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SMR 1331/11

AUTUMN COLLEGE ON PLASMA PHYSICS

8 October - 2 November 2001

Magnetic Holes

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These are preliminary lecture notes, intended only for distribution to participants.

Generation Mechanism for Magnetic Holes in the Solar Wind

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Abstract

A new mechanism for generation of magnetic holes in the solar wind is presented. In the high speed solar wind , large-amplitude right-hand polarized Alfvénic wave packets propagating at large angles to the ambient magnetic field are shown to generate magnetic holes (MHs). Characteristics of these holes crucially depend on plasma β (β being the ratio of kinetic pressure to magnetic pressure) and the ratio of electron temperature T_e to proton temperature T_i. Proton temperature anisotropy is found to be favorable but not essential for the development of MHs. From our simulations we observe MHs with microstructures bounded by sharp gradients (magnetic decreases) in some cases. The holes generated by this process have thicknesses of hundreds of ion Larmor radii, typical of many of the solar wind hole observations, the depths of the holes are also comparable. The theory can explain the presence of MHs seen in the solar wind for those cases when anisotropies are not favorable for the development of the mirror mode instability.

Geophys. Res. Lett. 28 , 1355 (2001)

Introduction

The study of magnetic holes (depressions in the magnitude of the magnetic field) in the solar wind is attracting a lot of attention as more observations have recently been reported by a number of authors using Ulysses data (Smith and Marsden, 1995). Linear magnetic holes (MHs) in the interplanetary medium were reported by Turner et al. (1977). These are isolated structures embedded in uniform magnetic fields and have widths typically $\sim 2 \times 10^4$ km. These were found in regions of high proton β with very small changes in the direction of the magnetic field. Turner et al. conjectured that diamagnetic effects associated with localized inhomogeneities were responsible for creating these holes. Mirror mode depressions have also been reported in magnetosheaths of Earth, Jupiter and Saturn (Tsurutani et al., 1982) and also in the environment of comets (Russell et al., 1987).

Ulysses high resolution magnetic field and plasma data recently stimulated further investigations. Winterhalter et al. (1994) looked for linear magnetic holes in the solar wind, outside the Jovian environment and indeed found a large number satisfying the criteria they had used to identify them. Winterhalter et al. (2000) have extended their studies to determine the latitudinal distribution of solar wind magnetic holes. In the polar solar wind at Ulysses distances ($\sim 2AU$) magnetic depressions (MDs), apparent holes bounded by discontinuities, with scale sizes of $25 - 70$ proton gyroradii have been reported by Tsurutani et al. (1999). They found that often these depressions are bounded by tangential discontinuities . These depressions have been reported as solitary structures. Similar conclusions about the association of tangential discontinuities with MHs have been drawn by Fränz et al. (2000); they also used Ulysses data for magnetic field as well as plasma parameters.

The observed linear magnetic holes have been interpreted as mirror modes by Winterhalter et al. (1994). They claim that outside the holes the mirror mode instability criterion is not satisfied but within the holes the anisotropies and plasma β are large enough to satisfy the threshold for this instability. The widths of such holes are determined by the time for nonlinear saturation of the mirror mode instability to set in. According to Fränz et al. (2000) the proton anisotropy within the holes is not always higher than that outside the holes and the mirror mode instability criterion is satisfied only in $\sim 50\%$ of all cases. In such cases mirror mode is debatable as the source for MHs, so one should look for an alternative explanation. Baumgärtel (1999) has suggested that these holes are dark soliton solutions of the Derivative Nonlinear Schrödinger (DNLS) equation (Kennel et al., 1988; Hada et al., 1989; Buti, 1990) governing AW propagating at large angles to the ambient magnetic field. From simulation studies of the interaction of two propagating dark solitons, Baumgärtel claimed that, unlike the bright solitons, the dark solitons are stable. In his simulation the interaction of two propagating DNLS solitons was carried out neglecting the evolution of density fluctuations altogether. We would like to emphasize that in the derivation of the DNLS (cf. Kennel et al. 1988) equation only linear density perturbations are retained. Consequently Baumgärtel's results do not guarantee a real stability. Inclusion of density fluctuations could drastically change the stability properties. We would like to stress that one should exercise great caution in accepting Baumgärtel's claim that MHs are DNLS dark solitons for the simple reason that the DNLS equation is strictly valid only for parallel / quasi parallel propagation (Kennel et al., 1988). Earlier DNLS was considered to be valid strictly for parallel propagation and Kennel et al. (1988) had shown that it is valid for quasiparallel propagation also but Baumgärtel wrongly claims that Kennel et al. (1988) had shown its validity even for propagation at very large angles. We have done hybrid simulations to investigate the evolution of dark as well as bright solitons appropriately incorporating nonlinear density fluctuations as well as kinetic effects. We find that in both cases even a single soliton, let alone two interacting solitons, is not stable. Bright solitons propagating at large angles are found to evolve into MHs. Since the present paper deals with MHs, discussion of stability of dark solitons is out of place here; results for dark solitons will be reported in a separate publication.

