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ABSTRACT

This paper discusses major and minor disruptions in tokamaks. A number of models and numerical simulations of disruptions based on resistive MHD are reviewed. A discussion is given of how disruptive current profiles are correlated with the experimentally known operational limits in density and current. It is argued that the q_a=2 limit is connected with stabilization of the m=2/n=1 tearing mode for $2 < q_a \lesssim 2.7$ by resistive walls and mode rotation. Experimental and theoretical observations indicate that major disruptions usually occur in at least two phases, first a "predisruption", or loss of confinement in the region $1 < q < 2$, leaving the $q \approx 1$ region almost unaffected, followed by a final disruption of the central part, interpreted here as a toroidal n = 1 external kink mode.

1. INTRODUCTION

The major disruption is probably the most important phenomenon limiting the range of operation for tokamaks. Disruptions manifest themselves as violent MHD activity on timescales ranging from those of resistive instabilities to the ideal MHD instability timescale. A disruption is termed major if it leads to the termination of the discharge after broadening the current profile, whereas disruptions from which the plasma recovers are termed minor.

Apart from the purely theoretical interest in understanding such dramatic behaviour, there are clear practical motivations why detailed understanding is desirable. First, disruptions set operational limits, hence it is important to understand what can trigger a disruption in order to be able to extend the range of operation and improve on tokamak performance. Secondly, in particular for large tokamaks at high current operation, very large amounts of energy are released in disruptions, leading to significant erosion of limiters and walls and giving rise to large mechanical stresses on coils, etc. Thus, from the viewpoint of experiment (or reactor) lifetime, it is important to minimize the number of hard disruptions, either by avoiding disruptions altogether or by finding some way to "soften" them.

This paper is an attempt to give an at least partial and necessarily somewhat biased review of the present theoretical understanding of disruptions. Section 2 reviews experimental results concerning both the factors leading to disruption and the observed dynamical behaviour. Section 3 gives a basic account of MHD stability for the cylindrical tokamak and Section 4 reviews a number of numerical simulations and theoretical models of disruptions. Recent results on the influence of resistive walls on the q_a-limit and concerning the final phase of major disruptions are discussed.

2. EXPERIMENTAL RESULTS

It is well known that tokamaks disrupt when the density or current exceed certain values. This restricts the operating regime to lie within a triangle of the so-called Hugill diagram, where the Murakami parameter[1] nR/B_T is plotted vs. $1/q_a$.

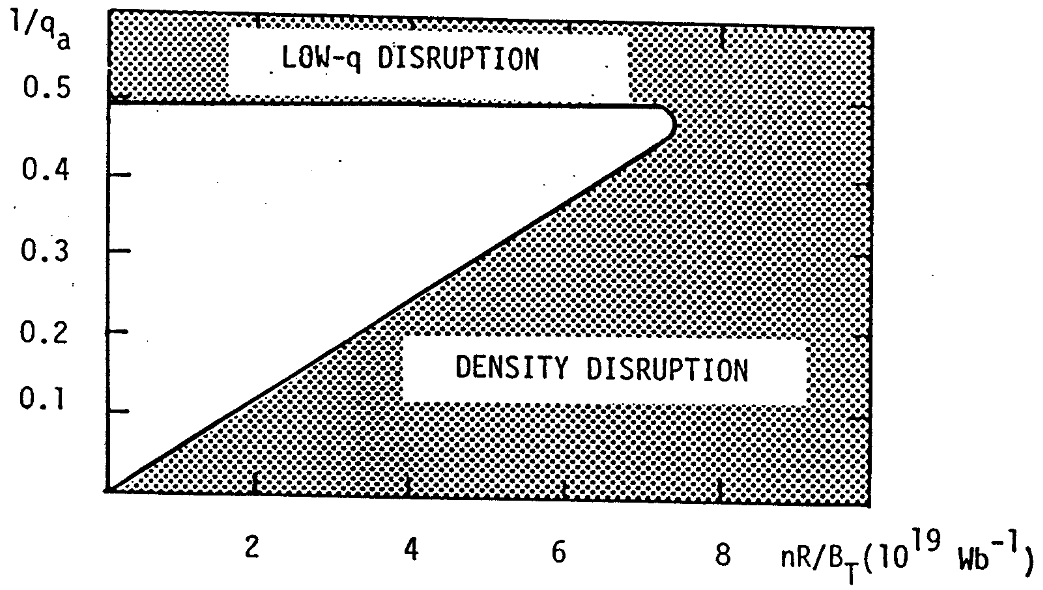

Figure 1: Schematic Hugill Diagram

The current limit is given by $q_a > 2$ and the density limit, which is proportional to the average current density, by $n < CB_T/Rq_a$, where C is a slightly machine-dependent number usually between $1 \cdot 10^{20}$ and $2 \cdot 10^{20}$ Wb^{-1} for ohmically heated plasmas. The Hugill diagram drawn in Fig. 1 is very schematic and, for example, ignores difficulties often encountered in crossing $q_a = 3$, but the basic picture is similar for most tokamaks.[2–5] The density limit can be increased somewhat by the addition of neutral beam injection[6,7] (which increases the amount of heating) and pellet fuelling[7] (which reduces the impurity fraction).

