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Demonstration for Novel Self-organization Theory by Three-Dimensional Magnetohydrodynamic Simulation

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It is demonstrated by three-dimensional simulations for resistive magnetohydrodynamic (MHD) plasmas with both "spatially nonuniform resistivity η" and "uniform η" that the attractor of the dissipative structure in the resistive MHD plasmas is given by $\nabla \times (\eta \mathbf{j}) = (\alpha/2)\mathbf{B}$ which is derived from a novel self-organization theory based on the minimum dissipation rate profile. It is shown by the simulations that the attractor is reduced to $\nabla \times \mathbf{B} = \lambda \mathbf{B}$ in the special case with the "uniform η" and no pressure gradient.

Keywords : self-organization, eigenfunction, dissipative dynamics operator, attractor, dissipative structure, 3-D MHD simulation

It has been widely believed that the self-organization phenomena in dissipative magnetohydrodynamic (MHD) plasmas and the resultant profiles such as the reversed field pinch (RFP) [1–4] and the spheromak [5] are explained by the energy-relaxation theory and the force-free field, $\nabla \times \mathbf{B} = \lambda \mathbf{B}$ with a constant profile of λ, which were proposed and derived by J. B. Taylor [6, 7] with use of the energy principle and the conjecture on "the time invariant of the total magnetic helicity". Taking account of the experimental RFP plasma which has the finite pressure gradient and satisfies the boundary condition that the current density $\mathbf{j} = \mathbf{0}$ at the wall, one of the authors (Y.K.) had introduced the partially relaxed state model (PRSM) [8–10] and developed numerical codes for the RFP equilibria and for the mode transition point of the relaxed states by introducing the energy principle with partial loss of magnetic helicity in the boundary region [11–15]. The experimental data of the RFP plasma in the TPE-1RM15 device [4] have been shown to be fitted well by the numerical results of the PRSM [11, 15]. However, the following five facts of (A) - (E), which do not agree basically with the Taylor theory [6, 7], are clarified recently by experiments and by three-dimensional (3-D) MHD simulations: (A) Having no concern with the total magnetic helicity conservation (for example, without depending on the magnetic helicity loss up to about 20 percent or on the negligible helicity loss), the MHD plasma system with the spatially uniform resistivity still relaxes to the state of $\nabla \times \mathbf{B} = \lambda \mathbf{B}$, as was demonstrated clearly by the 3-D MHD simulations shown in Fig.3 in ref.[16]. (B) The simple toroidal Z pinch without any initial magnetic helicity in the ZP-2 device [17] has been clarified experimentally to relax close to the state of $\nabla \times \mathbf{B} = \lambda \mathbf{B}$ with finite total helicity within a few tens of μs (i.e. the total helicity is not the time invariant during the relaxation phase), as was reported in refs.[17, 18]. (C) The relaxed state of the MHD plasma without pressure gradient (i.e. $\nabla p = 0$) does depend rather explicitly on the spatial profile of the resistivity η, as was demonstrated clearly by the 3-D MHD simulations in Fig.7 in ref.[19] (i.e. the spatially uniform resistivity leads to the relaxed state of $\nabla \times \mathbf{B} = \lambda \mathbf{B}$ and the nonuniform resistivity leads to the deviation from this so-called "Taylor state"). (D) It has been demonstrated that the recent merging experiments of two spheromaks realize the field-reversed configuration (FRC) plasma as a branch of the relaxed states, which is completely different from $\nabla \times \mathbf{B} = \lambda \mathbf{B}$ and has no connection with the total magnetic helicity invariant, as was shown in Fig.2 in ref.[20]. (E) The relaxed state of the MHD plasma with the pressure gradient does deviate obviously from the force-free state of $\nabla \times \mathbf{B} = \lambda \mathbf{B}$, as was demonstrated clearly by the 3-D MHD simulations with the heat conductivity in Figs.5, 6, and 8 in ref.[21]. As to the phenomenon of the self-organization, there is no reason to distinguish among these self-organized relaxed states with the uniform resistivity, with the nonuniform resistivity and/or with and without the pressure gradient. Taking account of those recent facts obtained by the experiments and the 3-D MHD simulations mentioned above, one of the authors (Y.K.) has proposed a novel self-organization theory based on the minimum dissipation rate profile and has shown that the attractor of the dissipative structure for MHD plasmas with resistivity η is given by the eigenfunction for the dissipative dynamics operator $-\nabla \times (\eta \mathbf{j})$, i.e. $\nabla \times (\eta \mathbf{j}) = (\alpha/2)\mathbf{B}$ [22–25]. We report here new results by 3-D MHD simulations demonstrating this novel self-organization theory.