In this paper, we present an alternative model for the generation of MHs. Taking clues from Ulysses observations showing the continuous existence of very large amplitude ($\delta B/B \sim 1-2$) Alfvén waves (Smith et al., 1995) accompanying the fast solar wind, we decided to explore the possibility of these waves as a potential source for the observed MHs. In principle large amplitude Alfvén waves should evolve into solitary waves (Kennel et al., 1988) but curiously enough according to our knowledge such solitary Alfvénic structures have not been observed in the solar wind. From the dispersive MHD simulations, Buti et al. (1998) and Velli et al. (1999) have shown that these structures get disrupted due to coupling of magnetic field fluctuations and density fluctuations. This coupling plays a very significant role in plasmas with large $\beta(\sim 1)$. In fact, the neglect of this coupling by Baumgärtel (1999) is the reason for the misleading conclusion drawn by him. Inhomogeneities in the system can also destroy these solitary structures (Buti et al., 1999). In the solar wind, plasma kinetic effects also play a very important role. Buti et al. (2000a, b) have incorporated the latter effects in their hybrid simulations of Alfvénic wave packets (AWP) and found interesting new features like collapse of left hand polarized wave packets and generation of ion density holes. All these investigations were restricted to parallel propagation of wave packets in isotropic plasmas. Here we have extended our high resolution hybrid simulations to investigate the evolution of right-hand polarized (RHP) AWP propagating at very large angles ($\sim 80°$) to the ambient magnetic field in anisotropic plasmas similar to the high speed solar wind (large β, $T_e < T_i$ and $T_{i\perp} \neq T_{i\|}$). These wave packets introduce localized inhomogeneities in an otherwise homogeneous plasma. For plasma parameters consistent with regions where MHs have been observed our simulations show that these packets evolve into MHs with characteristics similar to those observed in the polar solar wind. Propagation at large angles is a pertinent condition for the generation of MHs. However proton thermal anisotropy is not a must and it can be relaxed.

Simulation Model

We use a one-dimensional hybrid code (Winske and Leroy, 1985) to study the nonlinear dynamical evolution of large-amplitude Alfvénic wave packets propagating along the x-axis at an angle θ to the direction of the ambient magnetic field $\mathbf{B}_0 = (B_0 cos\theta,\ B_0 sin\theta, 0)$. Although there is only one spatial coordinate $(\partial/\partial y = \partial/\partial z = 0)$, all three components of velocity and other fields are retained in the calculations. In our simulations, electrons are treated as an isothermal fluid whereas the protons are treated as particles having an initial biMaxwellian distribution modified by the localized density inhomogeneity. For the high-resolution studies, we take the simulation box length as 860 ion - inertial lengths ($L_i \equiv V_A/\Omega_i$) with 2048 grid points and 200 particles in each cell. Simulations are carried out using periodic boundary conditions. For the initial condition, we take an AWP with,

