Dynamics of spherical tokamak plasma on the internal reconnection event

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Nonlinear magnetohydrodynamic (MHD) simulations are executed to investigate the dynamical behavior of the relaxation phenomenon observed in spherical tokamak (ST) plasma that is known as the Internal Reconnection Event (IRE). The simulation results successfully reproduce several key features of IRE, and the physical mechanisms are revealed. A sudden collapse of the pressure profile takes place as a result of a nonlinear time development of a pressure-driven instability. A magnetic reconnection induced between the internal and the external magnetic field is found to play a crucial role in determining the nature of the overall process, namely, the rapid expulsion of the plasma heat energy due to the pressure imbalance along the reconnecting field lines, and the large distortion in the overall shape. The resultant deformations in overall shape of the plasma are in good agreement with the experimental observations. © 2000 American Institute of Physics. [S1070-664X(00)02203-5]

I. INTRODUCTION

The spherical tokamak (ST) experiments have demonstrated their promising properties, such as excellent stability and confinement at high $\beta$.

The START (Small Tight Aspect Ratio Tokamak) device at Culham has achieved the highest average $\beta$ of 40% in 1999.

Experimentally, the ST plasma has revealed several unique features in its dynamical behavior. The Internal Reconnection Event (IRE) is the most noticeable relaxation phenomenon, which is characterized by a collapse and a subsequent resilience of the pressure profile, the appearance of a positive spike in the net toroidal plasma current trace, and a large elongating distortion in the overall shape of plasma. IRE has been clearly observed in a couple of devices.

It is differentiated from the major disruptions in conventional tokamaks in that the plasma current is not terminated by IRE, and from the sawtooth oscillations in that IREs are observed with an increase of plasma current, and in irregular intervals. It should be noted, however, it has been observed in more recent experiments in START that IRE can sometimes lead to occurrence of current-terminating disruptions, especially after the divertor coils were installed close to the core plasma.

The occurrence of IRE needs to be carefully considered when we aim to realize a fusion reactor by using the ST concept, partly because each time IRE occurs about 30% of the stored heat energy is lost, and partly because it can sometimes lead to an occurrence of major disruptions. The physical mechanisms of IRE have not been clarified well enough yet. To know them is an important and an urgent research topic for exploring larger scale ST devices.

The advanced diagnostics has provided many clues for understanding the mechanisms of IRE. Before describing the simulation results, we review the phenomenology of IRE. As is shown in Fig. 10 of Ref. 12, which is the temporal variation in electron temperature and density profiles through an IRE measured by Thomson scattering in START, the central values fall in a short time scale of around 100–150 $\mu$sec, and the radial profiles are broadened to the edge during an IRE. The experimental observations for the temporal change in the net toroidal current and the soft x-ray signal are plotted in Fig. 1. An IRE usually proceeds in three stages in START experiments, namely, thermal quench, current increase, and current quench stages, which are referred to as I, II, and III in Fig. 1(b), respectively. Furthermore, the thermal quench stage usually proceeds in two steps, which are indicated by IA and IB, respectively, where the signal falls more slowly in the first (IA) step than the other (IB).

In the START experiment, the overall structure of the plasma is monitored by using a CCD (charge coupled device) camera. Much fruitful information about the large-scale nonlinear behavior of plasma can be derived from the images. Some of the snapshots taken during an IRE are shown in Fig. 2, in which the brighter region corresponds to relatively cold plasma. Figure 2(a) is taken by a monochromatic CCD camera, which has a higher time-resolution of about 30 $\mu$sec, at a moment corresponding to an early stage of IRE. A bright helical structure is observed to extend in the periphery from the top down to the bottom of the torus. It is noted that the structure exists only in a narrow region and is localized in the toroidal direction. Similar localized structures are often observed during IREs.

