electron neoclassical non-ambipolar currents.

In the above discussion electron viscosity is usually neglected, for $\Pi_e/\Pi_i \sim \sqrt{m_i/m_e} \ll 1$ for equal ion and electron temperatures. While this approximation may be correct in ion-root plasmas, we need to consider electron viscosity and neoclassical flux to model our low-density electron-root plasmas. In this parameter region, the ion and electron neoclassical fluxes display a similarly strong $E_r$ dependence around its equilibrium value [11] and therefore both contribute to the return to the neoclassical ambipolarity.

The projection of the neoclassical pressure tensor is approximated in terms of the flow components:

$$\langle e_0 \cdot \nabla \cdot \Pi_i \rangle = \mu_p^\epsilon E_r + \mu_p^\circ A_r,$$

(9)

where $\mu_p^\epsilon$ and $\mu_p^\circ$ are the poloidal and bootstrap viscosities, respectively. We refer the reader to [15] for further details. Analytical expressions for the viscosities coefficients in the plateau regime (the relevant regime for the edge of TJ-II plasmas) can be found in [16]. These expressions were found to be in good agreement with the numerical solution of the drift-kinetic equation with DKES in LHD. A similar agreement has been found in TJ-II [17]. In this reference it is also shown that the poloidal viscosity term in equation (9) is dominant in the edge of TJ-II. Using this approximation in equation (8) we obtain

$$m_{in} \langle e_0 \cdot e_0 \rangle \frac{\partial E_i}{\partial t} = -(\mu_p^\epsilon + \mu_p^\circ)\langle E_i - E_{i0} \rangle + R.$$

(10)

Here, $E_{i0} = (2\pi/n_e)(p_i^\epsilon + p_i^\circ)\mu_p^\epsilon/\langle \mu_p^\epsilon + \mu_p^\circ \rangle$ represents the ambipolar ion poloidal flow obtained by zeroing the time derivative of equation (8). $R$ represents other possible momentum fluxes (i.e. Reynolds stress) not captured in equation (9). Therefore, our estimate of neoclassical damping rate is given by

$$\nu_p = \frac{\mu_p^\epsilon + \mu_p^\circ}{m_{in} \langle e_0 \cdot e_0 \rangle}.$$

(11)

A usual approximation of the inertia term $\langle e_0 \cdot e_0 \rangle$ for large aspect-ratio and circular cross-section is $\langle e_0 \cdot e_0 \rangle \approx r^2(1 + 2q^2)$. In the case of TJ-II this approximation is only correct within a factor of 2.

The poloidal viscosity is related to the non-ambipolar radial diffusion and was computed with DKES [18]: the mono-energetic neoclassical coefficients were calculated following [19] and then convolved with a Maxwellian distribution describing the equilibrium plasma. Infact, it can be seen that equation (9) is equal to the equations proposed by Sugama and Nishimura [19] for poloidal mass flow balance neglecting the parallel and poloidal heat fluxes (i.e. keeping only the particle flux terms). More precise calculations are underway. Also note that the calculated viscosities correspond to a fixed value of $E_r$, namely the ambipolar electric field calculated self-consistently [20]. We, therefore, assume that the viscosity nonlinearity is not very strong, as is expected in the low poloidal Mach number conditions discussed here. Finally, we must again note that the relatively fast ($\sim 10 \mu s$) time scales involved in the rise and decay of the potential structures would require a time-dependent treatment of the viscosity (i.e. without neglecting the explicit time derivative of the distribution function in the drift-kinetic equation [13, 4]).

A comparison of the estimated NC decay times with the experimental values is shown in figure 5. Despite the aforementioned approximations, reasonable (within a factor $\times 2$) agreement is found. A more systematic comparison would require local temperature and density measurements (to be used as an input for the viscosity calculations) and the extension of these studies to magnetic configurations with different damping characteristics, as those shown in figure 5. This is left to a future work.

