Particle simulation of bursting Alfvén modes in JT-60U\textsuperscript{a)}

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(Received 31 October 2006; accepted 29 January 2007; published online 3 April 2007)

The results of particle-in-cell simulations of a negative neutral beam heated Alfvén-mode experiment in the Japan Atomic Energy Research Institute Tokamak-60 Upgrade (JT-60U) [H. Ninomiya \textit{et al.}, Fusion Sci. Technol. \textbf{42}, 7 (2002); A. Kitsunezaki \textit{et al.}, \textit{ibid.} \textbf{42}, 179 (2002)] are presented. They seem to match quite well the dynamics of the abrupt large-amplitude event (ALE) experimentally observed in the reference JT-60U discharge. The time scale and frequency spread of the ALE are well reproduced too. The issue of the weaker Alfvén fluctuation phase following the ALEs, characterized by fast frequency sweeping modes, is also investigated and an interpretation of the full JT-60U bursting-mode phenomenology is presented. Finally, the simulation tool is exploited by \textit{ad hoc} synthetic diagnostics on the fast ion distribution function to get a deeper insight into the ALE nonlinear dynamics. The underlying fast-growing energetic particle mode saturates as resonant energetic ions are scattered out of the resonance region and displaced outwards. The radially displaced ions resonate with outer Alfvén modes and enhance their local drive, consistently with the “avalanche” paradigm for mode nonlinear dynamics and energetic ion transports. © 2007 American Institute of Physics. [DOI: 10.1063/1.2710208]

I. 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 evidence of rapid transport of energetic ions related with fluctuations in the Alfvén-mode frequency range has been observed in plasmas heated by different auxiliary power systems,\textsuperscript{1} as reported, e.g., in the Japan Atomic Energy Research Institute Tokamak-60 Upgrade (JT-60U),\textsuperscript{2,3} in connection with the so-called abrupt large-amplitude events (ALEs).\textsuperscript{4,6} A review of experimental observations of collective mode effects on fast ions is given in Ref. 1, 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)\textsuperscript{7}] has been found in particle-simulation studies.\textsuperscript{8–15}

In this paper we compare particle simulation results with experimental observations of Alfvén-mode dynamics in JT-60U discharges. We will show that simulations performed by the hybrid magnetohydrodynamics gyrokinetic code (HMGC)\textsuperscript{8,16,17} are in fair qualitative and quantitative agreement with those observations. We observe that the linear-growth phase is dominated by a fast-growing mode, localized at half radius, near the maximum of the energetic-ion pressure gradient\textsuperscript{18,19} and with a significant coupling with the Alfvén continuum, showing the energetic-particle-driven nature of the mode. As the nonlinear effects become important, a macroscopic outward displacement of the energetic ions is observed, producing a significant reduction of their density in the central region. Later on, the saturated phase exhibits weaker outer and inner modes. These simulation results seem to match quite well the dynamics of the ALE experimentally observed in the reference JT-60U discharge. The time scale and frequency spread of the ALE are well reproduced too. The issue of the weaker Alfvén fluctuation phase following the ALEs, characterized by fast frequency sweeping (hereafter “fast FS”) modes, is also investigated. We show that it is not possible to explain such a phase as the consequence of the mere relaxation of the energetic-ion density profile produced by the ALE. The decrease of the instability level of the system can, however, be fully accounted for by properly taking into account the nonlinear modification of the energetic-ion distribution function in the velocity space as well. An interpretation of the full JT-60U bursting-mode phenomenology is then presented. Finally, we exploit the simulation tool by \textit{ad hoc} synthetic diagnostics on the fast ion distribution function to get a deeper insight into the nonlinear dynamics underlying the ALE. We find that the fast-growing mode saturates as resonant energetic ions are scattered out of the resonance region.\textsuperscript{20} We show that such a scattering displaces the energetic ions both in the velocity and in the configuration space. The radially displaced ions resonate with outer Alfvén modes and enhance their local drive, consistent with the “avalanche” paradigm discussed in Refs. 20–23.

