ACCELERATION OF PLASMA FLOWS DUE TO REVERSE DYNAMO MECHANISM

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ABSTRACT

The “reverse dynamo” mechanism—the amplification/generation of fast plasma flows by microscale (turbulent) magnetic fields via magneto-fluid coupling—is recognized and explored. It is shown that macroscopic magnetic fields and flows are generated simultaneously and proportionately from microscopic fields and flows. The stronger the microscale driver is, the stronger the macroscopic products will be. Stellar and astrophysical applications are suggested.

Subject headings: acceleration of particles — instabilities — ISM: magnetic fields — magnetic fields — MHD — stars: atmospheres — stars: chromospheres — stars: coronae — stars: magnetic fields

Online material: color figures

1. INTRODUCTION

The generation of macroscopic magnetic fields from (primarily microscopic) velocity fields defines the standard dynamo mechanism. The dynamo action seems to be a very pervasive phenomenon; in fusion devices as well as in astrophysics (in stellar atmospheres and MHD jets) one sees the emergence of macro-scale magnetic fields from an initially turbulent system. The relaxation observed in the reverse field pinches is a vivid illustration of the dynamo in action. The search for interactions that may result in efficient dynamo action is one of the most flourishing fields in plasma astrophysics. The myriad phenomena taking place in the stellar atmospheres (heating of the corona, the stellar wind, etc.) could hardly be understood without knowing the origin and nature of the magnetic field structures weaving the corona.

The conventional dynamo theories concentrate on the generation of macroscopic magnetic fields in charged fluids. With time, the dynamo theories have invoked more and more sophisticated physics models—from the kinematic to the magnetohydrodynamic (MHD) to, more recently, the Hall MHD (HMHD) dynamo. In the latter theories the velocity field is not specified externally (as it is in the kinematic case) but evolves in interaction with the magnetic field. Naturally, both MHD and HMHD dynamo theories encompass, in reality, the simultaneous evolution of the magnetic and the velocity fields. If the short-scale turbulence can generate long-scale magnetic fields, then under appropriate conditions the turbulence could also generate macroscopic plasma flows. In this context, the results of Socas-Navarro & Manso Sainz (2005; see also Bellot Rubio et al. 2001; Blackman 2005) are rather pertinent: the structures/magnetic elements produced by the turbulent amplification are destroyed/dissipated even before they are formed completely, creating significant flows or leading to heating.

If the process of conversion of microscale kinetic energy to macroscopic magnetic field (flow) independent of the mix of the microscale energy (magnetic and kinetic).

Within the framework of a simple HMHD system, we demonstrate in this paper that the dynamo and reverse dynamo processes operate simultaneously: whenever a macroscopic magnetic field is generated, there is a concomitant generation of a macroscopic plasma flow. Whether the macroscopic flow is weak (sub-Alfvénic) or strong (super-Alfvénic) with respect to the macroscopic field depends on the composition of the turbulent energy. We derive the relationships between the generated fields and the flows and discuss the conditions under which one or the other process is dominant. In § 2 we display an analytical calculation based on the conversion of microscale magnetic and kinetic energy into macroscopic fields and flows. In particular, we dwell on the reverse dynamo mechanism: the permanent dynamical feeding of the flow kinetic energy through an interaction of the macroscopic magnetic field structures with weak flows (seed kinetic energy). In § 3 we illustrate that the theoretically derived processes do indeed take place by presenting simulation results from a general two-fluid code that includes dissipation.

2. THEORETICAL MODEL ANALYSIS

The physical model exploited for flow generation/acceleration is simplified HMHD—a minimal model that contains two interacting scales that can be quite disparate; the macroscopic scale of the system is generally much larger than the ion skin depth, the intrinsic microscale of HMHD at which ion kinetic inertia effects become important (Mahajan & Yoshida 1998; Mahajan et al. 2001; Ohsaki et al. 2002; Yoshida et al. 2004). In HMHD the ion ($v_i$) and electron ($v_e = v - f/e$) velocity flows are different even in the limit of zero electron inertia. In its dimensionless form, HMHD comprises

$$\frac{\partial b}{\partial t} = \nabla \times [(v - \alpha_0 \nabla \times b) \times b], \quad (1)$$

