Interaction of fast particles and Alfvén modes in burning plasmas

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Abstract. In this paper we study the interaction of fast particles with Alfvénic instabilities in Tokamak plasmas, with reference to present-day experiments that exploit strong energetic particle heating (namely, JT-60U [1]) and the consistency of proposed ITER burning plasma scenarios [2]. Concerning JT-60U, two different types of bursting modes have been observed by MHD spectrography in auxiliary heated (NNB) discharges. One of these modes has been dubbed fast frequency sweeping (fast FS) mode. It is characterized by a timescale of the order of few milliseconds and frequencies branching upwards and downwards. The other mode, called the abrupt large-amplitude event (ALE), has shorter timescale (order of hundred microseconds) and larger amplitude. On the occurrence of ALEs, a significant reduction of the neutron emission rate in the central plasma region is observed. Such a change has been attributed to a redistribution of the energetic ions, with a marked reduction of their on-axis density. We present an interpretation of these experimental observations, based on the results of nonlinear particle simulations performed by the Hybrid MHD-Gyrokinetic Code HMGC [3].

Concerning ITER, monotonic-q (scenario 2) and reversed-shear (scenario 4) equilibria are considered. Also an ITER hybrid scenario is examined and quantitatively compared with the previous ones. The transition from the low-amplitude Alfvénic instability saturation to the secondary excitation of a stronger mode is addressed, and its effect on the energetic particle transport analyzed.

Keywords: fast particles, Alfvén modes, burning plasmas, Hybrid MHD-Gyrokinetic Code, JT-60U, abrupt large-amplitude event, ALE, ITER, ITER scenarios

PACS: 52.35.Bj, 52.35.Mw, 52.55.Pi, 52.65.Rr, 52.65.Tt, 52.65.Ww

1. INTRODUCTION

In a burning plasma, alpha particles are expected to transfer their energy via Coulomb collisions to the thermal plasma, thus providing a nuclear self-heating mechanism and a route to ignition. The fusion reactions generate energetic particles characterized by velocities in the super-Alfvénic range, and it is well known that they can resonate with, and possibly destabilize, shear Alfvén modes. The mutual nonlinear interaction of shear Alfvén and energetic particles can, in turn, affect the energetic ion transport and confinement properties. Experimental evidences of rapid transport of energetic ions related with fluctuations in the Alfvén-mode frequency range have been observed in plasmas heated by different auxiliary power systems [4] as reported, e.g., in the JT-60U tokamak [1, 5, 6] in connection with the so called Abrupt Large amplitude Events (ALEs). A review of experimental observations of collective mode effects on fast ions is given in Ref. [4] and references therein. Meanwhile, evidence that energetic-ion redistribution can take place because of fast growing Alfvén modes (e.g., Energetic Particle driven Modes (EPMs) [7]) has been found in particle-simulation studies [8, 9,
Despite these numerous experimental and numerical evidences, supported by theory predictions, the operation scenarios for next-step proposed burning-plasma experiments, such as ITER-FEAT [2, 15], are usually obtained from equilibrium and transport codes that do not include the physics required to describe shear Alfvén modes. Thus, the possibility that these modes are excited and, eventually, produce macroscopic transport of alpha particles themselves is neglected. A basic question concerning the consistency of such envisaged burning-plasma scenarios may arise, in particular with reference to the alpha-particle radial profile and, therefore, to the fusion power density.

In this paper we will first test the particle simulation approach versus the experimental observations of the Alfvén mode dynamics in JT-60U discharges. We will show that simulations performed by the Hybrid MHD-Gyrokinetic Code HMGC [3, 8, 16] yield results in fair agreement with those observations. Then, we will apply the same tool to the investigation of the consistency of the proposed ITER scenarios with the nonlinear evolution of unstable Alfvén modes.

The paper is organized as follows. Section 2 describes the model adopted in our simulations. The bursting modes observed in the JT-60U device will be analyzed in Sect. 3. Section 4 is devoted to the analysis of ITER scenarios, with respect to both their linear stability and the nonlinear effects of the Alfvén modes on the alpha-particle transport.

