SPONTANEOUS FORMATION OF SURFACE MAGNETIC STRUCTURE FROM LARGE-SCALE DYNAMO IN STRONGLY STRATIFIED CONVECTION

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Received 2015 December 28; revised 2016 April 10; accepted 2016 April 18; published 2016 May 2

ABSTRACT

We report the first successful simulation of spontaneous formation of surface magnetic structures from a large-scale dynamo by strongly stratified thermal convection in Cartesian geometry. The large-scale dynamo observed in our strongly stratified model has physical properties similar to those in earlier weakly stratified convective dynamo simulations, indicating that the \(\alpha^2\)-type mechanism is responsible for the dynamo. In addition to the large-scale dynamo, we find that large-scale structures of the vertical magnetic field are spontaneously formed in the convection zone (CZ) surface only in cases with a strongly stratified atmosphere. The organization of the vertical magnetic field proceeds in the upper CZ within tens of convective turnover time and band-like bipolar structures recurrently appear in the dynamo-saturated stage. We consider several candidates to be possibly be the origin of the surface magnetic structure formation, and then suggest the existence of an as-yet-unknown mechanism for the self-organization of the large-scale magnetic structure, which should be inherent in the strongly stratified convective atmosphere.

Key words: convection – dynamo – magnetohydrodynamics (MHD) – Sun: magnetic fields – sunspots

1. INTRODUCTION

A longstanding goal of solar interior physics is to self-consistently reproduce active regions, composed mainly of sunspots, from magnetic fluxes generated in the solar interior. We now approach the subject from two different theoretical perspectives: one perspective focuses on the emergence and organization processes of the magnetic flux in the uppermost part of the convection zone (CZ), and the other explores the flux generation and maintenance processes, i.e., the dynamo process, operating deeper down.

Several leading-edge numerical studies that have focused on the uppermost part of the solar CZ have succeeded in simulating spontaneous formations of concentrated magnetic structures reminiscent of active regions (e.g., Cheung et al. 2010; Stein & Nordlund 2012; Rempel & Cheung 2014; Käpylä et al. 2016). In these studies, the solar surface convection and its nonlinear interaction with the magnetic field were simulated in a more or less realistic manner, with the steep density gradient just below the photosphere, and/or the radiative transfer with the ionization in Cartesian domains. However, since some sort of the large-scale seed magnetic field has been inconsistently assumed to be an initial or boundary condition, the dynamo mechanism and its connection to the formation process of the active region were beyond the scope of these studies.

A growing body of evidence is demonstrating that solar-like cyclic large-scale magnetic fields are organized in global spherical-shell convections (e.g., Ghizaru et al. 2010; Käpylä et al. 2012; Masada et al. 2013; Augustson et al. 2015; Yadav et al. 2015). Despite some differences in the numerical setup and method, there is a common outcome for the convective dynamo in these studies: diffuse magnetic flux extending over the CZ and/or the tachocline instead of magnetic flux tubes expected in the standard solar dynamo paradigm (e.g., Charbonneau 2010, and the references therein). Although the flux emergence-like event from distributed magnetic flux has been occasionally observed in some models (Nelson et al. 2013; Fan & Fang 2014), its universality or feasibility in the Sun is still a matter of considerable debate.

There is still a large gap between the dynamo in the interior and the active region formation at the surface. Our study in this Letter would be a first step for bridging the gap between them. By advancing our previous works on weakly stratified MHD convection (Masada & Sano 2014a, 2014b, hereafter MS14a, MS14b), we perform a convective dynamo simulation in a strongly stratified atmosphere resembling the solar interior in Cartesian geometry. The spontaneous formation of large-scale magnetic structures in the CZ surface self-consistently from the large-scale convective dynamo is reported.

2. NUMERICAL SETUP

A convective dynamo system is solved numerically in a Cartesian domain. Our model covers only the CZ of depth \(d_{cz}\) (0 \(\leq z \leq d_{cz}\)), omitting a stably stratified layer below it, where the \(x\)- and \(y\)-axes are taken to be horizontal, and the \(z\)-axis is pointing downward. We set the width of the domain to be \(W = 4d_{cz}\).

