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Partial reconnection in the sawtooth collapse

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Abstract

The effect of the electron pressure gradient in Ohm's law on the nonlinear development of the internal kink mode is investigated. While pressure fluctuations have a destabilizing effect, the average pressure gradient, giving rise to diamagnetic flows, is stabilizing. If the latter is strong enough it leads to saturation at finite island size. The results are compared with recent measurements of the sawtooth collapse on the Textor tokamak and alternative concepts of the collapse are discussed.

Keywords : Internal kink mode, sawtooth collapse, partial reconnection

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I. Introduction

The sawtooth oscillation is one of the most common phenomena in tokamak plasmas, observed in practically all devices under rather general conditions. In spite of its clear and seemingly simple signature the interpretation of the sawtooth oscillation, in particular the rapid collapse, has been a major challenge in plasma theory, and up to now no generally valid model has been presented. There seems to be little doubt that the collapse is in some way associated with the internal kink mode, but the details how this mode causes the observed flattening of the temperature and density profiles is still unclear. A stringent theoretical constraint arises from the observation of the central $q-$value remaining clearly below unity. The early observations by Soltwisch on the Textor device [1] have been refined and similar results have been obtained in various other machines using different observational techniques, all confirming the original Textor claim. This behavior essentially invalidates Kadomtsev's full reconnection model [2] as well as Wesson's quasi-interchange model [3], which seemed to be the natural extension of the former to the high temperature regime eliminating the slow time-scale problem. The full reconnection model was hence replaced by the idea of only partial reconnection, as earlier observations, e.g. on TFR [4], had already suggested. Here the reconnection process is assumed to be halted at some finite island size, while the electron energy diffuses out of the plasma core owing to stochastization of field lines caused mainly by toroidal coupling of magnetic perturbations on different magnetic surfaces. This idea has recently been studied in detail by Lichtenberg et al. [5], who also give some arguments how to explain the rapid density profile flattening. Nevertheless the actual dynamics of the collapse remains unexplained, in particular, why the reconnection process is halted. In fact numerical simulations, including toroidal geometry [5], seem to lead always to full reconnection once the dynamical process has started, with no indication of a slowing down until reconnection is complete. Generalizing the MHD model by including collisionless effects in Ohm's law, in particular electron inertia [6] and electron pressure [7], the situation tends to be even worse, since the reconnection process accelerates nonlinearly.

This fast quasi-collisionless process has recently been shown to allow a completely different solution of the problem [8]. Since a finite fraction of the magnetic energy released in the reconnection process is converted in to plasma flow energy, instead of being completely dissipated as in the resistive case, plasma inertia drives the system back to a $q_0 < 1$ state. Hence the collapse consists of a two-step reconnection process, a Kadomtsev-type full reconnection phase followed by a secondary process, where the helical flux pushed out of the core in the first phase, is (at least partly) reinjected. The process is found to be particularly efficient [9], if the initial state has reduced magnetic shear in the vicinity of the $q = 1$ surface, corresponding to a shoulder in the current profile as observed experimentally [10]. Since this process takes place on a fast time scale ($\sim 100\mu s$), the diagnostics used up to now are too slow to distinguish between the two-stage process and genuine partial reconnection.

Very recently, however, the Textor group [11] has been able to obtain information about the behavior of the magnetic field during the collapse using one fast Faraday rotation channel across the center in combination with detailed measurements of the density collapse. The still preliminary results are quite surprising. The magnetic signal, which can be related to the shift of the magnetic axis, increases with neutral beam power. However the shift appears to be rather small, in the case of only Ohmic heating less than 10% of the $q = 1$ radius, in the most strongly neutral beam heated case less than 50%. If these observations and their interpretation are valid, this rules out any collapse model involving a full reconnection phase. Hence we have thus to look for some mechanism, which can prevent full reconnection leading to finite amplitude saturation of the kink mode.

In a previous paper [12] I presented a nonlinear theory of the internal kink mode taking into account diamagnetic effects. It has been found that if the latter are sufficiently strong, the mode is stabilized nonlinearly at small island size. Recently this theory has been extended to include toroidal effects [13], in order to explain the absence of sawtooth oscillations in supershot TFTR discharges though q_0 is measured to be below unity. It is in fact found that for plasma parameters corresponding to those in the TFTR discharges, nonlinear stabilization occurs at small amplitudes, qualitatively confirming the results of Ref. [12].