2. THE MODEL

The plasma model adopted in the HMGC code consists of a thermal plasma and an energetic particle population [3, 8, 16]. The former is described by reduced $O(\varepsilon^3)$ MHD equations [17, 18] in the limit of zero pressure ($\varepsilon$ being the inverse aspect ratio of the torus), including resistivity and viscosity terms. The reduced MHD equations expanded to $O(\varepsilon^3)$ allow us to investigate equilibria with shifted circular magnetic surfaces. The energetic particle population is described by the nonlinear gyrokinetic Vlasov equation [19, 20], expanded in the limit $k_\perp \rho_H \ll 1$ (with $k_\perp$ being the perpendicular component of the wave vector to the magnetic field, and $\rho_H$ the energetic-particle Larmor radius), and solved by particle-in-cell (PIC) techniques. The coupling between energetic particles and thermal plasma is obtained through the divergence of the energetic-particle pressure tensor, which enters the vorticity equation [21]. Numerical simulations are performed retaining, for each case, the relevant thermal-plasma quantities -- the on-axis equilibrium magnetic field, major and minor radii ($R_0$ and $a$, respectively), the safety-factor $q$, the plasma density $n_e$, the electron temperature $T_e$, the anisotropy of the energetic particle distribution function and the ratio between their pressure and the magnetic one $\beta_H$.

Energetic particles are loaded in the $(\mu, v_\parallel, \psi)$ space, ($\mu$ and $v_\parallel$ being the magnetic moment and the parallel velocity of the energetic particles, and $\psi$ the magnetic flux function, respectively) according to a slowing-down (SD) distribution function. For the JT-60U device the SD distribution function is strongly anisotropic with an upper energy cut-off given by the injection energy of the (almost tangential) Negative Neutral Beam (NNB) ions; for the ITER scenarios, the SD distribution function is isotropic with an
upper energy cut-off given by the fusion energy of the alpha particles. Moreover, the critical energy $E_c$ is given by the Stix expression [22]. In the model for ITER scenarios we will assume, for simplicity, $n_D = n_T = n_i/2$ and $n_i = n_e$ (thus the critical energy is well approximated by the expression $E_c \simeq 33.0 T_e$), whereas the experimental quantities for bulk ion (deuterium) and electron densities will be used for modeling the JT-60U device. Here, $n_D$, $n_T$ and $n_i$ are the deuterium, tritium and (total) bulk-ion densities, respectively.

For simplicity we neglect the nonlinear mode-mode coupling among different toroidal mode numbers, then limiting the analysis to the evolution of a single toroidal mode $n$, while keeping fully nonlinear dynamics for energetic particles. Fluid nonlinearities are not expected indeed to considerably alter shear-Alfvén mode dynamics [8]. Note, however, that considering a single toroidal mode could underestimate the nonlinear effects on energetic-particle transport, e.g. neglecting resonance overlap and orbit stochastization in phase-space.

In order to have a sensitivity study and a more complete picture of the Alfvén mode dynamics for the considered ITER scenarios, we will extend our investigation to values of the on-axis $\beta_H$ ($\beta_{H0}$) larger than the reference value ($\beta_{H0,\text{scenario}}$), while keeping the normalized profile, $\beta_H/\beta_{H0}$, unchanged.

### 3. BURSTING MODES IN JT-60U

Two different types of bursting modes, both in the frequency range of Alfvén modes, have been observed by MHD spectroscopy in auxiliary heated (NNB) discharges in the JT-60U tokamak. The first one, the ALE [1], is characterized by a time scale of the order of hundred microseconds, a relatively large fluctuating magnetic field level ($\tilde{B}_\theta/B_\theta \sim 10^{-3}$ at the first wall) and large growth rates (corresponding to wide frequency spectra $\Delta \omega \sim \omega \sim 50 \text{ kHz}$). The ALEs are followed by a significant drop of the neutron emission rate in the central plasma region, which has been interpreted as a radial redistribution of the energetic ions, producing a marked reduction of their on-axis density. The repetition rate of ALEs appears to be related with the intensity of the event itself, as measured by the amount of such reduction. In the relatively quiescent phase between two ALEs, bursting modes of a second type, the fast frequency sweeping (fast FS) modes [1], are observed. Such modes are characterized by longer time scales (few milliseconds) when compared with ALEs, lower growth rate and fluctuating field level, and much smaller redistribution of the energetic ions. Most of them consist of bifurcating branches, chirping up and down in frequency.

