DYNAMO ACTIVITIES DRIVEN BY MAGNETOROTATIONAL INSTABILITY AND THE PARKER INSTABILITY IN GALACTIC GASEOUS DISKS

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ABSTRACT

We carried out global three-dimensional magnetohydrodynamic simulations of dynamo activities in galactic gaseous disks without assuming equatorial symmetry. Numerical results indicate the growth of azimuthal magnetic fields non-symmetric to the equatorial plane. As the magnetorotational instability (MRI) grows, the mean strength of magnetic fields is amplified until the magnetic pressure becomes as large as 10% of the gas pressure. When the local plasma β (= pgas/pmag) becomes less than 5 near the disk surface, magnetic flux escapes from the disk by the Parker instability within one period of rotation disk. The buoyant escape of coherent magnetic fields drives dynamo activities by generating disk magnetic fields with opposite polarity to satisfy the magnetic flux conservation. The fluctuation of the azimuthal magnetic flux from the disk and the subsequent amplification of disk magnetic field by the MRI drive quasi-periodic reversal of azimuthal magnetic fields on a timescale of 10 rotation periods. Since the rotation speed decreases with radius, the interval between the reversal of azimuthal magnetic fields increases with radius. The rotation measure computed from the numerical results shows symmetry corresponding to a dipole field.

Key words: dynamo – Galaxy: disk – magnetohydrodynamics (MHD)

Online-only material: color figures

1. INTRODUCTION

The origin and nature of magnetic fields in spiral galaxies is an open question. Magnetic fields in spiral galaxies have been intensively studied through observations of linear polarization of synchrotron radiation (Sofue et al. 1986; Beck et al. 1996). The typical magnetic field strength is a few μG in the Milky Way near the Sun and in spiral galaxies. The mean ratio of gas pressure to magnetic pressure is close to unity (Rand & Kulkarni 1989; Crocker et al. 2010), suggesting that magnetic fields in spiral galaxies are strong enough to affect the motion of the interstellar matter in gaseous disks.

Since the galactic gas disk rotates differentially, the magnetic field is amplified faster near the galactic center than in the outer disk, and the magnetic field strength can exceed 0.1 mG in the galactic center. Observational evidence of such strong magnetic fields has been obtained in the radio arc near the galactic center, where the field strength of about 1 mG is estimated from synchrotron radiation. Many threads perpendicular to the galactic plane are observed, indicating a strong vertical field (Yusef-Zadeh et al. 1984; Morris 1990). On the other hand, observations of near infrared polarization suggested that horizontal fields of 0.1–1 mG are dominant inside the disk (Novak et al. 2000; Nishiyama et al. 2009).

Taylor et al. (2009) presented results of the all-sky distribution of rotation measure (RM), which indicate that the magnetic fields along lines of sight are roughly consistent with the dipole-like models in which the azimuthal magnetic fields are anti-symmetric with respect to the equatorial plane, rather than quadrupole-like models in which the azimuthal magnetic fields show symmetry with respect to the equatorial plane. They also showed that the RM distribution has small-scale variations, which would be linked to the turbulent structure of the gaseous disk. Han et al. (2002) showed a view of RM distributions from the north Galactic pole using the results of pulsar polarization measurements. They indicated that the magnetic field direction changes between the spiral arms (Han et al. 2002; Han & Zhang 2007).

We carried out global three-dimensional magnetohydrodynamic simulations of dynamo activities in galactic gaseous disks assuming an axisymmetric gravitational potential and study the dynamo activities of galactic gas disks. In this paper, we present the results of global three-dimensional (3D) magnetohydrodynamic (MHD) simulations of galactic gas disks assuming an axisymmetric gravitational potential and study the dynamo activities of galactic gas disks.

In order to take into account the back-reaction of magnetic fields on the dynamics of gas disks, we have to solve the MHD equations in which the induction equation is coupled with the equation of motion. This approach is called the dynamical dynamo. Nishikori et al. (2006) presented results of 3D MHD simulations of galactic gaseous disks. They showed...
that, even when the initial magnetic field is weak, the dynamical dynamo amplifies the magnetic fields and can attain the current field strength observed in galaxies. They also showed that the direction of mean azimuthal magnetic fields reverses quasi-periodically, and pointed out that the reversal of the magnetic field is driven by the coupling of the MRI and the Parker instability. When the magnetic field is amplified by the MRI, the timescale of buoyant escape of magnetic flux due to the Parker instability becomes comparable to the growth time of the MRI and limits the strength of magnetic fields inside the disk.

