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A model for the turbulence in the scrape-off layer of tokamaks

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A MODEL FOR THE TURBULENCE IN THE SCRAPE-OFF LAYER OF TOKAMAKS

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ABSTRACT. The fluid theory of electrostatic perturbations in a scrape-off layer (SOL) plasma is analysed. The main difficulty is that the edge is theoretically found to be stable, while experimentally it is unstable. A possible explanation relies on the fact that the commonly used ballooning representation is not correct in the SOL. An alternative representation is proposed which reproduces the instability of the edge in several simple configurations and explains many experimental features.

1. INTRODUCTION

In both divertor and limiter configurations, tokamak plasmas are surrounded by a region with open field lines — the scrape-off layer (SOL). Direct measurements of the plasma density or potential with Langmuir probes show that this region is highly turbulent, i.e. $\tilde{n}/n \approx 50\%$ [1–4]. Theoretical models greatly fail to find linearly unstable modes, since the free motion of electrons along flux lines stabilizes the rippling interchange modes [5]; collisional drift modes are also found to be stable [6]. Several models have been developed which include specific edge effects to explain the large level of turbulence, for example impurity radiation, the Z_{eff} gradient or ionization and charge exchange [7, 8]. Other attempts have been made to take into account non-linear effects due to the large amplitude of the turbulence [9, 10].

However, theoretical models of edge turbulence are often expressed using a geometry of closed field lines. The fluid equations relative to this case are summarized in Section 2 and the rippling/interchange mode is briefly discussed. In fact, the field lines in the SOL are open — a feature which drastically changes the behaviour of modes. Nedospasov [11] studied a model which takes account of open field lines, showing that flute modes could be unstable in the SOL owing to the boundary conditions of the limiters. However, the exact spatial mode structure and the stabilizing effect of the polarization current were not taken into account in this work. In Section 3, a full study of mode stability is presented in the framework of fluid theory, demonstrating that interchange modes are unstable in the SOL. The consequences of this model are demonstrated in Section 4 and it is shown that some experimental facts could be explained with this model.

2. STANDARD FLUID EQUATION FOR EDGE TURBULENCE

We consider the plasma edge in a toroidal equilibrium (major radius R, minor radius a), characterized by its temperature T or its Spitzer resistivity η and its density n. Since modes with scales larger than the ion Larmor radius are experimentally observed and the collisionality is larger at the edge, it seems reasonable to use a fluid model. Moreover, the mode is restricted to an electrostatic perturbation, since the associated vector potential fluctuation plays only a minor role in the case of typical edge plasma parameters. The response to a potential perturbation $\tilde{\Phi}(r, \theta) \exp(\gamma t - im\phi)$ is deduced using the following standard equations:

Mass conservation

$$\frac{\partial \tilde{n}}{\partial t} + \mathrm{div}\,(n\tilde{\vec{V}}_E) = 0 \tag{1}$$

with

$$\tilde{\vec{V}}_E = \frac{-\nabla\tilde{\Phi} \times \vec{B}}{B^2}$$

and hence

$$\frac{\tilde{n}}{n} = \frac{V_e^*}{\gamma}\,\frac{e}{T}\,\frac{\partial\tilde{\Phi}}{r\partial\theta}$$

where

$$V_e^* = \frac{T}{eB}\,\frac{dn}{ndr} = -\frac{T}{eBL_N}$$

is the diamagnetic velocity of electrons and e is the algebraic electron charge.

Energy conservation

Neglecting perpendicular conduction and radiation, the energy conservation can be expressed as an equation for temperature or rather for resistivity evolution, multiplying by $d\eta/dT$:

$$\frac{\partial \tilde{T}}{\partial t} + \mathrm{div}\,(T\tilde{\vec{V}}_E) = \chi_\parallel \nabla_\parallel^2 \tilde{T}$$

or

$$\frac{\tilde{\eta}}{\eta} = \frac{V_{\eta e}^*}{\gamma + k_\parallel^2\chi_\parallel}\,\frac{e}{T}\,\frac{\partial\tilde{\Phi}}{r\partial\theta} \tag{2}$$

where $\chi_\parallel$ is the parallel thermal conductivity, $k_\parallel$ is the parallel wave number and

$$V_{\eta e}^* = -\frac{T}{eBL_\eta}$$

is the diamagnetic velocity associated with the resistivity gradient length L_η.