A brief description of the novel self-organization theory [22–25] is as follows: A general dissipative nonlinear dynamics system with n elements of $\mathbf{q}(t, \mathbf{x}) = \{q_1(t, \mathbf{x}), q_2(t, \mathbf{x}), \cdots, q_n(t, \mathbf{x})\}$ may be described by $\partial q_i/\partial t = L_i^N[\mathbf{q}] + L_i^D[\mathbf{q}]$, where $L_i^N[\mathbf{q}]$ and $L_i^D[\mathbf{q}]$ denote respectively the nondissipative and the dissipative dynamics operators. The dissipation rate of global auto-correlations $W_{ii} = \int q_i q_i \mathrm{d}V = \int q_i^2 \mathrm{d}V$

over the space volume of the system, which represent usually the total energy of the system, is written as $\partial W_{ii}/\partial t = 2\int q_i(\partial q_i/\partial t)\mathrm{d}V = 2\int q_i L_i^D[\mathbf{q}]\,\mathrm{d}V$, by using the definition for the nondissipative operator $\int q_i L_i^N[\mathbf{q}]\,\mathrm{d}V = 0$. When the dynamics system has some unstable regions, the nondissipative dynamics operator $L_i^N[\mathbf{q}]$ may become dominant and lead to the rapid growth of perturbations there and to the turbulent phases. This process would yield consequently spectrum transfers toward larger mode number (wavenumber) region in q_i distributions, like as the normal energy cascade demonstrated by 3-D MHD simulations in [16]. When the larger mode number becomes a large fraction of spectrum, the dissipative dynamics operator $L_i^D[\mathbf{q}]$ may act as dominant operators to yield the higher dissipations for the larger mode number components in W_{ii}, like as the selective dissipation reported in [16]. Through this rapid dissipation phase which is far from equilibrium, the unstable regions in the dynamics system are considered to vanish to lead to a stable configuration again. After these relaxation processes, the dynamic system is considered to attain peculiar spatial distributions such that yield the minimum dissipation rate of W_{ii} for the instantanious value of W_{ii}, and to spend the longest time there in a quasi-steady state of equilibria $L_i^N[\mathbf{q}] = 0$. We therefore come to the following self-organization theory by Kondoh [22–25] based on the minimum dissipation rate profile. The self-organized state of the general dissipative dynamics system after the turbulent and nonlinear relaxation phase is given by

$$min \mid \frac{\partial W_{ii}}{\partial t} \mid \; state \text{ for a given value of } W_{ii}. \tag{1}$$

This is a typical problem of the variational calculus with respect to the spatial variable x to find the spatial profiles of q_i^* such that satisfy eq.(1). Since $\partial W_{ii}/\partial t$ has usually negative value, by defining a functional F as $F \equiv -\partial W_{ii}/\partial t - \alpha W_{ii}$, the mathematical expressions for eq.(1) are written as $\delta F = 0$ and $\delta^2 F > 0$, where δF and $\delta^2 F$ are the first and the second variations of F, and α is the Lagrange multiplier. In the case of resistive MHD plasmas with $L_i^D[\mathbf{q}] = -\nabla \times (\eta \mathbf{j})$, we obtain the following Euler-Lagrange equation and the associated eigenvalue problem respectively from $\delta F = 0$ and $\delta^2 F > 0$ [22–25]:

$$\nabla \times (\eta \mathbf{j}^*) = \frac{\alpha}{2} \mathbf{B}^* . \tag{2}$$

$$\nabla \times (\eta \nabla \times \delta \mathbf{B}_i) - \frac{\mu_o \alpha_i}{2} \delta \mathbf{B}_i = 0. \tag{3}$$