$$\frac{\partial v}{\partial t} = v \times (\nabla \times v) + (\nabla \times b) \times b - \nabla \left( \frac{p + v^2}{2} \right), \quad (2)$$
with the standard normalizations: the density $n$ is normalized to $n_0$, the magnetic field to the same measure of the ambient field $B_0$, and velocities to the Alfvén velocity $V_{A0}$. We assign equal temperatures to the electrons and the protons so that the kinetic pressure $p$ is given by $p = p_e + p_n \simeq 2nT$, where $T = T_e \simeq T_n$. We note that the Hall current contributions become significant when the dimensionless Hall coefficient $\alpha_0 = \lambda_0/R_0$ ($R_0$ is the characteristic scale length of a system, and $\lambda_0 = c/\omega_0$ is the collisionless skin depth) satisfies the condition $\alpha_0 \gg \eta$, where $\eta$ is the inverse Lundquist number for the plasma. For a typical solar plasma, in the corona, the chromosphere, and the transition region, this condition is easily satisfied \[ \frac{\alpha_0}{\alpha_0} \approx 10^{-7} \] for densities within $10^{14} \sim 10^{18}$ cm$^{-3}$ and $\eta = c^2/(4\pi V_{A0} R_0 \sigma) \approx 10^{-14}$, where $R_0$ is solar radius and $\sigma$ is the plasma conductivity. In such circumstances, the Hall currents modify the dynamics of the macroscopic flows and fields could have a profound impact on the generation of macroscopic magnetic fields (Mininni et al. 2003) and fast flows (Mahajan et al. 2002, 2005).

In the following analysis, $\alpha_0$ is absorbed by choosing the normalizing length scale $\lambda_0$. Let us now assume that our total fields are composed of some ambient seed fields and fluctuations about them,

$$b = H + b_0 + \tilde{b}, \quad v = U + v_0 + \tilde{\nu},$$

where $b_0$ and $v_0$ are the equilibrium fields, $H$ and $U$ are the macroscopic fluctuations, and $\tilde{b}$ and $\tilde{\nu}$ are the microscopic fluctuations.

Note that our ambient fields are allowed to have a component at a microscopic scale. For analytical work, we choose for the ambient fields a special class of equilibrium solutions to equations (1) and (2). These solutions, also known as the double Beltrami (DB) pair (Mahajan & Yoshida 1998), come into existence because of the interaction of flows and fields; the Hall term is essential for their formation. The DB configurations are known to be robust and accessible, through a variational principle, for a variety of conditions including inhomogeneous densities. Non-constant-density cases display many interesting phenomena (Mahajan et al. 2002, 2005), but the dynamo and reverse dynamo actions can be very adequately described by the analytically tractable constant-density system. We therefore choose the DB pair (obeying the concomitant Bernoulli condition $V(p_0 + \rho_0^2/2) = \text{const}$; Ohsaki et al. 2001, 2002)

$$\frac{b_0}{a} + \nabla \times b_0 = v_0, \quad b_0 + \nabla \times v_0 = dv_0$$

as a representative ambient state. The general solution is expressible in terms of the single Beltrami fields $G_{\pm}$ that satisfy $\nabla \times G(\lambda) = \lambda G(\lambda)$:

$$b_0 = C_+ G_+(\lambda_+) + C_- G_-(\lambda_-),$$

$$v_0 = (a^{-1} + \lambda_+)C_+ G_+(\lambda_+) + (a^{-1} - \lambda_-)C_- G_-(\lambda_-).$$

Here $C_{\pm}$ are arbitrary constants, and the parameters $a$ and $d$ are set by the invariants of the equilibrium system: the magnetic helicity $h_{00} = \int (A_0 \cdot b_0) d^3x$ and the generalized helicity $h_{20} = \int (A_0 + v_0) \cdot \nabla \times (A_0 + v_0) d^3x$ (Mahajan & Yoshida 1998; Mahajan et al. 2001); here $A_0$ is the vector potential of the ambient field. The inverse scale lengths $\lambda_+$ and $\lambda_-$ are fully determined in terms of $a$ and $d$: $\lambda_{\pm} = \frac{1}{2} \left( (d - a^{-1}) \pm \sqrt{(d - a^{-1})^2 - 4} \right)^{1/2}$. As the DB parameters $a$ and $d$ vary, $\lambda_{\pm}$ can range from real to complex values of arbitrary magnitude.