In order to provide a possible explanation of this phenomenology in the framework of the nonlinear behavior of Alfvén modes, we perform particle simulations of a typical NNB-heated JT-60U discharge (# E036378) with the HMGc code [3]. In particular, we simulate the dynamics of energetic ion interactions with Alfvén modes, assuming the radial profiles experimentally observed immediately before an ALE, and postulating a strongly anisotropic SD distribution function in velocity space for the energetic deuterons injected at 0.397 MeV by the NNB auxiliary heating system. Modes with toroidal number $n = 1$ have been considered, to compare simulation results with the experimental observations. The linearly unstable phase is dominated, in this “ALE” sim-
ulation, by a fast-growing mode ($\gamma/\omega \approx 0.5$), located around the maximum energetic-ion pressure gradient ($r/a \approx 0.5$) and characterized by a significant coupling with the Alfvén continuum, showing the energetic-particle-driven nature of the mode (see Fig. 1 left, where the contour plot of the power spectrum of the fluctuating scalar potential is shown). Once the mode saturates, a complex phenomenology appears. At an earlier stage, the configuration is dominated by a Toroidal Alfvén Eigenmode (TAE)-like mode, localized around $r/a \approx 0.8$, on the lower part of one throat of the Alfvén continuum. The original mode, localized at half radius, is replaced by a couple of nearly degenerate modes. A weak core-localized mode also appears, with frequency well localized in the toroidal gap (Fig. 1 centre). At later times, the core-localized mode becomes the dominant one, while the external mode still persists at lower amplitudes (Fig. 1 right).

Figure 2 shows the contour plot in the energetic-ion $(\hat{\mu}, \hat{v}_{||})$ space of the wave-particle power transfer averaged over the radial annulus $0.36 < r/a < 0.67$, where the most unstable mode is localized during the linear phase. Here, $\hat{\mu} \equiv \mu \Omega_{H0}/m_H v_H^2$ and $\hat{v}_{||} \equiv v_{||}/v_H$ are normalized magnetic moment and parallel velocity ($m_H$, $v_H$ and $\Omega_{H0}$ are the energetic-ion mass, injection velocity and on-axis Larmor frequency, respectively). The continuous and dashed lines correspond to trapped-to-circulating particle transition at the boundaries of the annulus ($r/a = 0.36$ and $r/a = 0.67$, respectively); the dotted line indicates the injection energy of the NNB ions. The resonant drive is given by circulating energetic ions, consistently with the nearly tangential beam injection. The resonance region appears to be quite broad in energy, in agreement with the experimental observations [23].

As the nonlinear effects become important, a macroscopic outward displacement of the energetic particles is observed, producing a significant reduction of their density in the central region (Fig. 3 left), which compares fairly well with the observed energetic particle density profile after the ALE (black curve in Fig. 3 left). The fluctuating poloidal magnetic field $B_\theta$ vs. $t$ (at a radial position close to plasma boundary) is plotted in Fig. 3 centre, while Fig. 3 right shows the power spectrum of the same quantity in the plane $(t, \omega)$. The time scale $\Delta t$ of the “burst” is $\Delta t \approx 200\omega_{A0}^{-1}$, which corresponds to

![Contour plots of the power spectrum of the fluctuating scalar potential in the $(r/a, \omega/\omega_{A0})$ plane, for a typical JT-60U discharge (#E036378), during the linear-growth phase (left) and the early (centre) and late (right) saturated phases. Modes with toroidal number $n = 1$ are considered. The Alfvén continuum spectrum is also plotted by the black (solid) line. Here, $\omega_{A0} \equiv v_{A0}/R_0$ is the on-axis Alfvén frequency, $v_{A0}$ the on-axis Alfvén velocity and $R_0$ the major radius.](image-url)
FIGURE 2. Contour plot in the energetic-ion ($\hat{\mu}, \hat{v}_||$) space of the wave-particle power transfer, averaged over the radial annulus $0.36 < r/a < 0.67$, during the linear phase. The solid and dashed lines correspond to trapped-to-circulating particle transition at the boundaries of the annulus ($r/a = 0.36$ and $r/a = 0.67$, respectively); the dotted line indicates the injection energy of the NNB ions.

