Nonlinear Dynamics and Complex Behaviors in Magnetized Plasmas of Fusion Interest

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Abstract. Complexity and self-organization in burning plasmas are consequence of the interaction of energetic ions with plasma instabilities and turbulence; of the strong nonlinear coupling that will take place between fusion reactivity profiles, pressure driven currents, MHD stability, transport and plasma boundary interactions, mediated by the energetic particle population; and finally of the long time scale nonlinear (complex) behaviors that may affect the overall fusion performance and eventually pose issues for the stability and control of the fusion burn. These issues are briefly discussed in this work, with a view on their potential applications to other research areas.

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INTRODUCTION

A burning plasma is a complex self-organized system, where among the crucial processes to understand there are (turbulent) transport and fast ion/fusion product induced collective effects. Complexity and self-organization are intrinsic to the very nature of burning plasmas, where the self-sustainment of fusion reactions for efficient power production requires that stationary conditions are achieved when, in D-T plasmas, (almost) the whole power density balance to compensate losses is provided by heating from fusion alphas. Meanwhile, fast ions in the same (MeV) energy range as that of fusion alphas will be used to heat and fuel the thermal plasma, to provide rotation and to drive current, mainly by Ion Cyclotron Resonance Heating (ICRH) and Negative Neutral Beam Injection (NNBI). Together with fusion produced alphas, these fast ions are a potential free energy source for driving collective plasma oscillations, which may induce or enhance transport processes. These phenomena are peculiar to the presence of MeV energetic ions in plasmas with temperatures of the order of 10 keV. Thus, the physics of energetic ions encompasses issues related with both charged fusion products as well as fast ions due to externally injected power. Complexity and self-organization in burning plasmas are consequence of the interaction of energetic ions with plasma instabilities and turbulence; of the strong nonlinear coupling that will take place between fusion reactivity profiles, pressure driven currents, MHD stability, transport and plasma boundary interactions, mediated by the energetic particle population; and finally of the long time scale nonlinear (complex) behaviors that may affect the overall fusion performance and eventually pose issues for the stability and control of the fusion burn.
Reactor relevant conditions require fast ion (MeV energies) and charged fusion products good confinement; thus, identification of burning plasma stability boundaries with respect to energetic ion collective mode excitations and understanding of their nonlinear dynamic behaviors above the stability thresholds is necessary, for this has obvious impact on the reactor operation-space boundaries, because collective losses may lead to significant wall loading and damaging of plasma facing materials in addition to degrading fusion performance. Meanwhile, mutual interactions between collective modes and energetic ion dynamics with drift wave turbulence and turbulent transport should not deteriorate thermonuclear efficiency via long time-scale nonlinear behaviors, which they generate and are typical of self-organized complex systems.

These physics are truly peculiar to burning plasmas of fusion interest and require a conceptual step with respect to the way phenomena are currently investigated in present day experiments. For example, energetic ion power density profiles and characteristic wavelengths of the collective modes in reactor relevant plasmas will be different, while MeV energy ion tails introduce dominant electron heating and different weighting of the electron driven micro-turbulence. Furthermore, plasma operation scenarios will reflect different plasma edge conditions and plasma wall interactions at high density and low collisionality. For these reasons, among others, important roles will be played by predictive capabilities based on numerical simulations [1] as well as by fundamental theories for developing simplified yet relevant models, needed for gaining insights into the basic physics processes, which determine the long time scale nonlinear dynamics of burning plasmas. Experiments have a key role in this respect and provide experimental evidences for modeling verification and validation. In the perspective of the International Thermonuclear Experimental Reactor (ITER) it is crucial to investigate these physics, exploiting mutual positive feedbacks between experiment, numerical simulation and theory, and integrating the largest number of aspects that are important for complexity in reactor relevant plasmas. Present day experiments, even the largest ones, investigate the various issues separately and in operation regimes of limited interest. Nonetheless, on the basis of present understanding, it is possible to conceive plasma experiments, which will be capable of investigating integrated burning plasma dynamics without the use of Tritium and the connected complications of a nuclear environment [2, 3]. In this way, some aspects of complexity in reactor relevant plasmas could be studied experimentally, providing precious feedbacks for theory and modeling and, thus, helping preparing for ITER operations and early DEMO design (DEMO is the conventional name for the first demonstrative fusion reactor). However, exploration of self-organization will be beyond reach even in those experiments, since for such studies Deuterium-Tritium fuel is necessary together with a significant fraction of fusion alpha power production (in DEMO, the fusion alpha power density is in excess of 90% of the total).