Our primary interest is the creation of macrofields from the ambient microfields. Somewhat later we assume, for simplicity, that our zeroth-order fields are wholly at the microscopic scale. This allows us to create a hierarchy in the microfields: the ambient fields are much greater than the fluctuations at the same scale ($|\tilde{b}| \ll |b_0|$ and $|\tilde{\nu}| \ll |v_0|$). Following Mininni et al. (2003), we derive the evolution equations,

$$\frac{\partial b}{\partial t} = U \times (\nabla \times U) + \nabla \times H \times H + (v_0 \times (\nabla \times \tilde{\nu})) = (\nabla \times (\nabla \times \tilde{\nu})) - \frac{1}{2}(U^2),$$

$$\frac{\partial v}{\partial t} = (H \cdot \nabla) v_0 - (U \cdot \nabla) b_0,$$

$$\frac{\partial H}{\partial t} = \nabla \times (v_0 \times b_0) + v_0 \times \tilde{b} + \nabla \times (U \times (\nabla \times H) \times H),$$

where the angle brackets denote the spatial averages and $v_{0,0} = v_0 - \nabla \times b_0$. This set of equations can be regarded as a closure model of the Hall MHD equations, which are now general in two respects: (1) it is a closure of the full set of equations, since the feedback of the microscale is consistently included in the evolution of both $H$ and $U$, and (2) the role of the Hall current (especially in the dynamics of the microscale) is also properly accounted for (see Mininni et al. 2003, 2005 for details).

We now choose the constants $a$ and $d$ so that the two Beltrami scales become vastly separated (since these constants reflect the values of the invariant helicities, it is through $a$ and $d$ that the helicities control the final results). In the astrophysically relevant regime of disparate scales (the size of the structure is much greater than the ion skin depth), we deal with two extreme cases: (1) $a \sim d \gg 1$ and $(a - d)/ad \ll 1$ ($\lambda_+ \sim \mu$ and $\mu \sim (a - d)/ad$), and (2) $a \sim d \ll 1$ and $(a - d)/a \gg 1$ ($\lambda_+ \sim a - a^{-1}$ and $\mu \sim d - a$). At this time, we would like to draw the reader’s attention to the origin of scale separation in the original equilibrium system—it is the Hall term that imposes the macroscopic (ion skin depth) on the macroscopic MHD equilibrium.

Consistent with the main objectives of this paper, we now assume that the original equilibrium is predominantly microscale (the condition applicable for many astrophysical systems), i.e., the basic reservoir from which we generate macroscale fields is, indeed, at a totally different scale. Neglecting the macroscale component altogether, the assumed equilibria become simpler, with the velocity and magnetic fields linearly related as

$$v_0 = b_0 (\lambda + a^{-1}),$$

leading to

$$v_{0,0} = v_0 - \nabla \times b_0 = b_0 a^{-1},$$

$$\tilde{b} = [H - (\lambda + a^{-1})U] \cdot \nabla b_0.$$
Note the preponderance of nonlinear terms in the evolution equations for $U$ and $H$. One would expect that these terms will certainly play a very important part in the eventual saturation of the macroscopic fields, but in the early acceleration stage, when the ambient short-scale energy is much greater than the newly created macroscopic energy, these terms will not be significant. Deferring the fully nonlinear terms to a later stage, we limit ourselves to a "linear" treatment here. Neglecting the nonlinear terms and manipulating the system of equations, we readily derive (after "solving" for and eliminating the short-scale fluctuating fields)

$$
\ddot{H} \simeq \left(1 - \frac{\lambda}{a} - \frac{1}{a^2}\right)(\nabla \times (H \cdot \nabla)b_0 \times b_0),
$$

(15)

$$
\ddot{U} \simeq \left(\frac{\lambda}{a} + \frac{1}{a^2}\right)(\dot{\lambda} \hat{b} - \nabla \times \hat{v}) - \left(\frac{\lambda}{a} + \frac{1}{a^2}\right)(\nabla (b_0 \cdot \hat{v}) \times b_0)
- \left(\frac{\lambda}{a} \hat{b} - \nabla \times \hat{b}\right) \times b_0),
$$