Here, $\mathbf{B}^*$ denotes the field distributions such that satisfy eq.(1), the boundary conditions for the eigenvalue problem eq.(3) are $\delta \mathbf{B} \cdot \mathrm{ds} = 0$ and $(\eta \delta \mathbf{j} \times \delta \mathbf{B}) \cdot \mathrm{ds} = 0$ at the boundary, and α_i and $\delta \mathbf{B}_i$ denote the eigenvalue and the eigensolution, respectively. We find from eq.(2) that the profiles of the self-organized quasi-steady relaxed state with the minimum dissipation rate are given by the eigenfunction for the dissipative dynamics operator of $-\nabla \times (\eta \mathbf{j})$ [i.e. $-\nabla \times (\eta \nabla \times \mathbf{B}/\mu_o)$]. Since the present dissipative dynamics operator $-\nabla \times (\eta \mathbf{j})$ satisfies the the self-adjoint property of $\int \delta \mathbf{B} \cdot [\nabla \times (\eta \nabla \times \mathbf{B})] \, dv = \int \mathbf{B} \cdot [\nabla \times (\eta \nabla \times \delta \mathbf{B})] \, dv$ for the boundary conditions of $\delta \mathbf{B} \cdot \mathrm{ds} = 0$, and $(\eta \delta \mathbf{j} \times \delta \mathbf{B}) \cdot \mathrm{ds} = 0$ at the boundary, the eigenfunctions, $\mathbf{a}_k$, for the associated eigenvalue problem of eq.(3) form a complete set and the appropriate normalization is written as $\int \mathbf{a}_k \cdot [\nabla \times (\eta \nabla \times \mathbf{a}_j)] \, dv = \int \mathbf{a}_j \cdot [\nabla \times (\eta \nabla \times \mathbf{a}_k)] \, dv = (\mu_o \alpha_k/2) \int \mathbf{a}_j \cdot \mathbf{a}_k \, dv = (\mu_o \alpha_k/2) \, \delta_{jk}$, where $\nabla \times (\eta \nabla \times \mathbf{a}_k) - (\mu_o \alpha_k/2) \, \mathbf{a}_k = 0$ is used. After the spectrum transfers towards larger mode number region by instabilities which are followed possibly by the saturation of perturbation growth, it is considered that the nondissipative operator becomes less dominant and the dissipative operator comes into the more dominant operator. The field distribution $\mathbf{B}$ can then be expressed as $\mathbf{B} = \mathbf{B}^* + \sum_{k=1}^{\infty} c_k \mathbf{a}_k$, by using the eigensolution $\mathbf{B}^*$ for the boundary value problem of eq.(2) for the given boundary value and also by using orthogonal eigenfunctions $\mathbf{a}_k$ for the eigenvalue problem of eq.(3) with the boundary condition of $\mathbf{a}_k = 0$ and $\nabla \times \mathbf{a}_k = 0$ at the boundary. We obtain the following field equation

$$\frac{\partial \mathbf{B}^*}{\partial t} + \sum_{k=1}^{\infty} \frac{\partial (c_k \mathbf{a}_k)}{\partial t} = \nabla \times (\mathbf{u} \times \mathbf{B}) - \frac{\alpha}{2}\mathbf{B}^* - \sum_{k=1}^{\infty} \frac{\alpha_k}{2} c_k \mathbf{a}_k \, , \tag{4}$$

where $\nabla \times (\mathbf{u} \times \mathbf{B})$ acts now as a less dominant operator, the eigenvalues α_k are positive, and α_1 is the smallest positive eigenvalue. It is seen from eq.(4) that the components of $\mathbf{B}^*$ and $c_k \mathbf{a}_k$ decay approximately by the decay constants of $\alpha/2$ and $\alpha_k/2$, respectively, in the present short time interval. This selective dissipation indicated by eq.(4) gives us a detailed physical picture for the self-organization of the present dissipative system approaching the state of the basic mode $\mathbf{B}^*$ and also for the bifurcation of the dissipative structure from the basic mode $\mathbf{B}^*$ to the mixed mode with $\mathbf{B}^*$ and $\mathbf{a}_1$ which takes place at $\alpha = \alpha_1$.