In this paper interest is in finite amplitude saturation connected with partial reconnection in the sawtooth collapse. Here the island size should not be too small in order to allow field line stochastization and fast transport. In this amplitude range the approximations made in both analytical theory and simulations in Ref. [13] are not valid, and we have to resort to a fully nonlinear numerical treatment. The model equations are discussed in section II. In section III I present the results of numerical computations. It is found that the electron pressure term in Ohm's law may have both a destabilizing effect as discussed in Ref. [7] or a stabilizing one leading to finite island size saturation depending on the density scale length r_n. The strong ion flows generated along the separatrix tend to become Kelvin-Helmholtz unstable leading to turbulent ion viscosity which acts as an additional stabilizing factor. Section IV gives the conclusions.

II. Model equations

The linear stability of the internal kink mode depends on many different effects, toroidicity, pressure, dissipative processes, diamagnetic, finite Larmor radius and possibly kinetic effects due to the presence of high energy particles. Whether in the sawtooth rise phase the kink mode is linearly stable or nonlinearly stabilized at very low amplitude, may be a semantic question. In any case, particularly for precursorless sawteeth, the sudden onset of the collapse constitutes a major theoretical challenge, which is outside the scope of the present study. Instead I am concentrating on the strongly nonlinear phase of the collapse, which depends primarily on the nonlinear reconnection dynamics. Hence the dissipative or collisionless terms in Ohm's law are most important. One may therefore in a first approach neglect toroidicity restricting geometry to the lowest order tokamak approximation. The equations used are a simplified version of the four field model of Hazeltine et al.[14], neglecting parallel ion flow and ion pressure. Hence the dynamic variables are the helical flux function ψ, where $\vec{\hat{z}} \times \nabla\psi = \vec{B}$ is the so-called helical magnetic field, the stream function φ, $\vec{\hat{z}} \times \nabla\varphi = \vec{v}$ being the poloidal ion flow, and the electron pressure $p_e = nT_e$. Since the latter enters the equations only in the form $\nabla_{\|}p_e$, one can use the high parallel electron heat conductivity to assume T_e = const. along field lines $\nabla_{\|}p_e = T_e\nabla_{\|}n$, and for simplicity T_e = const. throughout :

$$\partial_t(\psi - d_e^2 j) + \vec{v}\cdot\nabla(\psi - d_e^2 j) = -\mu_e\nabla^2 j - \alpha\beta_p\vec{B}\cdot\nabla n \quad , \tag{1}$$

$$\partial_t\omega + \vec{v}\cdot\nabla\omega - \vec{B}\cdot\nabla j = \mu\nabla^2\omega \quad , \tag{2}$$

$$\partial_t n + \vec{v}\cdot\nabla n - \alpha\vec{B}\cdot\nabla j = D\nabla^2 n \quad . \tag{3}$$

Here $j = \nabla^2\psi + \text{const.}$ is the toroidal current density and $\omega = \nabla^2\varphi$ the vorticity. The equations are written in nondimensional form using as units the plasma radius a, the poloidal Alfvén velocity $v_A = B_0/\sqrt{4\pi n_0 m_i}$, B_0 = typical poloidal field, and n_0 a typical density; $\alpha = (c/\omega_{pi}a)(B_0/B_t) \simeq c/\omega_{pi}R$, R = major radius, since $B_t a/B_0 R = q = 1$; $d_e = c/\omega_{pe}a$; $\beta_p = 4\pi n_0 T_e/B_0^2$ = poloidal β. In equation (1) electron viscosity is chosen as the dominant dissipative process, since it controls the current density gradient, which in the absence of dissipation tends to become infinite (not the current density amplitude) due to the collisionless reconnection processes. Of the latter electron inertia and electron pressure are included in eq.(1). We have neglected the Hall term, which corresponds to a decoupling of the electrons from the ions inside a certain layer width. To include the Hall term the ion velocity in the convective term in eq.(1) has to be replaced by the electron velocity and an additional equation is needed to describe the fluctuation of the axial magnetic field. The Hall term is found to be the dominant reconnection process in the case of weak axial field $B_t \lesssim B_0$ [15], where it leads to an ion layer width of order c/ω_{pi}. In the low$-\beta$ case relevant to the sawtooth phenomenon this layer width has to be replaced by $(c/\omega_{pi})(B_0/B_t)$, i.e. α defined above, which for $\beta_p = O(1)$ is of the same order as the effective ion Larmor radius ρ_s, connected with the electron pressure contribution [16]. Since including the Hall term in the low$-\beta$ case would significantly increase the complexity of the equations without the expectation of a qualitative change of the results, the Hall term is neglected in this study.