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II. THE MODEL

The plasma model adopted in the HMGC code consists of a thermal (core) plasma and an energetic-ion population. The former is described by reduced $O(\epsilon^3)$ magnetohydrodynamic (MHD) equations\(^{24,25}\) in the limit of zero pressure ($\epsilon$ being the inverse aspect ratio of the torus), including resistivity and viscosity terms. The reduced MHD equations expanded to $O(\epsilon^3)$ allow us to investigate equilibria with shifted circular magnetic surfaces. The energetic-ion population is described by the nonlinear gyrokinetic Vlasov equation,\(^{26,27}\) expanded in the limit $k_\perp \rho_\parallel \ll 1$ (with $k_\perp$ being the perpendicular component of the wave vector to the magnetic field and $\rho_\parallel$ the energetic-ion Larmor radius), though fully retaining magnetic drift orbit widths and solved by particle-in-cell (PIC) techniques. The coupling between energetic ions and thermal plasma is obtained through the divergence of the energetic-ion pressure tensor, which enters the vorticity equation.\(^{28}\) Numerical simulations of experimental conditions are performed fitting 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 electron $n_e$ and ion $n_i$ plasma densities, the electron temperature $T_e$, the anisotropy of the energetic-ion distribution function, and the ratio $\beta_{HH}$ between fast ion and magnetic pressures.

Energetic ions are loaded in the ($\mu, v, \psi$) space ($\mu$ and $v$ being the magnetic moment and the parallel velocity, and $\psi$ the magnetic flux function, respectively) according to a strongly anisotropic slowing-down (SD) distribution function (see, later, the left-hand side of Fig. 9), with an upper energy cutoff given by the injection energy of the (almost tangential) negative neutral beam (NNB) ions and the critical energy $E_c$ given by the Stix expression.\(^{29}\) Experimental values are adopted for bulk ion (deuterium) and electron densities.

The code has been validated in Refs. 16 and 17 with respect to the linear dynamics (damping and energetic-particle drive mechanisms), and in Refs. 8 and 20 with respect to the nonlinear dynamics of the EPMs.

Here, 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 ions. Fluid nonlinearities are not expected indeed to considerably alter shear-Alfvén-mode dynamics for the cases under investigation.\(^{8,23}\)

III. 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,\(^4\) is characterized by a time scale of the order of 100 ms, a relatively large fluctuating magnetic-field level ($\Delta B/B \sim 10^{-3}$ at the first wall), and large growth rates (corresponding to wide frequency spectra: $\Delta \omega \sim 50$ kHz, $\Delta \omega/\omega \sim 1$). The ALEs are followed by a significant drop of the volume integrated neutron emission rate, $S_n$. This quantity is related with the number of energetic ions in the central region of the plasma; its decrease has then been interpreted as the effect of a radial redistribution of the energetic ions producing a marked reduction of their on-axis density. The ALE repetition rate appears to be inversely correlated 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 are observed: the fast FS modes.\(^4\) Such modes are characterized by longer time scale (a few milliseconds), lower growth rate and fluctuating field level, and much smaller redistribution of the energetic ions than ALEs.\(^5\) Most of them consist of bifurcating branches, chirping up and down in frequency. Figures 1 and 2 (Ref. 5) are representative of these experi-
mental results: Figure 1 presents the contour plot in the \((r/a, \omega_0/\omega)\) plane of the fluctuating poloidal magnetic field \(\hat{B}_\rho\) measured by Mirnov coils at the first wall, during a typical JT-60U NNB-heated discharge; the ALE, which occurs at \(t = 4345\) ms, can be easily distinguished from the fast FS modes. Figure 2 plots, for several bursting events, the mode amplitude (left) and the relative variation of the neutron emission rate, \(\Delta S_n/S_n\), induced by each bursting event (center) versus \(S_n\), and the burst period versus \(\Delta S_n/S_n\) (right).