We solve the fully compressible MHD equations in the rotating frame of reference with a constant angular velocity of \(\mathbf{\Omega} = -\Omega_0 \mathbf{e}_z\),

\[
\frac{\partial \rho}{\partial t} = -\nabla \cdot (\rho u),
\]

\[
\frac{\mathbf{D}u}{\mathbf{D}t} = -\nabla P + \frac{\mathbf{J} \times \mathbf{B}}{\rho} - 2\mathbf{\Omega} \times \mathbf{u} + \frac{\nabla \cdot \mathbf{\Pi}}{\rho} + \mathbf{g},
\]

\[
\frac{\mathbf{D}e}{\mathbf{D}t} = -\frac{P}{\rho} \nabla \cdot \mathbf{u} + Q_{\text{heat}},
\]

\[
\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{u} \times \mathbf{B} - \eta_0 \mathbf{J}),
\]
with the viscous stress $\mathbf{\Pi}$ and the heating term $Q_{\text{heat}}$ of

$$
P_{ij} = 2 \rho \nu_0 S_{ij} = \rho \nu_0 \left( \frac{\partial u_i}{\partial x_j} + \frac{\partial u_j}{\partial x_i} - \frac{2}{3} \frac{\partial P}{\partial x_i} \right),
$$

$$
Q_{\text{heat}} = \frac{\gamma \nabla \cdot (\kappa_0 \nabla \epsilon)}{\rho} + 2 \rho \nu_0 S^2 + \frac{\mu_0 \eta_0}{\rho} \epsilon,
$$

where $D/Dt$ is the total derivative, $\epsilon = c_v T$ is the specific internal energy, $J = \nabla \times B/\mu_0$ is the current density with the vacuum permeability $\mu_0$, and $g = g_0 e_z$ is the gravity of the constant $g_0$. The viscosity, magnetic diffusivity, and thermal conductivity are represented by $\nu_0$, $\kappa_0$, and $\kappa_0$, respectively. A perfect gas law $P = (\gamma - 1) \rho c^2$ with $\gamma = 5/3$ is assumed.

The initial hydrostatic state is polytropic stratification given by

$$
\epsilon = \epsilon_0 + \frac{g_0 z}{(\gamma - 1)(m + 1)}, \quad \rho = \rho_0 (\epsilon/\epsilon_0)^m,
$$

with the initial surface internal energy $\epsilon_0$, surface density $\rho_0$, and the polytropic index $m = 1.49$ providing the superadiabaticity of $\delta \equiv \nabla - \nabla_d = 1.6 \times 10^{-3}$, where $\nabla_d = 1 - 1/\gamma$ and $\nabla = (\partial \ln T/\partial \ln P)$.

Normalization quantities are defined by setting $d_{cz} = g_0 = \rho_0 = \mu_0 = c_p = 1$. The normalized pressure scale height at the surface, defined by $\xi = H_p/d_{cz} = (\gamma - 1) \epsilon_0 (g_0 d_{cz})$, controls the stratification level and is chosen here as $\xi = 0.01$, yielding a strong stratification with a density contrast between the top and bottom CZs of ~700.

Figure 1(a) shows the initial profiles of the density (solid) and temperature (dashed) of our model. The density profile in the range 0.71 \(\leq r/R \leq 0.991\) of the standard solar model is also shown in such a way as to fit the computational domain (dash–dotted; Model S; Christensen-Dalsgaard et al. 1996). Our model has a stratification that almost encompasses the solar CZ, except for its uppermost part.

All the variables are assumed to be horizontally periodic. Stress-free boundary conditions are used in the vertical direction for the velocity. Perfect conductor and vertical field conditions are used for the magnetic field at the bottom and top boundaries. A constant energy flux that drives the thermal convection is imposed on the bottom boundary, while the specific internal energy is fixed at the top boundary.