Since Nishikori et al. (2006) assumed a symmetric boundary condition at the equatorial plane of the disk, they could not simulate the growth of the anti-symmetric mode of the Parker instability, in which magnetic fields cross the equatorial plane (Horiuchi et al. 1988). The present paper is aimed at obtaining a realistic, quantitative model of the global magnetic field of the Milky Way based on numerical 3D MHD simulations of magnetized galactic gaseous disks. In Section 2, we show the initial setting of the simulation. We present the results of magnetized galactic gaseous disks. In Section 2, we show the initial setting of the simulation. We present the results of numerical simulations in Section 3. The result will be used to compute the distribution of the RM on the whole sky in Section 4.

2. NUMERICAL MODELS

2.1. Basic Equations

We solved the resistive MHD equations by using a modified Lax–Wendroff scheme with artificial viscosity (Rubin & Burstein 1967; Richtmyer & Morton 1967):

\[
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0
\]

(1)

\[
\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v} \cdot \nabla)\mathbf{v} = -\frac{1}{\rho} \nabla P - \nabla \phi + \frac{1}{4\pi \rho} (\nabla \times \mathbf{B}) \times \mathbf{B}
\]

(2)

\[
\frac{\partial}{\partial t} \left( \frac{1}{2} \rho \mathbf{v}^2 + \frac{B^2}{8\pi} + \frac{P}{\gamma - 1} \right) + \nabla \cdot \left[ \left( \frac{1}{2} \rho \mathbf{v}^2 + \frac{P}{\gamma - 1} \right) \mathbf{v} + \frac{c}{4\pi} \mathbf{E} \times \mathbf{B} \right] = -\rho \mathbf{v} \nabla \phi
\]

(3)

\[
\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{v} \times \mathbf{B} - \eta \nabla \times \mathbf{B}),
\]

(4)

where \( \rho, P, \mathbf{v}, \mathbf{B}, \) and \( \gamma \) are the density, gas pressure, velocity, magnetic field, and the specific heat ratio, respectively, and \( \phi \) is the gravitational potential proposed by Miyamoto & Nagai (1975). The electric field \( \mathbf{E} \) is related to the magnetic field \( \mathbf{B} \) by Ohm’s law. We assume the anomalous resistivity \( \eta \)

\[
\eta = \begin{cases} 
\eta_0 (v_d/v_c - 1)^2, & v_d \geq v_c, \\
0, & v_d \leq v_c,
\end{cases}
\]

(5)

which sets in when the electron-ion drift speed \( v_d = j/\rho \), where \( j \) is the current density, exceeds the critical velocity \( v_c \) (Yokoyama & Shibata 1994).

2.2. Simulation Setup

We assumed an equilibrium torus threaded by weak toroidal magnetic fields (Okada et al. 1989). The torus is embedded in a hot, non-rotating spherical corona. We assume that the torus has constant specific angular momentum \( L \) at \( \sigma = \sigma_0 \) and assume a polytropic equation of state \( P = K \rho^\gamma \), where \( K \) is constant. The density distribution is same as that for Equation (8) in Nishikori et al. (2006).

We consider a gas disk composed of one-component interstellar gas. Radiative cooling and self-gravity of gas are ignored in this paper.

We adopt a cylindrical coordinate system (\( \sigma, \varphi, z \)). The units of length and velocity in this paper are \( \sigma_0 = 1 \) kpc and \( v_0 = (GM/\sigma_0)^{1/2} = 207 \) km s\(^{-1} \), respectively, where \( M = 10^{10} M_\odot \). The unit time is \( t_0 = \sigma_0/v_0 \). Other numerical units are listed in Table 1.