Figures 2(b)–2(d) are taken by a color CCD camera which has relatively poor time-resolution of about 100 $\mu$sec, which is comparable to the time scale of the fall of the peak pressure. Although instantaneous structures may be smoothed out, characteristic conical structures, which consist of multiple layers, are clearly seen at the top and the bottom of the torus in Fig. 2(b) (denoted by A). A filamentation structure is observed as bright stripes in the middle of the image (denoted by B). An overall distortion is captured in Figs. 2(c) and 2(d), which appears in the slightly...
A model for describing IREs has been developed by Buttery et al. In Ref. 14, they reconstructed equilibrium using a free-boundary Grad-Shafranov code with the experimental data obtained in START. By comparing the pre- and the post-IRE states, they found out a conservation constraint in “modified helicity.” Scanning in various aspect ratios and falls in $\beta_p$ using this constraint, they successfully explained the distinctive features of IRE such as disruption immunity at low aspect ratio.

Such is an outline of experimental observations and theoretical analyses of IREs. It can be summarized that an IRE is a relaxation phenomenon which accompanies a drastic change in overall shape in a short time scale, and is a highly nonlinear and fully three-dimensional phenomenon. It is necessary for understanding these complex features to make use of a nonlinear numerical simulation. In our previous simulation, we have succeeded in reproducing several main features of an IRE such as a rapid fall in pressure and a large distortion in overall shape. In this paper, by advancing the analysis of the simulation result to a direct comparison with experimental results, the physical mechanisms of IRE are revealed.

There have been several nonlinear simulation works concerning disruption phenomena in a conventional tokamak. Kleva and Drake studied the density limit disruption by using a fluid model including the edge radiation. They explained the pressure degradation by an injection of cold bubble as a result of nonlinear development of the $q=1$ kink mode, and succeeding relaxation between multiple $q=2$ modes. Park et al. executed a nonlinear simulation to study high-$\beta$ disruptions. They showed that toroidally localized high-$n$ ballooning modes driven by a local pressure steepening due to evolution of low-$n$ kink modes make a local pressure bulge and destroy the flux surfaces. In contrast, Kleva and Guzdar claimed that a beta limit disruption can be driven by a dynamical process of high-$n$ ballooning modes. IRE is different from these disruptions that occur in a conventional tokamak. The main difference is a large distortion in overall shape, and experimentally the property of the strong resiliency. In this paper, we investigate the nonlinear development of low-$n$ modes in a simulation geometry in which the open external field is included to treat such a large deformation.

In Sec. II, the simulation model and the numerical method are introduced. The simulation results are presented in Sec. III. The comparisons between the simulations and the experiments, and some additional investigations, are discussed in Sec. IV. Finally, the results are summarized in the last section.

II. SIMULATION MODEL

The numerical simulations are executed in a three-dimensional full-toroidal geometry as shown in Fig. 3. The compressible nature in the governing equations is fully retained in the simulation model because of its importance in the nonlinear dynamics in this configuration. The governing
The equations are the full set of normalized nonlinear resistive compressible magnetohydrodynamic (MHD) equations as follows:

\[
\frac{\partial \rho}{\partial t} = -\nabla \cdot (\rho \mathbf{v}), \tag{2.1}
\]

\[
\frac{\partial}{\partial t}(\rho \mathbf{v}) = -\nabla \cdot (\rho \mathbf{v} \mathbf{v}) - \nabla p + \mathbf{j} \times \mathbf{B} \tag{2.2}
\]

\[
+ \mu \left( \nabla^2 \mathbf{v} + \frac{1}{3} \nabla (\nabla \cdot \mathbf{v}) \right), \tag{2.3}
\]

\[
\frac{\partial \mathbf{B}}{\partial t} = -\nabla \times \mathbf{E}, \tag{2.4}
\]

\[
\frac{\partial p}{\partial t} = -\nabla \cdot (p \mathbf{v}) - (\gamma - 1)(p \nabla \cdot \mathbf{v} + \eta \mathbf{j}^2 + \Phi), \tag{2.5}
\]

\[
\mathbf{j} = \nabla \times \mathbf{B}, \tag{2.6}
\]

\[
\mathbf{E} = -\mathbf{v} \times \mathbf{B} + \eta \mathbf{j}, \tag{2.7}
\]

\[
\Phi = 2\mu \left( e_{ij} e_{ij} - \frac{1}{3} (\nabla \cdot \mathbf{v})^2 \right), \tag{2.8}
\]

\[
e_{ij} = \frac{1}{2} \left( \frac{\partial v_i}{\partial x_j} + \frac{\partial v_j}{\partial x_i} \right), \tag{2.9}
\]

where \(\rho\) is the mass density, \(\mathbf{v}\) is the velocity, \(\mathbf{B}\) is the magnetic field, and \(p\) is the plasma pressure. The equations contain the resistivity \(\eta\) and the viscosity \(\mu\) as uniform constants. The ratio of specific heats \(\gamma\) is 5/3.