$\Delta t \approx 150\mu s$. The frequency spread is of order $\Delta \omega \approx 40kHZ$. Both these results are close to the experimental observations.

Given the fair agreement between numerical results and experimental observations, we investigate whether the nonlinear dynamics described by our simulations can also explain the relatively quiescent phase (low amplitude fluctuating fields with negligible effects on the energetic-ion distribution) between two successive ALEs. As a first step, we check to which extent the reduced instability level of the system after an ALE is due to the modification of the energetic-ion radial density gradient. Thus, we perform a simulation (“after ALE”) initializing the energetic-ion density profile equal to that experimentally observed after the ALE, while postulating, as before, an anisotropic SD distribution function in velocity space. The linear growth rate of the mode is lower than that found in the “ALE” simulation ($\gamma/\omega \approx 0.28$) but still very high (this could be expected observing that, in the region where the linear mode is driven, magnitudes
of the radial gradient of the energetic particle pressure in the two simulations are not very different from each other (see Fig. 3 left and compare the blue and black curves). Correspondingly, in spite of fairly little nonlinear effects on the spatial energetic-ion distribution (Fig. 4 left), we get quite large fluctuating field levels (Fig. 4 centre): approximately half of the values found in the “ALE” simulation, to be compared with the experimental observation of the fast FS-to-ALE amplitude ratio (∼ 0.2). This seems to be inconsistent with the observed periodic character of the ALEs. Indeed, with such a level of instability, even a partial reconstruction of the spatial energetic-ion density profile would generate a fast instability, which would essentially clamp the energetic-ion density profile to the “after-ALE” shape (this was confirmed by other simulations not reported here).

Numerical results and experimental observations can be reconciled accounting for the fact that, differently from what we assumed above, the velocity-space distribution function after an ALE is not a SD. In other words, retaining only part of the ALE nonlinear effects (the spatial redistribution), with no attention to the distortion produced in the velocity space, is not appropriate. To show this, we perform a further numerical experiment (“after ALE modified”), initializing the energetic particle distribution function according to the distribution function obtained from the “ALE” simulation, after the rapid dynamics of the event has taken place (\(t\omega_{A0} \approx 500\), see Fig. 3 centre); in such a way, we retain both the radial profile broadening and the distortions of the distribution function in the velocity space. Figure 5 left shows the normalized amplitude of the fluctuating poloidal magnetic field \(B_\theta\) vs. \(t\) (at a radial position close to plasma boundary).

FIGURE 4. “After ALE” simulation. Left: nonlinear modifications of the energetic ion density profile (blue curve refers to the initial profile, red curve, to the simulation relaxed profile). Right: fluctuating poloidal magnetic field \(B_\theta\) vs. \(t\) (at a radial position close to plasma boundary).

On the basis of these results, a possible interpretation of the experimental observations could be the following: given the energetic-ion source provided by the beam injection, the free energy reconstruction rate is set by the need of rebuilding the velocity space distribution function. In this way, the intermediate configurations between two successive...
ALEs are characterized by a lower drive than that of a SD distribution with the same energetic-ion density radial profile. Low growth rates and amplitude modes, such as the fast FS modes, are then excited and have a weaker effect than that of ALEs in contrasting the density profile reconstruction. As soon as the combined restoration of the configuration and velocity space distribution provides enough drive for a fast growing Alfvén mode, a new ALE occurs. For this process, the 2D character of velocity space and its interplay with the radial profile variation of the fast ion source play a fundamental role.

4. ITER SCENARIOS

Given the encouraging agreement obtained in comparing our model for nonlinear energetic-particle dynamics with Alfvén modes with the experimental observations in JT-60U discharges, we now apply the simulation approach to the proposed ITER operation scenarios, in order to test their consistency with the potential presence of Alfvénic turbulence.