The density limit has been correlated with cooling of the outer region of the plasma, leading to contraction of the current profile so that little or no current flows outside the q = 2 surface.[5] In terms of the current density at q=2, cooling of the outer regions gives rise to profiles that are similar to those arising at low-q operation where $q_a \approx 2$. Ashby and Hughes[5] showed that a limit in nRq_a/B_T arises from the condition of 100% energy loss in a radiating layer at the edge of the plasma (assuming that the impurity content and plasma temperature are kept fixed). Perkins and Hulse,[8] comparing radiation losses with ohmic heating for a flat q-profile, found that the radiative density limit is almost independent of the electron temperature and should depend on the impurity content as

$Z_{eff}/(Z_{eff}-1)$. They also found strongly dependence on the impurity species, light impurities giving a higher density limit than heavy impurities.

The current - or q_a - limit is in practically all cases $q_a > 2$. There is some variation in how closely the $q_a = 2$ limit can be approached; $q_a \approx 2$ is reported from Doublet III,[3] Jet,[9] and TFR[10]; $q_a \approx 2.2$ from FT[2] and JIPP-II.[4] Experiments with $q_a < 2$ have been reported from DIVA[11] with a conducting shell, T-10[12] in the presence of large amounts of light impurities, and from TFR[10] under discharge cleaning conditions.

Concerning the mechanism for disruption, an almost universal feature is that the disruption is triggered by an m=2/n=1 mode becoming unstable.[2,4,9,13-19] In the case of high-density disruptions the reason for instability of the m=2 mode is often identified as radial contraction of the current profile caused by edge cooling.[4,9,13,14,16,18] Many experiments show excitation of an m=3/n=2 mode after the m=2 mode has grown to large amplitude.[2,4,17-19] In TOSCA, even high order modes such as m=5/n=3 have been observed.[17] With respect to the presence of the 2/1 and 3/2 modes, the high-density and low-q disruptions are similar, but the timescales are quite different, the low-q disruptions occuring much faster.[9,16] This difference becomes very pronounced in a large machine like JET, where the density disruptions can have precursors lasting for about 1 second.[9] A typical feature in the early phase of disruptions is the slowing down of the m=2/n=1 mode rotation as the mode starts to grow.[16,20]

While the initial phase of a disruption is similar on different machines, there is less agreement on the final phase. TFR[16] reports strong coupling between the m=2/n=1 and m=1/n=1 modes leading to a large drop in the central electron temperature. In JIPP-II[19] the coupled 2/1 and 3/2 modes were found to lead to only partial disruption, leaving the central region unaffected, while the final termination of the discharge occured via coupled 1/1 and 2/1 modes. It is suggested in Ref. 19 that the ability of the plasma to survive partial disruptions depends on the ability of the control system to maintain the plasma position. Observations of a sequence of events, namely, shrinking of the discharge, "predisruption" by an m=2 instability, "mixing" by m=1 activity, and a final stage with m=2, 3 and 4 modes were presented already in 1976 by Mirnov and Semenov.[14] An interesting observation concerns the angular distribution of tracks in the T-10 liner[21] which gives evidence that the plasma touches the limiter in the form of an m = 1/n = 1 helix.

3. BASIC MHD STABILITY CONSIDERATIONS

The disruption models discussed in this paper are all based on resistive MHD calculations, usually in the large aspect ratio, zero β approximation. In this limit, the only remaining MHD instability is the resistive tearing mode (or, if the electrical conductivity is sufficiently low at the resonant surface, the external kink mode). A common feature of all these models is the dominant role of the m = 2/n = 1 tearing mode in triggering the disruption.[22-31]

Before discussing the different models it is useful to recall some basic stability considerations for resistive modes in the cylindrical tokamak.[32] The tearing mode is known to be stable if and only if[33]

$$\Delta' \equiv \lim_{\delta \to 0^+} \frac{1}{\psi(r_s)} \left[\frac{d\psi}{dr}(r_s+\delta) - \frac{d\psi}{dr}(r_s-\delta)\right] < 0. \qquad (1)$$

Here, ψ denotes the magnetic flux perturbation $\psi = rB_r$ in the external, ideal MHD region, and r_s is the radius of the resonant surface where $q(r_s) = m/n$. In the external region, ψ satisfies the force balance condition $\hat{z} \cdot \nabla \times (\vec{J}\times\vec{B})_1 = 0$, or

$$\frac{1}{r}\frac{d}{dr}\, r\, \frac{d\psi}{dr} - \frac{m^2}{r^2}\psi - \frac{m}{B_\theta(m-nq)}\frac{dj_z}{dr}\psi = 0. \qquad (2)$$

