Turbulence control by energetic particles

D. Zarzoso, Y. Sarazin, X. Garbet, R. Dumont, A. Strugarek[1], J. Abiteboul,
T. Cartier-Michaud, G. Dif-Pradalier, Ph. Ghendrih, V. Grandgirard, G. Latu and O. Thomine
CEA, IRFM, F-13108 Saint-Paul-lez-Durance, France.
[1] *Laboratoire AIM Paris-Saclay, CEA/Irfu Universite Paris-Diderot CNRS/INSU, 91191 Gif-sur-Yvette, France*

Analysis of the turbulent transport in present tokamaks is essential to predict the performance of the next step fusion devices. Although it is well-known that a low-frequency zonal flow (LFZF) is a possible way to reduce the turbulence level, the role of the oscillatory component of ZFs, namely the geodesic acoustic modes (GAMs), has been marginally analysed in the literature (see e.g. [1] and [2]). In the context of core-turbulence suppression, the role of GAMs is not evident for they are highly Landau damped. However, a population of energetic ions can drive them unstable. These so-called EGAMs have been predicted theoretically [3], and observed experimentally [4] and numerically [5] in gyrokinetic simulations using the full-f and global 5D GYSELA code [6]. The possibility to efficiently excite EGAMs in gyrokinetic simulations opens the way to the possible turbulence control by external ways. In this paper, first numerical evidence of the excitation of EGAMs in a fully developed ITG turbulence is reported as well as the associated improved confinement due to the turbulence reduction.

GYSELA solves the standard gyrokinetic equation in the electrostatic limit for the ion guiding-centre distribution function in a simplified magnetic topology consisting of toroidal flux surfaces with circular poloidal cross-sections. The self-consistent system is composed by the gyrokinetic equation coupled to the quasi-neutrality equation

$$\frac{\partial F}{\partial t} + \mathbf{v}\cdot\nabla F + \dot{v}_{\parallel}\frac{\partial F}{\partial v_{\parallel}} = \mathscr{C}(F) + \mathscr{D}(F) + S \tag{1}$$

$$\phi - \langle\phi\rangle - \frac{1}{n_{\mathrm{eq}}}\nabla_{\perp}\cdot\left(\frac{n_{\mathrm{eq}}}{B^2}\nabla_{\perp}\phi\right) = \frac{n - n_{\mathrm{eq}}}{n_{\mathrm{eq}}} \tag{2}$$

where electrons are assumed adiabatic, $\mathscr{C}(F)$ is the collision operator, $\mathscr{D}(F)$ a dissipation term which is non vanishing in narrow buffer regions at radial boundaries, the electrostatic potential is normalized to the thermal energy, the subscript *eq* stands for equilibrium quantities, the brackets $\langle\rangle$ represent a flux-surface average, B is the magnetic field and n the ion guiding-centre density. The source term S on the right-hand side of Eq. 1 is essential in flux-driven simulations on energy confinement time scales, for it drives equilibrium gradients and therefore neoclassical and turbulent fluxes. Notice that the heating source term can lead to the inversion of the negative slope of the distribution function with respect to the energy, which clearly represents an out-of-equilibrium state. Thus, the system tends to relax to a Maxwellian unless it is forced by additional external sources aiming at inverting the population of the initial Maxwellian distribution. The source S is decomposed into two terms, $S = S_{\text{th}} + S_{\text{fp}}$, where S_{th} is a source of energy localized in the inner region and S_{fp} is centered around the mid radial position $r_{\text{mid}} = (r_{\text{min}} + r_{\text{max}})/2$. The energy dependence of S_{fp} is such that it brings the distribution function out of the equilibrium by creating a positive slope in energy, i.e. it represents the energetic particles generation. This source is built by using projections onto Laguerre and Hermite polynomial basis. The coefficients are chosen in order to inject only parallel energy. For symmetry reasons, the source is written as $S_{\text{fp}} = \frac{1}{2} S_{\text{fp},0} S_r(r)(S_+ + S_-)\, \mathrm{e}^{-\bar{\mu} B(r,\theta)}$ where $S_{\text{fp},0}$ is the amplitude of the source, S_r is the normalized radial envelope, B is the magnetic field and $\bar{\mu} = \mu / T_{s\perp}$. Here μ is the adiabatic invariant and $T_{s\perp}$ is a temperature, both normalized to the thermal energy. $T_{s\perp}$ is set to 1 in the simulations presented here. $S_{\pm}$ includes the decomposition onto the specified polynomial basis

$$S_{\pm} = \mathrm{e}^{-\left(\bar{v}_{\parallel} \pm \bar{v}_0\right)^2} \sum_{h,l} c_{hl} \mathscr{H}_h\left(\bar{v}_{\parallel} \pm \bar{v}_0\right) \mathscr{L}_l\left(\bar{\mu} B(r,\theta)\right) \tag{3}$$