$$B\ (x,\ t=0)\ =\ \frac{(2^{1/2}-1)^{1/2}\ B_p\ e^{i\phi\ (x)}}{\sqrt{(2^{1/2}\ \cosh\ (2\ V_p\ x) - 1)}}, \tag{1}$$

where $B = B_y + iB_z$, B_p is the amplitude of the wave packet, ϕ is the phase given by,

$$\phi\ (x)\ =\ V_p\ x\ -\ \frac{3}{8(1\ -\ \beta)}\ \int_{-\infty}^{2x} \mid B \mid^2 dx' \tag{2}$$

and V_p is the speed in the wave frame of reference and is defined by,

$$V_p\ =\ \frac{(2^{1/2}-1)\ B_p^2}{8\ (\beta - 1)}. \tag{3}$$

Throughout, we use the normalized variables e.g., $\mathbf{B}$ is normalized to B_0, ρ to ρ_0, $\mathbf{v}$ to $V_A = B_0/(4\pi\rho_0)^{1/2}$ (V_A being the Alfvén velocity), t to inverse of Ω_i, the ion cyclotron frequency and l to L_i. The subscript '0' refers to the equilibrium quantities. Note that Eq. (1) is an RHP super-Alfvénic soliton solution of the DNLS (for parallel propagating Alfvén waves) equation in the wave frame of reference (Verheest and Buti, 1992). We call it an AWP simply because, unlike a soliton, AWP does not remain stable as a result of evolution. β appearing in Eqs.(2) and (3) is the plasma β i.e., $\beta = \beta_e + \beta_i$ with T_i defined as $T_i = (T_{i\|} + 2T_{i\perp})/3$. Specific reasons for selecting soliton solution as initial condition are elaborated in Buti et al. (1999). We study the spatio-temporal evolution of this wave packet as it propagates along the x-axis. Eq.(1) represents a localized magnetic field pulse in an otherwise homogeneous plasma. Correspondingly the initial transverse magnetic field $\mathbf{b} = (b_y, b_z)$ is given by,

$$b_y(x, t = 0) = B_0 sin\theta + \mid B \mid cos\phi \tag{4}$$

and

$$b_z(x, t = 0) = \mid B \mid sin\phi. \tag{5}$$

For strong pulses with $\mid B \mid \geq 1$ (as observed by Tsurutani et al., 1999b) density perturbations, driven by the pondermotive force of magnetic field fluctuations ($\delta\rho \propto |B|^2$), become very significant. To incorporate these density perturbations appropriately, for the initial density we have considered the following three different cases:

$$\rho(x, t = 0) = 1, \tag{6}$$

$$\rho(x, t = 0) = 1 + \frac{1}{2(\beta - 1)} \mid B(x, t = 0) \mid^2, \tag{7}$$

$$\rho(x, t=0) = 1 + \frac{1}{(1+\beta_e/\beta_i)} \mid B(x, t=0) \mid^2 . \tag{8}$$

It may be worth pointing out that the sign of the second term on the right hand side of Eq. (7) is compatible with the kinetic derivation of the DNLS equation whereas Eq.(8) represents the density fluctuations corresponding to the solution obtained from the drift kinetic treatment (Inhester, 1990). Even when we initially take uniform plasma i.e., Eq.(6), density perturbations are generated by the pondermotive force during the evolution (see Fig.3b). For the initial velocities we use the appropriate relations governing Alfvén waves namely $v_{y,z} = -B_{y,z}$. The code of Winske and Leroy (1985) has been modified to impose correct periodic boundary conditions on the vector potential in the presence of the localized inhomogeneities introduced by our specific initial conditions.