By partial integration of (2), Δ' can be expressed as

$$r_s\Delta'\psi_s^2 = -\int_0^a \left\{\left(\frac{d\psi}{dr}\right)^2 + \left[\frac{m^2}{r^2} + \frac{m}{B_\theta(m-nq)}\frac{dj_z}{dr}\right]\psi^2\right\} rdr$$

$$- m\Lambda\psi_a^2 \;, \qquad (3)$$

where

$$\Lambda = \frac{1 + (a/b)^{2m}}{1 - (a/b)^{2m}} = - \frac{a}{m} \psi_a'/\psi_a \tag{4}$$

and where b is the radius of a conducting wall. Equation (3) shows that the standard current gradient in tokamaks ($dj_z/dr < 0$) is destabilizing when it occurs inside the resonant surface for $q < m/n$, and is stabilizing on the outside. Furthermore, current gradients located where $q \approx m/n$ have a strong effect on the linear stability because of the $m/(m-nq)$ weighting factor. The $-m^2\psi^2/r^2$ term makes high-m modes stable unless a large part of the current gradient is confined to a region just inside the resonant surface such that $0 < m-nq < 1$. A conducting wall is strongly stabilizing if it is located close to the resonant surface.

If we now consider the stability of the m=2/n=1 mode in a tokamak, it is clear that when the current density is low for $q > 2$, the 2/1 mode tends to be unstable because the main part of the current gradient is located inside the resonant surface. If the current density at the $q = 2$ surface is constrained to a low value, the best solution for stability of the 2/1 mode is to move the current gradient as far inward as possible, away from $q = 2$. Thus, it is desirable to minimize q in the centre. However, at least in the cylindrical model, q_0 cannot be less than unity. Then, in order for the 2/1 mode to remain stable, if $j_z = 0$ at $q = 2$ (and $j_z(r)$ is non-negative), the current profile must approach a stepfunction with $j_z = 2B_T/R\mu_0$ = constant in the central region where $q_0 = 1$ and $j_z = 0$ for $q > 1$. The stability of the stepfunction current profile (for which tearing modes become equivalent to external kinks) was analyzed by Shafranov.[34] For $q_0 = 1$ and in the absence of a conducting shell, the Shafranov profile is marginal to the modes with $m = n$ and $m = n+1$ and stable to all other modes. The stabilizing effect of a perfectly conducting wall can be seen from the range of q_0-values for which a current channel of radius r_0 is unstable to the m/n external kink

$$(1/n)\left[m - 1 + (r_0/b)^{2m}\right] < q_0 < m/n \quad . \tag{5}$$

We see that in the absence of a close-fitting wall, q_0 must be close to unity in order for the 2/1 mode to be stable.

It is emphasized that if the current density is finite at $q = 2$, as would be the case if $q_a > 2$ and in the absence of excessive edge cooling, much more relaxed profiles stable to all tearing modes can be constructed.[35] Although the stepfunction current profile in a sense is optimal for stability, it is clearly in conflict with the effects of transport on the equilibrium. Consequently, decreasing the current density at $q = 2$, either by increasing the total plasma current so that $q_a \to 2$ or by cooling the edge plasma, produces a conflict between transport and MHD stability, unless wall stabilization is effective. The influence of wall stabilization will be discussed in some detail in Sec. 4.D.

4. DISRUPTION MODELS

4.A Profiles Unstable to the m=2/n=1 Tearing Mode

In a nonlinear, single helicity study of the m=2/n=1 tearing mode,[22] White, Monticello, and Rosenbluth found that if q_0 is well above unity and the main part of the current gradient occurs inside q=2, the saturated width of the q=2 island is very large. The q-profiles used in Ref. 22 and many subsequent studies of disruptions were parametrized as

$$q(r) = q_0[1 + (r/r_0)^{2\lambda}]^{1/\lambda} \quad , \quad (a/r_0)^{2\lambda} = (q_a/q_0)^{\lambda} - 1. \qquad (6)$$

For q_0=1.37, r_0=0.6a, and λ=4, giving a current profile whose gradient is concentrated just inside q=2, White et al[22] found that the q=2 island encompasses virtually the whole plasma cross-section, the saturated island width being $w \simeq 0.7a$. They also pointed out the sensitive dependence of the island width on the central q-value, with $q_0 = 1.1$, the island width was reduced to $w \simeq 0.4a$.

The conclusions of Ref. 22 are in agreement with those of the linear theory in Sec. 3. One aspect that becomes clear in nonlinear computations is that low shear inside the resonant surface makes possible very large islands nonlinearly.[22,27] This has also been observed in ideal MHD simulations of the so-called vacuum bubbles[39,40] due to external kink modes. In linear theory, shear has two opposing effects on the tearing mode. First, low global shear inside the resonant surface tends to increase the driving energy Δ', because of the m/(m-nq) weighting of the current gradient in Eq.(3). On the other hand, the dynamics of the mode, taking place inside the resistive layer, is sped up by shear, and the linear growth rate[33] scales as $q'^{2/5}\Delta'^{4/5}$. In the nonlinear phase, when the island is wider than the resistive layer (the so-called Rutherford regime[41]), the growth only depends on the resistivity η and a nonlinear Δ',[41,42]

$$\frac{dw}{dt} = 1.66\,\eta\,[\,\Delta'(w) - \alpha w\,] \quad , \qquad (7)$$

where α depends on the resistivity profile inside the island.