The numbers of grid points we used are \( (N_\sigma, N_\varphi, N_z) = (250, 128, 640) \). The grid size is \( \Delta \sigma/\sigma_0 = 0.05, \Delta \sigma/\sigma_0 = 0.01 \) for \( 0 \leq \sigma/\sigma_0 \leq 6.0 \), and \( |z|/\sigma_0 \leq 2.0 \), and otherwise they increase with \( \sigma \) and \( z \), respectively. The grid size in the azimuthal direction is \( \Delta \varphi = 2\pi/128 \). The outer boundaries at \( \sigma = 56\sigma_0 \) and \( |z| = 10\sigma_0 \) are free boundaries where waves can be transmitted. We applied periodic boundary conditions in the azimuthal direction. An inner absorbing boundary condition is imposed at \( r = \sqrt{\sigma^2 + z^2} = r_{\text{in}} = 0.8\sigma_0 \), since the gas accreted onto the central region of the galaxy will be converted to stars or swallowed by the black hole. To illuminate the effect of the boundary condition at the equatorial plane, we calculate two models: one is model PSYM, which imposes a symmetric boundary condition at the equatorial plane. This model is identical to that in Nishikori et al. (2006). The other is model ASYM, in which we include the equatorial plane inside the simulation region, and imposed no boundary condition at the equatorial plane. In this paper, we adopted model parameters \( \sigma_0 = 10\sigma_0, b_0 = 100 \) at \( (\sigma, z) = (\sigma_0, 0) \), \( \gamma = 5/3, v_c = 100v_0 \), and \( n_0 = 0.1v_0/\sigma_0 \). The thermal energy of the torus is parameterized by \( E_{\text{th}} = c_{\text{th}}^2/(\gamma v_0^2) = K \rho_n^{-1} v_0^2 = 0.05 \) where \( c_{\text{th}} \) and \( \rho_n \) are the sound speed and density at \( (\sigma, z) = (\sigma_0, 0) \), respectively. The initial temperature of the torus is \( T \sim 2 \times 10^5 \) K. This mildly hot plasma mimics the mixture of hot \( (T \sim 10^6 \) K) and warm \( (T \sim 10^4 \) K) components of interstellar matter, which occupy a large volume of the galactic disks.

3. NUMERICAL RESULTS

3.1. Time Evolution of the Galactic Gaseous Disk

Figure 1(a) shows the time evolution of the mass accretion rate measured at \( \sigma/\sigma_0 = 2.5 \). After \( t/t_0 \sim 200 \), the gas starts to inflow and the mass accretion rate increases linearly. Subsequently \( (t/t_0 > 600) \) the accretion rate saturates and becomes roughly constant. The evolution of the mass accretion rate in model ASYM (black) is similar to that in model PSYM (gray), where the mass accretion rate for model PSYM is

| Physical Quantity | Symbol | Numerical Unit |
|------------------|--------|----------------|
| Length           | \( \sigma_0 \) | 1 kpc          |
| Velocity         | \( v_0 \) | 207 km s\(^{-1} \) |
| Time             | \( t_0 \) | 4.8 \times 10^9 \) yr |
| Density          | \( \rho_0 \) | 1.6 \times 10^{-24} \) g cm\(^{-3} \) |
| Temperature      | \( T_0 \) | 5.15 \times 10^4 \) K |
| Magnetic field   | \( B_0 \) | 26 \( \mu \)G |

Table 1 Units of Physical Quantities
Figures 1(b)–(d) show the time evolution of the plasma $\beta$ and the magnetic energy of each component averaged in the region where $2 \leq \sigma / \sigma_0 < 5$, $|z|/\sigma_0 < 1$, and $0 \leq \varphi \leq 2\pi$, respectively. It is clear that plasma $\beta$ decreases in the linear stage as the magnetic energy increases. In both models, the averaged magnetic energy first increases exponentially. After that, it saturates and becomes roughly constant. The time evolution of magnetic fields is similar between model ASYM and model PSYM. The time evolutions of the growth and saturation in the mass accretion rate are quite similar to those of the averaged magnetic energy. This means that the magnetic turbulence driven by the MRI is the main cause of the angular momentum transport which drives mass accretion. In the remaining part of this section, we discuss the results of model ASYM.