The fundamental equations are solved by the second-order Godunov-type finite-difference scheme that employs an approximate MHD Riemann solver (Sano et al. 1998). The magnetic field evolves with the Consistent MoC-CT method (Evans & Hawley 1988; Clarke 1996). Non-dimensional parameters $Pr = 20$, $Pm = 2$, and $Ra = 6 \times 10^7$, an angular velocity of $\Omega_0 = 0.5$, and the spatial resolution of $(N_r, N_z, N_p) = (256, 256, 256)$ are adopted, where the Prandtl, magnetic Prandtl, and Rayleigh numbers are defined by

$$
Pr = \frac{\nu_0}{(\kappa_0/\rho c_p)}, \quad Pm = \frac{\nu_0}{\eta_0}, \quad Ra = \frac{g_0 d_{cz}^2 \delta}{\chi_0 \nu_0 H_p},
$$

where $\rho$, $\delta$, and $H_p$ are evaluated at $z = d_{cz}/2$.

In the following, the volume-, horizontal-, $x$-, and $y$-averages are denoted by single angular brackets with subscripts “v,” “h,” “x,” and “y,” respectively. The time average of each spatial mean is denoted by additional angular brackets. The relative importance of convection to the rotation is measured by the Rossby number $Ro = (u_{cz} k)/2(\Omega_0)$, where $u_{cz} = \sqrt{\langle (\kappa_0 \nabla \epsilon)^2 \rangle}$ is the mean convective velocity, and $k_2 = 2\pi/d_{cz}$. The global convective turnover time and the equipartition field strength are defined by $\tau_{cv} \equiv 1/(u_{cw} k_2)$ and $B_{eq}(z) \equiv \sqrt{\langle (\mu_0 \eta_0 a^2) \rangle}$. Note that $B_{eq}(z)$ is evaluated from the local convective energy and thus has a depth-dependence.

Since the sound speed in the deep CZ becomes very large in the strongly stratified model and imposes a strict limit on the time-step, a long thermal relaxation time is required in our fully compressible simulation. To alleviate it, we first construct a progenitor model, in which the convection reaches a fully developed state and the system becomes thermally relaxed, by evolving a non-rotating hydrodynamic run for 800$\tau_{cv}$.

Shown in Figures 1(b) and (c) are the vertical profile of $\langle (\kappa_0 \nabla \epsilon)^2 \rangle$ (dashed) and the distributions of $u_z$ in the $y$–$x$ plane at $z/d_{cz} = 0.04, 0.23, 0.78$ in the equilibrated state of the progenitor model. The black (orange) tone denotes down-(up)-flows. The multi-scale convection with the strong up-down asymmetry, i.e., the slower and broader upflow cell surrounded by networks of faster and narrower downflow lanes, is developed in the progenitor model (e.g., Spruit et al. 1990; Miesch 2005). The dynamo run is started by adding the rotation and a seed weak horizontal field to the progenitor model.

### 3. SIMULATION RESULTS

#### 3.1. Basic Properties of Convection and Dynamo

The temporal evolutions of $\epsilon_K \equiv \langle (\kappa_0 \nabla \epsilon)^2 \rangle$, $\epsilon_M \equiv \langle B_z^2 \rangle$ (the energy of the mean magnetic components) are shown by solid, dashed, and dash–dotted lines in Figure 2(a). The evolution of the $\epsilon_K$ of the progenitor run before starting the dynamo run is also shown. Note that, from the horizontal symmetry and div $B = 0$, $(B_z)_h$ and $(B_x)_h$ are zero, independent of time.

After a short relaxation time, the convective kinetic energy reaches a quasi-steady state at $t \approx 50\tau_{cv}$. The magnetic energy is gradually amplified by the convection and is saturated at $t \approx 120\tau_{cv}$. The mean values evaluated there are $u_{cw} = 5.9 \times 10^{-3}$ and $B_{eq} = \sqrt{\langle (\mu_0 \eta_0 a^2) \rangle} = 7.2 \times 10^{-2}$, providing $\tau_{cv} = 54$ and $Ro = 0.02$. The vertical profile of $\langle (\kappa_0 \nabla \epsilon)^2 \rangle$ (solid) and the distributions of the $u_z$ on the horizontal planes in the dynamo-saturated stage are also shown in Figures 1(b) and (c). Since the rotation gives rise to the Coriolis force acting on the convective motion, the convective cell shrinks and thus the scale separation becomes larger in the rotating system. Since the mean kinetic helicity, which is a prerequisite for exciting the large-scale dynamo, arises as a natural consequence of the rotation, the dynamo-generated magnetic field also affects the convective motion shown in Figure 1(c).