The equations are solved in a cylindrical \((r, \theta, z)\) geometry with a rectangular poloidal cross section (see Fig. 3). The spatial derivatives of the equations are represented by the fourth-order finite differences and the time integration is calculated by the fourth-order Runge-Kutta method. The grid numbers are chosen to be \((N_r, N_{\theta}, N_z) = (128 \times 64 \times 128)\), which is large enough to treat the large-scale phenomena. The numerical soundness of the choice of the grid number and the time step has been confirmed by comparing results computed with coarser grid numbers or different time step. The initial condition is given by solving the Grad-Shafranov equation numerically under an assumption of polynomial pressure and current profiles. It should be noted that the computation region includes both inside and outside the separatrix. An open magnetic field extends in the external region (see Fig. 3). The plasma pressure gradient exists only inside the separatrix in the initial condition. The external region is initially filled with a low temperature plasma, which equals 10% of that at the magnetic axis. The initial mass density is assumed to be uniform. The boundary condition is put as a perfect-conducting and no-slip wall at all boundaries of the computation region.

The variables are treated as a normalized form. Length, magnetic field, and mass density are normalized to the major radius of the geometrical center of the computation region, the toroidal field at the magnetic axis, and the initial mass density, respectively. Therefore, the unit of time equals the Alfvén transit time \((\tau_A)\) encircling the magnetic axis divided by \(2\pi\).

### III. SIMULATION RESULTS

The simulation starts from a linearly unstable configuration which causes initial tiny perturbations to grow spontaneously. The perturbation is introduced on the plasma velocity field at \(t=0\) as a random white noise. We have tried to execute the simulations under several conditions. A typical case, which well reproduces the properties of IRE, is presented in this section. The parameters used in the calculations are the aspect ratio \(A=1.5\), the elongation \(\kappa=1.9\), the central beta \(B_0=44\%\), the full volume averaged beta \(\langle \beta \rangle = (\langle \rho \rangle / \langle B^2 \rangle / 2) = 8.1\%\), the volume averaged beta with a commonly used definition \(\beta_T=\langle \rho \rangle / B_0^2/2 = 12\%\), where \(B_0\) is the vacuum field at the magnetic axis, and the central safety factor \(q_0=0.91\). \(\mu\) and \(\eta\) are assumed to be uniform values of \(4 \times 10^{-4}\) and \(1 \times 10^{-5}\), respectively.

Shown in Fig. 4(a) is the time development of the perturbation in magnetic energy for each toroidal Fourier component. In the early stage of the growth (until \(t=170\tau_A\)), each of the components grows exponentially with its linear growth rate and phase independent of other components. Hereafter, we call it the “linear stage.” The subsequent time development may be called the “nonlinear stage.”

![FIG. 4. (a) Time development of perturbation in the magnetic energy for each toroidal Fourier mode in the simulation result. (b) Growth rate of the linear instability for each toroidal Fourier mode.](image-url)
A. Linear stage

The growth rate for each toroidal mode number \((n)\) in the linear stage is plotted in Fig. 4(b). It can be seen that multiple numbers of eigenmodes are excited simultaneously and the growth rates come out to be larger for lower \(n\) modes. Particularly, the \(n=1\) and \(n=2\) modes have almost the same large growth rate. Since the initial perturbations include every mode with the same amplitude (white noise), the two modes dominate the growth with similar amplitude.

Figure 5 shows the poloidal mode structure for the \(n=1\) and the \(n=2\) modes. The variation of the color indicates the amplitude of the fluctuations in the plasma pressure. It can be clearly seen that the dominant poloidal components are \(m=1\) for the \(n=1\) mode and \(m=2\) for the \(n=2\) mode. The position of the \(q=1\) rational surface drawn together in Fig. 5 suggests that these modes are resonant with it. The flow patterns are also drawn in Fig. 5. It can be seen that \(m\) pairs of vortices are formed on the poloidal cross section.