Charge conservation

The current response to the potential perturbation is the sum of three terms:

(1) $\tilde{\vec{J}}_{curv}$ is the response due to the curvature drift of electrons ($\vec{V}_{ge}$) and ions ($\vec{V}_{gi} = -\vec{V}_{ge}$):

$$\mathrm{div}\,\tilde{\vec{J}}_{curv} = \mathrm{div}\,(2\tilde{n}e\vec{V}_{ge}) = \frac{2ne^2 V_e^* V_{ge}}{T\gamma r}$$

$$\times \left[\sin\theta\,\frac{\partial^2}{\partial\theta\partial r} - \frac{\sin\theta}{r}\,\frac{\partial}{\partial\theta} + \frac{\cos\theta}{r}\,\frac{\partial^2}{\partial\theta^2}\right]\tilde{\Phi} \tag{3}$$

$$V_{ge} = -\frac{T}{eBR}$$

(2) $\tilde{\vec{J}}_{pol}$ is the polarization current (ion contribution):

$$\mathrm{div}\,(\tilde{\vec{J}}_{pol}) = \mathrm{div}\left(-\frac{\gamma n m_i}{B^2}\,\nabla\tilde{\Phi}\right)$$

$$= \frac{-\gamma n m_i}{B^2}\left[\frac{\partial^2}{\partial r^2} + \frac{1}{r}\,\frac{\partial}{\partial r} + \frac{1}{r^2}\,\frac{\partial^2}{\partial\theta^2}\right]\tilde{\Phi} \tag{4}$$

where m_i is the ion mass.

(3) $\tilde{\vec{J}}_\parallel$ is the parallel current response which can be expressed as a function of the unperturbed loop voltage V_{loop}:

$$\mathrm{div}\,(\tilde{\vec{J}}_\parallel) = \frac{V_{loop}}{2\pi R\eta^2}\,\nabla_\parallel\tilde{\eta} - \frac{\nabla_\parallel^2\tilde{\Phi}}{\eta} \tag{5}$$

The charge conservation is given by

$$\mathrm{div}\,(\tilde{\vec{J}}_{curv}) + \mathrm{div}\,(\tilde{\vec{J}}_\parallel) + \mathrm{div}\,(\tilde{\vec{J}}_{pol}) = 0 \tag{6}$$

This problem is generally made unidimensional assuming radial invariance. The perturbation is expressed using the ballooning representation:

$$\tilde{\Phi} = \text{periodified}\,(\hat{g}(\theta))\, e^{im(q(r)\theta-\phi)} e^{\gamma t}$$

$$= \sum_{\ell} \tilde{f}(r - r_\ell)\, e^{-im\phi + i\ell\theta}$$

where $\hat{g}(\theta/sk_\theta)$ is the Fourier transform of $\tilde{f}(r)$ and r_ℓ is the radius of the resonant surface, $q = \ell/m$. A simple case is to consider highly localized modes, i.e. $\theta \approx 0$. This tends to overestimate the growth rate. The charge conservation appears as a Schrödinger equation in a parabolic potential well for $\tilde{f}$:

$$\frac{\partial^2 \tilde{f}}{\partial x^2} + \frac{1}{\rho_{thi}^2} Q(x)\tilde{f} = 0 \tag{7}$$

where $x = r - r_\ell$ and

$$Q(x) = -k_\theta^2\rho_{thi}^2 + \frac{4k_\theta^2 V_{ge} V_e^*}{\gamma^2} - \frac{m_i}{m_e}\frac{k_\theta^2 s^2 V_{thi}^2}{\nu_{coll}\gamma q^2 R^2}x^2$$