In order to check the magnetic field structure, we analyze the time evolution of mean azimuthal magnetic fields. The mean fields are computed by the same method as Nishikori et al. (2006). Figure 2(a) shows the time evolution of the mean azimuthal magnetic fields. The black curve shows $B_\varphi$ averaged in the region where $5 < \sigma / \sigma_0 < 6$, $0 \leq \varphi \leq 2\pi$, and $0 < z/\sigma_0 < 1$, and the gray curve shows $B_\varphi$ averaged in the region where $5 < \sigma / \sigma_0 < 6$, $0 \leq \varphi \leq 2\pi$, and $1 < z/\sigma_0 < 3$. Black and gray curves correspond to the disk region and halo region, respectively. The azimuthal magnetic fields reverse their direction on a timescale $t \sim 300 t_0 \sim 1.5$ Gyr. This timescale is comparable to that of the buoyant rise of azimuthal magnetic fields. The halo fields also reverse their directions on the same timescale as the disk component.

The radial distribution of the mean azimuthal magnetic fields at $t = 1000 t_0$ are shown in Figure 2(b). Black and gray curves show the averages in the disk ($0 \leq \varphi \leq 2\pi$ and $0 < z/\sigma_0 < 1$) and in the halo ($0 \leq \varphi \leq 2\pi$ and $1 < z/\sigma_0 < 3$), respectively. Since the magnetic fields in the disk become turbulent, the direction of azimuthal magnetic fields changes frequently near the equatorial plane. On the other hand, long wavelength magnetic loops are formed in the halo region.

Figure 3(a) shows the radial distribution of the specific angular momentum at the equatorial plane. Although the rotation is assumed only inside the initial torus (gray curve in Figure 3(a)), the torus deforms its shape from a torus to a disk by redistributing angular momentum (black curve in Figure 3(a)). Since the outflow created by the MHD-driven dynamo transports angular momentum out of the model, the torus deforms. The rotation is transported outward by the outflow and turns the torus into a disk.

doubled since model PSYM includes only the region above the equatorial plane.
Figure 2. Evolutions of the mean azimuthal magnetic fields averaged in the region where $5 < \sigma / \sigma_0 < 6, \ 0 \leq \psi < 2\pi$, and $0 < z / \sigma_0 < 1$ (black), and in the region where $5 < \sigma / \sigma_0 < 6, \ 0 \leq \psi < 2\pi$, and $1 < z / \sigma_0 < 3$ (gray). (a) and (b) show time evolution and radial distribution, respectively.

Figure 3. (a) Radial distribution of the specific angular momentum ($l = \sigma v_\psi / (\sigma_0 v_0)$) at the equatorial plane. The black curve shows the distribution at $t = 1000 t_0$ and the gray curve shows the initial condition. The light gray dashed line shows the position of the initial density maximum. (b) Spatial distribution of the azimuthally averaged rotation speed at $t = 1000 t_0$.

3.2. Evolution of the Magnetic Field Structure

We show the 3D structure of the mean magnetic fields in Figure 4. The colored surfaces represent isodensity surfaces, and the curves show magnetic field lines. Light brown curves show field lines passing through the plane at $z / \sigma_0 = 5$ and $\sigma / \sigma_0 < 1$. The vertical magnetic fields around the rotation axis are produced by the magnetic pressure-driven outflow near the galactic center. Color on the curves depicts the direction of the azimuthal magnetic fields. Red curves show a positive direction of the azimuthal magnetic fields (counterclockwise direction) and blue shows a negative direction (clockwise). The figure indicates that the azimuthal field reverses the direction with the radius.

Figure 4 indicates that long wavelength magnetic loops are formed around the disk–halo interface. Therefore, we discuss the formation mechanism of the magnetic loops. Horizontal magnetic fields embedded in a gravitationally stratified atmosphere become unstable against long wavelength undular perturbations. This undular mode of magnetic buoyancy instability, called the Parker instability, creates buoyantly rising magnetic loops (Parker 1966). Nonlinear growth of the Parker instability in gravitationally stratified disks was studied by MHD simulations by Matsumoto et al. (1988). They showed that magnetic loops continue to rise when $\beta < 5$. In a weakly magnetized region where $\beta > 5$, the Parker instability only drives nonlinear oscillations (Matsumoto et al. 1990).