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 (#E039672) with the HMGC code.\(^\text{14}\) 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, with on-axis density \(n_H = 0.12 \times 10^{19} \text{ m}^{-3}\). Other relevant parameters are \(a = 0.95\ m\), \(R_0 = 3.3\ m\), toroidal magnetic field \(B_T = 1.2\ T\), on-axis electron density \(n_e = 2.5 \times 10^{19} \text{ m}^{-3}\), on-axis ion density \(n_I = 1.74 \times 10^{19} \text{ m}^{-3}\) (main ion species: deuterium), on-axis ion temperature \(T_I = 2.1\ \text{keV}\), on-axis safety factor \(q = 1.3\). Modes with toroidal number \(n = 1\) are considered, to compare simulation results with the experimental observations. The linearly unstable phase is dominated in this “ALE” simulation, by a fast-growing mode \((\gamma/\omega = 0.5)\), located around the maximum energetic-ion pressure gradient \((r/a = 0.5)\) and characterized by a significant coupling with the Alfvén continuum, showing the EPM nature of the mode [see left-hand side of Fig. 3, where the contour plots of the power spectrum of the fluctuating scalar potential are shown in the \((r, \omega)\) plane]. Once the mode saturates, a complex phenomenology appears. At an earlier stage, the configuration is dominated by a mode that progressively adjusts its frequency in order to minimize its coupling with the Alfvén continuum and eventually localizes, as a toroidal Alfvén eigenmode (TAE)-like mode, at \(r/a \approx 0.8\), above the lower continuum. The original mode, localized at half radius, is replaced by a couple of nearly degenerate modes. A weak core-localized mode\(^\text{30,31}\) also appears, with frequency well localized in the toroidal gap (Fig. 3, center). At later times, the core-localized mode becomes the dominant one, while the external mode still persists at lower amplitudes (right-hand side of Fig. 3).

Figure 4 shows the wave-particle power exchange (averaged over the poloidal and toroidal angles) in the energetic-ion \((\hat{E}, \alpha)\) plane during the linear phase at \(r/a = 0.5\), where the EPM is peaked. Positive sign corresponds to ions pumping energy into the wave. Here, \(\hat{E}\) and \(\alpha\) are defined by \(\hat{E} = \hat{\mu} + \hat{v}_r^2/2\), \(\alpha = \arccos(v_t/(2\hat{\mu} + \hat{v}_r^2)^{1/2})\), with \(\hat{\mu} = \mu\Omega_{00}/m_H\hat{v}_r^2\) and \(\hat{v}_r = v_r/v_H\) being the normalized magnetic moment and parallel velocity, \(m_H\) and \(\Omega_{00}\) the energetic-ion mass and on-axis Larmor frequency, respectively, and \(v_H = (E_{\text{peaks}}/m_H)^{1/2}\). The resonant drive is given by circulating energetic ions, consistent with the nearly tangential beam injection; it appears to be quite broad in energy, in agreement with estimates based on the experimental measurements.\(^\text{32}\)

As the nonlinear effects become important, a macroscopic outward displacement of the energetic ions is observed, producing a significant reduction of their density in the central region (left-hand side of Fig. 5), which matches reasonably well the energetic-ion density profile after the ALE inferred from the measurements of the radial neutron emission profiles (black curve in Fig. 5, left). The fluctuating poloidal magnetic field \(\hat{B}_\rho\) vs \(t\) (at a radial position close to the plasma boundary) is plotted in the center of Fig. 5, while the right-hand side of Fig. 5 shows the power spectrum of the same quantity in the plane \((t, \omega)\). The time scale \(\Delta t\) of the “burst” is \(\Delta t \approx 200\omega_0^{-1}\), which corresponds, for the considered parameters, to \(\Delta t \approx 150\ \mu s\). The frequency spread is of.

FIG. 3. (Color online) Contour plots of the power spectrum of the fluctuating scalar potential in the \((r/a, \omega_0/\omega)\) plane, for a typical JT-60U discharge (#E036378), during the linear-growth phase (left) and the early (center) 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, \(\omega_0 = \omega_0/R_0\) is the on-axis Alfvén frequency and \(v_0\) the on-axis Alfvén velocity. A rainbow scale has been used, ranging from violet (zero) to red (maximum value relative to each frame).
order $\Delta \omega \approx 40$ kHz. Both these results are close to the experimental observations.