Other problems are of great importance for reactor relevant fusion plasmas, starting from those related with fusion technology. The issues mentioned above, however, are peculiar to burning plasmas for, in a sense, they stem from the complex self-organized nature of these systems itself. In this brief work, we address two of the issues mentioned so far, i.e. the energetic ion induced collective behaviors and fast ion transport in burning plasmas and the mutual interactions between collective modes and energetic ion dynamics with drift wave turbulence and turbulent transport. Thus, we focus on two important aspects concerning the physics of the core of reactor relevant fusion plasmas,
neglecting edge conditions and plasma boundary interactions. General references on the first issue are provided by the ITER physics basis [4] and its recent update [5], while recent and sufficiently comprehensive reviews can be found in [6, 7, 8]. On the second issue, experimental investigation is in the very early development stage, so only fairly recent theoretical reviews are available [6, 8, 9].

In the following, examples of the broader applications to other fields of plasma physics and beyond of fundamental problems that appear in studies of complex behaviors of burning plasmas are also given, to demonstrate that fusion plasmas are an exciting research field of practical importance and intellectually challenging.

**COLLECTIVE BEHAVIORS AND FAST ION TRANSPORT**

The possible detrimental effects of collective shear Alfvén (s.A.) fluctuations [10, 11] as well as of lower frequency MHD modes [12, 13, 14] on energetic ion confinement properties were recognized in the early years of fusion research and ever since the stability properties of Alfvénic and MHD fluctuations have been an important subject of the field. Shear Alfvén waves are of particular relevance for two main reasons: (i) they are nearly incompressible (easier to excite) and (ii) their group velocity is directed along the equilibrium confining $B$ field, so that resonant particles moving parallel to $B$ can stay in resonance with the s.A. wave-packet.

In nonuniform magnetized plasmas, s.A. waves are characterized by a continuous spectrum, whose local (singular) fluctuations decay as $(1/r)\exp(-i\Omega_{\psi}(\psi)t)$ because of phase mixing [15], with

$$\omega^2 = k_{\parallel}^2(\psi)v_{A}^2(\psi) = \Omega_{\psi}^2(\psi),$$  \hspace{1cm} (1)

$k_{\parallel}$ representing the parallel (to $B$) wave vector, $v_{A} = B/\sqrt{4\pi\rho}$ the Alfvén speed, $\rho$ the plasma mass density and $\psi$ indicating the magnetic flux function, i.e. representing a radial-like flux coordinate, $r$, since $\psi = \psi(r)$. If the plasma is perturbed at frequency $\omega_0$, the shear Alfvén continuous spectrum is resonantly excited at the radial position $\psi(r_0) = \psi_0$ where $\Omega_{\psi}(\psi_0) = \omega_0$ and energy is locally absorbed at the rate $\gamma_\psi \simeq |\Delta\rho \partial_\psi \Omega_{\psi}(\psi_0)|$, representing continuum damping [16], with $\Delta\rho$ the perturbation radial extent.

In toroidal geometry, common to magnetized plasmas of fusion interest such as tokamaks (like ITER), the magnetic field intensity varies along $B$, creating a one-dimensional periodic lattice structure for s.A. wave propagation. Similar considerations can be made for other toroidal magnetized plasma equilibria, such as stellarators, or – more generally – for any general periodic modulation of the $B$ field strength. As a consequence of this lattice symmetry breaking, frequency gaps are formed in the s.A. continuous spectrum, similar to forbidden energy bands for electrons moving in a one-dimensional crystal lattice [6, 8, 9]. The best known example is the toroidicity induced frequency gap [17], however each equilibrium geometry modulation generates its own particular structures in the s.A. continuous spectrum [6, 8, 9].