It is possible to identify the driving sources of these eigenmodes by estimating the kink and the interchange terms of the potential energy integral \((\delta W)\) appearing in the energy principle. These terms are expressed as

\[
\delta W_{\text{kink}} = \frac{1}{2} \int d^3x \mathbf{J}_0 \cdot \mathbf{B}_0 \frac{\mathbf{B}_0 \times \mathbf{\xi} \cdot \mathbf{\bar{B}}}{|\mathbf{B}_0|^3}, 
\]

\[
\delta W_{\text{interchange}} = - \int d^3x \langle (\mathbf{\xi} \cdot \nabla p_0)(\mathbf{\xi} \cdot \mathbf{\kappa}) \rangle,
\]

where \(\mathbf{\xi}\), \(\mathbf{\kappa}\), and \(\mathbf{\bar{B}}\) denote the displacement vector, the normal curvature of the equilibrium magnetic field, and the perturbation of the magnetic field, respectively. All the equilibrium quantities are denoted by suffix 0. Substituting the eigenfunction obtained by the simulation for \(\mathbf{\xi}\), we can estimate each term as listed in Table I with other stabilizing terms. Both of the kink and interchange terms are negative for both \(n=1\) and \(n=2\) modes, hence, they are driven by the combination of the current and the pressure gradient. However, the interchange term is much larger than the kink one in absolute value for each mode. Therefore, these modes are identified as pressure-driven modes. In addition, judging from the poloidal mode patterns shown in Fig. 5, which are not localized in the poloidal direction, it is concluded that these modes are not categorized into the ballooning mode but into the interchange mode.

B. Nonlinear stage

Until the amplitude for each mode reaches a certain level (around \(1 \times 10^{-6}\) for the magnetic energy), the instabilities keep growing linearly with the growth rates and the mode structures fixed, as described in the previous subsection. The first sign of the nonlinearity appears as an abrupt change in the growth rate for the higher \(n\) modes. As seen in Fig. 4(a), the growth rate of the \(n=4\) mode suddenly increases at \(t = 130 \tau_A\). It is caused by the development of the higher harmonics for the \(n=2\) mode and by the nonlinear couplings between the \(n=1\) and the \(n=3\) modes. In the early nonlinear stage, other higher \(n\) modes are also excited in the similar way. The growth of these nonlinear modes is so rapid that the fluctuation levels soon become significant enough to affect the global configuration. On the other hand, the growth rate of the \(n=1\) and \(n=2\) modes is saturated at around \(t = 190 \tau_A\) as seen in Fig. 4(a).

The time evolution of the pressure profile in the early nonlinear stage is indicated in Fig. 6. The displacement becomes a visible scale at \(t = 165 \tau_A\) [Fig. 6(b)]. The pressure profile is elliptically elongated in the poloidal cross section, reflecting the convection motion of the \(m=2n=2\) mode. At around \(t = 190 \tau_A\) [Figs. 6(c)–6(d)], the deformation develops into a large bulge. The plasma pressure at the core region is transported from the core to the edge by such convection motions in a short time scale. The rapid fall of the central pressure observed in the experiments is explained by this mechanism. It is worthy of notice that although in the linear stage the instabilities have the nature of the internal modes, in the nonlinear stage they grow so much that even the out-
The outermost surface of the torus is deformed like an external mode. The three-dimensional view of the deformation of the torus at $t=195\tau_A$ is shown in Fig. 7 by a plot of an iso-surface of the pressure near the periphery of the torus. The structure of the surrounding magnetic field is also plotted in the figure by picking up typical external field lines. It can be clearly seen that toroidally localized deformation is formed as it pushes aside the neighboring magnetic field lines.

Due to the deformation of the surface, a large topological change in the configuration, namely, the magnetic reconnection, is induced on the surface of the torus. Shown in Fig. 8 is a magnetic field line at a different time traced from a fixed point on the top boundary, and connects directly into the torus at $t=195\tau_A$.

FIG. 6. (Color) Time development of pressure profile on a poloidal cross section in the simulation result for times equal to (a) 0, (b) 165, (c) 185, and (d) 200 $\tau_A$.