Given the fair qualitative and quantitative 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. A possible explanation of such a weaker fluctuation phase could be that the modification of the energetic-ion radial density gradient reduces the instability level of the system after an ALE. To test this hypothesis, we perform a simulation (“after ALE”) assuming an initial energetic-ion density profile equal to that experimentally observed after the ALE, while postulating, as before, an anisotropic SD distribution function in the velocity space. The linear growth rate of the mode is lower ($\gamma/\omega = 0.28$) than that found in the ALE simulation, but still very high (this could have been indeed expected observing that, in the region where the linear mode is driven, magnitudes of the initial radial gradient of the energetic-ion pressure in the two simulations are not very different: see left-hand side of Fig. 5 and compare the blue and black curves). Correspondingly, in spite of modest nonlinear effects on the spatial energetic-ion distribution (left-hand side of Fig. 6), fluctuating field levels are quite large (right-hand side of Fig. 6): 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 ($\sim 0.2$). 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 is confirmed by a third simulation (“intermediate”), characterized by an intermediate initial energetic-ion density profile (namely, the average between the two initial profiles adopted in the ALE and after-ALE simulations); the final profile is substantially coincident with that found at the end of the two previous cases (Fig. 7).

The reason for the above numerical results can be traced back to the fact that, in assuming a SD for the initial distribution function in the after-ALE simulation, we have retained only part of the ALE nonlinear effects (the spatial redistribution), with no attention to the distortion produced in the velocity space. In order to check this conjecture, we perform a further numerical experiment (“after-ALE modified”) in which we initialize the energetic-ion distribution function according to that obtained in the ALE simulation, after the rapid dynamics of the event has taken place ($t\omega_{\phi} \approx 500$, see right-hand side of Fig. 5); in such a way, we retain both the radial profile broadening and the distortion of the distribution function in the velocity space. As expected, the left-hand side

FIG. 4. (Color online) Wave-particle power exchange (averaged over the poloidal and toroidal angles) in the energetic-ion ($E, \alpha$) plane during the linear phase, at $r/a \approx 0.5$, where the most unstable mode is peaked. The positive sign corresponds to ions pumping energy into the wave. The solid lines contain the trapped particle region. Here and in the following colored contour plots, a rainbow scale has been used, with negative values represented by colors ranging from violet (minimum value) to light blue, and positive values by colors from light green to red (maximum value).

FIG. 5. (Color online) “ALE” simulation. (Left) Nonlinear modifications of the energetic-ion density profile (the blue curve refers to the initial—simulation and experimental—profile and the red curve to the simulation relaxed profile; the black curve corresponds to the experimentally inferred profile just after the ALE and is plotted here for comparison). (Center) Fluctuating poloidal magnetic field $\vec{B}_\phi$ vs $t$ (at a radial position close to plasma boundary). (Right) Power spectrum of the same quantity in the $(t, \omega)$ plane.
of Fig. 8 still indicates that only a negligible further relaxation of the energetic-ion density profile is obtained in the new simulation. Meanwhile, Fig. 8, center, shows the normalized amplitude of the fluctuating poloidal magnetic field close to the plasma edge, which is now close to the experimental observations (−0.2 the level found in the ALE simulation). Moreover, the characteristic time scale of the mode becomes longer, consistent with the weaker drive, with the power spectrum traces bifurcating in two branches, chirping up and down in frequency (Fig. 8, right): the frequency sweeping of the weaker fluctuation supports its identification with a fast FS mode.