$$+ \frac{eV_{loop}}{T}\frac{m_i}{m_e}\frac{V_{thi}^2}{q^2R^2}\frac{V_{\eta e}^*}{\gamma\nu_{coll}\left[\gamma + \chi_\parallel\left(\frac{k_\theta sx}{qR}\right)^2\right]}$$

$$\times \frac{k_\theta^2 sqx}{2\pi} \tag{8}$$

where $k_\theta = \ell/r$ is the poloidal wave number, the parallel wave vector has been replaced by its expansion, $k_\parallel = (k_\theta s/qR)(r - r_\ell)$ (s is the shear parameter, equal to 2 at the edge) and ν_{coll} is the collision frequency. The four terms represent the inertia (stabilizing), the interchange effect (destabilizing since $V_{ge}V_e^* > 0$), the parallel current (stabilizing) and the rippling term (destabilizing). For typical edge parameters the conclusions that can be drawn from Eq. (8) are the following:

(a) The parallel current response is so large that the third term represents a very deep parabolic well. The corresponding solutions in $\tilde{f}(x)$ are very localized (much less than an ion Larmor radius) and hence are stabilized by the finite Larmor radius effect. Even neglecting the latter effect, moderate growth rates are found.

(b) The classical growth rate of rippling modes (see, for instance, Ref. [9]) can be found from Eq. (8):

$$\frac{\gamma R}{v_i} \approx \left\{\frac{m_i}{m_e}\frac{v_{thi}}{\nu_{coll}R}\left(k_\theta\rho_{thi}\frac{q}{s}\right)^2\left(\frac{eV_{loop}}{T}\frac{R}{L_\eta}\right)^4\right\}^{1/5} \tag{9}$$

is of order one in usual conditions. However, the corresponding mode is characterized by a radial scale δ:

$$\delta \approx \left\{\frac{m_e}{m_i}\frac{\nu_{coll}R}{v_{thi}}\frac{\gamma R}{v_{thi}}\left(\frac{q}{s}\right)^2\left(\frac{1}{k_\theta\rho_{thi}}\right)^2\right\}^{1/4}\rho_i$$

which is much smaller than the ion Larmor radius ρ_i, except at low temperatures, $T < 10$ eV. The ion Larmor radius effects should therefore stabilize the modes.

In conclusion, when fluid theory is used within the ballooning formalism, very localized and weakly unstable modes are found unless the plasma is very collisional. This is due to the large parallel heat conductivity or the current response which are strongly stabilizing for finite $k_\parallel$.

3. STABILITY OF OPEN FLUX LINES

The standard fluid theory summarized in Section 2 cannot explain the edge turbulence. The rippling term can be enhanced by adding specific edge effects such as an effective charge gradient or radiation [7, 8] and non-linear effects [8–10]; enhanced instability is generally found. However, it has been shown [12] that the rippling modes are stabilized in the SOL by line tying effects. This is not true for interchange modes because of the specific boundary condition for the current on a limiter [11]. Moreover, according to the models based on rippling modes, the edge turbulence should be very sensitive to the wall material and the plasma parameters (determining Z_{eff}), and to the limiter geometry which determines the current density in the SOL (through a shadowing effect). Moreover, Thayer et al. [12] predict that the edge turbulence should be stabilized during current drive experiments since the loop voltage is zero. All these effects have not been reported.

We propose here to take into account a more trivial linear effect, based on the following points:

— It is not realistic to assume ballooning invariance within the SOL;

— The field lines are open and the modelling of the parallel response is wrong; Equation (5) and (7) have to be modified to take into account this important feature.