3.4. Density profile modulations and comparison with density compression by the $E \times B$ flow

In [1] we reported on the correlation between the ZF amplitude and the $H_\parallel$ monitors as a proxy of the outward flux of particles. Positive ZF amplitude (corresponding to a positive increment of $E_r$) was seen to be dynamically correlated with decreased levels of $H_\parallel$ radiation. This correlation is shown in figure 6(a) (inset) for the shot no 26971. In this shot the 2D probe was operated with columns 1, 2 and 4 (see figure 1) operated in sat) mode. The extraction method outlined above was applied to the floating pins together with the distant probe potential to obtain the $\phi_t$ modulations shown in figure 6(a) with the estimated radial electric field $\langle -\frac{\partial E_i}{\partial \phi_t} \rangle$ and $I_{sat}$ modulations shown in figures 6(b) and (c). The line-averaged density is somewhat higher than for the shots shown in figure 2 and the average floating potential profile displays more negative values approaching the electron-to-ion root transition.

The ion saturation current depends mainly on the plasma density and electron temperature $I_{sat} \propto n\sqrt{T_e}$, so its
with the second term in the rhs being a factor $n_1/n_0$ smaller than the first. Compression of $v_E = (B \times \nabla \phi(\psi)/B^2)$ is due to toroidicity $\nabla \cdot v_E \approx -2v_E \cdot \nabla \ln B \approx -2v_E \cdot \kappa$, where $\kappa$ is the field line curvature and the approximations are good for low-$\beta$ plasmas. Note that the compression of the mean $E \times B$ velocity is compensated by the Pfirsch–Schlueter parallel flow, so we only consider the compression of the ZF velocity field.

The inset in figure 6(c) shows the values of $2\kappa_s \equiv 2(\nabla \psi \times B/B |\nabla \psi|) \cdot \kappa$, which according to the above discussion relate to the increments of density by

$$\frac{\Delta n}{n_0} \approx 2\kappa_s \frac{E_{ZF}^2}{B} \Delta t,$$

where $E_{ZF}^2 = -d\phi_{ZF}/dr$. The values of $2\kappa_s$ are positive in the outer probe region $\sim 0.5 \text{ m}^{-1}$, become zero around $\rho = 0.85$ and turn slightly negative further inside. Figure 6(c) shows a comparison of the actual density modulation with the $E \times B$ compression estimate given by equation (12) with $\Delta t = 20 \mu s$ (see figure 4) shown in grey. This estimate casts considerably smaller relative modulations (a few per cent) compared with the observed $\sim 25\%$. It is well known (see [23]) that the amplitudes of density perturbations in geodesic acoustic modes are much smaller than their electric potential counterpart. Another prominent disagreement in this comparison is that the location of maximum $E_r$ modulation and expected density compression coincides with the radial position where $I_{sat}$ hardly varies. All the above observations indicate that the observed density modulations are not due to the compression of the background density by the $E \times B$ flow, but rather due to the radial transport of particles in agreement with our previous findings [1].

4. Summary and conclusions

In this work we have studied the dynamics of the zonal-flow like-structures responsible for the long-range floating potential correlation observed in the TJ-II stellarator under certain conditions [24]. We made use of a previously introduced [1] ZF extraction method to discuss the possible drive (Reynolds stress) and damping (neoclassical viscosity, geodesic transfer) mechanisms for the associated $E \times B$ velocity. We show that (a) while the observed turbulence-driven forces can provide the necessary perpendicular acceleration, a causal relation could not be firmly established, possibly because of the locality of the Reynolds stress measurements. (b) The calculated neoclassical viscosity and damping times are comparable to the observed zonal-flow relaxation times. The configuration dependence of this decay will be addressed in a future work. (c) Although an accompanying density modulation is observed to be associated with the zonal flow, it is not consistent with the excitation of pressure sidebands, as those present in geodesic acoustic oscillations, caused by the compression of the $E \times B$ flow field.

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