The latter results show that the effects of the nonlinear saturation of the EPM on the velocity distribution function can indeed play a fundamental role in explaining the ALE-fast FS dynamics. We can address this point in a more quantitative way, by comparing the initial velocity-space distribution functions adopted in the above after-ALE and after-ALE-modified simulations (Fig. 9). The former—a SD distribution function—is plotted in the (\(\hat{E}, \alpha\)) plane in the left frame; the latter—corresponding to the distribution function obtained in the ALE simulation at \(t_{\tau_{AE}}=500\) (late saturated phase)—is reported in the central frame. Both functions refer to the radial position \(r/a=0.5\), where the EPM is peaked, and are normalized to the local energetic-ion density. We see that the nonlinear ALE dynamics (fully retained only in the after-ALE-modified case) causes both a spread of the distribution function in the low-energy region and a depletion of its high-energy portion. Such a redistribution can be better seen in the right-hand side of Fig. 9, which shows the difference between the two distribution functions. It is not surprising that the instability of the system is significantly reduced in the after-ALE-modified framework, as the initial energetic ion distribution function is apparently depleted in the region of the velocity space where the drive is expected to be larger. This can be appreciated by weighting the difference reported in the right-hand side of Fig. 9 by the power transfer plotted in Fig. 4. The resulting quantity, which represents an estimate of the local variation of the drive in the velocity space, is shown in Fig. 10: the overall effect is definitely negative.

We expect that, until the resonance region in the velocity space has been refilled, the system is stable enough to preserve the energetic-ion density profile from radial redistribution. To check the validity of such expectation, we perform one more simulation ("after-ALE–peaked density" simulation), in which we initialize the energetic-ion distribution function assuming a fully reconstructed density profile (equal to that adopted in the ALE simulation) and a fully distorted velocity-space distribution function (as in the after-ALE-modified simulation). Such a simulation shows (Fig. 11) that even a large energetic-ion density gradient could be maintained against instabilities in the presence of a depleted velocity-space distribution function.

On the basis of these results, a possible interpretation of the experimental observations is the following: given the energetic-ion source provided by the beam injection, the free-energy reconstruction rate is mainly set by the need for rebuilding the velocity-space distribution function after the occurrence of an ALE. 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-rate and am-
plitude modes, such as the fast FS modes, are then excited with 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 an EPM, a new ALE occurs.

IV. ALE SATURATION

In the previous section we observed that considering the nonlinear distortion of the velocity-space distribution function is needed for a satisfactory interpretation of the observed ALE-fast FS phenomenology in terms of Alfvén-mode dynamics. In the present section we investigate more accurately the details of the saturation of the EPM in the ALE simulation.

To this aim we consider three radial shells: $0 \leq r/a \leq 0.33$ (left), $0.33 \leq r/a \leq 0.67$ (center), and $0.67 \leq r/a \leq 1$ (right). Each shell is centered, in fact, around one of the three main modes that appear in the different stages of the simulation: the weak core-localized mode characterizing the late saturation phase, the EPM dominating the linear phase, the outer TAE-like mode emerging in the early saturation phase (cf. Fig. 3). Figure 12 shows the contour plots of the volume-averaged wave-particle power transfer in the energetic-ion $(\hat{E}, \alpha)$ plane for the three shells, during the linear phase. As observed in the previous section, the EPM dominating this phase is strongly coupled with the Alfvén continuum. Then, it can only exist if the energetic-ion drive is sufficient to overcome the large continuum damping; that is, if the mode finds an efficient compromise, in choosing its frequency, between minimizing the damping and maximizing the resonant drive. From Fig. 12 we see that such a condition is clearly satisfied in the two outer shells, where almost all the resonant particles pump energy into the mode. In the inner shell, the structure of the resonance corresponds to draining energy from the lower parallel velocity particles (larger $\alpha$ values), while transferring it to the higher velocity ones (lower $\alpha$), with a little net contribution to the drive. This is consistent with the mode structure observed in the left-hand side of Fig. 3, which extends preferentially over the two outer shells.

Figure 13 presents, for the same three radial shells, the volume-averaged variation of the energetic-ion distribution function in the $(\hat{E}, \alpha)$ plane induced by the saturation of the

FIG. 8. (Color online) “After-ALE-modified” simulation. (Left) Nonlinear modifications of the energetic-ion density profile. The black curve refers to the initial profile (equal to the relaxed one in Fig. 5, left). The red curve represents the profile obtained at the end of the simulation. The initial profile used in the ALE simulation is also reported for comparison (blue curve). (Center) Fluctuating poloidal magnetic field $\tilde{B}_p$ vs $t$ (at a radial position close to plasma boundary). (Right) The power spectrum of the same quantity in the $(t, \omega)$ plane. Note that the two quantities have been reported in the same scales adopted in Fig. 5, center and right, respectively.