In differentially rotating disks, the MRI couples with the Parker instability. The growth timescale of the non-axisymmetric the MRI is $t_{MRI} \sim 1/(0.1 \Omega) \sim 10 H / c_s$ where $c_s$ is the sound speed. On the other hand, the growth timescale of the Parker instability is $t_{PI} \sim 5H / v_A \sim 5\sqrt{\beta} \gamma / 2(H / c_s)$ where $v_A$ and $\gamma$ are the Alfvén speed and the specific heat ratio, respectively. This timescale becomes comparable to the growth time of the MRI when $\beta \sim 5$. Since magnetic turbulence...
produced by the MRI limits the horizontal length of coherent magnetic fields, the growth of the Parker instability can be suppressed in a weakly magnetized region inside the disk. On the other hand, in regions where $\beta < 5$, the nonlinear growth of the Parker instability creates magnetic loops buoyantly rising to the disk corona (Machida et al. 2000, 2009).

Figure 5(a) shows an example of a magnetic loop at $t = 1000t_0$. Magnetic loops are identified by the same algorithm as that reported in Machida et al. (2009). Figures 5(b)–(d) show the distribution of the vertical velocity, plasma $\beta$, and density, respectively, along the magnetic field line depicted in Figure 5(a). A black circle shows the starting point of the integration of a magnetic field line. The positive vertical velocity indicates that the magnetic loop is rising. The value of plasma $\beta$ decreases from the footpoint of the loop where $\beta \sim 100$ to the loop top where $\beta < 5$ (Figure 5(d)). The density also decreases from the footpoints to the loop top. These are typical structures of buoyantly rising magnetic loops formed by the Parker instability.

From the numerical results presented so far, the initial magnetic field was found to be amplified and transported to the halo region after several rotation periods. In order to investigate the time evolution of the azimuthal magnetic field, we show a butterfly diagram of the azimuthal magnetic fields at $\sigma / \sigma_0 = 2$ (Figure 6(a)). The horizontal axis shows the time $t / t_0$, and the vertical axis shows the height $z / \sigma_0$. Color denotes the direction of the azimuthal magnetic fields. The white curve shows the isocontour where $\beta = 5$ above the equatorial plane. We should note that there is no magnetic flux at $\sigma / \sigma_0 = 2$ initially. As the time elapses, the accretion of the interstellar gas driven by the MRI transports magnetic fields from $\sigma / \sigma_0 = 10$ to $\sigma / \sigma_0 = 2$. The amplified azimuthal magnetic fields change their directions quasi-periodically. Turbulent magnetic fields are dominant around the equatorial plane.

When the magnetic tension becomes comparable to that of the turbulent motion in the disk surface and the coherent length of the magnetic field line increases, the magnetic flux buoyantly escapes from the disk due to the Parker instability. We find that the instability significantly grows, when the plasma $\beta$ decreases to around 5. The averaged escape velocity of the flux is about 25 km s$^{-1}$, approximately equal to the Alfvén velocity.

The periodic reversals of azimuthal magnetic fields can also be seen in the correlation between the fields above and below the equatorial plane (Figure 6(b)). White shows the region where $B_\phi(z) \cdot B_\phi(-z) > 0$, and black shows the region where $B_\phi(z) \cdot B_\phi(-z) < 0$. We can see fine structures of turbulence near the equatorial plane. In the halo region where $z / \sigma_0 > 1$, the color of the outgoing flux changes quasi-periodically. This means that the disk magnetic topology changes between the symmetric state and the anti-symmetric state.

Finally, we check the spatial distribution of the azimuthal magnetic fields. The azimuthal magnetic fields in the $\sigma - z$ plane averaged in the region $0 \leq \varphi \leq 2\pi$ at $t / t_0 = 1000$ is shown in Figure 7(a). We can see that a few $\mu$G fields are distributed as high as at $z / \sigma_0 = 5$. Figures 7(b) and (c) show the distribution of the azimuthal magnetic fields in the $\sigma - \varphi$ plane, where the azimuthal magnetic fields are averaged in the halo ($1 < z / \sigma_0 < 1.5$) and in the disk ($0 < z / \sigma_0 < 0.5$), respectively. Since turbulent magnetic fields are dominant inside the disk, short magnetic filaments are formed (Figure 7(c)). Long and strong azimuthal magnetic sectors are formed in the halo region because the plasma $\beta$ in
Figure 5. (a) A typical magnetic loop at $t = 1000 t_0$. The black circle denotes the starting point of integration of the magnetic field line. The distribution of the physical quantities along the magnetic field line, (b) vertical velocity, (c) plasma $\beta$, and (d) density, respectively.