FIG. 7. (Color) Appearance of a localized deformation at $t=195\tau_A$ in the simulation result. (a) and (b) are viewed from the different angles of 180 deg apart from each other in the toroidal direction. The iso-pressure surface and the surrounding magnetic field lines are drawn.

FIG. 8. (Color) Magnetic reconnection between the internal and the external field. An identical magnetic field line and an iso-pressure surface near the separatrix are drawn for (a) $t=0$ and (b) $t=195\tau_A$. The magnetic field line is traced from a fixed point on the top boundary, and connects directly into the torus at $t=195\tau_A$.

FIG. 9. (Color) Schematic figure for the dynamics of the magnetic reconnection between the external and the internal field. The external (blue) and the internal (red) field lines reconnect each other in the localized deformation. A pair of strong plasma flows passing each other in the localized deformation is formed along the large pressure gradient [seen in (c)].
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FIG. 10. (Color) Three-dimensional structure of plasma pressure, flow, and the magnetic field at $t = 195 \tau_A$ in the simulation result. The semitransparent surface indicates the iso-contour of plasma pressure, the fat curves are the arbitrary traces of magnetic field lines, and the thin lines represent the plasma flows. The direction of the flow is represented by the variation of the color on the lines [from the blue (red) to the white]. The difference between the blue and the red is the sign of toroidal velocity. The magnetic lines are colored in green to red as the local pressure increases.

FIG. 11. (Color) Nonlinear time development of the three-dimensional structures of the plasma pressure represented by iso-pressure surfaces and the surrounding magnetic field lines. The time equals (a) 0, (b) 195, and (c) $250 \tau_A$.

FIG. 12. (Color) Time development of the pressure profile in the latter stage of the simulation result at (a) $t = 250$ and (b) $340 \tau_A$.

FIG. 13. Change in radial profiles of (a) pressure, (b) current, and (c) safety factor on the midplane. The broken line and the solid line [the cross markers in (c)] correspond to $t = 0$ and $t = 340 \tau_A$, respectively.
Though the system is ideally stable in this state for the modes which are linearly unstable in the initial equilibrium, the resistive modes become unstable due to the increase of the resistivity. The simulation is executed after $t = 500 \tau_A$ for two different resistivity terms, $\eta = 1 \times 10^{-5}$ and $2 \times 10^{-4}$. The time development of the kinetic energy is plotted in Fig. 14. It shows that the perturbations actually grow for larger resistivity. The growth may be considered as the second step in the whole process. The configuration becomes a distorted one like Fig. 14(b). The mode structure for the second step is shown in Figs. 14(c) with a bird’s-eye view of pressure fluctuations as well as poloidal mode patterns in the pressure. It can be seen that the $m = 2/n = 1$ mode is included as a dominant component for the second step. The location of the $q = 2$ surface is also indicated in Fig. 14(d). The mode structure is mostly resonant to the $q = 2$ surface, which shows the general nature of resistive modes.

**IV. DISCUSSION**

**A. Comparison with experiments**

In this section we compare the simulation results with the experimental observations of IREs. Figure 15 shows a temporal change in the radial pressure profile on the midplane in a fixed poloidal cross section, which can be compared with the experimental observations of the electron temperature or the density profiles reported in several papers. The time scale of $1 \tau_A$ for the simulation roughly corresponds to 0.5 µsec for the experiments of START. It can be seen that the pressure profile suddenly collapses at around $t = 200 \tau_A$, while there is no visible change in the linear stage. The time scale of the collapse is around $100 \tau_A$. The pressure profiles are broadened to the edge during the collapse. These features are in good agreement with Fig. 10 of Ref. 12, which is the temporal variation in electron temperature and density profiles through an IRE measured using Thomson scattering.

The second step is also plotted in Fig. 15 for $\eta = 2 \times 10^{-4}$. It can be seen that the excitation of the resistive mode leads to a two-step relaxation in the overall process, which agrees with the experimental observation as is denoted by IA and IB in Fig. 1.