All these geometry dependent structures are of crucial importance in determining the linear as well as nonlinear plasma behaviors (see discussions below). For example toroidal geometry influences both mode structures of Toroidal Alfvén Eigenmodes
As well as their resonant absorption due to the nonlocal coupling with the s.A. continuum and there exists an entire “zoology” for their classification based on empirical observations [21]. Here, we want to stress that all these fluctuations can be classified, according to one single unified theoretical framework, as either Alfvén Eigenmodes (AE) or Energetic Particle Modes (EPM) [6, 8, 13, 22, 23, 24, 25]. In fact, perturbations of the s.A. wave spectrum generally consist of singular (inertial) and regular (ideal MHD) structures. For this reason, via asymptotic analyses it is always possible to derive a general “fishbone-like” dispersion relation in the form [6, 8, 13, 22, 23, 24, 25] 

\[ i\Lambda(\omega) = \delta \tilde{W}_f + \delta \tilde{W}_k \]  

(2)

Here, \( i\Lambda(\omega) \) is the inertial layer contribution due to thermal ions, while the right hand side comes from background MHD and energetic particle contributions in the regular ideal regions. Alfvén Eigenmodes are resonant cavity modes of the plasma, which are weakly damped by the coupling with the s.A. continuum since they exists within frequency gaps, where \( \partial_{\Omega_M}(\psi) \simeq 0 \) due to equilibrium/geometry effects and \( \gamma_c \) is minimized; EPMs are instead excited at the relevant energetic ion characteristic frequency, which generally falls in the s.A. continuum, when the energetic ion drive exceeds continuum damping. The s.A. continuous spectrum is described by [26] \( \Lambda^2 = k_{||}^2 q^2 R_0^2 \). Thus, the s.A. frequency gap is given by the condition \( \text{Re}\Lambda^2 < 0 \) and, on the basis of Eq. (2), AE exist for \( \text{Re}\Lambda^2 < 0 \) while EPM [23] correspond to \( \text{Re}\Lambda^2 > 0 \). The combined effect of \( \delta \tilde{W}_f \) and \( \text{Re}\delta \tilde{W}_k \), which determines the existence conditions of AE by removing the degeneracy with the s.A. wave accumulation point \( \Lambda = 0 \), depends on the plasma equilibrium profiles. Thus, various effects in \( \delta \tilde{W}_f + \text{Re}\delta \tilde{W}_k \) can lead to AE localization in various gaps, i.e. to different species of AE [6, 8]. In the case of EPM, meanwhile, \( \omega \) is set by energetic ion characteristic frequencies and mode excitation requires the drive exceeding a threshold due to continuum damping; i.e., \( \text{Im}\delta \tilde{W}_k > \text{Re}\Lambda \) [13, 22, 23] in Eq. (2). Thus, toroidal geometry is very important for the linear dispersive properties of both AE and EPM, for geometry and details of the plasma equilibria affect structures of the s.A. continuum with gaps as well as energetic ion own frequencies, i.e. transit frequency \( \omega_{tH} \) for circulating particles freely traveling along magnetic field lines, or bounce frequency \( \omega_{bH} \) and toroidal precession frequency \( \omega_{dH} \) for particles that are trapped between magnetic mirror points, due to parallel modulation of the magnetic field strength, and move toroidally at a precession rate \( \omega_{dH} \) because of their bounce averaged magnetic drift motion. Here, we have used the conventional definition of toroidal coordinates \((r, \theta, \phi)\), where \( r \) is the radial-like flux variable introduced above, \( \theta \) is the angular variable increasing by \( 2\pi \) when encircling the magnetic axis and \( \phi \) is the toroidal angle, with respect to which the system is rotationally symmetric. In general, wave-particle resonance conditions can be expressed as \( \omega = k_\parallel v + \ell \omega_{tH} (\ell \in \mathbb{Z}) \) for circulating particles and \( \omega = \omega_{dH} + \ell \omega_{bH} (\ell \in \mathbb{Z}) \) for trapped particles. The fact that \( \omega_{dH} \) depends on energy and not mass has recently attracted significant attention due to experimental observations of collective fluctuations driven by energetic electrons via the \( \omega = \omega_{dH} \) resonance, which may provide useful insights into trapped energetic ion behaviors in fusion plasmas due to their small dimensionless orbit size (normalized to the system size), which better reproduces reactor relevant plasma conditions than fast ions in present day experiments.
characterized by much larger dimensionless orbits due to limited plasma current [25].