Figure 6. (a) Time evolution of the azimuthal component of the magnetic field averaged in the azimuthal direction. The white curve shows the location where the plasma $\beta = 5$ above the equatorial plane ($z > 0$). (b) Correlation of the direction of $B_\phi$ below the equatorial plane and above it. White shows the region of positive correlation where the direction of the azimuthal magnetic fields above and below the disk is identical and black is the region with negative correlation.

The halo region becomes lower than that in the disk due to the density decrease (Figure 7(b)). Therefore, the magnetic tension of buoyantly rising magnetic loop exceeds the ram pressure by the turbulent motion around the disk surface and in the halo. These magnetic loops buoyantly escape from the disk by the Parker instability and create the rising magnetic fluxes in the butterfly diagram shown in Figure 6(a).

3.3. Dependence on the Azimuthal Resolution

In order to study the dependence of the numerical results on azimuthal resolution, we carried out simulations in which we used 64 (model ASYM64) and 256 (model ASYM256) grid points in $0 \leq \phi \leq 2\pi$. Other parameters are the same as those in model ASYM.

Figure 8(a) shows the time evolution of the magnetic energy averaged in the same region as Figure 1(c). Black, dashed, and gray curves show models ASYM, ASYM256, and ASYM64, respectively. This panel indicates that the saturation level of the magnetic energy near the equatorial plane slightly decreases as the resolution increases because shorter wavelength turbulence inside the disk can be resolved better in the model with higher azimuthal resolution, so that more magnetic energy dissipates in the nonlinear stage. The time evolution of the azimuthal magnetic fields averaged in the region where $5 < \sigma / \sigma_0 < 6$, $0 < z / \sigma_0 < 1$, and $0 \leq \phi \leq 2\pi$ is shown in Figure 8(b). Colors are the same as in Figure 8(a). The amplitude of the azimuthal magnetic fields and the timescale of the field reversal are almost similar in models ASYM and ASYM256. In order to resolve the most unstable wavelength of the MRI, we need more than 20 grid points per disk thickness for azimuthal direction (Hawley et al. 2011). Although ASYM does not satisfy this condition, the qualitative behavior of the results for ASYM is identical to that of ASYM256, which satisfies the condition (see Figure 8(b)). For example, the ratio of the radial magnetic energy to the
azimuthal magnetic energy ($B_{\phi}^2 / B_z^2$) is 0.12 in model ASYM and 0.13 in model ASYM256. Hence, we conclude that the numerical results of model PSYM and ASYM in this paper obtained by simulations with 128 grid points in the azimuthal direction do not significantly differ from results with higher azimuthal resolution in the nonlinear stage.

4. SUMMARY AND DISCUSSION

We carried out 3D MHD simulations of galactic gaseous disks. The numerical results indicate that magnetic fields are amplified by the MRI. The amplified azimuthal magnetic fluxes buoyantly escape from the disk by the Parker instability. Azimuthal component of the mean magnetic field reverses its...
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Figure 9. Schematic drawing of the mechanism of the MHD dynamo.

Figure 10. Rotation measure distribution obtained from numerical results. (A color version of this figure is available in the online journal.)

direction quasi-periodically. The timescale of the reversal is about 300\(t_0\) which corresponds to about 1.5 Gyr.

The numerical results also show that galactic gaseous disks become turbulent and that averaged plasma \(\beta\) stays at around \(\beta \sim 10\). Plasma \(\beta\) becomes lower near the disk surface due to the density stratification. Since rotational speed decreases from the disk to the halo (see Figure 3(b)), turbulent magnetic fields are stretched around the disk surface, forming ordered fields. When the magnetic pressure of the ordered fields becomes as large as \(\beta \sim 5\), magnetic flux escapes from the disk due to the buoyancy.

Figure 9 schematically illustrates the mechanism of the galactic dynamo obtained from numerical simulations. Let us consider a differentially rotating disk threaded by weak seed fields, which rotates counterclockwise. As the MRI grows in the gaseous disk, positive and negative fields whose strengths are comparable are amplified. When the averaged magnetic pressure becomes about 10% of the gas pressure in the disk, local magnetic pressure near the surface exceeds 20% of the gas pressure and the Parker instability creates magnetic loops near the disk surface. Since the magnetic field strength parallel to the seed field grows faster than that of the opposite field, the counterclockwise fields satisfy the condition of the Parker instability earlier and buoyantly escape from the disk. Subsequently, clockwise magnetic fields which remain inside the disk are amplified by the MRI, and mean azimuthal magnetic fields inside the disk reverse their direction. The reversed fields become a seed field for the next cycle.