(16)

where the spatial averages are yet to be performed. We use the standard isotropic ABC solution of the single Beltrami system,

$$
b_{0x} = \frac{b_0}{\sqrt{3}} (\sin \lambda y + \cos \lambda z),
$$

$$
b_{0y} = \frac{b_0}{\sqrt{3}} (\sin \lambda z + \cos \lambda x),
$$

$$
b_{0z} = \frac{b_0}{\sqrt{3}} (\sin \lambda x + \cos \lambda y),
$$

(17)

to compute the spatial averages. After some tedious but straightforward algebra, we arrive at the final acceleration equations,

$$
\ddot{U} = \frac{\lambda b_0^2}{3} \nabla \times \left\{ \left[\left(\frac{\lambda}{a} + \frac{1}{a^2}\right)^2 - 1\right] U - \lambda H \right\},
$$

(18)

$$
\ddot{H} = -\frac{\lambda b_0^2}{3} \left(1 - \frac{\lambda}{a} - \frac{1}{a^2}\right) \nabla \times H,
$$

(19)

where $b_0^2$ measures the ambient microscale magnetic energy (also the kinetic energy, because of eq. [11]). The coefficients in these equations are determined by $a$ and $d$ ($\lambda = \lambda(a,d)$).

We see that, to leading order, $H$ evolves independently of $U$, but the reverse is not true: the evolution of $U$ does require knowledge of $H$.

In the dynamo context, the Hall currents in the microscale are known to modify the $\alpha$-coefficient so that it survives the standard cancellation of the kinetic and magnetic contributions for Alfvénic perturbations (Mininni et al. 2002). It is also known that, depending on the state of the system, the Hall effect (by replacing the bulk kinetic helicity by the electron flow helicity) can cause a large enhancement or suppression of the dynamo action as compared to the standard MHD (Mininni et al. 2005).

Writing equations (18) and (19) as

$$
\ddot{H} = -r(\nabla \times H), \quad \ddot{U} = \nabla \times (sU - qH),
$$

(20)

where

$$
r = \frac{\lambda b_0^2}{3} (1 - \lambda a^{-1} - a^{-2}), \quad s = \frac{\lambda b_0^2}{6} \left[\left(\lambda + a^{-1}\right)^2 - 1\right],
$$

$$
q = \frac{\lambda^2 b_0^2}{6},
$$

(21)

and performing a Fourier analysis, one obtains

$$
-\omega^2 H = -ir(k \times H), \quad -\omega^2 U = ik \times (sU - qH),
$$

(22)

yielding the growth rate,

$$
\omega^4 = r^2 k^2, \quad \omega^2 = -|r|(k),
$$

(23)

at which $H$ and $U$ increase. The growing macrofields are related to each other by

$$
U = \frac{q}{s + r} H.
$$

(24)

We now show how a choice of $a$ and $d$ fixes the relative amounts of microscopic energy in the ambient fields and consequently in the nascent macroscopic fields $U$ or $H$. We continue with our two extreme cases.

1. For $a \sim d \gg 1$, the inverse microscale $\lambda \sim a \gg 1$, implying $v_0 \sim ab_0 \gg b_0$, i.e., the ambient microscale fields are primarily kinetic. These types of conditions may be met in stellar photospheres, where the turbulent velocity field at some stage can be dominant, although some $b_0$ is present as well. For these parameters, it can easily be seen that the generated macrofields have precisely the opposite ordering, $U \sim a^{-1} H \ll H$. This is an example of the straight dynamo mechanism. Microscale fields with kinetic dominance create, preferentially, macroscale fields that are magnetically dominant—super-Alfvénic "turbulent flows" lead to steady flows that are equally sub-Alfvénic (remember that we are using Alfvénic units). It is extremely important, however, to emphasize that the dynamo effect (dominant in this regime) must always be accompanied by the generation of macroscale plasma flows. This realization can have serious consequences for defining the initial setup for the later dynamics in the stellar atmosphere. The presence of an initial macroscale velocity field during the flux emergence processes is, for instance, always guaranteed by the mechanism exposited above. The implication is that all models of chromosphere heating/particle acceleration should take into account the existence of macroscopic primary plasma flows (even weak ones) and their self-consistent coupling (see Mahajan et al. 2001, 2005; Ohsaki et al. 2002 and references therein).