The remarkable changes in the overall shape are compared with the CCD camera images introduced in Sec. I. The helical bright structure seen in Fig. 2(a) may be an instant phenomenon, since it is captured only when it is taken by a high time-resolution camera. The localized deformation described in Sec. III B corresponds to it with respect to the shape and the duration time scale. The conical layers appearing in Fig. 2(b) resemble the structure which appears right after the reconnection. To compare our simulation results with the CCD camera observations, in which the bright part mainly corresponds to the low temperature region, we reconstruct images using a volume rendering scheme to visualize the line-integrated intensity of plasma radiation. The radiation from the plasma is modeled in a simple way, in which the intensity of radiation per unit volume is assumed to be peaked when the local value of the pressure is around a certain low level. The images shown in Fig. 16 are thus created.
by using the simulation results corresponding to Fig. 2. Figure 16(a) is drawn with the variables at \( t = 250 \tau_A \). The conical structures are clearly recognized on the image as a bright part on the top and the bottom of the torus, and are in good agreement with Fig. 2(b). The bulk plasma, on the other hand, forms a spiral structure winding around the center rod as seen in Fig. 16. It also agrees with Fig. 2(b).

The \( m=2n=1 \) distortion which grows in the second step is represented in Figs. 16(b) and 16(c). The only difference between them is the view angle, 90 deg. apart in toroidal angle. The tilt distortion and the fatness in only one side, which are seen in Figs. 2(c) and 2(d), are well reproduced in Figs. 16(b) and 16(c), respectively.

With regard to the toroidal current, there is no remarkable change in our simulation, although the system has experienced such a drastic change. This simulation is corresponding to the thermal quench phase, which is indicated with stage I in Fig. 1, and the toroidal current is not so increased in this stage in experiments, either.

Although the simulation cannot continue up to the current increase and the current quench phase, there is fairly good agreement between the simulation and the experiments. It supports the validity of our modeling.

**B. Spontaneous phase alignment**

The appearance of the localized bulge described in Sec. III.B is a key point of this phenomenon in that it causes the collapse in pressure profile and the expulsion of the heat energy through the magnetic reconnection induced on it. Here we consider how the localized deformation is generated.

The most dominant mode in the linear stage is the \( n=2 \) mode; however, the deformation appears only in one side of the torus in the simulation result, as shown in Fig. 7. The reason why the number of the ridge of the deformation is reduced in the nonlinear stage is found to be the existence of the accompanying \( m=1/n=1 \) mode. We execute another simulation run to check the role of the \( n=1 \) mode in the nonlinear stage. This run is realized by artificially suppressing the growth of the \( n=1 \) mode. The result shows that two bulges originating from the \( n=2 \) mode are formed. It proves the necessity of the accompanying \( n=1 \) mode for the production of the single localized deformation.

The \( n=1 \) mode plays an additional role to the growth of the \( n=2 \) mode. Shown in Fig. 17 is the three-dimensional mode structure of the \( n=1 \) and the \( n=2 \) mode, respectively, where a pair of iso-surfaces for the negative (blue) and the positive (red) fluctuation components in the plasma pressure is plotted. It can be seen that the localized deformation observed in Fig. 7 is formed at the toroidal direction where the phases are aligned between the \( n=2 \) and the \( n=1 \) modes and the positive fluctuations are additional.

This phase-alignment mechanism among the excited modes, as is observed in Fig. 17, is a universal nature in the development of this phenomenon. Namely, the phases are spontaneously aligned to particular modes as a result of nonlinear interaction among them. In order to confirm the nature of the spontaneous phase alignment, we execute a simulation, in which the toroidal phase of the \( n=1 \) mode is initially shifted from that of \( n=2 \). In an axisymmetric system, the toroidal phase of each linear mode is determined by the perturbations given initially. Figure 18 is a top view of the torus and shows the time development of the fluctuations in the pressure on the midplane for \( n=1 \) and \( n=2 \). It is observed that the \( n=2 \) mode never changes the direction, while the \( n=1 \) mode is eventually split in two [Fig. 18(d)], and the outer one spontaneously rotates till it is aligned to the \( n=2 \) mode [Figs. 18(g) and 18(j)]. In this way, the localized deformation always appears in spite of the initial discrepancy among the phases.
C. Dependence on parameters

The linear instability described in Sec. III A is quite sensitive to the initial profiles. The existence of the $q = 1$ surface especially causes the growth of the low $n$ modes resonant to it such as $m = 1/n = 1$ and $m = 2/n = 2$. Here we investigate the nonlinear development when the $q = 1$ surface does not exist. We have executed another simulation run, which we call hereafter case B, starting from an equilibrium where $q_0 = 1.06$, and other parameters are almost the same as those described in Sec. III, which is called case A.