Similar quasi-periodic reversal of azimuthal fields have been reported in local 3D MHD simulations of a differentially rotating disk (Shi et al. 2010). Although the MRI can grow so long as \(\beta \gtrsim 0.1\) (Johansen & Levin 2008), the saturation level of plasma \(\beta\) is about \(\beta \sim 1–10\) when the simulations include the effect of the vertical stratification. The growth of the MRI saturates when the growth rate of the Parker instability is largest (\(\beta \sim 1–10\)). This indicates that the saturation levels of the magnetic energy inside the disk are determined by the Parker instability around the surface layer.

Our finding about the periodical reversal of the direction of the halo azimuthal field would have an impact on the observation of the all-sky RM map. There are a number of reports about all-sky RM distributions (Taylor et al. 2009; Braun et al. 2010; Han & Zhang 2007). The study of all-sky RM distribution is essential not only to understand the global structure of galactic magnetic fields but also to explore the intergalactic magnetic fields in the intergalactic medium (Akahori & Ryu 2010, 2011, and references therein). For instance, Taylor et al. (2009) suggested that the existence of the halo poloidal field is consistent with a dipole field rather than a quadrupole field.

We calculated the distribution of the RM obtained from numerical results (Figure 10). The position of the observer is assumed to be at \(r = 8\) kpc at \(t = 1000t_0\). The direction of the rotational velocity and azimuthal magnetic fields is counterclockwise in our numerical simulation. Since the angular velocity observed in the Milky Way is clockwise, we calculated the RM distribution by reversing the \(z\)-direction (i.e., north is the \(z\)-direction). From the characteristics of MHD, if the initial magnetic field directions were reversed, we can obtain the same result except for the direction of the magnetic fields. Thus we reversed the color of the RM from the original one. The scale of the RM is arbitrary in this paper. We found that magnetic loops produced by the Parker instability formed multiple reversals of the sign of the RMs along the latitude. These reversals reveal that magnetic fields with opposite polarity emerge from the disk periodically. The snapshot shown in Figure 10 indicates that the distribution of the RM is point symmetric with respect to the galactic center, and that this feature remains in a timescale of \(\sim 1.5\) Gyr (Figure 2(a)).

Braun et al. (2010) observed RM distributions in nearby galaxies. They also pointed out that the field topology in the upper halo of galaxies is a mixture of an axisymmetric spiral quadrupole field in thick disks and a radially directed dipole field in halos. They suggested that the origin of the dipole components might be a bipolar outflow. A mixed structure was also suggested by recent studies of the all-sky RM map (see, e.g., Figure 8 of Pshirkov et al. 2011). The RM distribution in the halo region of our numerical results shows an anti-symmetric distribution corresponding to the dipole field.

Outflows are produced by emerging magnetic fluxes driven by the Parker instability. The time interval of the buoyant escape of the magnetic flux by the Parker instability is about 10 rotational periods, which corresponds to the growth timescale of the MRI. Since the buoyant escape of magnetic flux takes place about 10 cycles during the simulation for 4 Gyr at \(\sigma/\sigma_0 = 2\), the disk becomes turbulent enough, and the influence of the initial field...
Our numerical calculation showed that an initially non-rotating halo rotates after several rotation periods because the angular momentum was supplied from the disk to the halo by rising magnetic loops formed by the Parker instability. In fact, H\textsc{i} observations of the late-type spiral galaxy NGC 6503 indicated that there were two components of the H\textsc{i} disk, a thin dense disk and a thick low-density disk, and that the thick low-density disk rotates slower than the thin disk (Greisen et al. 2009). The rotation speed of the ionized gas in the thick disk is slower than that of H\textsc{i} (Bottema 1989).

In this paper, we assumed that the interstellar gas is a one-component gas with a temperature of $2 \times 10^5$ K in order to compare the result with that by Nishikori et al. (2006). However, in a realistic galaxy, the interstellar gas has multi-temperature, multi-phase components. In future works, we would like to include the energy input by supernovae, multi-temperature structure of the interstellar gas, and cosmic rays.

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