where $\mathscr{H}_h$ and $\mathscr{L}_l$ are the Hermite and Laguerre polynomials of degree h and l respectively, $\bar{v}_{\parallel} = v_{\parallel}/\sqrt{2T_{s\parallel}}$ and $\bar{v}_0 = v_0/\sqrt{2T_{s\parallel}}$. The parallel velocities $v_{\parallel}$, v_0 and the temperature $T_{s\parallel}$ are normalized to the thermal energy. In the context of EGAM excitation, the parameters v_0 and $T_{s\parallel}$ are essential, since they enable the inversion of the slope in energy space. The effect of EGAMs on the turbulent transport is analysed by comparing two simulations, characterized by the following dimensionless parameters: the collisionality has been chosen $\nu_\star = 5 \cdot 10^{-3}$ (banana regime), the q-profile is parabolic with a value of $q(r_{\text{mid}}) \approx 2.7$, the normalized Larmor radius $\rho_\star = 1/75$ and the initial normalized temperature gradient $R/L_T = 6.5$. Both simulations are identical in the time window $0 < \omega_c t < 1.96 \cdot 10^5$. Within this interval, only the source S_{th} has been used in order to provide initial ambient turbulence (AT) spectrum in which EGAMs will be excited. This spectrum is shown in Fig. 1 by the grey region. At $\omega_c t = 1.96 \cdot 10^5$, the source S_{fp} is switched on. In one of the sim-

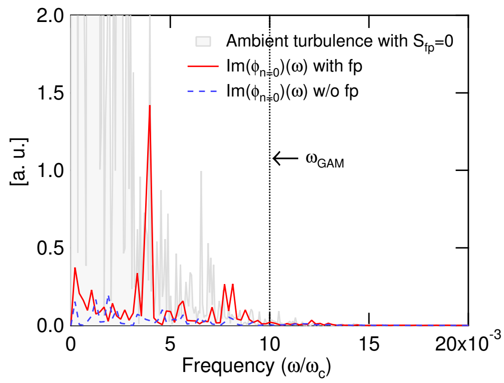


Figure 1: Time Fourier transform of the up-down asymmetric components of the axisymmetric electrostatic potential for both simulations with (solid red line) and without (dashed blue line) fast particles. The spectrum of all the fluctuations is given by the grey region.

ulations the parameters are $v_0 = 2$ and $T_{s\|} = 0.5$. In the other simulation, the parameters are $v_0 = 0$ and $T_{s\|} = 1$. The former one will be referred to as the simulation *with energetic particles* and the latter one as the simulation *without energetic particles*. In both simulations, the total injected power is the same. The only difference comes from the excitation of EGAMs by means of the convenient modification of the distribution function in the velocity space. The existence of the energetic mode is observed in Fig. 1, where the time Fourier transform of the poloidally up-down asymmetric component of the modes $n = 0$ is represented. For comparison, a vertical dashed line indicates the position of the standard GAM frequency.

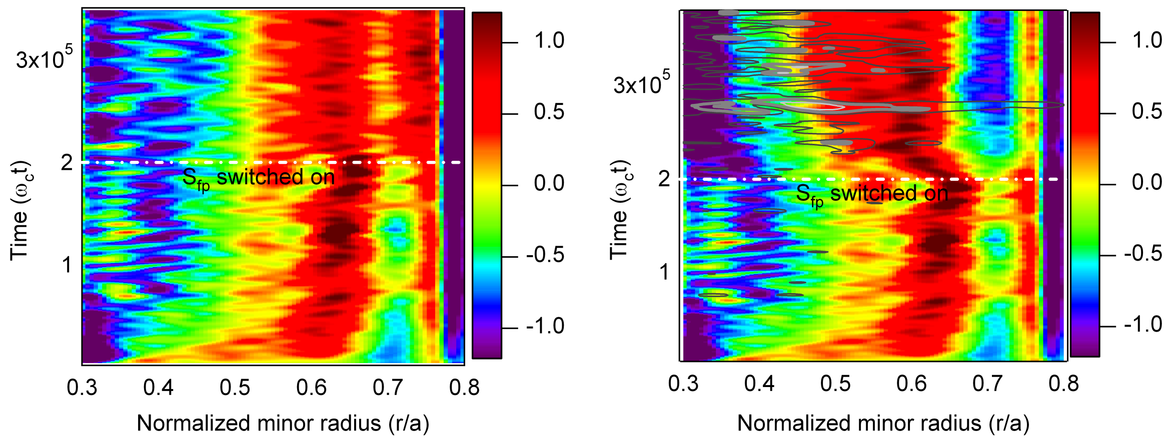
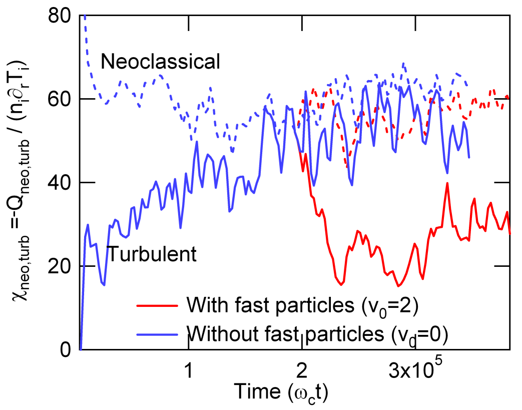