Sykes and Wesson,[23] using an initial current profile with q_0=1.5 and low current density at q=2, found that the 2/1 mode is destabilized if the q=2 island makes contact with the limiter. Limiter contact cools the island and hence decreases the current density near the O-point and this enhances the growth of the island. (However, as the island grows mainly inward in cylindrical geometry,[42] q_a must be very near 2 for the island to touch the limiter.) It has been pointed out that other physical effects cooling the interior of the island, such as increased radiation losses,[43] can enhance island growth. Conversely, magnetic islands can be suppressed by the application of radiofrequency waves[44,45] so as to (a) heat the interior of the island to increase the ohmic current at the O-point, (b) heat the plasma outside the resonant surface[46] to move the current gradients outward and reduce Δ', or (c) drive current non-inductively within the island.

4.B Multiple-Helicity Interactions

A significant element in the theoretical understanding of disruptions was found in three-dimensional MHD simulations carried out at Oak Ridge.[28-30] Waddell et al[28] showed that when the 2/1 tearing mode produces a large island it can strongly destabilize the m=3/n=2 mode. The simultaneous presence of resistive modes with different helicities leads to field line stochasticity and destruction of the magnetic surfaces. If the islands created by the 2/1 and 3/2 modes are large enough to overlap radially, large scale stochasticity and loss of confinement results. Refs. 28-31 showed that such overlap also leads to very violent MHD activity.

Waddell et al[28] presented a simulation starting from an equilibrium (6) with q_0=1.38, q_a=4, and λ=4 (the same as used by White et al[22]). For this equilibrium, the 2/1 and 3/2 modes are both highly unstable and, if each is evolved in a single-helicity computation, their islands are large enough to overlap. When overlap occurs in the multiple-helicity calculation, the consequences are dramatic. The 3/2 mode is strongly destabilized and its growth rate is increased by a factor of almost 4, at a Lundquist number $S=\tau_r/\tau_A=1.3\cdot 10^5$. The destabilization is dependent on the coupling to nonlinearly driven modes such as m/n=5/3 and 1/1.

The simulations of Refs. 28-31 showed that nonlinear interaction between modes with different helicities could lead to rapid instability, on timescales comparable to those observed in major disruptions, and considerably faster than single-helicity calculations would predict. Furthermore, the 3-D nonlinear interaction of resistive modes is a mechanism that results in stochastic magnetic fields and loss of confinement over a large region of the plasma.

The violently unstable multiple-helicity computations[28-31] stimulated a large interest in MHD turbulence and several different mode-coupling theories were presented to explain the numerical simulations. Carreras, Rosenbluth, and Hicks,[47] in a third-order random-phase calculation found non-linear growth on the ideal-MHD timescale, with a growthrate independent of resistivity but proportional to the fluctuation level. Tetrault[48] derived an anomalous resistivity due to the turbulent fields, whereas Biskamp and Welter[49] proposed a model in which small-scale turbulence acts on the large-scale fields as a negative resistivity. Diamond and coworkers[50] derived renormalized equations for the large-scale fields, in which the short wavelength turbulence gives rise to an anomalous resistivity by turbulent fluid convection and an anomalous vorticity damping by turbulent magnetic stresses.

In Refs. 51 and 52, a simpler question was asked, namely, what is the ideal and resistive stability of equilibria with large q=2 islands. These studies point at the destabilizing influence on the 3/2 mode when the q=2 island grows to large size and pushes the main part of the current gradient inside the q=3/2 surface. Kleva, Drake, and Bondeson showed that a significant fraction of the increase in the growthrate of the resistive 3/2 mode could be accounted for by the steepening of the current gradient by the island together with the relative displacement of the current gradient and the q = 3/2 surface, so that the current gradient became highly localized just inside q = 3/2. Reference 51 showed that even ideal instability can occur when the q=3/2 surface becomes sufficiently distorted by the inner q=2 separatrix. In contrast with the random phase prediction,[47] however, this ideal instability has a large threshold in island width.

Recently, the accuracy of the Oak Ridge computations[28–30] has been questioned, both with regard to the time-stepping, which is explicit and unstable in the absence of dissipation,[53] and the spatial discretization, where aliasing may cause nonlinear, numerical, but non-physical, instabilities if the radial resolution is insufficient.[54] While several independent simulations[31,37,55] have verified the destabilization of the 3/2 and higher order modes and also find generation of short-wavelength turbulence, the final catastrophic growth of the 3/2 mode at a rate independent of the resistivity reported in Ref. 50 is not observed in other simulations.[31,37] A clear indication of numerical difficulty in Ref. 30 is the temperature profile, which, in a minute fraction of the ohmic heating time develops sharp local extrema, although the temperature should mainly be diffusing on this timescale. Biskamp and Welter[31] stressed that the required resolution increases sharply with decreasing resistivity and viscosity.

4.C Transport-MHD Simulations

It has been emphasized throughout this survey that the current profile plays a crucial role in triggering disruptions. To a large extent, the temperature and current profiles are determined by microscopic transport processes in the plasma. However, when large scale MHD activity occurs, such as sawteeth or q = 2 islands, the magnetic field structure is affected and the transport changes. Thus, there is an intimate coupling of transport and instability, or, to cite Turner and Wesson[25] "it is clear that neither transport nor the behaviour of instabilities can be treated in isolation." This is particularly true if we want to understand the conditions leading to disruption.