Figure 2(b) shows the time-depth diagrams of $(B_z)_h$ and $(B_y)_h$ normalized by $B_{eq}(z)$. Note that the turbulent magnetic component is eliminated by taking a horizontal average. It is found that the oscillatory large-scale horizontal magnetic component is spontaneously organized in the bulk of the CZ. It has a peak with the super-equipartition strength in the mid-part of the CZ and propagates from there to the top and base of the CZ. Since there exists a phase difference of $\pi/2$ between $(B_z)_h$ and $(B_y)_h$, the mean horizontal magnetic flux, defined by $B_h \equiv \sqrt{(B_z)_h^2 + (B_y)_h^2}$, has a quasi-steady vertical profile.
The large-scale dynamo observed here in the strongly stratified model has physical properties similar to those in the weakly stratified convective dynamo simulations (e.g., Käpylä et al. 2013; MS14a). Because of the horizontal symmetry and thus no differential rotation in our system, the turbulent electromotive force would be solely responsible for the dynamo (see MS14b for a mean-field $\alpha^2$-dynamo model that can quantitatively reproduce the DNS results). Our intriguing finding in this Letter, which has not been observed in the weakly stratified model with similar boundary conditions, is a spontaneous formation of large-scale magnetic structures in the CZ surface, which will be reported in the following sections.

3.2. Spontaneous Formation of Surface Magnetic Structure

A series of snapshots where the distribution of the $B_z$ at different times on the horizontal cutting plane at $z/d_{cz} = 0.04$ is shown in the top panel of Figure 3(a). The darker (lighter) tone denotes positive (negative) $B_z$. While $B_z$ has a small-scale tangled structure with a typical size comparable to the convective cell in the initial evolutionary stage [(a1)–(a2)], it evolves as time passes to organize the large-scale structure to have spatial-scale much larger than the convective cell [(a3)–(a4)]. The surface magnetic structure has a dynamically important strength comparable to $B_{eq}(z)$ and recurrently appears in the dynamo-saturated stage [(a5)], implying that it should be
Figure 2. (a) The temporal evolutions of $\epsilon_k$ (solid), $\epsilon_M$ (dashed), and $\epsilon_{Mn}$ (dashed–dotted). (b) Time-depth diagrams of $\langle B_x \rangle_n / B_{eq}(z)$ and $\langle B_y \rangle_n / B_{eq}(z)$.

Figure 3. (a) A series of snapshots for the horizontal distribution of $B_z$ at $z/d_z = 0.04$ and the vertical distribution of the $\langle B_y \rangle_n / B_{eq}(z)$ or $\langle B_z \rangle_n / B_{eq}(z)$. (b) The temporal evolution of the 2D Fourier spectrum of the $B_z^2$ in the upper CZ, where $k_c = 2\pi/W$. 

The Astrophysical Journal Letters, 822:L22 (7pp), 2016 May 10

Masada & Sano
Figure 4. (a) The growth rate of the Parker instability evaluated from the simulation result. (b) Comparison between the growth time of the fastest-growing mode of the Parker instability ($\tau_{\text{pk,m}}$) and the local convective turnover time ($\tau_{\text{cv}}$) as a function of the depth.