Simulation Results

Recently Buti et al. (2000 a, b) had looked into the spatio-temporal evolution of large amplitude AWP propagating parallel to the direction of the ambient magnetic field in isotropic plasmas. In our present simulations, we instead take the localized AWP propagating at $\theta \sim 75° - 80°$. Simulation parameters used are compatible with the regions where MHs are observed (Winterhalter et al., 1994; Fränz et al., 2000). The spatio-temporal evolution of the transverse magnetic field b ($= (b_y^2 + b_z^2)^{1/2}$) for $\beta = 4$, $T_i/T_e = 3$, $A \equiv T_{i\perp}/T_{i\parallel} = 2$ and propagation angle $\theta = 80°$ are shown in Fig.1. Note that the average value of A in polar solar wind is ~ 1

(Marsch et al., 1982) but just outside (pre event) an MH, A can be ≥ 2 (Neugebauer et al., 2001). For Figs.1 and 2, thermal noise in the simulation data is removed by averaging over $10L_i$. However for Table 1. no averaging has been done. The initial amplitude of the packet B_p in case of Fig.1 is 2. We see a very quick conversion of the compressive wave packet into an MD bounded by sharp gradients (see Fig.1b). Later on, around $t = 800$ ion gyroperiods (see Fig.1c), this MD evolves into a deep magnetic hole with depth ($= b_{max} - b_{min}$) ~ 0.6 (0.7) B_0 and width (shown by dotted lines) ~ 350 L_i i.e., $200R_L$ (= ion gyroradii). From Fig.1d we see that at $t = 1000$ the hole attains depth of ~ 0.65 (0.8) B_0 and width ~ 400 L_i. The numbers in parenthesis for depths correspond to data with no averaging. In the solar wind for proton density ($N_P = 0.5$) $L_i \sim 360$ kms. The change in the direction of magnetic field across the hole in this case is only 2.7°. The corresponding density fluctuations (N) (see Fig.3a)) are always compressive i.e., b and N are anticorrelated as in the case of observed MHs (see Fig.6 of Fränz et al.,2000). Fig.3b shows the evolution of N for the case $N(t = 0) = 0$. Fig.3 clearly shows the effect of nonlinear forces due to magnetic field fluctuations on the particle dynamics.

We can roughly understand the conversion of magnetic field from compression to depression as follows. A Hamiltonian approach for the study of *stationary* nonlinear Alfvén waves governed by the DNLS equation showed (Hada et al., 1989; Buti, 1990) that these waves are governed by a pseudopotential with saddle point at $b = B_{0y}$. The two constant energy (or equipotential) contours passing through this point represent compressive ($b > B_{0y}$) and rarefactive ($b < B_{0y}$) solitons. In this case there is no possibility of overlapping / crossing of these contours. However, the evolution of density fluctuations in our model (see Fig.3) and the coupling of these fluctuations with the magnetic field fluctuations provide a source of perturbation that allows the overlapping and consequently the conversion of a compressive wave packet to a MH. For the simple reason that the saddle point is very sensitive even to small perturbations, we observe quick transitions to MHs. Further this argument justifies the absence of MHs for parallel propagation simply because b can not be < 0 and the saddle point for $\theta = 0$ is at $b = 0$.

Table 1 shows how the properties of MHs change with plasma β, the initial amplitude of the wave packet (B_P) and with anisotropy (A). From the summary of the simulation results shown in this table, we infer that the holes do form even in an isotropic case unlike the mirror mode model for MHs. However the size of the hole, both depth and width, are smaller in an isotropic case compared to the case of $A = 2$. Similar results hold good for $B_p = 2$ as well as 0.8 but the decreases in depth and width are much larger for the latter case. For all the cases shown in the Table, the angle of propagation is 80°. It is worth noting that the change in magnetic field direction across the hole in all the cases is $\leq 6°$. Moreover the holes reported here are solitary structures like the 'linear holes' but have microstructures in some cases. The depth of MH decreases but the width increases when plasma β is changed from 4 to 2.4 whereas both depth and width decrease when β_i/β_e is changed from 3 to 1 keeping all the other parameters constant.