The single and general dispersionless relation, Eq. (2), describes the whole AE/EPM “zoology” [21], i.e. it accounts for the resonant excitation of the s.A. frequency spectrum by energetic and thermal ions in the range \( \omega_{pi} \approx \omega_t \ll \omega \ll \omega_d [24] \), i.e. from the low Kinetic Ballooning Mode (KBM) [22, 27, 28] and Beta induced Alfvén Eigenmode (BAE) [29, 30] frequency up to the higher frequency typical of TAEs [18]. Here, \( \omega_{pi} = (Te/eB)(k \times b) \cdot \nabla P_i/P_i, b \equiv B/B, k \) is the wave-vector, \( \omega_t \) is the thermal ion transit frequency, \( \omega_d = v_A/qR_0, q = \delta(d \theta/2\pi)B \cdot \nabla \phi/(B \cdot \nabla \theta) \) is the safety factor, \( R_0 \) the torus major radius and other symbols are standard. Thermal ion excitations of AE are connected with wave-particle resonant interactions in the inertial layer and are possible only in the presence of finite \( \nabla T_i \), so that this type of fluctuations have been dubbed Alfvén Ion Temperature Gradient modes (AITG) [26, 31], being the Alfvénic counterpart of the well known ITG micro-instability. More detailed discussion on AITG will be given in the next Section. As a final remark, we note that, since singular structures characterizing the s.A. spectrum are independent of the mode number except for a scale factor, Eq. (2) applies to macroscopic MHD modes as well, including \( (n/m) = (1, 1) \) fishbone oscillations [12], with \( n/m \) the toroidal/poloidal mode number. In fact, according to the present classification scheme, diamagnetic fishbones [14] can be considered as AE while precessional fishbones [13] are the first notable example of EPM [6, 8, 23]. Extensions of Eq. (2), which include finite Larmor radius and finite orbit width effects [32] as well as plasma resistivity [33, 34], are readily derived. In the banana regime, \( |\omega| \ll \omega_{bi} \), with \( \omega_{bi} \) the (magnetically trapped) thermal ion bounce frequency, \( \Lambda \) is given by [35]

\[
\Lambda^2 = \left( \omega^2/\omega_{bi}^2 \right) \left[ 1 - \omega_{pi}/\omega \right] \left[ 1 + \left( 1.6(R_0/r)^{1/2} + 0.5 \right) q^2 \right],
\]

\[
\Lambda^2 = \frac{\omega^2}{\omega_{BAE}^2} \left[ 1 + \frac{\omega_{BAE}^2 (46/49) + (32/49)(T_e/T_i) + (8/49)(T_e/T_i)^2}{(1 + (4/7)(T_e/T_i))^2} \right],
\]

with \( \omega_{BAE} = q\omega_t (7/4 + T_e/T_i)^{1/2} \) being the asymptotic expression of the BAE [29, 30] accumulation point in the fluid limit and \( \omega_t = (2T_i/m_i)^{1/2}/(qR_0) \). The general kinetic expression of \( \Lambda \) (\( \Lambda \) is generally complex) for \( \omega_{bi} < |\omega| \ll \omega_d \) is given in Ref. [26, 36].