3.3. What Makes Surface Magnetic Structure Organized?

The so-called negative magnetic pressure instability (NEMPI) has been proposed as a mechanism for the self-assembly of the magnetic flux near the CZ surface (e.g., Kleeorin et al. 1996; Brandenburg et al. 2010). Although its presence has been confirmed numerically in the forced MHD turbulence (Brandenburg et al. 2011, 2013; Warnecke et al. 2013; Mitra et al. 2014; Jabbari et al. 2014), it would not play a significant role in organizing the surface magnetic structure in our simulation. This is because relatively rapid rotation with $\text{Ro} = 0.02$ is assumed here, but according to Losada et al. (2012), $\text{Ro} \gtrsim 5$ is required to excite the NEMPI. Here two possible alternatives are discussed as a cause of the magnetic structure formation: one is the Parker instability (Parker 1979) and the other is the flux expulsion that accompanies the strong downflow (e.g., Weiss 1966; Kitiashvili et al. 2010; Stein & Nordlund 2012; Käpylä et al. 2016).

Figure 4(a) shows the growth rate of the Parker instability obtained from the WKB dispersion equation

\[
(c^2_s + v_A^2)\gamma^2 + v_A^2 \left[ 2(c^2_s + v_A^2)k_h^2 + \frac{g_0}{D_z} \ln \left( \frac{B_h}{\rho} \right) \right] \gamma^2 + k_h^2 v_A^2 \left( k_h^2 c_s^2 + \frac{g_0 D_z}{D_z} \ln B_h \right) = 0
\]

(e.g., Gilman 1970), where $\gamma$ is the growth rate, $k_h$ is the horizontal wavenumber, $B_h$ is the horizontal magnetic flux, $c_s = \sqrt{P/\rho}$, and $v_A = B_h/\sqrt{\mu_0 \rho}$. For deriving the depth-dependent growth rate, we adopt $P = \langle P \rangle_h$, $\rho = \langle \rho \rangle_h$, and $B_h = \sqrt{\langle B_x^2 \rangle_h + \langle B_y^2 \rangle_h}$ evaluated from the simulation model. The different line-type denotes the growth rate at different depths. The vertical and horizontal axes are normalized by $1/\tau_{\text{cv}}$ and $k_c = 2\pi/W$. Note that the instability is inhibited in the range $k/k_c < 1$ due to the box-width constraint.

The dynamo-maintained magnetic flux is unstable to the Parker instability in the span $0 \lesssim z/d_{\text{cz}} \lesssim 0.4$. Since the typical growth time of the Parker instability is comparable to or a bit smaller than $\tau_{\text{cv}}$, it has sufficient time to grow during the simulation. However, the most unstable mode has a smaller wavelength compared with the box-width, indicating the difficulty of explaining the dominance of the box-sized surface magnetic structure in our simulation.

In addition to the mismatch of the magnetic spatial scales, the growth of the Parker instability itself may be inhibited by vigorous convective motions. In Figure 4(b) the growth time of the most unstable mode of the Parker instability ($\tau_{\text{pk,max}}$)
and the local convective turnover time defined by $\tau_{cv,2}(z) = H_p(z)/\langle u_z \rangle_{k_h}^{1/2}$ (dashed) are compared as a function of the depth, where $H_p \equiv dz/d \ln \rho$. In the upper CZ, the condition $\tau_{cv,1} \ll \tau_{pk,max}$ is always satisfied. Since, in such a situation, the small-scale convective motion violently disturbs the coherency of the magnetic flux, we would have to say that the Parker instability would not be responsible for the large-scale structure formation observed in our simulation.

Next, the large-scale flow and its association with the surface magnetic structure are analyzed. For casting light on the large-scale pattern, the small-scale structures with $k/k_c \gtrsim 8$ are eliminated by applying Fourier filtering (e.g., Warnecke et al. 2016; Jabbari et al. 2016). A series of snapshots where $\langle B_z \rangle_{k_h}$ and $\langle u_z \rangle_{k_h}$ on the horizontal plane at $z/d_{cz} = 0.04$, are shown in Figure 5(a), where $\langle \cdot \rangle_{k_h}$ denotes Fourier filtering. The overplotted arrows are the velocity vectors composed of $\langle u_x \rangle_{k_h}$ and $\langle u_y \rangle_{k_h}$. Additionally, 2D spectra of $B_h^2$, $B_z^2$, $\rho v_h^2$, and $\rho v_z^2$ in the upper CZ are also shown in Figure 5(b). The spectrum at each depth is spatially averaged over the normalized depth from 0.0 to 0.25 and is temporally averaged over $10\tau_{cv}$ around the corresponding reference time.