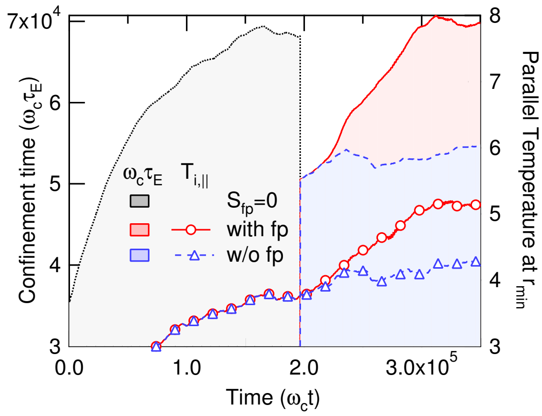
Figure 2: Colour plot of the $E \times B$ shearing rate without (left) and with (right) fast particles. In the figure with fast particles, isocontours of the mode $\Im(\phi_{1,0})$ at the EGAM frequency are given.

The time evolution of the $E \times B$ shearing rate, namely $\gamma_E \approx \frac{r}{q}\partial_r\left(\frac{q}{r}\partial_r\phi_{00}\right)$, is illustrated in Fig. 2, normalized to the estimated maximum value of the linear growth rate γ_{lin} of the ITG modes. It can be observed that the $E \times B$ shearing rate is slightly reduced in the absence of fast particles. However, when energetic particles and EGAMS are present, the shearing rate is shifted inwards. The isocontours of the mode $\Im(\phi_{1,0})$ at EGAM frequency are overlapped on the right-hand side of Fig. 2. It turns out that the radial modification of the $E \times B$ shearing rate follows the radial location of the maximum EGAM amplitude, indicating a correlation between EGAMs and ZFs. This modification leads to a reduction of the turbulent transport for radial positions $r/a > 0.5$. This is quantified by the effective turbulent transport diffusivity $\chi_{\text{turb}} = -\frac{Q_{\text{turb}}}{n_i|\nabla T_i|}$, where $Q_{\text{turb}} = \left\langle \int \mathbf{v}_E \cdot \nabla\chi\,(v_\|^2/2 + \mu B)\,J_0 \cdot F\,\mathrm{d}^3\mathbf{v} \right\rangle$. In these expressions, $\mathbf{v}_E$ is the electric drift velocity, $\nabla\chi$ is a vector orthogonal to the flux surface, J_0 is the gyro-average operator, n_i the ion density and T_i the ion temperature. The time evolution of χ_{turb} for both simulations together with the evolution of the neoclassical diffusivity χ_{neo} is illustrated in Fig. 3a.

Finally, the ratio between the energy content of the plasma and the total injected power pro-



(a) Time evolution of $\chi_{\rm turb}$ and $\chi_{\rm neo}$ for both simulations.



(b) Confinement time and ion parallel temperature at the inner radial position.

vides an estimate of the confinement time $\omega_c\tau_{\rm E} = \frac{W_{\rm int}}{P_{\rm add}}$. The time evolution of the parallel temperature at the inner radial position is also plotted. As observed in figure 3b, in the presence of EGAMs, the inner parallel temperature is increased by $\sim 25\%$ with respect to the case with no fast particles, due to the transient reduction of the turbulent transport. The overall increase of $W_{\rm int}$ leads to an increase of $\sim 30\%$ in the confinement time with respect to the case with no EGAMs.

In conclusion, we have reported on the first evidence of turbulence reduction by excitation of EGAMs in full-f gyrokinetic simulations with GYSELA code. The reduction of the effective turbulent transport diffusivity, leads to an improved core confinement.

References

[1] T. Hahm *et al.*, Physics of Plasmas **6**, 922 (1999)

[2] G. Conway *et al.*, *Phys. Rev. Lett.* **106**, 65001 (2011)

[3] G. Fu, *Phys. Rev. Lett.* **101**, 185002 (2008)

[4] R. Nazikian *et al.*, *Phys. Rev. Lett.* **101**, 185001 (2008)

[5] D. Zarzoso *et al.*, *Phys. Plasmas* **19**, 022102 (2012)

[6] V. Grandgirard *et al.*, *Comm. Nonlinear Science Num. Sim.* **13**, 81 (2008)