Thus, the above representation is not well suited for edge turbulence. Since the high parallel conductivity prevents any variation along the magnetic field, realistic perturbations tend to be uniform along field lines. Hence, we shall look for solutions of the form $\tilde{f}(r)\, e^{im(q(r)\theta-\phi)} e^{\gamma t}$ (since the field lines are open, there are no problems of multiple determination when θ varies as 2π). This representation has also the advantage that it is well suited for an integral formulation of quasi-neutrality (Eq. (6)); the problem is simplified by expressing the quasi-neutrality of a flux tube rather than using a local expression as in Eq. (7). In the following, we shall neglect the rippling effect, which is stabilized by the parallel thermal conductivity, even with open field lines [12], and by the current density limitation due to obstacles in the SOL.

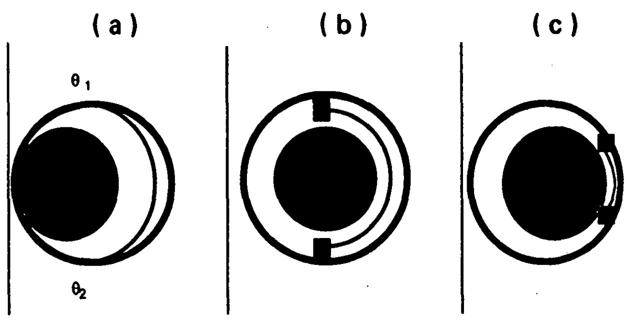

FIG. 1. Three axisymmetric configurations with unfavourable curvature. (a) Inner wall operation, (b) low field side in a divertor double-null configuration (the short stabilizing part in the divertor chamber is not represented), and (c) outer part in a belt limiter configuration.

3.1. Charge conservation in a flux tube

For the sake of simplicity, we consider only the case of an axisymmetric configuration (i.e inner wall operation, belt limiter, standard divertor). Each field line in the SOL intersects the wall or the limiter at two poloidal angles, θ_1 and θ_2, as shown in Fig. 1. Its length is $L = qR(\theta_2 - \theta_1)$. The charge conservation in a flux tube of vanishing cross-section can be expressed as

$$\int_{\theta_1}^{\theta_2} \mathrm{div}\,(\tilde{\tilde{J}}_{curv} + \tilde{\tilde{J}}_{pol})qRd\theta + \tilde{J}_\parallel(\theta_1) + \tilde{J}_\parallel(\theta_2) = 0 \quad (10)$$

$\tilde{J}_\parallel(\theta_{1,2})$ represents the parallel current through the sheath at both ends of the flux tube. When the plasma is in equilibrium, the potential adjusts itself to a Φ_0 value so that the losses are ambipolar:

$$neV_{thi} = neV_{the} \exp\left(-\frac{e\Phi_0}{T}\right)$$

where $v_{the} = \sqrt{2T/m_e}$ is the electron thermal velocity. When Φ_0 is perturbed at the end of the flux tube by $\tilde{\Phi}$, the number of reflected electrons is slightly perturbed and the resulting current is

$$\tilde{J}_\parallel(\theta_{1,2}) = -ne^2V_{thi}\frac{\tilde{\Phi}(\theta_{1,2})}{T} \quad (11)$$

Note that $\tilde{\tilde{J}}_\parallel(\theta_{1,2})$ is generally much smaller than the standard parallel current response discussed in Section 2, $k_\parallel\tilde{\Phi}/\eta$, by a factor $\sqrt{m_e/m_i}\, k_\parallel V_{the}/\nu_{coll}$.

3.2. Stability of a flux tube

Equation (10) implies the differential equation in $\tilde{f}(r)$:

$$\frac{1}{r}\frac{d}{dr}\left(r\frac{d\tilde{f}}{dr}\right) + \left(\frac{\mu}{r^2} + \nu\right)\tilde{f} = 0 \quad (12)$$

where

$$\frac{\mu}{r^2} = -\frac{k_\theta^2[G(\theta) - AH(\theta)]_{\theta_1}^{\theta_2}}{[\theta]_{\theta_1}^{\theta_2}}$$

$$\nu = \frac{-4\omega_{ci}^2}{q\gamma RV_{thi}[\theta]_{\theta_1}^{\theta_2}}$$

with $G(\theta) = s^2\theta^3/3 + \theta$ and $H(\theta) = (1 + s)\sin\theta - s\theta\cos\theta$.