We repeated the simulations for other large angles of propagation ($70° < \theta < 80°$); the results (not shown here) were qualitatively the same but showed slight quantitative differences. However, for parallel propagation no holes whatsoever were observed. Instead only turbulence/noise was observed. In all the cases discussed so far, the initial density fluctuations taken for the simulations were according to Eq. (7). In order to make sure that the MH formation was not due to this specific initial condition, we repeated the calculations using Eq.(6) as an initial condition i.e., $N(t=0)=0$; the results are shown in Fig.2a. Similar results for the density profile given by Eq.(8) are shown in Fig.2b. On comparing Figs.2a and 2b with Fig.1d showing the evolution at $t=1000$, we see some quantitative but no qualitative differences in the properties of the MHs generated. This confirms that the source of MHs in our model is the localized AWP independent of any specific *initial* localized density inhomogeneity. We would like to emphasize however that the coupling of density fluctuations generated during the evolution (see Fig.3) with the magnetic field fluctuations plays a crucial role in the development of MHs. If this coupling was absent there would be no evolution of the density fluctuations whatsoever.

Discussion and Conclusions

Localized inhomogeneities introduced by large amplitude Alfvénic wave packets lead to MHs in the fast solar wind. We would like to stress that the initial magnetic pulse given by Eq.(1) is very different from the dark soliton solutions of the DNLS (for AW propagating at large angles) studied by Baumgärtel. As pointed out by Kennel et al. (1988), validity of DNLS equation for propagation of Alfvén waves at large angles is questionable. Unfortunately Baumgärtel had misinterpreted our results (Kennel et al., 1988) to justify his use of DNLS equation for large angle propagation. Moreover in his stability analysis of two counterstreaming dark solitons, density *evolution* has been ignored whereas in the present paper we have shown the crucial role played by *evolving* density perturbations in the nonlinear evolution of AWP. In summary the mechanism proposed here is altogether different from Baumgärtel's model. In the initial stages of the evolution, we observe MDs bounded by sharp gradients (see Fig.1b) and later on these MDs evolve into MHs with thicknesses of hundreds of Larmor radii. Most of the holes produced in our model seem to be solitary in nature and have shapes resembling either 'V' or 'W'. For generation of these holes, propagation at very large angles is found to be essential whereas the restriction on proton temperature anisotropy can be relaxed. Consequently our model can also explain the presence of observed MHs in the solar wind for the cases when anisotropies are not favorable for the development of the mirror mode instability. The initial localized density inhomogeneities are favourable but not a must for the generation mechanism outlined in the present paper. However this does not mean that the evolution of the density can be ignored. As shown in Fig.3b, density perturbations are generated during the evolution even for the case of zero initial density perturbation. Properties of MHs e.g., depth, width etc. depend on plasma β, ion-temperature anisotropy as well as the initial amplitude of the wave packet (see Table 1).

Acknowledgments. The research conducted at the Jet Propulsion Laboratory, California Institute of Technology, was performed under contract to the National Aeronautics and Space Administration. BB would like to acknowledge some helpful discussions with E.J. Smith, M. Velli and D. Winterhalter.

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Received ____________________

Figure Captions

Fig. 1 shows spatio-temporal evolution of magnitude of the magnetic field (b) of an RHP wave packet for $\beta = 4$, $B_p = 2$, $A = 2$, $\theta = 80°$ and $\beta_i/\beta_e = 3$ at $t = 0, 100, 800$ and 1000 ion-gyroperiods.

Fig. 2 shows evolution of b at $t = 1000$. a) shows MH similar to the one in Fig. 1d but with zero initial density fluctuations; b) same as Fig. 1d but with initial density given by Eq.(8).

Fig. 3 Stack plots for density fluctuations (N) for $\beta = 4$, $B_p = 2$, $A = 2$, $\theta = 80°$ and $\beta_i/\beta_e = 3$. Temporal evolutions shown are for $t = 0, 80$ and $(80 + \Delta t)$ with $\Delta t = 40$ ion-gyroperiods. Initial density profiles used for Figs.3a and 3b respectively are given by Eqs. (7) and (6).