Chandrasekhar and Woltjer [26] derived $\nabla \times (\nabla \times \mathbf{B}) = \lambda^2 \mathbf{B}$ in order to lead to $\nabla \times \mathbf{B} = \lambda \mathbf{B}$ and to explain the cosmic magnetic fields, by introducing two principles of "Maximum magnetic energy state for a given mean-square current density" and "Minimum dissipation state for a given magnetic energy with the assumption of the spatially uniform resistivity". The present self-organization theory by one of the authors (Y.K.) has been motivated by the experimental fact of $\mathbf{j} = 0$ near the wall where the resistivity η goes to infinity [23, 24], without knowing the work by Chandrasekhar and Woltjer [26]. It has been clarified theoretically in ref.[24] that the configurations of the FRC plasma can be derived from the relaxed state of $\nabla \times (\eta \mathbf{j}) = (\alpha/2)\mathbf{B}$. This theoretical result agrees with the fact of (D) mentioned in the introduction.

Essential difference between $\nabla \times (\eta \mathbf{j}) = (\alpha/2)\mathbf{B}$ by the present theory and $\nabla \times \mathbf{B} = \lambda \mathbf{B}$ by the Taylor theory is that the former contains explicitly the resistivity η. The latter has been shown to be the special case for the former such that has spatially uniform η and no pressure gradient (i.e. $\nabla p = 0$) [24]. Substituting $\mathbf{j} =$

$\mathbf{j}_{||} + \mathbf{j}_{\perp}$ and $\mathbf{j}_{||} = f(\mathbf{x})\mathbf{B}$ into $\nabla \times (\eta \mathbf{j}) = (\alpha/2)\mathbf{B}$, we obtain the following approximate solution, $\mathbf{j}_{||}^*$, for the parallel component $\mathbf{j}_{||}$ at the self-organized relaxed state,

$$\mathbf{j}_{||}^* \simeq \sqrt{\frac{\mu_0 \alpha}{2\eta}}\, \mathbf{B} \,, \tag{5}$$

where the subscripts $||$ and $\perp$ denote respectively the parallel and the perpendicular components to the field $\mathbf{B}$. Using spatially nonuniform resistivity and comparing the approximate solution $\mathbf{j}_{||}^*$ of eq.(5) with numerical results obtained by 3-D MHD simulations, we can demonstrate the novel self-organization theory shown above.

The basic equations for compressible, zero-β ($\nabla p = 0$), dissipative MHD plasmas without viscousity [19] to be solved here are written as

$$\frac{\partial \rho}{\partial t} = -\nabla \cdot (\rho \mathbf{u}) \,, \tag{6}$$

$$\frac{\partial \rho \mathbf{u}}{\partial t} = -\nabla \cdot (\rho \mathbf{u}\mathbf{u}) + \mathbf{j} \times \mathbf{B} \,, \tag{7}$$

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{u} \times \mathbf{B}) - \nabla \times (\eta \mathbf{j}) \,, \tag{8}$$

$$\mathbf{j} = \frac{1}{\mu_0} \nabla \times \mathbf{B} \,. \tag{9}$$

The explicit and second order Richtmyer method is used for the difference scheme in our 3-D MHD simulation code. The simulation domain is a rectangular column which is periodic along z axis with a periodic length $2\pi R$ and enclosed by perfectly conducting wall at $x = y = \pm a$, just as the same with ref.[16]. Normalization method of variables and boundary conditions are all the same as those in [16]. Two types of resistivity profiles are used here, one is a "uniform η" given by $\eta = \eta_c$, and the other is a "nonuniform η" (a parabolic type) given by $\eta = \eta_c \{ 1.0 + 19.5[(x/a)^2 + (y/a)^2]\}$, where η_c is the value of η at the center of $x = y = 0$. Numerical results for the magnetic Reynolds number of $R_\eta = \mu_0 a v_A / \eta_c = 10^4$ are shown here, where v_A is the initial value of the mean Alfvén velocity, and the time t is measured by the Alfvén time t_A defined by $t_A = a/v_A$.