Turner and Wesson[25] formulated a one-dimensional model where the tearing modes were treated according to Rutherford theory,[41] obeying Eq. (7) with $\alpha=0$. The temperature was evolved selfconsistently including thermal conduction, ohmic heating and radiation losses. When a magnetic island developed, the temperature was flattened across the island and the effect of the sawteeth was modelled by a large increase in the thermal conductivity in the region where $q<1$. The current profile was evolved according to a one-dimensional diffusion equation with Spitzer resistivity, $\eta \propto T^{-3/2}$. Despite its simplicity, this model includes the essential element of coupling transport processes and MHD instability.

When q_a was reduced from 8 to 4 and q_0 went below 2, the $m=3,2/n=1$ tearing modes gradually became unstable, while still $q_0 > 1$. (The current profile is generally more peaked in a tokamak. The peaking depends on the thermal conductivity profile, which was taken as a constant in Ref. 25. Neoclassical corrections[56] to the resistivity also increase the peaking.) With the transport model of Turner and Wesson, a $q=1$ surface appeares in the plasma when $q_a \approx 3.4$. Further increase of the plasma current makes the $q=2$ surface move outward and forces the main part of the current gradient to occur inside $q=2$. As the maximum current density is constrained by $q > 1$, lowering q_a also forces the current gradient outward, closer to the $q=2$ surface, where it destabilizes the 2/1 mode producing an increasingly large magnetic island. Turner and Wesson concluded that as q_a drops to about 2.8, the combination of a growing $q=1$ region in the centre and a large island at $q=2$ finally removes the isolating layer between $q=1$ and $q=2$ and leads to disruption. They could also produce repeated soft disruptions for q_a above 3 if the current density at $q=2$ was reduced by impurity radiation. An important feature of all simulations described in this subsection is that they used no[25,36] or a very distant[37] conducting wall so that there was no wall stabilization.

Hopcraft and Turner[36] refined the 1-D model, including the 3/2 mode and its corresponding island. By ramping the plasma current so that q_a went from above 3 to about 2.5, they found a series of mini-disruptions and finally catastrophic interaction where the $q=2$ and 3/2 islands spanned the whole region from the plasma edge to the

$q \simeq 1$ region. According to the assumed transport model, with flattening over these regions, confinement was then completely lost.

A major contribution of Refs. 25 and 36 was to relate the disruption clearly to the experimental operating conditions. The interaction of different tearing modes in this model does not include mode-coupling, but only occurs through the quasi-linear modification of the current profile.

The simulations of the present author[37] include transport and MHD phenomena within the framework of the three-dimensional reduced MHD equations.[28–31] The thermal conductivity is highly anisotropic, with a large conductivity along the fieldlines. Although the disruptions simulated with the 3-D model are more distinct than those of the 1-D model, the 3-D simulations give a similar condition for the occurrence of disruptions: $q_a < 2.6$. As this limit is approached, first weak oscillations occur in the equilibrium profiles, followed by soft, and later harder disruptive events.

The initial motivation for these simulations was to see whether in a self-consistent 3-D simulation, with slowly varying external parameters, distinct disruptions would occur, or whether the plasma would respond in a "soft" way, for example, by increasing the level of turbulence. It is clear that the carefully selected, highly disruptive initial conditions used in the previous 3-D studies[28–31] could hardly arise in practice, simply because they are too unstable.

Disruptive behaviour indeed occurred[37] as a result of multiple-helicity interactions of resistive modes and their effects on the temperature profile. However, as these disruptions start from effective "initial conditions" that are much less unstable than those of the nonselfconsistent simulations,[22,28–31] for example by having $q_0 \simeq 1$, their subsequent evolution differs considerably. The selfconsistent disruptions are initiated by the 2/1 mode becoming unstable at a time when the 3/2 mode is stable, and the 3/2 mode becomes unstable primarily as a result of the profile modifications due to the growing 2/1 mode. The destabilization of the 3/2 and subsequent higher-m modes takes the character of a "shock" moving inward from the q=2 surface tearing successive flux surfaces. In front of the shock, the flux surfaces are intact but behind it the magnetic fields are stochastic and the plasma is turbulent. In the case of a disruption, the shock proceeds past the q=3/2 region, destabilizing modes with successively lower q=m/n (and therefore higher m) to remain in resonance with the background magnetic field, as the front propagates toward the q=1 region. Figure 2 shows the equitemperature lines in a poloidal cross-section at an early (2a) and late (2b) phase of such a disruption.