Table 1. Properties of Magnetic Holes from evolution at t = 800 gyroperiods

β	β_i / β_e	B_P / B_0	A	Depth (B_0)	Width (R_L)	$\lvert\Delta\phi\rvert$ (degrees)
4	3	2	2	0.70	200	2.7
4	3	2	1	0.25	70	3.7
4	1	2	2	0.35	100	2.7
4	3	0.8	2	0.55	215	5.6
4	3	0.8	1	0.40	175	1.6
2.4	3	0.8	2	0.35	260	0.1
2.4	2	0.8	1.4	0.30	160	0.7

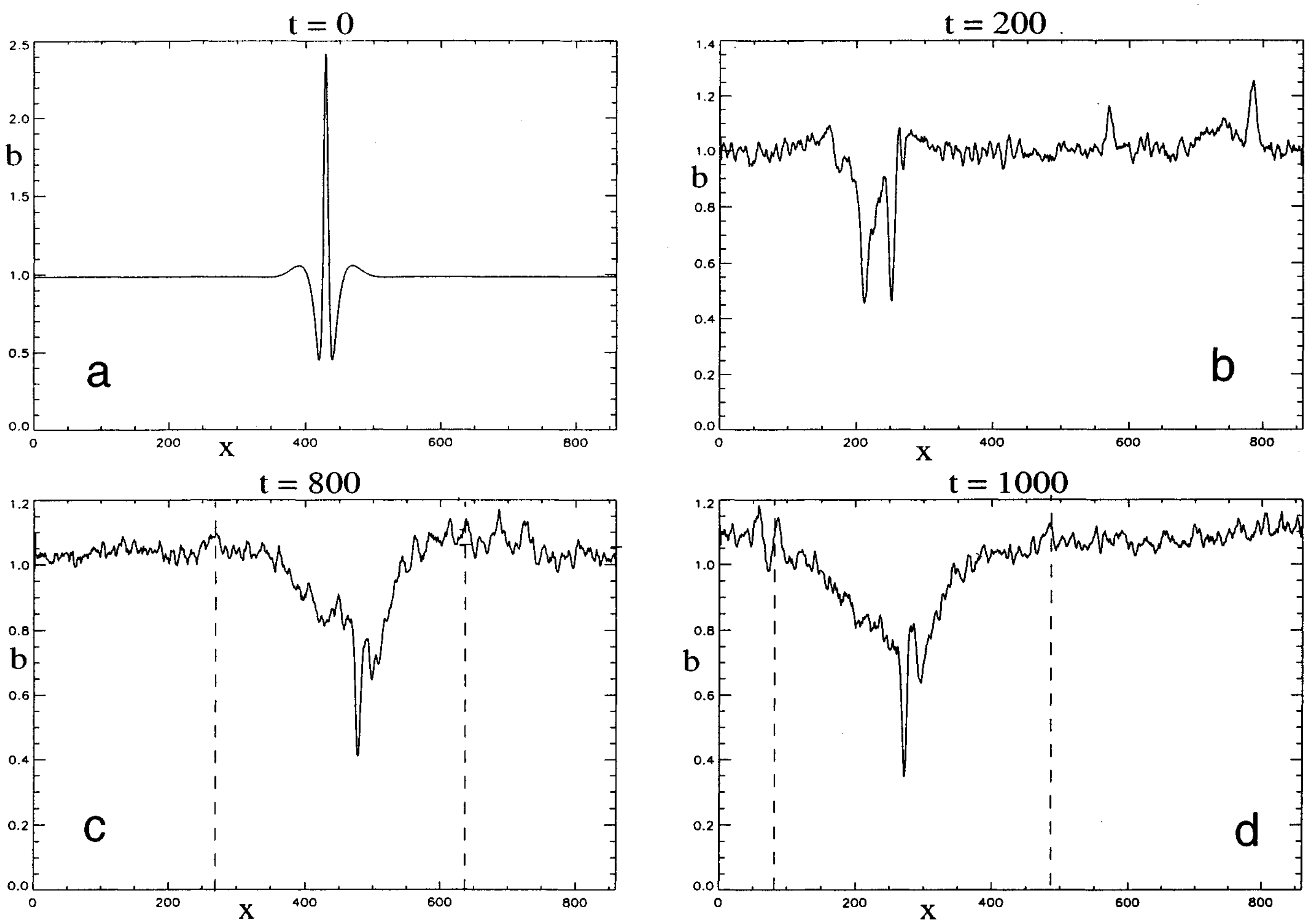

Fig. 1

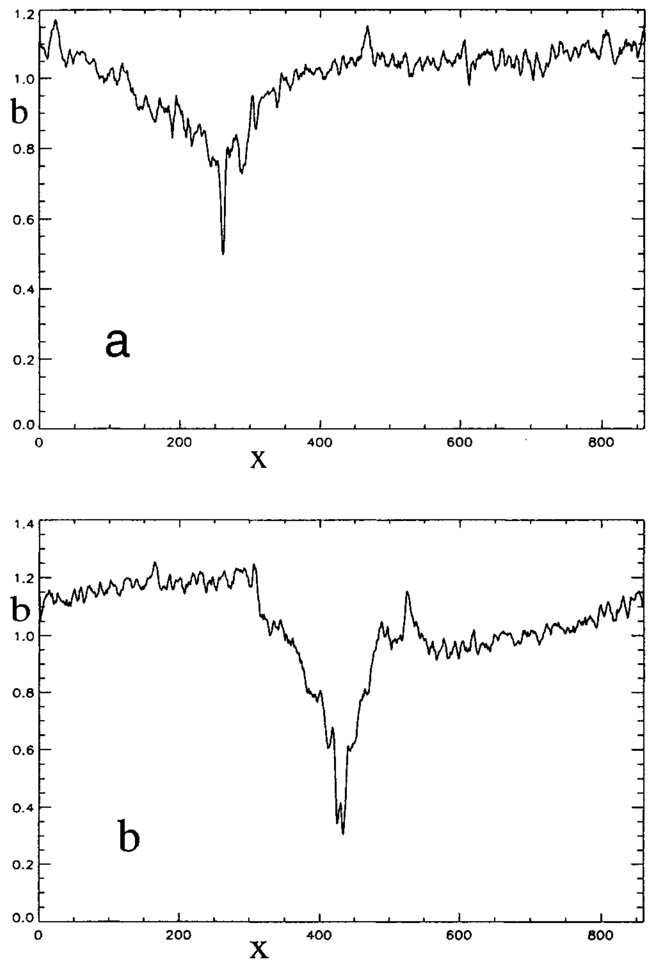

Fig. 2

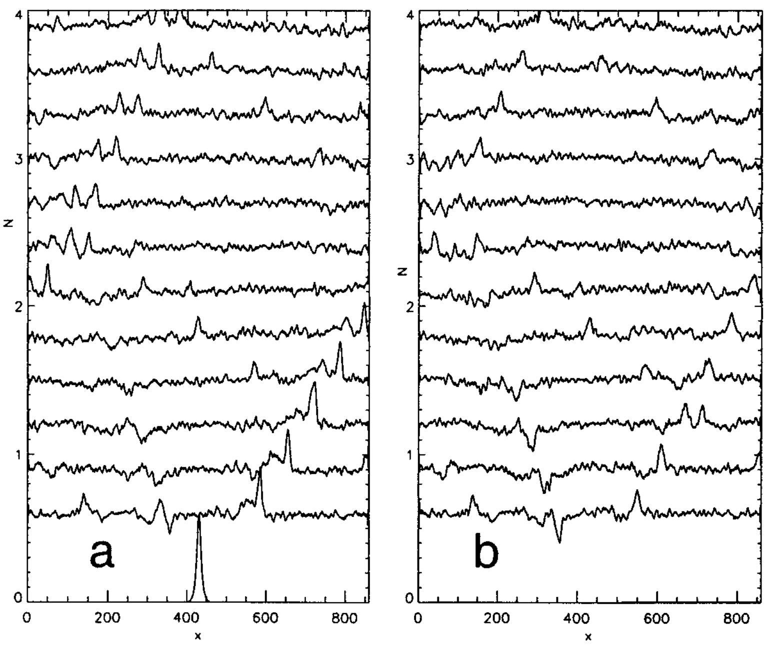

Fig. 3