Figure 1 shows a typical time evolution of three-dimentional display of the toroidal magnetic flux contours for the case of the "nonuniform η". From Fig.1, we recognize the same type of the self-organization processes with those reported in ref.[16] through the magnetic field reconnection at around $t = 4.0\ t_A$.

Figure 2 shows typical time evolutions of the toroidal flux contours on the mid-xy-plane for the two cases of (a) "nonuniform η" and (b) "uniform η". We recognize from both Fig.2 (a) and (b) that the spectrum transfers toward larger wavenumber region are taking place at around $t = 4.0\ t_A$ by instabilities , and the helical deformation still remains at around $t = 18.0\ t_A$ in the inner plasma region (cf. Fig.1).

Figure 3 shows typical time evolutions of spatial distributions of $|\ \mathbf{B}\ |$, $\mathbf{j}_{\|}$ and $\mathbf{j}_{\|}^*$ on the midhorisontal line of $y = 0$ for the two cases of (a) "nonuniform η" and (b) "uniform η". Here, the mark □ and the black one denote the data of $|\ \mathbf{B}\ |$ and $\mathbf{j}_{\|}$, respectively. The approximate solution $\mathbf{j}_{\|}^*$ by eq.(5) is denoted by the bold lines, and the both data of $\mathbf{j}_{\|}$ and $\mathbf{j}_{\|}^*$ are normalized so that the values at the center at $t = 34.0\ t_A$ coincide with the data of $|\ \mathbf{B}\ |$ in Fig.3. We recognize from Fig.3 (a) with "nonuniform η" that the distribution of the data $\mathbf{j}_{\|}$ has high wavenumbers at around $t = 4.0\ t_A$, and becomes to converge gradually around the bold line of $\mathbf{j}_{\|}^*$ at $t = 34.0\ t_A$, not around the data of $|\ \mathbf{B}\ |$ expected by $\nabla \times \mathbf{B} = \lambda \mathbf{B}$ which is obtained from the Taylor theory. In the case of Fig.3 (b) with "uniform η", $\mathbf{j}_{\|}^*$ coincides with $|\ \mathbf{B}\ |$, and the data of $\mathbf{j}_{\|}$ is recognized to converge gradually around the bold line of $\mathbf{j}_{\|}^*$, in the same way with the case of "nonuniform η" in Fig.3 (a). The distributions of $|\ \mathbf{B}\ |$, $\mathbf{j}_{\|}$ and $\mathbf{j}_{\|}^*$ on other xy-plane show the same features with those in Fig.3 (a) and (b). It is seen from comparison between (a) and (b) of Fig.3 that the attractor of the dissipative structure in the present system is given by $\nabla \times (\eta \mathbf{j}) = (\alpha/2)\mathbf{B}$, and the attractor is reduced to $\nabla \times \mathbf{B} = \lambda \mathbf{B}$ in the special case with the spatially uniform resistivity and no pressure gradient.

In conclusion, the 3-D MHD simulations presented here demonstrate that the attractor of the dissipative structure in the resistive MHD plasmas is given by $\nabla \times (\eta \mathbf{j}) = (\alpha/2)\mathbf{B}$, and the attractor is reduced to $\nabla \times \mathbf{B} = \lambda \mathbf{B}$ in the special case with the spatially uniform resistivity and no pressure gradient. We note here that the present novel self-organization theory [24, 24] is agreeable with all of the five facts of (A) - (E) mentioned in the introduction .

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Figure captions

Fig.1. Typical time evolution of three-dimentional display of the contours of the toroidal magnetic flux for the case of "nonuniform η".

Fig.2. Typical time evolutions of the contours of the toroidal flux on the mid-xy-plane. (a) nonuniform η. (b) uniform η.