2. For $a \sim d \ll 1$, the inverse microscale $\lambda \sim a - a^{-1} \gg 1$. Consequently, $v_0 \sim ab_0 \ll b_0$, and the ambient energy is mostly magnetic. These conditions might pertain to certain domains in the photospheres or chromospheres, where a turbulent velocity field may exist, but the turbulent magnetic field is the dominant component. This macroscale magnetically dominant initial system creates macroscale fields $U \sim a^{-1} H \gg H$ that are kinetically abundant. The situation is fully reversed from the one discussed in the previous example—starting from a strongly sub-Alfvénic turbulent flow, the system generates a strongly super-Alfvénic macro-scale flow; this mode of conversion could be called the reverse dynamo mechanism. In the region of a given astrophysical system where the fluctuating/turbulent magnetic field is initially dominant, the magneto-fluid coupling induces efficient/significant acceleration, and part of the magnetic energy will be transferred to steady plasma flows. The eventual product of the reverse dynamo mechanism is a steady super-Alfvénic flow—a macro flow accompanied by a weak magnetic field. (Compare this with Blackman & Field [2004] for a magnetically driven dynamo. In this study magnetic field growth on much larger scales and significant velocity fluctuations with finite volume-averaged kinetic
helicity are found.) It is tempting to stipulate that the reverse dynamo may be the explanation for observations that fast flows are generally found in weak-field regions of the solar atmosphere (Woo et al. 2004).

This simple analysis has led to what we believe are several far-reaching results. (1) The dynamo and reverse dynamo mechanisms have the same origin: they are manifestations of the magneto-fluid coupling. (2) The proportionality of $U$ and $H$ implies that they must be present simultaneously, and the greater the macroscale magnetic field is (generated locally), the greater the macroscale velocity field will be (generated locally). (3) The growth rate of the macroscale fields is defined by DB parameters (hence, by the ambient magnetic and generalized helicities) and scales directly with the ambient turbulent energy $C^2_k u^2_0$. Thus, the larger the initial turbulent (microscopic) magnetic energy is, the stronger the acceleration of the flow will be. We believe that these novel results will surely help in advancing our understanding of the evolution of large-scale magnetic fields and their opening up, with respect to the fast particle escape from the stellar coronae. This effect may also have an important impact on the dynamical and continuous kinetic energy supply of plasma flows observed in various astrophysical systems. We add here that in this study both the initial and final states have finite helicities (magnetic and kinetic). The helicity densities are dynamical parameters that evolve self-consistently during the process of flow generation. It is also important to notice that the end product of the reverse dynamo action is a macroscopic flow (produced from a microscopic helical magnetic field), while for the inverse dynamo (Blackman & Field 2004) it is still the macroscopic magnetic field, but produced from a velocity field with helicity.

We end this section with a remark on the nonlinear terms in equations (7) and (10) that do not appear later in our analysis. It is amazing that the linear solution given in equations (22)–(24) makes the nonlinear terms strictly zero. Thus, the solution discussed in this section is an exact (in a special class) solution of the nonlinear system and thus remains valid even as $U$ and $H$ grow to larger amplitudes. This interesting but peculiar property that a basically linear solution solves the nonlinear problem pertains to both MHD and HMHD. In MHD, for example, it manifests itself as Waleń’s nonlinear Alfvén wave (Waleń 1944, 1945), while in HMHD it is revealed through the recently discovered solution of Mahajan & Krishan (2005).

3. A SIMULATION EXAMPLE

In order to strengthen and support the conclusions of the simple analytical model, we now present some representative results from our 2.5 dimensional numerical simulation of the general two-fluid equations in Cartesian geometry (Mahajan et al. 2001). For a description of the code, Mahajan et al. (2005) should also be consulted. The simulation system is somewhat different because of the existence of an ambient embedding macroscale field. We find that when such a field is present, the basic qualitative features of the dynamo and reverse dynamo mechanisms do not change much, but the algebra is considerably more complicated and will be presented in a longer paper.

The simulation system contains several effects not included in the analysis; it has, for instance, dissipation and heat flux in addition to the vorticity and the Hall terms. The plasma is taken to be compressible and embedded in a gravitational field; this provides an extra possibility for microscale structure creation. Transport coefficients for heat conduction and viscosity are taken from Braginski (1965).