III. Discussion of the equations and numerical results

Including the electron pressure gives rise to two distinctly different effects, which can already be seen by linearizing the equations, in particular the pressure term,

$$\widetilde{\vec{B} \cdot \nabla n} = \vec{B}_0 \cdot \nabla \tilde{n} + \tilde{B}_r n_0' \quad . \tag{4}$$

The first term in (4) leads to the kinetic Alfvén wave,

$$\omega^2 = k_\parallel^2 v_A^2 (1 + k^2 \rho_s^2) \quad , \tag{5}$$

neglecting the electron inertia effect since ρ_s is usually larger than the electron skin depth c/ω_{pe}. Neglecting the mean density gradient $n_0' = 0$, equations (2) and (3) become essentially identical, hence $n = \alpha\omega$, and $\alpha\beta_p \vec{B} \cdot \nabla n =$

$\rho_s^2 \vec{B} \cdot \nabla \omega$ in eq.(1). This model has recently been studied [17] for the case of flux bundle coalescence. In that paper the pressure term has been found to dramatically increase the reconnection rate E for small resistivity compared with the purely resistive case, E being essentially independent of ρ_s and η (though the asymptotic behavior for ρ_s, $\eta \to 0$ is difficult to assess.) This effect is the same as the accelerated reconnection process in the nonlinear kink mode observed by Aydemir [7], where the growth rate increases nonlinearly from the linear value $\gamma_0 \sim \rho_s$ (more accurately $\rho_s^{2/3} d_e^{1/3}$) to a much larger value. For not too large amplitude this behavior is consistent with a similarity solution $E \propto (t_0 - t)^{-\lambda}$, $\lambda = O(1)$, which depends on the initial amplitude (through the constant t_0), but not on the smallness parameter ρ_s.

If n_0' is large, such that $r_n = n_0/n_0' \lesssim a$, wave propagation is dominated by the drift frequency $\omega_* \sim c_s/r_n$, which is tokamak plasmas is comparable to the growth rate γ, $\omega_* \sim \gamma_0 > \gamma$. References [12], [13] have focused attention on this range. The linear eigenfuncions are no longer symmetric, since the terms $\propto n_0'$ destroy the parity of the mode, which is now dominated by the differential diamagnetic rotation. For sufficiently large ω_*/γ_0, the growth rate becomes small, but the system remains linearly unstable. To decide, whether the kink mode is stabilized at some finite amplitude or island size, or full reconnection takes place, equations (1) $-$ (3) are solved numerically using a split finite difference $-$ spectral method as in the previous studies [8], [9] with up to 500 radial grid points and 96 Fourier modes. The parameters are chosen in a rather ad hoc way, since considering the approximate character of our model equations there is no advantage of using so-called realistic values, in particular since the effective values of μ_e and D are not well known. In most runs we have used $d_e = 10^{-2}$, $\alpha = 5 \times 10^{-2}$, $\mu_e = 10^{-9}$, $\mu = D = 10^{-6}$; β_p and r_n are varied in the range $0 \leq \beta_p \leq 4$, $0.7 \leq r_n \lesssim 2$, and $r_n = 10$ corresponding essentailly to $r_n \to \infty$.