The linear instability for case B is dominated by higher $n$ modes than case A, such as $m = 12/n = 8$. These modes are identified as the ballooning modes by following the same procedure as described in Sec. III A. The figures corresponding to Fig. 4(b) and Fig. 5 for case B are shown in Fig. 19. As seen in Fig. 19(b), the poloidal mode structures are localized in the bad curvature region. The growth rate of those unstable modes for case B is much smaller than that of the dominant modes for case A.

It is worthy of notice that the nonlinear development for case B shows a quite similar behavior to that of case A. Namely, the localized deformation also appears in case B in the early nonlinear stage. In case B, however, the localized deformation consists of several small bulges, not a single one, as shown in Fig. 20. The occurrence of the spontaneous phase alignment described in Sec. IV B is difficult to recognize clearly in case B because of the small discrepancy in phases among the modes due to the high-$n$ nature.

The development after that resembles case A. Namely, the magnetic reconnection between the internal and the external field occurs at a couple of large bulges, and the characteristic conical shape appears at the top and the bottom of the torus. It should be noted that the magnetic reconnection occurs simultaneously at a couple of large bulges, as is seen in Fig. 21. After expelling part of plasma pressure out of the separatrix, the system becomes ideally stable. The stability for the resistive modes is similar to case A, that is, the $m = 2/n = 1$ distortion grows when we use a larger resistivity.

Hence, there is no significant difference in the nonlinear development between the cases A and B, except for the number of localized deformation. That is to say, an IRE can occur under such condition where $q_0$ is greater than unity.

V. CONCLUDING REMARKS

Nonlinear MHD simulations are performed to reveal the dynamical behavior of STs. It is found that multiple mode numbers of pressure-driven instabilities can be excited simultaneously in ST equilibria with the central safety factor close to unity. The nonlinear evolution of these instabilities shows the collapse of pressure profile in a short time scale by a convection motion. Localized deformation, which appears in the early nonlinear stage in consequence of couplings among the linear modes, leads the magnetic reconnection between the internal and the external magnetic field. The heat energy stored in the higher pressure region inside the separatrix is quickly ejected to the external region by strong plasma flow, which is induced along the reconnected field lines, much like a rubber ball which gets a hole on its surface. The torus is stabilized due to the change in the profiles, in particular, the safety factor profile and the decrease in plasma beta. However, a resistive $m = 2/n = 1$ mode still remains unstable.
The results are intensively compared with the experimental observations of START, and are in good agreement with them in several points as follows.

(1) The temporal and the spatial scales of the thermal quench.
(2) Large distortion of overall shape.
(3) Thermal quench proceeds in two steps, separated from increase in the current.

The second point is confirmed by experimental CCD images which show (1) appearance of a toroidally localized bulge, (2) formation of conical structures at the top and the bottom of the torus, and (3) growth of \( m = 2n = 1 \) type helical distortion.

The results shown in this paper suggest the importance of the self-consistent dynamic analysis of the coupling between the internal and external region of a torus in understanding the mechanism of collapsing phenomena including the major disruptions for a conventional tokamak. Our conjecture is that when a torus has a mechanism to expel excessive heat energy quickly to the external region, the process can have a nature of resiliency like IRE, while the process becomes disruptive when the excessive heat energy remains inside long enough to destroy the configuration. Another topic that is raised by this simulation study is the necessity of a detailed analysis of the dynamics of the driven reconnection when there is a large pressure imbalance along reconnected field lines. In this simulation, the physics of heat conductivity along the magnetic field is not included. The parallel heat transport by electron conductivity is pretty large for high temperature plasma, while the perpendicular conductivity is negligible. The expulsion of stored energy to the external region would be enhanced if the conductivity were included in the model. The self-consistent treatment of such heat transport is a subject for future work.

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