Three different ITER-FEAT scenarios have been considered: the reference monotonic-\(q\) scenario (“scenario 2”, SC2), the reversed shear scenario (“scenario 4”, SC4) and an intermediate (“hybrid”), recently proposed, flat-\(q\) scenario (“scenario H”, SCH). Data corresponding to the two former scenarios are available at the (ITER) Joint Work Site [24]; those related to SCH have been obtained by the package of simulation codes CRONOS [25].

The SC2 is the reference monotonic-\(q\) inductive, 15 MA scenario, with 400 MW fusion power and fusion yield \(Q \simeq 10\). The SC4 is a steady state, 9 MA, weak-negative shear scenario, with about 300 MW fusion power and \(Q \simeq 5\); the \(q\) profile is characterized by \(q_{\text{min}} \simeq 2.4\) and \(r_{q_{\text{min}}}/a \simeq 0.68\). The SCH is a steady state, 11.3 MA weak-positive shear scenario, with about 400 MW fusion power and \(Q \simeq 5\). Main parameters for the three equilibria (SC2, SC4, SCH) are: major/minor radii \((m) R_0/a = 6.195/2.005, 6.34/1.859, 6.3734/1.8017\); \(q_{95\%} = 3.14, 6.34, 6.3734\); on-axis toroidal magnetic field \(B_{T0}(T) = 5.3, 5.183, 5.3\); on-axis electron density \(n_{e0} (10^{20} \text{m}^{-3}) = 1.02, 0.73, 0.72\),
ratio between on-axis alpha-particle and magnetic pressures $\beta_{H0} (%) = 1.10, 0.92, 1.37$. The radial profiles of the relevant quantities are also shown in Fig. 6. The bulk ion density profiles (not shown in Fig. 6) are almost flat (for SC2 up to $r/a \simeq 0.95$, for SC4 up to $r/a \simeq 0.65$ and for SCH up to $r/a \simeq 1$) rapidly decreasing at the plasma edge. Note that, while $\beta_{H0}$ varies within a factor 1.5 among the different scenarios (with SCH corresponding to the largest value, and SC4 to the smallest one), the maximum value of the local alpha particle drive $\alpha_{H,max} \equiv \max\{-R_0q^2\beta_H'\}$ (with “prime” denoting the radial derivative) varies by more than a factor 7 (with SC4 and SC2 having the highest and lowest values, respectively).

![Figure 6](image)

**FIGURE 6.** Radial profiles of safety factor ($q$), alpha-particle ($n_H$) density, electron temperature ($T_e$), and alpha-particle local drive ($\alpha_H$), for ITER-FEAT SC2 (black solid curves), SC4 (red dashed curves) and SCH (blue dot-dashed curves) scenarios.

### 4.1. ITER: Linear dynamics

In this Section, we examine the linear dynamics of the considered ITER-FEAT scenarios, as it emerges from the first (low field amplitude) phase of the simulations. Several toroidal mode numbers have been analyzed, up to $n = 8$; in Fig. 7 *top*, *left* we report the growth rates obtained at different values of $\beta_{H0}$ (with fixed normalized profile, $\beta_H/\beta_{H0}$) for the most unstable toroidal mode of each considered scenario.