If r_n is sufficiently large, $r_n \gtrsim 2$, the pressure term is strongly destabilizing. Figure 1a gives the growth rates defined by the evolution of the kinetic energy $\gamma = \partial_t E^K / 2E^K$ with $E^K = \frac{1}{2}\int v^2 dV$, for $r_n = 10$ (essentially infinity) and different values of the parameter β_p. While the linear growth rate γ_0 increases moderately with β_p, the most conspicuous feature is the accelerated growth in the nonlinear phase, as observed by Aydemir [7]. I have therefore plotted γ/γ_0 as function of $\gamma_0 t$, which shows that the nonlinear enhancement of the growth rate increases with β_p. If on the other hand r_n is small, $r_n < 1$, where the diamagnetic effect is dominating, increasing β_p leads to a reduction of the nonlinear growth rate as shown in Fig.1b, which for sufficiently high β_p can lead to finite amplitude saturation, $\beta_p \gtrsim 2$ in this figure. Hence there is a line $\beta_p(r_n)$ separating states leading to full reconnection from those leading to finite island saturation, as shown qualitatively in Fig.2. This figure illustrates the behavior of the average nonlinear growth rate γ in the β_p, r_n^{-1} plane. γ is positive in the lower left part, implying full reconnection. It increases toward a broad maximum in the hedged region, from where it drops rather abruptly. The region above the full line corresponds to saturated island states, i.e. partial reconnection. Figure 3 illustrates a saturating system $\beta_p = 2$, $r_n = 0.7$, the corresponding γ being given in Fig.1b. While at $\gamma_0 t = 8.78$ there is still a strong kink mode-like plasma flow toward the $X-$point (somewhat off-set), the flow at $\gamma_0 t = 9.34$ is weak and turned by 90° away from the $X-$point. The island does not grow further, reconnection has stopped.

From the different runs performed we can obtain a semi-quantitative condition for a saturated state to form :

$$\omega_* > \gamma_{\max} \quad , \tag{6}$$

where ω_* is the linear mode frequency and $\gamma_{\max}$, the maximum (nonlinear) growth rate. If $\omega_* < \gamma_{\max}$, full reconnection occurs, even if ω_* exceeds the linear growth rate. This generalizes previous criteria [11], [12], which relate ω_* to γ_0. Practically speaking (6) is not a very useful condition to predict saturation, since the nonlinear enhancement of γ is only known a posteriori.

A special feature is the development of strong turbulence in the island, which is caused by the Kelvin-Helmholtz instability of the strong, collimated flow out of the reconnection region. This turbulence provides an efficient damping of the plasma flow. Kelvin-Helmholtz instability occurs for system parameters slightly above the critical line in Fig.2. For fully reconnecting systems the outflow from the reconnection region is not collimated into a narrow channel, since the outflow cone is wide corresponding to fast reconnection, hence Kelvin-Helmholtz instability does not arise. For systems saturating at small island size, the outflow is collimated, but the velocity amplitude is too small for Kelvin-Helmholtz instability.

IV. Conclusions

We have seen that including the parallel electron pressure gradient in Ohm's law gives rise to a variety of new phenomena in the nonlinear development of the internal kink mode. The pressure effect is characterized by two parameters, β_p, related to the poloidal beta, and r_n, the width of the average pressure profile or density profile assuming $T_e = \text{const}$. For large r_n the electron pressure term is strongly destabilizing leading to the enhanced nonlinear growth observed by Aydemir. For small r_n the diamagnetic effect $\beta_p n_0' \propto \beta_p / r_n$ dominates, which has a stabilizing effect and for sufficiently large β_p leads to island saturation. Close to the critical conditions saturation is assisted by strong turbulence arising in the island due to Kelvin-Helmholtz instability of the strong collimated flow originating from the reconnection region. Relating these results to the sawtooth collapse, we have found that diamagnetic effects allow partial reconnection to occur. Conditions for partial reconnection depend on r_n and β_p in a nontrivial way, since an increase in β_p has a destabilizing effect for larger r_n and a stabilizing effect for smaller r_n. The recent observation on the Textor device, seem at first sight, to contradict the present results, since they indicate an increase of the final island size with increasing β_p, while the results given above predict that within in the partial reconnection range, increasing β_p leads to smaller island width. One possible explanation is that in addition to the increase of β_p there is an increasingly strong local flattening of n_0' near the resonant surface (due for instance to subordinate sawtooth activity), which according to Fig.2 is a strongly destabilizing effect. The more probable explanation, however, is that the increase of the ideal linear growth rate, considered in Ref.[13], has to be taken into account, which is outside the scope of the present model.