FIG. 9. (Color online) Plot in the energetic-ion $(\hat{E}, \alpha)$ plane of the slowing-down distribution function (left) and that for the distribution function obtained, in the ALE simulation, in the late saturated phase (center). Both functions refer to the radial position $r/a = 0.5$, where the EPM is peaked, and are normalized to the local energetic-ion density. (Right) Plot of the difference between the two normalized functions.
EPM nonlinear dynamics. The time considered corresponds to the early saturated phase (Fig. 3, center). Consistent with the result reported in Fig. 9, right, the EPM saturation is accompanied, in the radial shell where the mode is peaked, by a scattering of energetic ions out of the resonance region (Fig. 13, center). A similar effect involves the energetic ions in the inner shell (Fig. 13, left). We indeed observe that structured variations of the distribution function become appreciable simultaneously in this shell and in the central one. The energetic-ion scattering has then to be ascribed to the interaction with the EPM (the core-localized one being, at the considered time, too weak to appreciably influence particle orbits). Note that the depletion in the inner shell is relatively less pronounced in the low-energy region; this seems consistent with the fact that less energetic ions have weaker wave-particle interactions as well as narrower orbits and do not explore the radial region where the EPM is peaked.

The variation of the distribution function in the central shell occurs simultaneously to the beginning of the resonance destruction. This can be seen from Fig. 14, which plots the wave-particle power transfer at the same time considered in Fig. 13 (note that the time evolution of this quantity would show that, in the central shell, not only the definite signature, but also the time coherence is lost). The power transfer gets unstructured in the inner shell as well, in spite of the fact that the core-localized mode is clearly visible in the considered phase (Fig. 3, center), eventually dominating the later one (Fig. 3, right). This is, however, not surprising. Such a mode rises, indeed, in a region characterized by a little energetic-ion pressure gradient. Its existence condition is then determined, at the leading order, by the need for avoiding the potentially large continuum damping. This condition can be reached only if its frequency localizes in the toroidal gap, where the coupling with the Alfvén continuum is negligible. Different from the EPM case, the drive intensity anyway small, because of the little pressure gradient) plays no role in determining, at the leading order, the frequency of the mode; it is rather determined itself by the given frequency and mode structure. In this situation, on one side, an unstructured resonance pattern can be obtained (as it is, in fact); on the other side, even a little net power transfer from the ions to the wave is sufficient to make the mode grow.

For the outer shell, though the signature of the distribution function modification is not very different from that in...
the central shell, the results obtained in this phase have specific characteristics. The velocity-space distribution function shows indeed (Fig. 13, right) a smaller depletion of the large-energy resonant population than that observed for the central shell, while the increase in the relatively low-energy region, where the local linear-phase resonance was less important (Fig. 12, right), is more evident in this case. Moreover, a clear resonance structure can still be observed (Fig. 14, right), and its variation with respect to the linear-phase structure appears to be related with the distribution function modification. This indicates that the “low-energy” energetic ions are able to resonate with the outer mode. A natural interpretation of these results could be the following: the central EPM grows; resonant ions are displaced outwards; the EPM resonance pattern is destroyed and the mode saturates; the displaced ions affect the resonance pattern in the outer shell, driving the TAE unstable. Such a picture is consistent with the “avalanche” paradigm, based on the mode particle pumping and discussed in Refs. 20–23: a radial mode structure that evolves nonlinearly on the same time scale of the convective energetic-ion transport, following the unstable moving front.33 To test the validity of this interpretation, we have to show that the central-shell depletion and the outer-shell increase of resonant ions are abridged by a radial displacement, rather than being two local, mutually independent phenomena. To this aim, we separate, for each shell, the displacement of resonant ions and, hence, modifies the distribution function, not in the sedentary-particle one. We can then conclude that there is an intrinsically different cause-effect relationship between the mode growth and the local distribution-function modification in the various radial regions. While the mode growth in the central region causes the outward radial displacement of resonant ions (and, hence, modifies the distribution function), in the outer region it is the variation of the distribution function, due to energetic ions redistributed from the central region, which causes further local mode growth.