Nonlinear dynamics and fast ion transport induced by s.A. fluctuations reflect the different nature of AE and EPM. Alfvén Eigenmodes are predicted to have small saturation levels and yield negligible transport unless the stochasticization threshold in particle phase space is reached [37, 38]. This fact has been also confirmed by numerical simulations of alpha particle driven AE in ITER [39]. Analyses of phase space structures due to fast ion resonant interactions with one single AE [38, 40] demonstrated that the dominant “transient loss” mechanism below stochastic threshold is due to resonant drift motion across the orbit-loss boundaries in particle phase space and is \( \sim \delta B_r/B \). More precisely, it is expected to be that of scattering of barely counter-passing particles into unconfined (trapped particle) “fat” banana orbits [38, 40]; however, this loss mechanism is expected to be weak in ITER, due to the small ratio of the banana orbit width to the system size. Diffusive losses, scaling as \( \sim (\delta B_r/B)^2 \) are expected to occur above a
stochastic threshold, due to stochastic diffusion in phase space across orbit-loss boundaries. The stochastic threshold for a single AE is $\delta B_r/B \approx 10^{-3}$, although that may be greatly reduced in the multiple mode case ($\delta B_r/B \lesssim 10^{-4}$) [38, 40]. In the case of multi-mode wave-particle interactions, the nonlinear resonant particle characteristics may exceed the Chirikov resonance overlap criterion [41] only locally, producing fast ion avalanches [42] and eventually causing stochastic diffusion in particle phase space, possibly because of a domino effect [43]. In this context, the single transport events due to AEs (avalanches) [42] may exhibit characteristic aspects of sandpile physics involving Self Organized Criticality (SOC) [44, 45].

Strong energetic particle redistributions are predicted to occur above the EPM excitation threshold [46] in 3D Hybrid MHD-Gyrokinetic simulations [47]. Due to the intrinsic EPM resonant character and their localization at the radial position where the drive is strongest [48, 49, 50], EPMs rapidly redistribute energetic particles. Simulation results indicate that, above the linear stability threshold, strong EPM induced fast ion transport occurs via processes that can be identified with the convective amplification of an unstable front, which coherently propagates in the radial direction as the fast ion energy density gradient (free energy source) steepens before it eventually relaxes [51]. Such strong transport events occur on time scales of a few inverse linear growth rates (generally 100 – 200 Alfvén times, $\tau_d = R_0/v_{A0}$, with $v_{A0}$ the Alfvén speed on the magnetic axis evaluated at $B_0$) and have a ballistic character [52] that basically differentiates them from the diffusive and local nature of weak transport. This situation is depicted in Figure 1. The fluctuating fields are Fourier decomposed; e.g., the scalar potential is taken as $\phi = \exp(-i\omega t + in\phi) \sum m \phi_{m,n}(r,t) \exp(-im\theta)$. Since the functions $\phi_{m,n}(r,t)$ are radially localized with respect to equilibrium variation scales, one usually assumes

$$\phi_{m,n}(r,t) = A_{m,n}(t) \Phi(nq(r) - m,t) \simeq A_n(r,t) \Phi(nq(r) - m,t)$$