In addition a general problem, common to all partial reconnection concepts of the sawtooth collapse, has to be solved: How are temperature and density released in the short crash time of $100 - 200 \mu s$. While in the case of a strong magnetic perturbation, corresponding to island size or axis shift of a major fraction of the resonant radius, toroidal effects lead to complete stochastization of the core region and hence fast electron transport, no such explanation is possible in the Ohmic case, where the shift is so small that only a small region becomes stochastic corresponding at most to a subordinate collapse. In addition the density should respond on longer time scales, and no such difference between temperature and density collapse has been observed. Excitation of small-scale turbulence is a possible way out of the dilemma. A model based on small-scale interchange modes (Mercier modes) has recently been suggested [18], involving a rather sophisticated concept to overcome the problem of inherently slow transport by small-scale turbulence. A different challenging approach consists in allowing strong deviation from helical geometry. As recently proposed by Nagayama [19], reconnection occurs only in a small angular section in the toroidal direction, where the $X-$point of the helical perturbation is on the outer torus side, since on this side reconnection could be favored due to the compression of flux surfaces. Such local reconnection can be much faster than in th 2D case, and destruction of surfaces is more efficient, corresponding to a localized loss of equilibrium. However, to validate this concept fully 3D nonlinear computations are needed.

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Figure captions

Fig. 1 Normalized growth rates γ/γ_0 as functions of normalized time
(a) $r_n = 10$; $\beta_p = 0, 2, 4$
(b) $r_n = 0.7$; $\beta_p = 1.4, 2, 4$

Fig. Behavior of average nonlinear growth rate. Above the full line saturation occurs (partial reconnection), while below full reconnection occurs. The vertical and horizontal lines indicate the parameter range, where simulation runs have been performed.

Fig. 3 Evolution toward a saturated state of the case $r_n = 0.7$, $\beta_p = 2$ of Fig. 1b
(a) $t = 1668$
(b) $t = 1780$ (saturated state)
The excitation of the Kelvin-Helmholtz instability is clearly seen.