The quantity $A = 4V_e^*V_{ge}/(\rho_{thi}^2\gamma^2) \approx 2V_{thi}^2/\gamma^2RL_N$ represents the interchange effect ($L_N > 0$ is the characteristic length of the pressure gradient). Since there is no resonant layer, it is reasonable to look for radial variation in the eikonal form $\tilde{f}(r) = e^{ikr}$. The dispersion equation is $k^2 = \mu/r^2 + \nu$ or

$$\left([G(\theta)]_{\theta_1}^{\theta_2} + \frac{k^2[\theta]_{\theta_1}^{\theta_2}}{k_\theta^2}\right)\left(\frac{\gamma}{V_{thi}/qR}\right)^2 + \frac{4}{k_\theta^2\rho_{thi}^2}\left(\frac{\gamma}{V_{thi}/qR}\right) - \frac{2q^2R}{L_N}[H(\theta)]_{\theta_1}^{\theta_2} = 0 \quad (13)$$

The first term (stabilizing) containts the effect of the shear and inertia, and the second term (stabilizing) represents the end losses. The third term represents the curvature effect, which is stabilizing if $H < 0$. As in the standard interchange mode theory, there are real γ values only if H is positive, i.e. if the curvature is unfavourable. The maximum growth rate corresponds to $k = 0$. In practice, G is rather large and, thus, γ is only a weak function of k for $k \leq k_\theta$. By superposing several radial modes, it is possible to construct solutions localized in the SOL where this model is correct. The most unstable modes are found when the end loss term can be neglected. This happens if

$$k_\theta \rho_{thi} > \left(\frac{2L_N}{q^2 R[H(\theta)]_{\theta_1}^{\theta_2} [G(\theta)]_{\theta_1}^{\theta_2}} \right)^{1/4} \quad \text{(typically} \approx 0.1)$$

and the corresponding growth rate is

$$\gamma = \frac{V_{thi}}{R} \left(\frac{2R[H(\theta)]_{\theta_1}^{\theta_2}}{L_N [G(\theta)]_{\theta_1}^{\theta_2}} \right)^{1/2}$$

In the opposite case, the growth rate is

$$\gamma = k_\theta^2 q \rho_i V_e^* [H(\theta)]_{\theta_1}^{\theta_2}$$

a formula which is similar to the expression obtained by Nedospasov [11].

Three schematic cases corresponding to unfavourable curvature are considered (see Fig. 1). The growth rate for $k = 0$ has been computed, taking $R/L_N = 20$:

	θ_1, θ_2	$[G(\theta)]_{\theta_1}^{\theta_2}$	$[H(\theta)]_{\theta_1}^{\theta_2}$	$\gamma \frac{R}{V_{thi}}$
Inner wall	$-\pi, \pi$	90	13	2.4
Double- null low field side	$-\pi/2$	13	6	4
Belt limiter	$-\pi/12, \pi/12$	0.6	0.6	6

4. CONSEQUENCES AND COMPARISON WITH EXPERIMENT

The contact of field lines with the wall causes interchange modes to be unstable when the average curvature is unfavourable, i.e. in many common configurations. This could explain why the SOL is unstable in the case of inner wall operation and between two toroidal belt limiters (as in JET [13]), and why the external part of the SOL is unstable in the double-null configuration (as observed in ASDEX [14]). This explains also other observations:

— The large radial scale of turbulence, i.e. $k_\perp \rho_{thi} \approx 0.1$ [2];
— Collisionality plays no role, as observed in CALTECH [4];
— The large phase shift (60°) between density and potential fluctuation [2];
— The very large phase velocity along field lines [3].