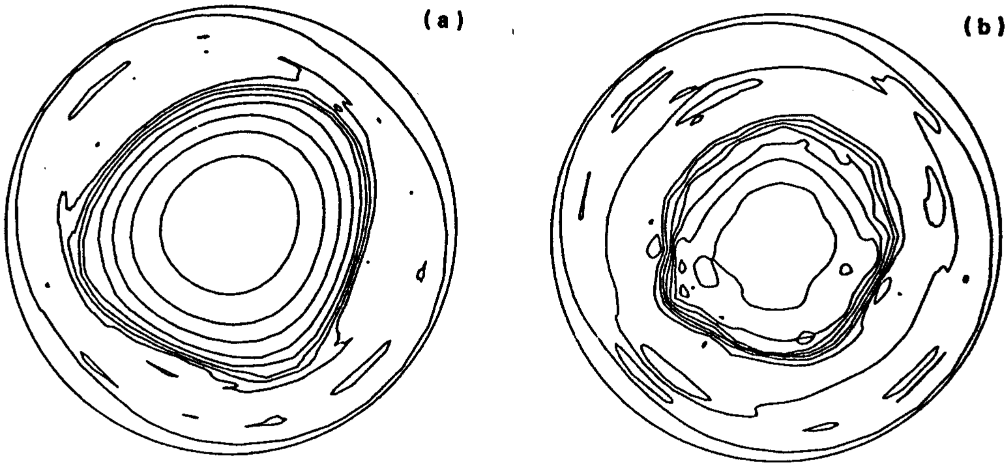

Figure 2. Equitemperature lines at (a) an early and (b) a late phase of disruption in the region 1 < q < 2.

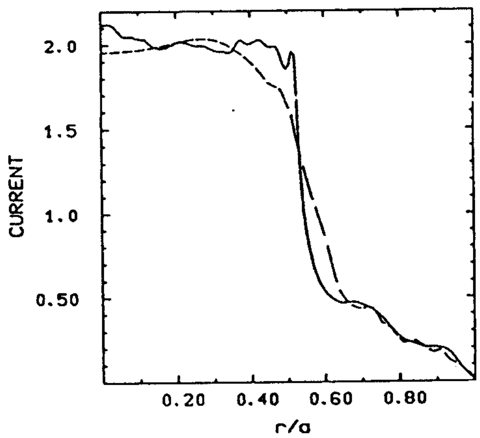

Figure 3. Current profiles at the end phase of a disruption in the region 1<q<2, (- -) before and (—) after giant sawtooth.

An unexpected result found in Ref. 37 was that the disruption fronts always stopped when the reached the q=1 surface. At this point, the current profile is highly contracted, in fact, it is almost a stepfunction with q_0 slighly below unity, an example of which is shown in Fig. 3. As discussed in Sec. 3, such a profile is stable to all external modes with m/n > 1 (although near marginal to the m=n+1 modes). Evidently, the shock stops because it runs out of driving energy and the q ≈ 1 region is not penetrated. The contraction ends with an internal sawtooth-like relaxation in which the central temperature drops sharply because the plasma outside $q \approx 1$ is cold. Following this internal relaxation, the calculation recovers and the profile broadens again, which eventually leads to another contraction.

Despite several similarities with the experimental observations, the results of Ref. 37 obviously disagree with the experimental findings in that there is no final termination of the discharge with a pronounced broadening of the current profile. In particular, the drop in the internal inductance found in Ref. 37 was small, at most a few percent, and it appears highly probable that the contractions in the external regions simulated in Ref. 37 correspond to what is often referred to as "predisruption".[14,19]

It will be argued here that toroidal effects are important for the final stage of a major disruption, and that the instability may be fairly ideal-like because the q=2 surface is cold and highly resistive. The n=1 stability of contracted profiles has been studied[58] within ideal MHD using the ERATO stability code.[59] For the sake of

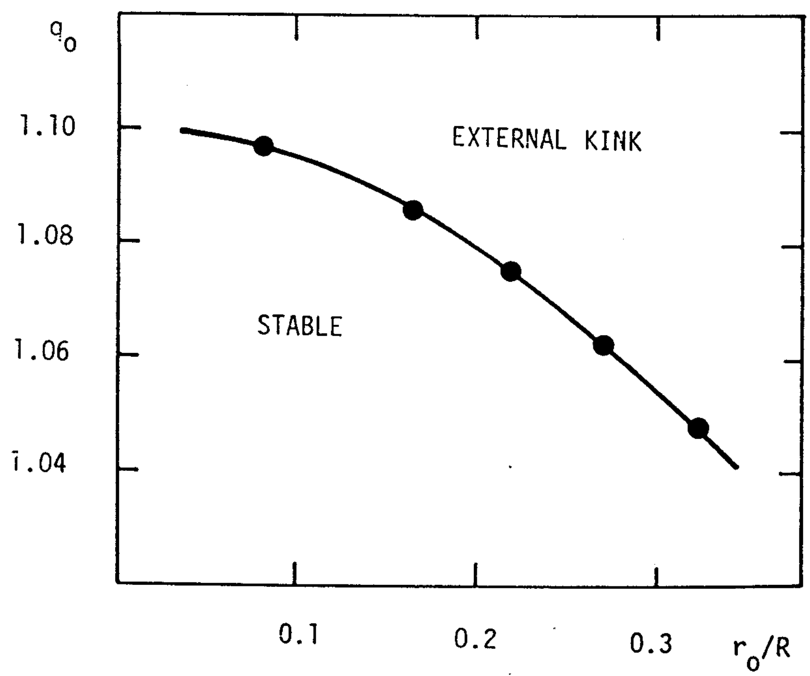

Figure 4. Stability diagram, q_0 vs. inverse aspect ratio for Shafranov equilibrium with a conducting wall at $b/r_0 = 1.78$. Solid curve: $q_0 = 1.1 - 0.505\,(r_0/R)^2$.