For the monotonic-$q$ scenario (SC2), the most unstable mode is $n = 2$ with $\gamma/\omega_{A0} \simeq 0.027$ (Fig. 7 *top*, *left*, black open circles), though the growth rate of the $n = 4$ is only slightly smaller ($\gamma/\omega_{A0} \simeq 0.023$). Figure 7 *top*, *centre* shows the power spectra of scalar-potential fluctuations in the $(r, \omega)$ plane for the $n = 2$ case: the mode structure
is radially localized around $r/a \simeq 0.3$, close to the position where the local drive $\alpha_H$ is maximum (see Fig. 6). We observe that the mode stays below the local minimum of the lower Alfvén continuum, in a large (because of the internally flat $q$ profile) damping-free region. On the basis of such feature, we expect that this mode would survive even without the contribution of energetic particles, revealing its original MHD nature, and this is indeed confirmed performing simulations with vanishing small energetic particle pressure ($\beta_{H0} \rightarrow 0$). That the maximum growth rates are found in the range $n = 2 \div 4$ can be traced back to the fact that the mode is mainly driven by trapped particles. Indeed the drive, that one would expect to scale as the energetic-particle diamagnetic frequency $\omega_*H$ (and thus with $n$), is limited, among other reasons (e.g., localization of the mode itself in order to minimize the continuum damping), by finite orbit width averaging. Thus, the maximum-drive $n$ is roughly given by $k_\perp \rho \sim 1$ [26, 27, 28, 29], with $k_\perp \approx nq(r)/r$ and $\rho$ being the typical resonant-particle orbit width. Trapped particles are particularly sensitive to these effects due to the $\propto n$ dependence of their toroidal precession frequency as well as to their characteristic orbit size $\rho \equiv \rho_B \gg \rho_d$ (with $\rho_B$ and $\rho_d$ being the banana and drift orbit widths, respectively). For trapped particle driven modes, if localized as
the one shown in Fig. 7 top, centre, the maximum-drive condition yields \( n_{\text{max}} \approx 4 \) (a wider discussion on the driving mechanisms of the \( \alpha \)-particles on shear-Alfvén modes can be found in Ref. [30]).

The toroidal mode number \( n = 2 \) is the dominant mode also for the reversed-shear scenario (SC4), with growth rate \( \gamma/\omega_{A0} \approx 0.027 \) (see Fig. 7 top, left, red open squares). For this scenario, the estimate \( k_{\perp} \rho_B \sim 1 \) yields \( n_{\text{max}} \approx 2 \). The localization in radius and frequency of the dominant mode is connected with the local maximum of the lower Alfvén continuum, corresponding to the minimum in the safety factor \( q \) (see Fig. 7 top, right): the mode resembles the so-called Cascade modes [31]. A second mode, slightly weaker than the dominant one, is also observed. It is localized around \( r/a \approx 0.42 \) and \( \omega/\omega_{A0} \approx 0.17 \), in the throat of the toroidal gap immediately inside the minimum-\( q \) surface. It has growth rate \( \gamma/\omega_{A0} \approx 0.019 \) (not reported in Fig. 7 top, left).

The ITER-FEAT hybrid scenario (SCH) presents some peculiar features. At the reference \( \beta_{H0} \) value, only modes very close to marginal stability are found. The dominant one (\( n = 4 \), see Fig. 7 top, left) has a radial localization similar to the one examined in the SC2 case. The similarity extends to the frequency localization (see Fig. 7 bottom, left), since both modes lie below the local minimum of the lower Alfvén continuum. In this case, \( \gamma/\omega_{A0} \approx 0.004 \) (other modes, which lie within the continuum gap, are even weaker). The analogies with the SC2 case are not surprising, due to the similarity of the internal \( (r/a \lesssim 0.5) \) \( q \) profiles and Alfvén-continuum structures in the two scenarios. What clearly distinguishes the two cases is the fact that, in the hybrid scenario, the energetic-particle drive is maximum around \( r/a \approx 0.6 \) (see Fig. 6, bottom, right), quite far from the mode localization. This motivates the lower growth rates obtained in the SCH than in the SC2 case, in spite of the higher value of \( \alpha_{H,\text{max}} \). This case is an excellent example to show how the mode localization results from the (global) competition between drive and damping mechanisms. If the damping that the modes would suffer around the maximum-drive surface is too large, the most unstable modes will grow around a different radial surface. This is indeed the case for the present scenario: the internal radial localization is more favorable than that around the \( \alpha_{H,\text{max}} \) surface, because the mode can better accommodate its frequency in such a way that large continuum damping is avoided. Most interesting, for this scenario, is the large-\( \beta_{H0} \) behavior of the system. When increasing the drive above the reference value, a different class of unstable modes is found. They are localized around the maximum drive radial position \( (r/a \approx 0.6) \), with frequencies deeply inside the Alfvén continuum (see Fig. 7 bottom, centre and right, where the \( n = 4 \) mode is shown). Different from the modes examined until now, these modes are not connected with MHD modes: they can exist only if the energetic-particle drive exceeds a certain threshold set by continuum damping, thus inducing us to identify them as EPMs. This is clear from Fig. 7 top, left, where two modes are shown using full symbols to distinguish them from the MHD-like modes: \( n = 4 \) and \( n = 8 \), this latter being the most unstable for \( \beta_{H0}/\beta_{H0,\text{SCH}} \gg 2 \). The growth rate of these modes increases with increasing \( n \) at a given \( \beta_{H0} \) (Fig. 7 top, left, comparing the full symbols for \( n = 4 \) and \( n = 8 \)). This means, of course, that the larger-\( n \) modes will have a lower threshold in \( \beta_{H0} \). For \( n = 8 \), such a threshold is \( \beta_{H0,\text{th}} \approx 1.6 \beta_{H0,\text{SCH}} \). Note that the strong dependence of the alpha-particle density on the electron and bulk-plasma ion temperatures \( n_H \sim T_i^2 T_e^{-3/2} \) could give these modes a practical relevance: an uncertainty
of few tens percent in these quantities, which is plausible in consideration of the approximation of transport models, can easily yield a corresponding quite large uncertainty in $\beta_{H0,SCH}$.