Though the limiter effect should be restricted to the SOL, it is expected that mode coupling (non-linear or toroidal) effects will ensure the continuity with the plasma bulk turbulence [15], in agreement with the radial continuity of $\tilde{n}/n$ observed experimentally [3, 4].

This model could have an impact on the design of many plasma facing components, for example divertor plates (asymmetric heat load) and the pump limiter (the pumping efficiency depends on the stability of the flux tube in the throat). The problem is not so simple in a non-axisymmetric geometry (as in the case of discrete limiters) where there is a mixture of stable and unstable field lines. This model could explain the existence of localized turbulent spots that are sensitive to the limiter configuration and the current direction, such as those observed in TEXT [16].

5. CONCLUSION

Up to now, standard models of turbulence have failed to explain the edge turbulence, since it is stabilized by the parallel dynamics of electrons. In the SOL, the field lines intersect the wall so that the stabilizing parallel currents are inhibited. When this effect is taken into account, the finite length of the lines makes interchange modes unstable when the radial scales are large, provided $k\rho_{thi} \approx k_\theta \rho_{thi} > (L_N/q^2R)^{1/4} \approx 0.1$. Thus, it is possible to explain the existence of turbulence in the SOL when the field line curvature is locally unfavourable. It is also possible to understand the radial scale of the turbulence as well as several observations related to the in-out asymmetry of the SOL.

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HIGHER HARMONIC DRIFT MODES IN TOKAMAKS

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ABSTRACT. The stability of higher harmonic electrostatic electron drift waves in toroidal geometry in studied for a representative plasma profile. Using a model driving function, it is shown that the dissipative trapped electron (DTE) mode admits solutions $\omega/\omega_e^* \gg 1$ (ω_e^* is the electron diamagnetic drift frequency) that may relate to density fluctuation measurements which have suggested the presence of such higher frequency modes. Comparison of analytic and numerical results for both outward and inward ballooning DTE modes gives good agreement for the former modes but generally poor agreement for the latter ones. Further, it is found from the numerical results that the inward ballooning modes may be more unstable than the outward ballooning modes, even though the trapped electron drive is maximized on the outboard midplane.

A number of recent theoretical works [1–3] on anomalous energy transport in tokamaks attribute confinement degradation in ohmically heated and beam heated tokamaks to the presence of electrostatic drift wave turbulence. It has been conjectured that low density neo-Alcator scaling of confinement time in ohmically heated tokamaks results from dissipative trapped electron (DTE) mode turbulence, while saturation of the confinement time in high density ohmically heated tokamaks and confinement degradation in beam heated tokamaks results from ion temperature gradient (ITG) mode turbulence. The simplified model of the DTE mode employed in these analyses has been shown [4] to apply only in the extreme collisionality regime, $\nu_{eff}/\omega_e^* \gg 10$, with ν_{eff} the effective trapped electron collision frequency and ω_e^* the electron diamagnetic drift frequency. Nevertheless, a more realistic model [4] of the DTE mode still preserves the essential scaling features of the energy confinement time. At present, these two components, the DTE and ITG modes, form the basis for the drift wave transport model, which attempts to explain bulk plasma confinement in tokamaks.

The DTE branch has received scant attention in recent years, with much more theoretical activity being devoted to the study of the ITG mode [5–10]. However, there are several important features of the DTE mode that bear directly on experimental results which have not been previously noted in the literature. Specifically, it is shown in the present work that in toroidal geometry it is possible to obtain higher harmonic DTE eigenmodes in the bulk of the plasma, $r/a \le 0.5$, with frequencies $\omega \gg \omega_e^*$ in the long wavelength range, $k_\theta \rho_s \le 0.3$ (k_θ is the poloidal wavenumber, $\rho_s = c_s/\Omega_i$, with c_s the sound speed, $c_s = (T_e/m_i)^{1/2}$, and Ω_i the ion cyclotron