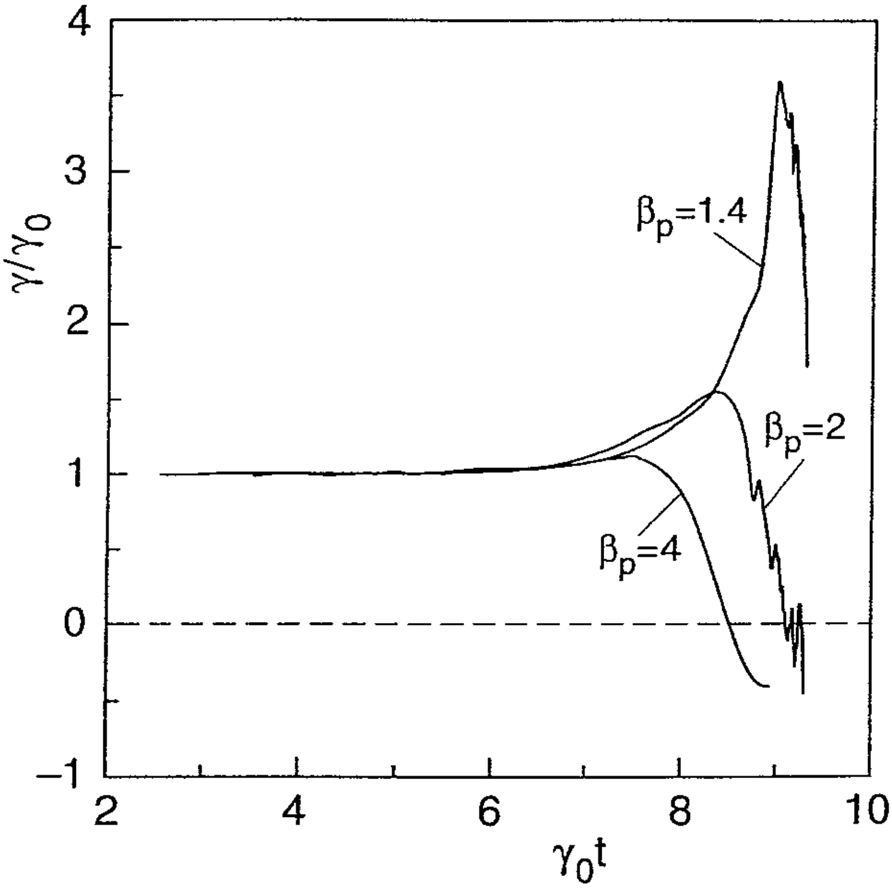

Fig. 1a

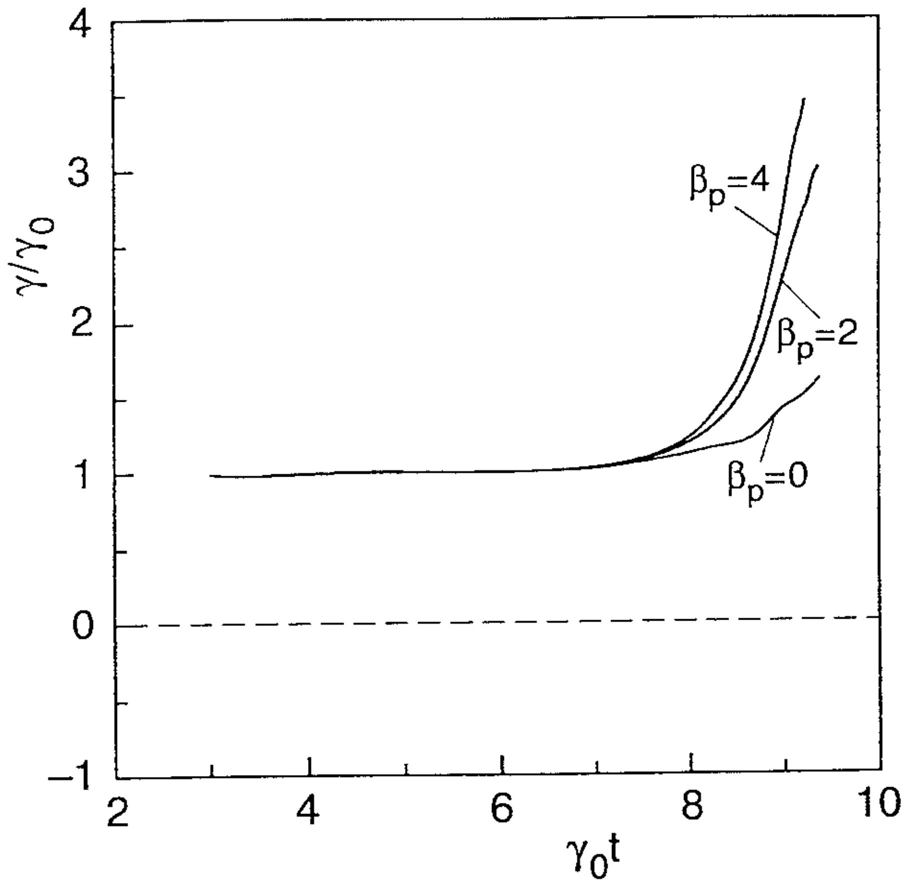

Fig. 1b