simplicity, the equilibrium was chosen to have a slightly rounded-off stepfunction for the current profile (j_ϕ = TT'/R, and $\langle j_\phi \rangle$ ≈ stepfunction), zero β, circular cross-section and a vacuum region for $q > q_s \approx 1.6$. With a conducting wall at r=b, a current channel of radius r_0 is stable to the m=2/n=1 mode in a cylinder if $1 < q_0 < 1+(r_0/b)^4$. In the toroidal calculation,[58] $b/r_0 = 1.78$ so that the corresponding cylindrical case is unstable if $q_0 > 1.10$. Figure 4 shows the stability boundary in q_0 as a function of inverse aspect ratio. Note that the Shafranov equilibrium is considerably more unstable in a low aspect ratio torus than in a cylinder. For $r_0/R=0.324$ the toroidal equilibrium is unstable for q_0 down to 1.048 to an n=1 mode that shows strong coupling of m=1, 2, and 3. It is proposed on the basis of these linear computations that toroidal external kinks (possibly further destabilized by finite pressure) are responsible for the final phase of major disruptions and that it is this instability that produces the large drop in internal inductance and the negative voltage spike. To make a more conclusive statement, nonlinear, fully toroidal simulation would be necessary.

The picture presented here of the major disruption thus involves a sequence of events. (a) Low current density at q=2 destabilizes the m=2/n=1 mode which, by nonlinear multiple helicity interactions, destroys the confinement in the external, q>1 region. (b) The contraction of the current channel makes the central q fall below unity after which an internal sawtooth-like instability strongly reduces the central temperature and makes $q_0 > 1$. (c) The contracted current profile with $q_0 > 1$ is unstable to the toroidal n=1 external kink which terminates the discharge. This scenario is in accord with recent data from JET,[60,61] and is also very similar to the early observations of Mirnov and Semenov[14] as well as to those from JIPP-II[19] and PLT.[57]

4.D Mode Locking, Stabilization by Resistive Walls and the q=2 Limit

It is usually observed experimentally that as the 2/1 mode begins to grow, its rotation frequency decreases[16,20,62] and the instant at which it locks to the wall signals the beginning of the disruption.

It is clear that when the rotation frequency of the mode is much larger that the inverse time-constant for the vessel, the wall looks conducting, and the mode may be wall stabilized. For instance, in JET, the 2/1 rotation frequency is typically between 1 and 10 kHz depending on operating conditions, and the time-constant of the wall is about 4 ms, thus the wall should appear nearly ideally conducting. Rutherford[45] and Nave and Wesson[63] derived equations for the slowing down of the plasma rotation due to eddy currents induced in the resistive wall. These authors point out the nonlinear, self-reinforcing effect of the wall locking, namely, once the mode starts to grow, the resulting torque slows down the plasma rotation, as a result the wall stabilization is reduced and the mode grows even faster, etc.

Reference 64 points at another aspect, namely that mode rotation may resolve the discrepancy between the disruption limit $2.6 \lesssim q_a \lesssim 2.8$ found in the cylindrical free-boundary computations[25,36,37] and the experimental result, $2.0 \lesssim q_a \lesssim 2.2$. As long as the mode rotates fast enough for the resistive wall to appear conducting, the 2/1 tearing mode becomes increasingly stable as q_a is lowered toward 2 (in fact, $\Delta' \to -\infty$ as $q_a \to 2$ if the wall is close-fitting). We conclude that tokamaks can be operated with $2 < q_a \lesssim 2.7$ because of stabilization by resistive walls and mode rotation. As long as the q=2 surface is inside the conducting plasma, the rotation frequency of the mode is determined by the plasma flow (and diamagnetic drifts, etc.) locally at the resonant surface. The important point is that if the q=2 surface moves out of the conducting plasma, the 2/1 mode is free to rotate (with moderate frequency) with respect to the plasma, as the mode is frozen into the plasma only at its resonant surface. Thus, in this case, the mode can lock to the wall and may be thought of as an external kink, slowed down by the resistive wall. A simple calculation shows that the growth time is $\tau = s\mu_0/\Delta'$, where $s = \sigma_w\delta_w$ is the surface conductance of the wall and Δ' is evaluated at the wall with $\psi \propto r^{-m}$ on the outside.