4.2. ITER: Nonlinear dynamics

We will consider, here, the nonlinear dynamics of Alfvén modes found unstable in the different ITER scenarios and their influence on alpha particle transport and confinement. This will reveal possible inconsistencies of the proposed scenarios with collective alpha-particle behaviors.

We have shown in Sect. 3 and also in previous papers [3, 9] that sufficiently unstable Alfvén modes saturate at large fluctuating-field amplitudes through a (few hundreds of Alfvén times) convection, generally yielding a sudden broadening of the energetic-particle pressure profile. This process is followed by a slower phase, during which the saturated electromagnetic fields act as a scattering source for the energetic particles, producing their further diffusion [32]. A different saturation regime takes place for lower growth rates, when modes saturate at lower field levels with very limited modifications of the energetic-particle pressure profile; this can be traced back to a softer mechanism, e.g., the trapping of resonant energetic ions in the potential well of the wave [33]. The transition from the weak-transport regime to strong transports dominated by avalanches has been theoretically analyzed in Ref. [34]. The numerical simulation approach, nevertheless, is necessary for a quantitative assessment of fast ion transport and confinement in the presence of collective mode excitations. Indeed, the quantitative effect of nonlinear Alfvén-mode dynamics depends, in experimental situations, on the competition of many factors, whose overall consequences are hard to predict on the basis of purely theoretical analyses. Transport of energetic particles can be enhanced, e.g., when the safety factor $q$ is larger, because of the larger typical particle-orbit width, as well as the stronger drive $\alpha_H$ [9] (this is the case, indeed, of the SC4 scenario); moreover, the avalanching process can be enhanced or prevented by a particular shape of the Alfvén continuum [34].

On the basis of the low growth rates observed in the simulations (see Sect. 4.1) we would expect low saturated field levels and quite negligible effects on the energetic-particle confinement. However, as noted above, the strong dependence of the equilibrium $\beta_H$ on the electron and bulk-ion temperatures can cause the drive to vary within a significantly large range; it is then worth extending the investigation up to a reasonable value of $\beta_{H0}$ above the reference value for each scenario.

Figure 8 left shows the effects of Alfvén wave nonlinear dynamics on the SC2 monotonic-$q$ scenario, at the reference $\beta_{H0}$ value (top) and at the artificially increased $\beta_{H0} \approx 2 \beta_{H0,SC2}$ value (bottom). This scenario, at the reference $\beta_{H0}$ value, is scarcely affected by the nonlinear saturation of the mode: such a result is consistent with the low growth rate and radial localization of the mode, which exists in a low-$q$ region of the discharge (thus, in a region where the alpha particles have small orbit size). We can thus affirm that this scenario is consistent with the nonlinear dynamics of shear-Alfvén modes. The artificially increased case ($\beta_{H0} \approx 2 \beta_{H0,SC2}$, Fig. 8 left, bottom),
however, indicates that some caution should be taken in extending this operation regime: the enhanced energetic particle case, indeed, shows a larger radial redistribution of the energetic pressure profile, that can be quantified in a decrease larger than 25% of the on-axis $\beta_H$ and a consistent broadening of the profile up to $r/a \lesssim 0.6$.