The simulation presented here deals with the trapping and amplification of a primary flow impinging on a single closed-line

![Fig. 1.—Top: Contour plot for the $y$-component of the vector potential $A_y$ (flux function) in the $x$-$z$ plane for an initial distribution of ambient arcade-like magnetic field. The field has a maximum $B_{max}(x_0 = 0, z_0 = 0) = 100$ G. Middle: Initial symmetric profiles of the radial velocity $V_r$ and density $n$. The respective maxima (at $x = 0$) are $\sim 2$ km s$^{-1}$ and $10^{12}$ cm$^{-3}$. Bottom: Time evolution of the initial flow: $V_r(t, z = 0) = V_{in} \sin (\pi t t_0), V_r(t > t_0) = 0$, and $t_0 = 100$ s. [See the electronic edition of the Journal for a color version of this figure.]
structure. The choice of initial conditions is guided by observational evidence (Aschwanden et al. 2001; Woo et al. 2004) of the self-consistent process of acceleration and trapping/heating of plasma particles in the finely structured solar atmosphere. The simulation begins with a weak symmetric upflow (initially Gaussian, \( n_0 \approx 10^3 \) cm\(^{-3} \)) and a vertically distributed magnetic field (location of field maximum \( B_{0z} = 100 \) G; see the top part of Fig. 1 for the vector potential [flux function] defining the two-dimensional arcade). This field was assumed to be initially uniform in time. The magnetic field is represented as \( B = \nabla \times \mathbf{A} + B_0 \mathbf{z} \), with \( \mathbf{A}(0, A_z, 0) \), \( \mathbf{b} = B_0 \mathbf{z} \), and \( \mathbf{b}_s(x, z \neq 0) \neq 0 \). From numerous runs on the flow field evolution, we have chosen to display the results corresponding to the initial and boundary flow parameters \( V_{0max}(x, z = 0) = V_{dc} = 2.18 \times 10^3 \) cm s\(^{-1} \), \( n_{0max} = 10^{12} \) cm\(^{-3} \), and \( T(x, z = 0) = \text{const} = T_0 = 10 \) eV. The background plasma density is \( n_0 = 0.2n_{0max} \). In simulations, \( n(x, z, t = 0) = n_0n_{0max} \) is an exponentially decreasing function of \( z \). Experience was a guide for imposing the boundary condition \( \partial_z K(x = \pm \infty, z, t) = 0 \), which was used with sufficiently high accuracy for all parameters \( K(A, T, V, B, n) \). The initial velocity field has a pulselike distribution (middle and bottom panels of Fig. 1) with a time duration \( t_0 = 100 \) s. We find the following:

1. The acceleration is significant in the vicinity of the magnetic field maximum (originally present or newly created during the evolution) with strong deformation of field lines and energy redistribution due to magneto-fluid coupling and dissipative effects.

2. Initially, a part of the flow is trapped in the maximum field localization area, accumulated, cooled, and accelerated (plots corresponding to \( t = 100 \) s). The accelerated flow reaches speeds greater than 10 km s\(^{-1} \) in less than 100 s (in agreement with recent observations; Schrijver et al. 1999; Seaton et al. 2001; Ryutova & Tarbell 2003 and references therein).

3. After this stage the flow passes through a series of quasi-equilibria. In this relatively extended era (~1000 s) of stochastic/oscillating acceleration, the intermittent flows continuously

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**Fig. 2.—Contour plots in the x-z plane at three time frames, \( t = 100, 1000, \) and 2500 s, for the dynamical evolution of \( A_y \) (left column), \( n \) (second column), \( |V| \) (third column), and \( T \) (right column) for flow–arcade field interaction. The realistic viscosity and heat flux effects as well as the Hall term (\( \alpha_B = 3.3 \times 10^{-6} \)) are included in the simulation. The primary flow (the type displayed in Fig. 1) is accelerated as it makes its way through the magnetic field with an arcade-like structure (Fig. 1). The primary flow, locally sub-Alfvenic, is accelerated, reaching significant speeds (\( \geq 100 \) km s\(^{-1} \)) in a very short time (\( \leq 100 \) s). Initially, the effect is strong in the strong-field region (center of the arcade). There is a critical time (\( \leq 1000 \) s) at which the accelerated flow bifurcates; the original arcade field is deformed correspondingly. After the bifurcation, strong magnetic field localization areas, carrying currents, are created symmetrically about \( x = 0 \). Postbifurcation daughter flows are localized in the newly created magnetic field localization areas. The maximum density of each daughter flow is the order of the density of the mother flow. Daughter flows have distinguishable dimensions (\( \geq 0.05 R_s \)). At \( t \geq 1000 \) s, the velocities reach (\( \geq 500 \) km s\(^{-1} \)) or even greater (\( \leq 800 \) km s\(^{-1} \)) values. The distance from surface where bifurcation happens is \( \geq 0.01 R_s \). In the regions of daughter flow localization there is significant cooling, while the nearby regions are heated. [See the electronic edition of the Journal for a color version of this figure.]
acquire energy (see Fig. 3 for the flow kinetic and magnetic energy maxima and also Fig. 2 results at $t = 1000$ s).