We are led to the conclusion that the q-limit is reached when the q=2 surface approaches the edge of the plasma where the conductivity is sufficiently low for a wall-locked mode to slip with respect to the plasma. This condition agrees very well with the experimental q_a-limit $2 \lesssim q_a \lesssim 2.2$, where the variations between different machines may be due to differences in the temperature profile near the plasma edge. The wall locking with $q_a \approx 2$ has been simulated[64] with a modified version of the code in Ref. 37. The mode rotation is due to mass flow of the plasma, and because of the reduced MHD ordering, only poloidal flows can be handled (but we expect similar behaviour in the case of toroidal rotation). When $q_a \approx 2$, the resistive layer formed at the imperfectly conducting wall coincides with the usual resistive layer at the resonant surface q=2. Under these circumstances, a variety of behaviours can be produced for the mode-locking by varying the flow speed and the conductance of the wall and the edge plasma. A consistent feature is that the locking is fast compared with the q > 2 cases analyzed by Nave and Wesson.[63] It is clear that the $q_a > 2$ scenario[63] is relevant to high-density disruption and that of Ref. 64 to low-q disruption. The difference in time-constant between the two cases is in excellent agreement with the experimental observations regarding precursor lengths for the high-density and low-q disruptions.[9,16,62] We note that for the $q_a \approx 2$ case, limiter contact and cooling in the center of the island, as proposed by Sykes and Wesson,[23] will be effective in enhancing the island growth. Figure 5 gives the time evolution of the mode amplitude and frequencies together with $\partial B_\theta/\partial t$ for a case where q_a goes through 2 and the conducting plasma extends to the resistive wall.[64]

The free boundary limit of $q_a \approx 2.6$ found in the cylindrical model would also be lowered by the stabilizing effect on the tearing mode of toroidicity and finite pressure.[65-68] However, this stabilization depends on the combined effects of high conductivity and finite β, therefore, it is at its weakest at the edge of the plasma and may only have a slight influence on the free-boundary q_a-limit. Thus, it seems justified to conclude that stabilization due to mode rotation and resistive walls plays a role in low-q operation of tokamaks and that the $q_a \approx 2$ disruption limit is due to loss of wall stabilization when the q=2 surface moves out of the conducting plasma.

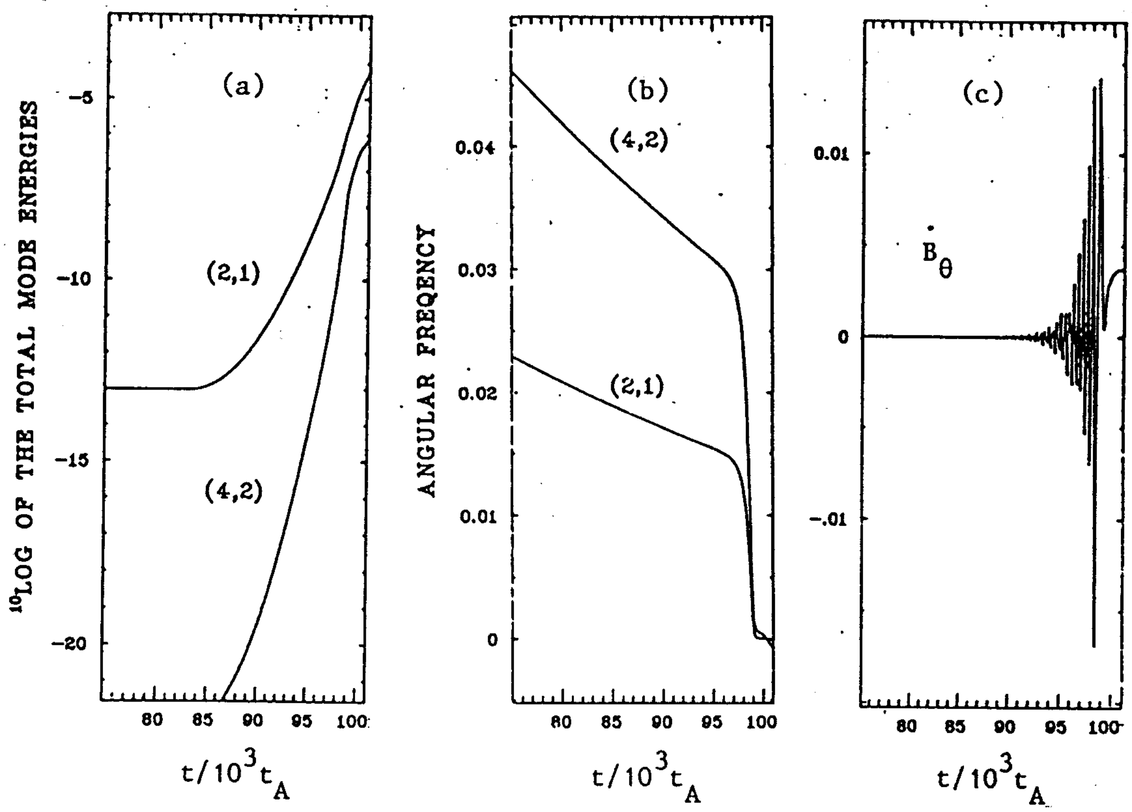

Figure 5. Time histories of (a) mode amplitude (b) mode frequency and (c) $\partial B_\theta/\partial t$ for mode-locking when $q_a \simeq 2$. ($dq_a/dt = -2\cdot 10^{-6}/t_A$ and $q_a=2$ at t/t_A=72 000.)

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