It is significant that, in the dynamo-saturated stage, the bipolar "band-like" structure elongated along the direction of the horizontal magnetic flux is predominant (see Figure 2(b)). Although the faster horizontal flow and stronger downflow can be found in/around the region with stronger $B_z$ before the dynamo-saturation (left column), a large-scale flow pattern is not necessarily associated with the magnetic structure in the dynamo-saturated stage (middle and right columns). In addition, we can find from the spectra that the energy contained in the large-scale magnetic components is much larger than that of the large-scale flow in the upper CZ. This suggests that large-scale flows are not the cause, but a consequence of the large-scale magnetic structures in the upper CZ.

Figure 5. (a) A series of snapshots of $\langle B_z \rangle_{k_h}$ and $\langle u_z \rangle_{k_h}$ on the horizontal plane at $z/d_{cz} = 0.04$. The overplotted arrows are the velocity vectors. (b) 2D spectra of $B_h^2$, $B_z^2$, $\rho v_h^2$, and $\rho v_z^2$ in the upper CZ around the reference time $t_r$. 
Overall our analyses indicate that there should be an as-yet-unknown mechanism for the self-organization of large-scale magnetic structures, which would be inherent in the strongly stratified atmosphere. The band-like magnetic structure observed in our simulation is similar to that observed in a large-scale dynamo of forced turbulence in a strongly stratified atmosphere (Mitra et al. 2014; Jabbari et al. 2016). This may imply that surface magnetic structure formation is a common universal feature of strongly stratified MHD turbulence, regardless of its details.

4. SUMMARY

In this Letter, we studied numerically MHD convections in a strongly stratified atmosphere resembling the solar interior. The large-scale dynamo observed in our simulation had physical properties similar to those observed in the weakly stratified model (see MS14ab): the oscillatory large-scale horizontal magnetic component with the dynamically important strength was spontaneously organized in the bulk of the CZ. Its spatiotemporal evolution strongly suggests that the \( \alpha \)-type mechanism is responsible for the MHD dynamo in our system.

Our intriguing finding, which has not been observed in the weakly stratified model, was the spontaneous formation of the large-scale \( B_z \) structure in the CZ surface. Small-scale tangled components of \( B_z \), which was seen in the earlier evolutionary stage, gradually evolved into a large-scale organized structure with a size much larger than the convective cell as time passes in the upper CZ. In the dynamo-saturated stage, the bipolar “band-like” \( B_z \) structure was predominant and appeared recurrently. Since the possible candidates, the NEMPI, Parker instability, and flux expulsion due to the strong downflow, could not explain the surface magnetic structure formation without difficulty, our results may suggest the existence of an as-yet-unrecognized mechanism that induces the spontaneous formation of the large-scale magnetic structure in the strongly stratified convective atmosphere.

The large-scale dynamo observed here maintains the quasi-steady magnetic flux and thus is different from an actual solar dynamo with a quasi-periodic flux modulation. Furthermore, since we adopt a faster rotation than that achieved in the actual Sun and ignore the global effects, such as differential rotation and meridional flows, our model still remains a long way from the solar interior. However, despite a lack of some solar elements, it would be interesting if the relatively shallow root of the surface magnetic structure, implied from our simulation, turns out to be compatible with some observations of the magnetic patches on the solar surface (e.g., Brandenburg 2005; Hara 2009). Understanding the self-organization mechanism of the surface magnetic structure observed in our study should deepen our knowledge of solar magnetism and is a higher-priority of our future work.

We acknowledge the anonymous referee for constructive comments. Computations were carried out on XC30 at NAOJ. This work was supported by JSPS KAKENHI grant number 15K17611 and the joint research project of the Institute of Laser Engineering, Osaka University.

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