V. SUMMARY AND DISCUSSION

In this paper we have presented the results of a numerical investigation, based on particle-in-cell simulations, of the bursting-mode phenomenology observed in NNB-heated JT-60U discharges. We have shown that the experimental observations can be interpreted as the effect of the nonlinear interaction between Alfvén modes and the energetic ions produced by the NNB injection. In particular, our investigation, related to $n=1$ modes, shows that an EPM localized around the maximum of the energetic-ion pressure gradient is driven unstable by the resonant interaction with such ions. Its saturation produces a radial displacement of energetic ions in fair agreement with the experimental findings related with ALEs. Simulation results demonstrate that the displaced ions resonate with the Alfvén modes in the outer region, causing a TAE-like mode to become dominant as the saturation of the EPM proceeds.

We also observe that the scattering due to the EPM is more effective on the resonant ions than the nonresonant ones. Besides the relaxation of the density profile, a distortion in the velocity-space distribution function is then produced. This fact can explain why a quieter phase, characterized by weaker bursting modes (the fast FS), is observed after an ALE, allowing for the system to restore the free
energy needed for a new ALE. In the absence of the velocity-space distortion, any reconstruction of the density profile would indeed generate relatively large amplitude modes: energetic ions would be further scattered by these modes and their density profile would be essentially clamped to the relaxed profile produced by the ALE. Once the phase-space distortion is fully taken into account, the free-energy reconstruction rate is instead set by the need for rebuilding both the density profile and the resonant part of the distribution function. The slow time-scale evolution of energetic-ion equilibria in intermediate configurations between two successive ALEs is then characterized by a lower drive than that corresponding to the unperturbed velocity-space distribution function, and the weak modes excited are less effective in contrasting the density profile reconstruction. Only when the combined restoration of the configuration and velocity-space distributions provides enough drive for a fast growing Alfvén mode, a new ALE occurs.

Several approximations were adopted in the model, but they can be justified a posteriori. In particular, treating the equilibrium as a circular-magnetic surface one seems to be fairly accurate for the main mode (the EPM), which is localized around the maximum fast-ion density gradient ($r/a = 0.5$) where the actual magnetic surfaces, for the specific discharge examined in the paper, are not very different from circular ones. Moreover, taking into account the noncircular shape would mainly affect the toroidal gap structure, with little impact on a strongly driven mode like the EPM (characterized by frequency relatively far from the gap and, hence, by an essentially cylindrical nature). Neglecting the finite-Larmor-radius effect is appropriate too, for the EPM: looking at the portion of the energetic-ion population that mainly contributes to the drive ($E \approx 0.5E_{\text{beam}}$), we indeed get $k_{\perp} \rho_{T} \approx 0.3$, which means that the guiding-center approximation affects the simulation results with a relative error of a few percent. Finally, neglecting the fluid nonlinearities should not prevent a satisfactory description of the nonlinear evolution of the EPM. Indeed, as observed in Ref. 23, EPM growth rate has a stronger $n$ dependence than Alfvén eigenmodes (e.g., TAE), producing typically a narrow toroidal-mode-number unstable spectrum. As a consequence, single-$n$ nonlinear dynamics dominates the rapid convective phase that characterizes the EPM saturation and is essentially dictated by the energetic-ion nonlinear response.

Further development of the present particle-simulation approach should include, eventually, the possibility of treating energetic-ion sources and Alfvén-mode nonperturbative nonlinear dynamics on the same footing. At the moment, this is beyond reach for all existing hybrid MHD codes and remains one of the challenging goals for integrated modeling of burning plasmas.

Note that the unstable moving front effect depends on the nature of the mode. For high toroidal mode number, as in Ref. 20, the wave-packet amplitude is convectively amplified along its secular radial motion. In the present $n=1$ case, it is the sequence of mode excitation at different radial locations that is synchronized by the coherent nonlinear interaction with the fast ions. No appreciable radial structure distortion is instead observed for the nonlinear fishbone cycle, discussed in Ref. 22, due to the very peculiar wave-field form of the internal kink mode.