Fig.3. Typical time evolutions of distributions of $\mid \mathbf{B} \mid$, $\mathbf{j}_{\parallel}$ and $\mathbf{j}_{\parallel}^{*}$ on the midhorisontal line of $y = 0$. (a) nonuniform η. (b) uniform η. The mark □ and the black one denote the data of $\mid \mathbf{B} \mid$ and $\mathbf{j}_{\parallel}$, respectively. The approximate solution $\mathbf{j}_{\parallel}^{*}$ is denoted by the bold lines.

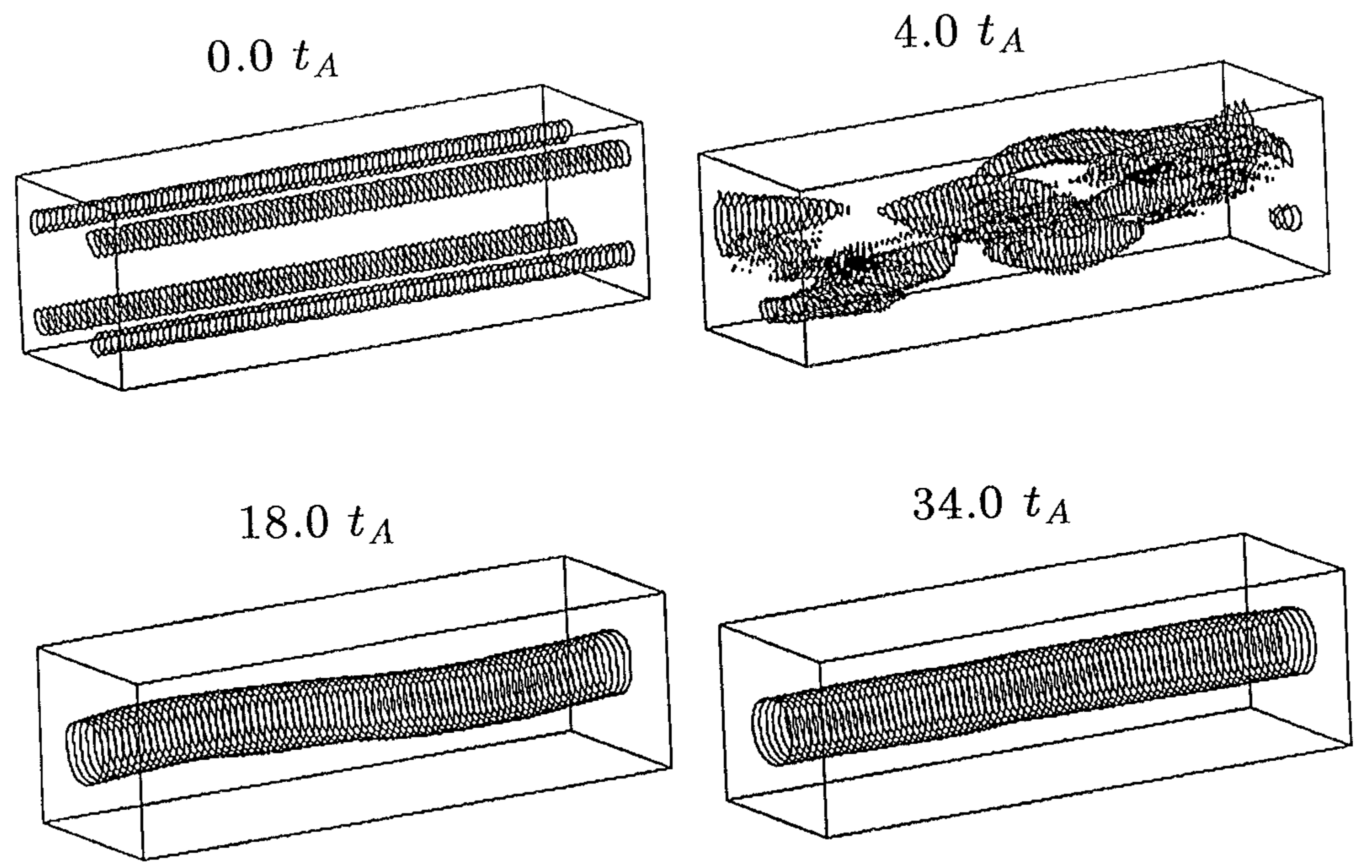

Fig.1

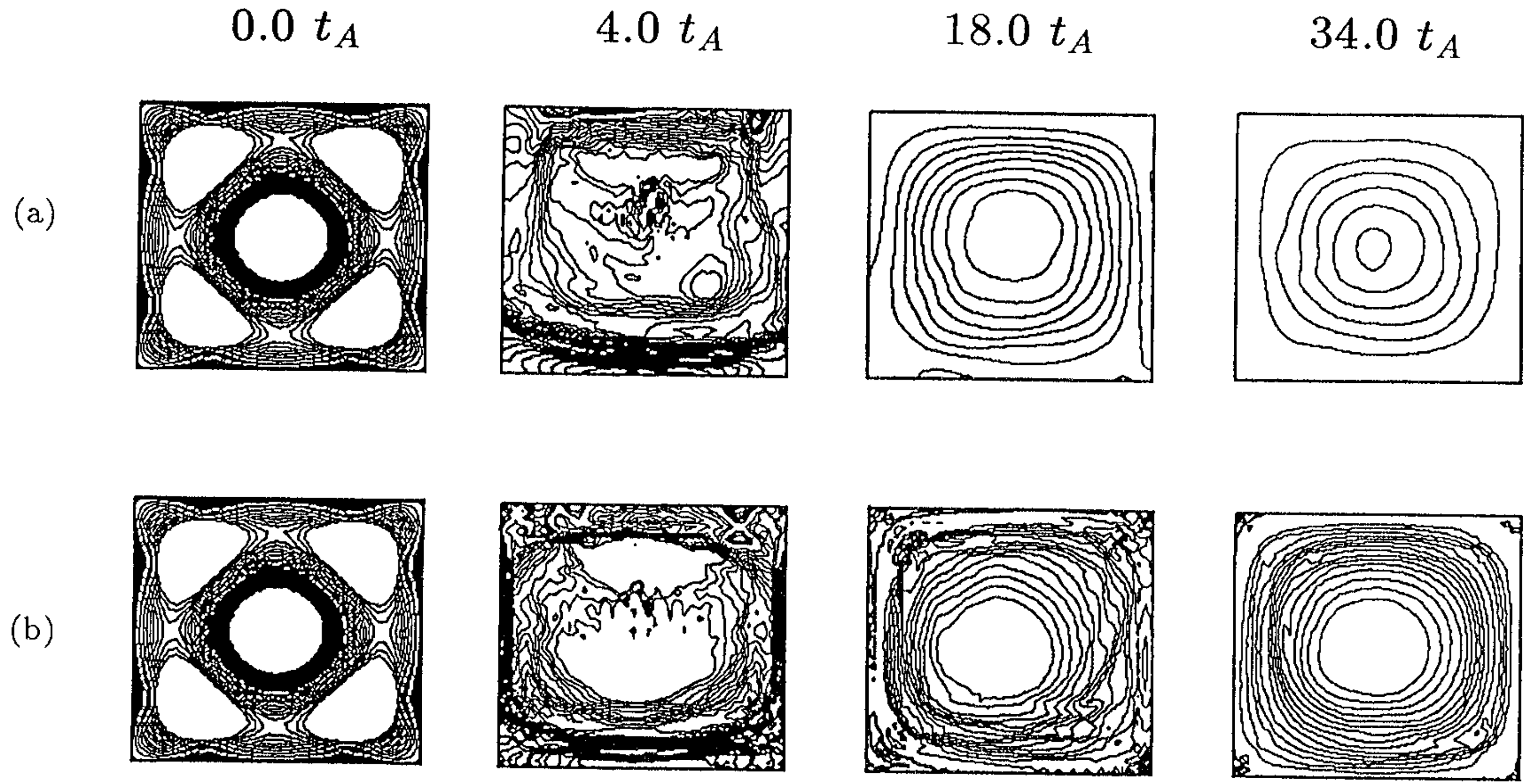

Fig.2

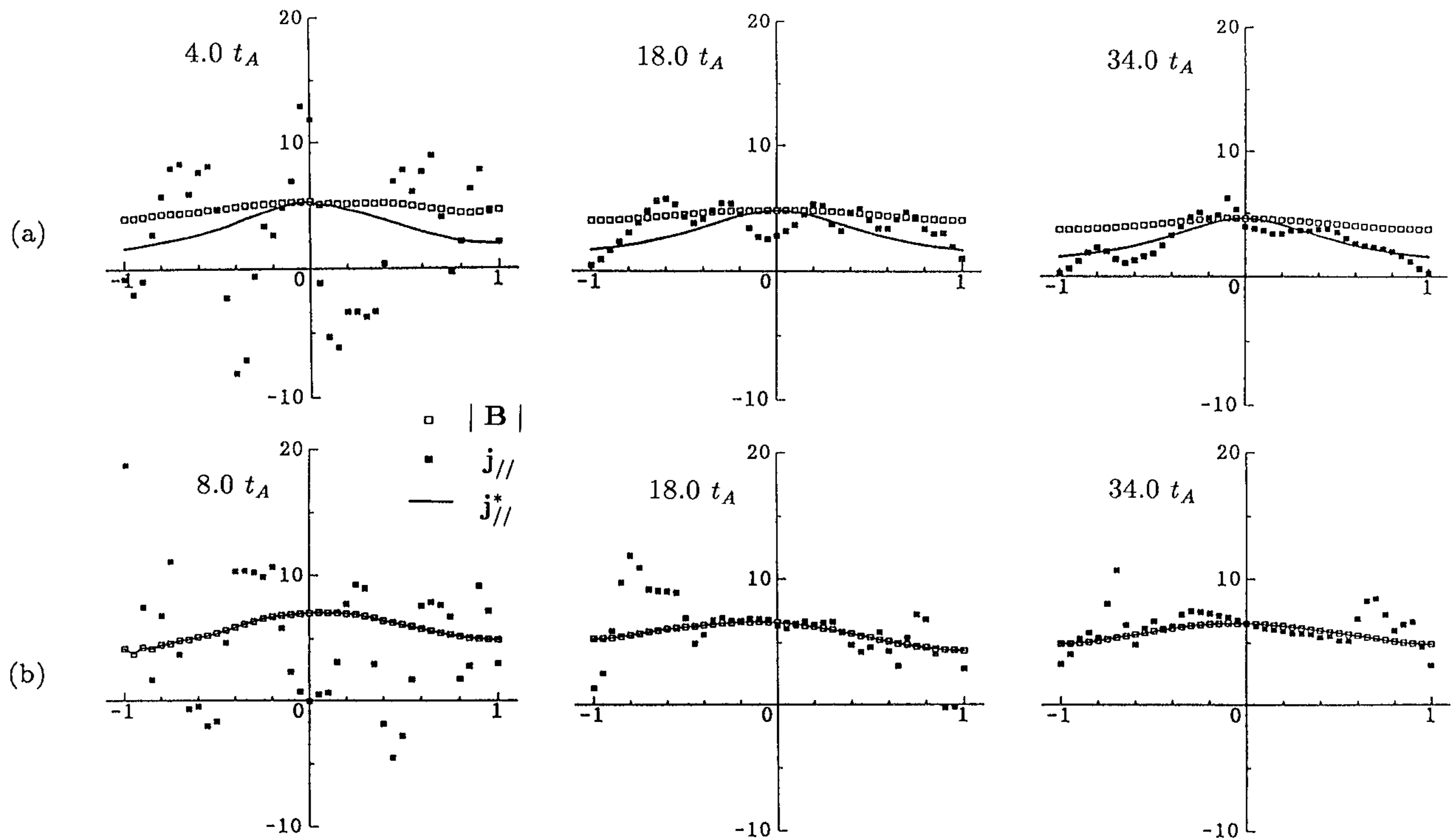

Fig.3

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