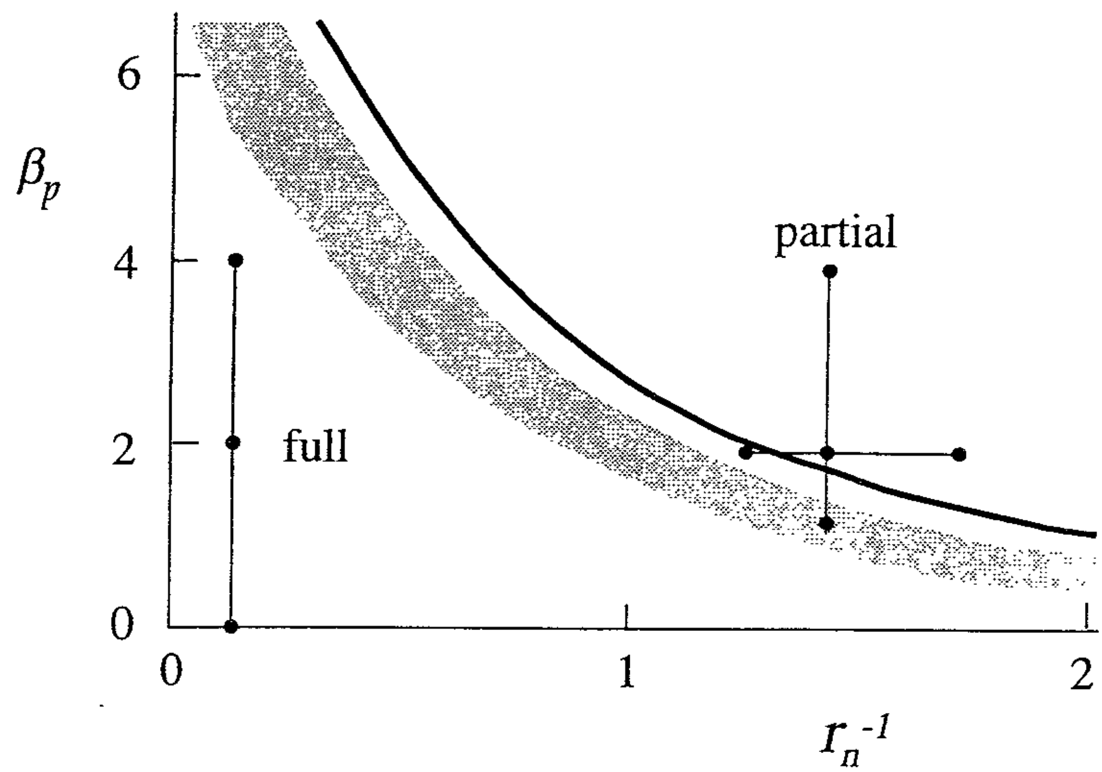

Fig. 2

$t = 1668$

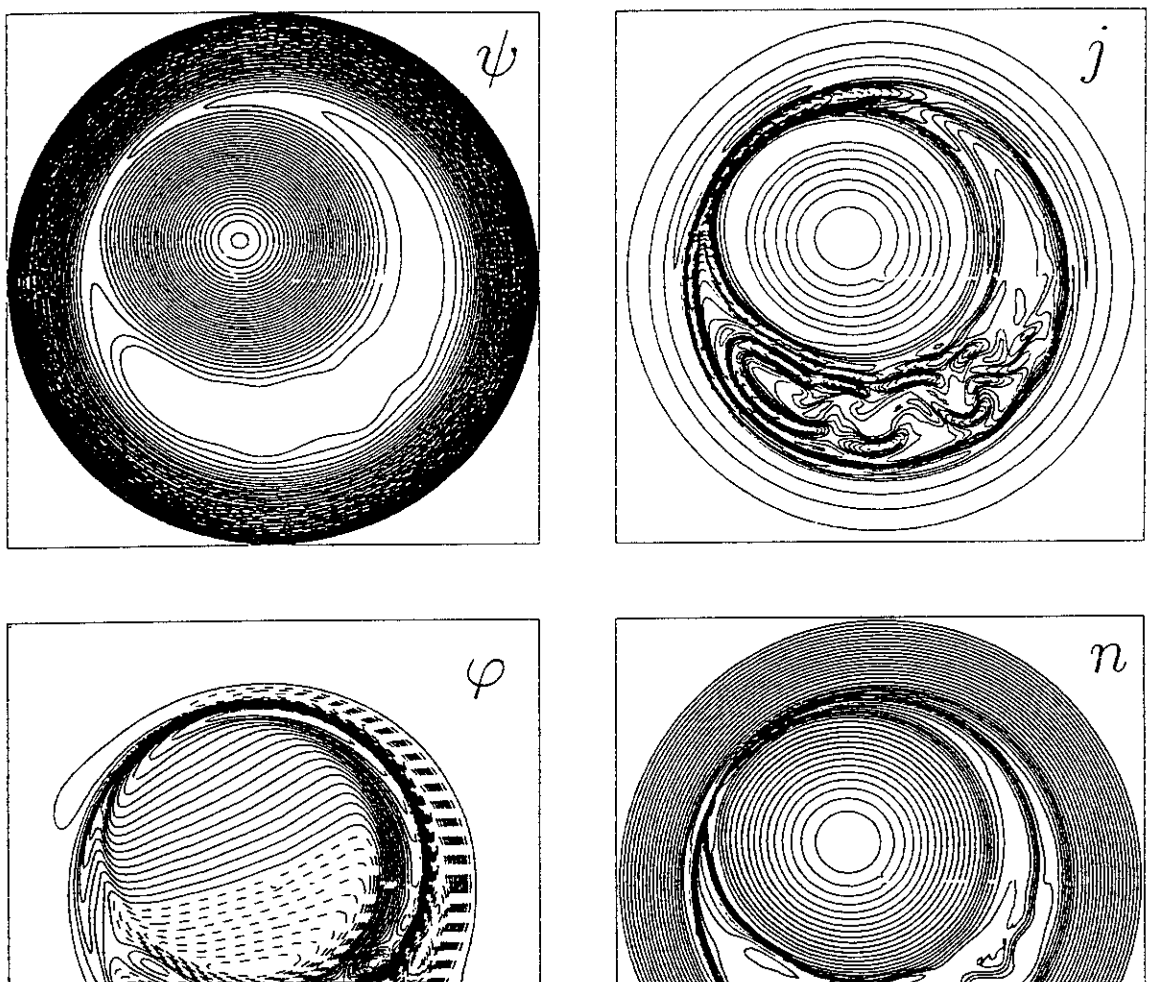

Fig. 3a