The reversed shear scenario, SC4, reveals nonlinear dynamics effects that are more significant than the SC2 case (see Fig. 8 centre). Indeed, already the reference $\beta_{H0}$ case (top) shows an energetic particle pressure profile modified, at saturation, on the whole radius, with a broadening and a pronounced decrease slightly inside the radial surface where the linear mode develops ($r/a \simeq 0.55$). It has to be noted that, despite the radial localization of the mode is fairly far from the axis, an evident decrease of the on-axis energetic particle pressure is observed, because of the large orbit size of the energetic particles in this high-$q$ region. A certain amount of inconsistency of this scenario with the nonlinear shear-Alfvén mode dynamics is then apparent. Moreover, it has to be noted the the outer localization of the mode represents an important risk factor for the global confinement of the alpha particles and, hence, for the damage of first wall of the device. The increased-drive case ($\beta_{H0} \simeq 2 \beta_{H0,SC4}$, Fig. 8 centre, bottom) shows an even more conspicuous effect.

The hybrid SCH scenario presents two antithetic features – the lowest growth rates at the reference $\beta_{H0}$ value, and the existence of fast growing EPMs for higher (than
reference) drive –, which give some specific reasons of interest to the analysis of the nonlinear behavior of this equilibrium. The results concerning the nonlinear saturation of the MHD-like modes at $\beta_{H0} = \beta_{H0, SCH}$ are very similar to those reported for the SC2 scenario, and even less relevant from the point of view of the alpha-particle confinement because of the lower growth rates involved. We then omit to show them explicitly.

We compare, instead, the results of the saturation for an over-driven MHD-like mode ($\beta_{H0} \simeq 2 \beta_{H0, SCH}$, Fig. 8 right, top), as for the previous scenarios, and the saturation of a well developed EPM ($\beta_{H0} \simeq 3.3 \beta_{H0, SCH}$, Fig. 8 right, bottom). The effects of the MHD-like mode ($n = 2$) are completely negligible; the EPM has instead some moderate effects on the global profile, but more limited that that observed in the over-driven SC4 case (Fig. 8 centre, bottom), in spite of the larger drive enhancement factor and the larger growth rate.

As a conclusive remark for this ITER Section, we can assert that the SC2 and SCH scenarios are consistent with the presence of nonlinear dynamics of shear-Alfvén modes. The SC4 scenario, instead, shows some broadening of the alpha-particle pressure profile, indicating a certain level of inconsistency of the scenario itself. Over-driven simulations (performed by artificially increasing the energetic particle pressure) show that a strong flattening of the alpha-particle pressure profile could occur in the inner plasma region for the SC2 case, while the global confinement of such particles is not significantly affected. The SC4 case presents more pronounced effects in the outer portion of the discharge, because of the outer localization of the modes, with less impact on the on-axis pressure value. In the SCH scenario, no significant consequences are observed below the EPM threshold. If such threshold is exceeded, effects on the alpha-particle pressure are similar to those in the SC4 case, but more limited for given growth rate. A wider investigation of the effects of shear-Alfvén modes on ITER scenarios can be found in Ref. [30].

ACKNOWLEDGMENTS

The authors are indebted with K. Shinohara, M. Ishikawa and M. Takechi for providing the experimental data concerning JT-60U and for several useful discussions, and to M. Schneider for providing the ITER-FEAT Hybrid scenario with the CRONOS codes.

REFERENCES

1. Shinohara K et al. 2001 Nucl. Fusion 41 603–612
2. ITER Physics Basis Editors, ITER Physics Expert Group Chairs and Co-Chairs and ITER Joint Central Team and Physics Integration Unit 1999 Nucl. Fusion 39 2137–74
24. ftp://itergps.naka.jaeri.go.jp/PF_control/EQDSK_files/...
28. Tsai S T and Chen L 1993 Phys. Fluids B 5 3284

263