where $\Phi(nq(r) - m,t)$ describes the local self-similar structure of the various poloidal Fourier harmonics and $A_n(r,t)$ describes the radial envelope of the mode structure. These functions are often assumed to weakly depend on time as well, although this is not strictly necessary. The panels from left to right refer to three subsequent times of the 3D Hybrid MHD-Gyrokinetic simulation discussed in [51], with the nonlinear evolution of the radial envelope of a single $n = 4$ EPM, shown as a function of the radial variable normalized to the plasma minor radius $a$ on top of the nonlinear distortion of the fast ion free energy source, $\delta\alpha_T = -(8\pi/B_0^2)R_0d^2(d/dr)(P_T - P_{T,\text{equil}})$, which is identically zero during the linear unstable phase, where profiles are those given at equilibrium initial conditions. Particle radial transport in this case is secular due to coherent non-linear interactions with the modes. No stochasticization of particle orbits leading to fast ion diffusion is produced in the phase space, unlike for the AE case discussed above [42]. In addition, the radial mode structure evolves on the same time scale as fast ion transport: thus, the regime is strongly non-perturbative. Since EPM induced transport events consist of a radially propagating unstable fronts producing convective fast ion radial redistribution, it is possible to speak of a single coherent avalanche [51]. This peculiarity is due to the fact that the EPM growth rate has a strong $n$ dependence, producing a narrow toroidal mode number unstable spectrum, in contrast to the AE case. Thus, single-$n$ nonlinear dynamics may be expected to dominate the initial
FIGURE 1. Nonlinear evolution of EPM radial envelope and free energy source for three subsequent times expressed in units of $\tau_d$. The vertical scale is amplified by a factor 10 in the second upper panel and by $\sim 30$ in the third upper panel for better visibility of the convective amplification of the unstable front. The lower panels show, from left to right, the nonlinear distortion of the fast ion free energy source.

rapid convective phase. The radial propagation of a single-$n$ nonlinear EPM localized mode structure takes place via couplings among poloidal harmonics and their interplay with nonlinear distortions of the fast ion source. Such dramatic transport processes due to bursty Alfvénic fluctuation activities have been observed experimentally, e.g. as Abrupt Large amplitude Events (ALE) in JT60-U [53] that have been successfully interpreted as nonlinear $n = 1$ EPM bursts in 3D Hybrid MHD-Gyrokinetic simulations, showing the convective nature of fast ion radial redistributions within a single ALE [54]. These conclusions, based on experimental observations and numerical simulations that are in close agreement, are further supported by formal theoretical analyses based on nonlinear extensions of Eq. (2) [6, 8, 9]. In fact, the possibility of writing Eq. (2) in its form relies on the asymptotic matching procedure between ideal region and singular layer, but it never requires the problem to be linear. Assuming that the EPM wave-packet grows out of thermal noise at the radial position where the drive is strongest [48, 49, 50], with given (complex) frequency $\omega_0$ locally satisfying Eq. (2), the nonlinear evolution equation for $A_n(r,t)$ is readily derived as

$$\Delta (\delta \vec{W}_f + \delta \vec{W}_k) A_n(r,t) - \delta \vec{W}_{k,NL}(r,t; |A_n|^2) A_n(r,t).$$

Here, $\Delta (\delta \vec{W}_f + \delta \vec{W}_k)$ accounts for the radial dispersiveness, which reflects both equilibrium geometry and equilibrium profile changes due to “slow” cumulative effects of nonlinear EPM dynamics, while $\delta \vec{W}_{k,NL}(r,t; |A_n|^2)$ is the nonlinear fast ion response, in which we have made the dependence on EPM intensity explicit [51]. Thus, EPM wave-packets grow most strongly when linear and nonlinear dispersions balance on the
r.h.s., describing the convective amplification of the unstable front accompanied by rapid frequency “tuning” to the wave-particle resonance condition and ballistic fast ion transport [51]. The structure of Eq. (6) suggests that, within this theoretical framework, it is possible to systematically generate a number of nonlinear partial differential equations, including nonlinear Schrödinger and Ginzburg-Landau equations, that are relevant for other fields of plasma physics research and beyond [9, 51, 55]. It is also possible to naturally derive the fractional derivative Ginzburg-Landau equation, which incorporates the key features of non-Gaussianity and long-range dependence in thresholded nonlinear dynamical systems [56] such as burning plasmas [9].

In summary, energetic ion transport in burning plasmas has two components: one associated with slow diffusive processes due to weakly unstable AEs and a residual component possibly due to plasma turbulence [57, 58]; and another one consisting of rapid transport processes with ballistic nature due to coherent nonlinear interactions with EPM and/or low-frequency long-wavelength MHD fluctuations [53, 54].