$t = 1780$

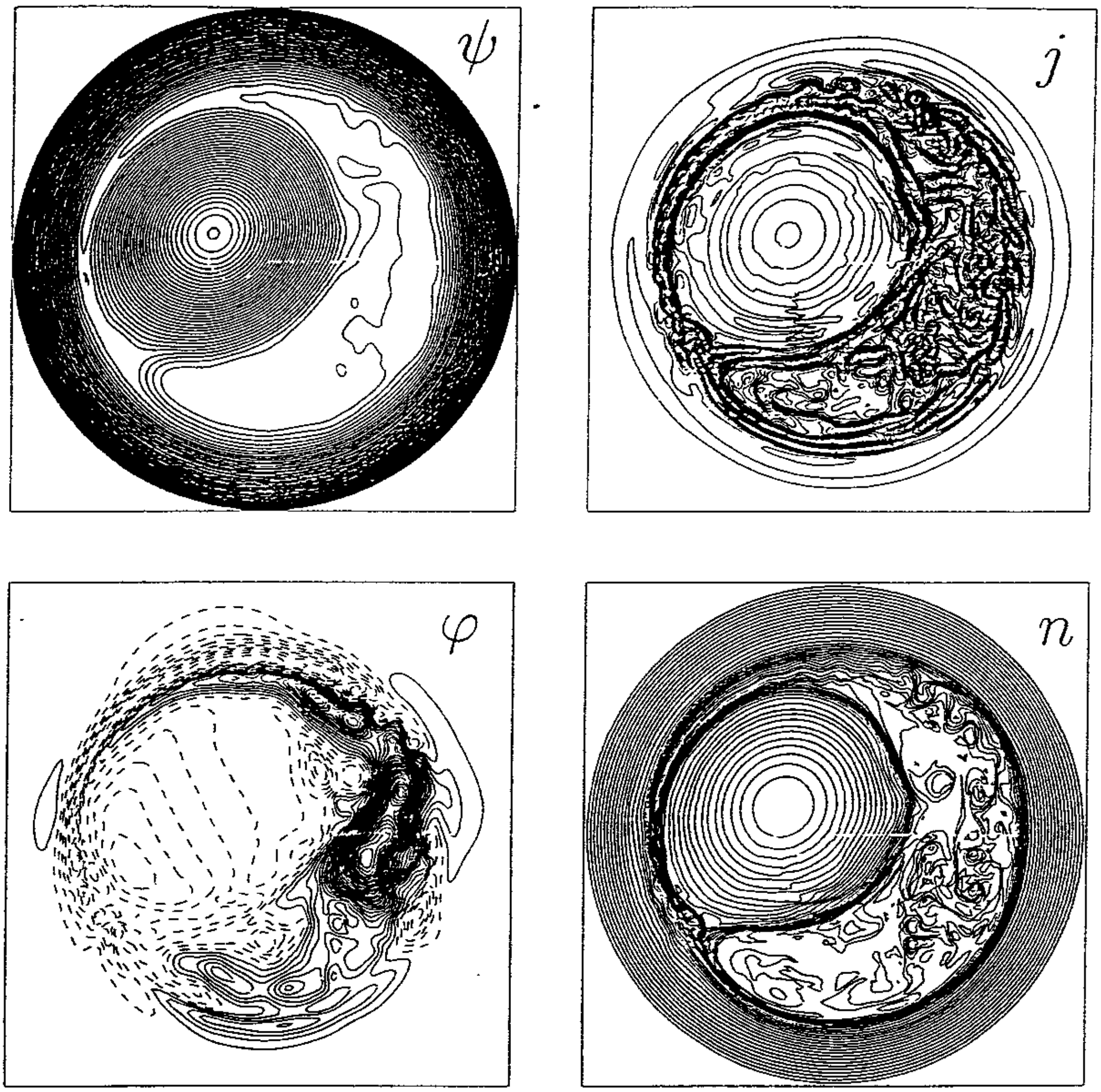

Fig. 3b

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