4. The flow starts to accelerate again (Figs. 3a–3c for the velocity field evolution). This process is completely consistent with the analytical prediction; the acceleration is highest in the strong-field regions (newly generated, Fig. 2). At this moment the accelerated daughter flows (macroscale) are decoupled from the mother

flow, carrying currents and modifying the initial arcade field, creating new $b_{\text{max}}$ localization areas that span the region between $\lesssim 0.05$ and $\sim 0.01$ $R_\odot$ from the interaction surface.

The extensive simulation runs also show that when dissipation is present, the Hall term (proportional to $\alpha_0$), through the mediation of microscale physics, plays a crucial role in the

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**Fig. 3.**—Evolution of maximum values of (a) $|V|$, (b) $|V_p| = (V_x^2 + V_y^2)^{1/2}$, (c) $V_z$, (d) $|b|$, (e) $|b_p| = (b_x^2 + b_y^2)^{1/2}$, and (f) $b_z$ in time. (a, d) It is shown that much of the transfer from magnetic field energy happens during the first, very fast ($\sim 100$ s) acceleration stage. (b, e) Later, the dissipation of perpendicular (toward height) magnetic field fluctuations leads to the maintenance of the quasi–equilibrium fast perpendicular flows for a period of $\sim 1000$ s, and then the effective acceleration of flow follows. (c, f) The maximum value of the magnetic field component along the height is not changed, and the radial component of the velocity field dissipates effectively. It should be emphasized that these maximum values of both field parameters change the localization dynamically and follow the relationship found analytically: fast flows (see Fig. 2) are observed in the regions of macroscale magnetic field maximum localization (initially given or later generated).
acceleration/heating processes. The existence of initial fast acceleration in the region of maximum localization of the original magnetic field and the creation of new areas of macroscale magnetic field localization (Fig. 2, first column for $A_y$) with simultaneous transfer of the magnetic energy (oscillatory, microscale) to flow kinetic energy (Fig. 2, third column for $|V|$, and Fig. 3 results) are manifestations of the combined effects of the dynamo and reverse dynamo phenomena. The maintenance of quasi-steady flows for a rather significant period is also an effect of the continuous energy supply from fluctuations (due to the dissipative Hall and vorticity effects). These flows are likely to provide a very important input element for understanding the finely structured atmospheres with their richness of dynamical structures as well as for the mechanisms of heating and possible escape of plasmas.

Note that in the simulation the actual magnetic and generalized helicity densities are dynamical parameters. Thus, even if they are not in the required range initially, their evolution could bring them into the range in which they could satisfy the conditions needed to efficiently generate flows. The required conditions could be met at several stages. This could, perhaps, explain the existence of several phases of acceleration. Dissipation effects could play a fundamental role in setting up these distinct stages; it could, for example, modify the generalized vorticity, which will finally lead to a modification of field lines and even to the creation of microscale fields (shocks or fast fluctuations).

4. CONCLUSIONS

From an analysis of the two-fluid equations, in this paper we have extracted the reverse dynamo mechanism—the amplification/generation of fast plasma flows in astrophysical systems with initial turbulent (microscale) magnetic fields. This process is simultaneous with and complementary to the highly explored dynamo mechanism. It is found (both analytically and numerically) that the generation of macroscale flows is an essential consequence of the magneto-fluid coupling and is independent of the initial and boundary conditions. The generation of micro-/macroscale magnetic fields and generation of macroscale flows go hand in hand; the greater the macroscale magnetic field (generated locally) is, the greater the macroscale velocity field (generated locally) will be. The acceleration due to the reverse dynamo is directly proportional to the initial turbulent magnetic energy. When the microscopic magnetic field is initially dominant, a major part of its energy transforms to macroscale flow energy; a weak macroscale magnetic field is generated along with it.

The reverse dynamo mechanism, providing an unfailing source for macroscale plasma flows, is likely to be an important mechanism for understanding a host of phenomena in astrophysical systems.

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