MUTUAL INTERACTIONS BETWEEN COLLECTIVE MODES AND ENERGETIC ION DYNAMICS WITH DRIFT WAVE TURBULENCE AND TURBULENT TRANSPORT

Collective oscillations excited by energetic ions in burning plasmas are characterized by a dense spectrum of modes with characteristic frequencies and spatial locations [6, 8]. Thus, both collective modes as well as plasma turbulence, represented by the Fourier series \( \delta \phi = \sum_{m,n} \phi_{m,n}(r,t) \exp(i\phi - im\theta) \) with \( \phi_{m,n}(r,t) \) given by Eq. (5), can be described by three degrees of freedom: the toroidal mode number \( n \), the parallel mode structure \( \Phi(nq(r) - m,t) \) reflecting the radial width of a single poloidal harmonic \( m \), and the radial mode envelope \( A_m(r,t) \). Correspondingly, nonlinear interactions can take the following three forms: mode coupling between two \( m \)s, distortion of the parallel mode structure and modulation of the radial envelope [59, 60].

It has been recognized that drift wave turbulence is the channel through which turbulent transport occurs; at the same time, however, in the description of turbulent transport processes, it is crucially important to account for the radial structures that are spontaneously generated by turbulence itself and regulate turbulence intensity and turbulent transport [61]. Among those structures, zonal flows (ZF) [62] play a major role in the nonlinear dynamics of drift waves. These are toroidally and poloidally symmetric flow patterns, due to the low frequency radial electric field nonlinearly generated by drift waves [62], which regulate turbulence intensity and turbulent transport, as demonstrated in numerical simulations [63]. Zonal flows are ubiquitous in plasmas and fluids, e.g. in atmospheric pressure systems where Coriolis forces drive Rossby wave turbulence [64], which is known to obey the same nonlinear partial differential equations as drift wave turbulence [62]. As a consequence, it is often said that zonal flows in burning plasmas are the counterpart of the Jupiter’s stripes, which are signatures of zonal flows driven by Rossby wave turbulence [65]. Similar to ZF, other toroidally symmetric flow patterns with more complicated poloidal structures and finite frequency, the geodesic acoustic modes (GAM) [66], are known to play a role in regulating plasma turbulence [61, 67].
Furthermore, magnetic field patterns can be generated as well, generically dubbed as zonal fields [61, 68, 69]. Zonal flows and fields are generated by nonlinear AE and EPM dynamics as well, depending on proximity to marginal stability [68, 69]. Strongly driven EPM cause radial modulations in energetic ion profiles [9], which may produce similar structures in the electron temperature profile and eventually alter the free energy source driving drift wave turbulence and transport, for nuclear heating by fusion reactions is dominated by fast ions slowing down on thermal electrons. Radial or zonal structures can be viewed as generators of nonlinear equilibria [70], which determine the long time scale nonlinear evolution of burning plasmas, affecting the reactor fusion performance.

Since collective modes are due to energetic particles and drift wave turbulence is due to thermal plasma components, there is an intrinsic separation of spatial scales in the corresponding free energy source, associated with characteristic orbit widths. Similarly, characteristic times/frequencies are different. Thus, mutual interactions between collective modes and energetic ion dynamics with drift wave turbulence and turbulent transport are generally possible when they are mediated by zonal structures such as zonal flows, fields and corrugation of the radial profiles; i.e. cross-scale couplings are most effective when they are caused by the generation of neighboring nonlinear equilibria [9, 70]. Another case of particular interest is when frequencies of micro- and meso-scale fluctuations become comparable such as for AITG, BAE and GAM [6, 9, 25, 71]. Both BAE and AITG are described by the same Eq. (2), as discussed above, the only difference among them being given by the drive, dominated by fast ions at long wavelength for BAE and by thermal ions at short wavelength for AITG. Furthermore, BAE and GAM are degenerate in frequency, with \( \omega \) obtained from Eq. (4) and \( \Lambda^2 = 0 \). This fact poses issues of general interest that have been recently discussed [71].

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