Heat conduction probably depends on the magnetic field configuration, which plays a decisive role in most theoretical models of the particle acceleration. 

Michelangelo et al. 1987; Blasi 2000; Ohno et al. 2002; Fujita et al. 2003; Brunetti et al. 2004), where the magnetic fields play a definite role in various aspects of ICM. Some observational evidence indicates that the ICM is magnetized. For example, some clusters have diffuse nonthermal synchrotron radio emission that is called radio halos or relics, which shows that there are magnetic fields as well as relativistic electrons in the intracluster space (Giovannini et al. 1999; Kemper & Sarazin 2001). In addition, Faraday rotation measurements of polarized radio sources such as radio lobes and active galactic nuclei (AGNs) behind and/or in clusters indicate that the magnetic field structures are quite random (Clarke et al. 2001; Vogt & Ensslin 2003; Govoni et al. 2006).

Comparing the synchrotron radio flux with the hard X-ray one (or its upper limit) due to the inverse Compton scattering of cosmic microwave background (CMB) photons, we are able to estimate the volume-averaged magnetic field strength (or its lower limit). Typically, a strength of ~0.1 $\mu$G is obtained with this method (Fusco-Femiano et al. 1999, 2005), although those detections of nonthermal hard X-ray components are still controversial (Rossetti & Molendi 2004; Fusco-Femiano et al. 2007). On the other hand, somewhat higher values (~several $\mu$G) tend to be obtained with the Faraday rotation measure method (Clarke et al. 2001; Vogt & Ensslin 2003; Govoni et al. 2006), although these results depend on the detailed magnetic field structures that are not fully understood (Ensslin & Vogt 2003; Murgia et al. 2004).

Although the magnetic energy density is typically ~ a few percent of the thermal one in the intracluster space, it is believed that the magnetic fields play a crucial role in various aspects of the ICM. It is likely that nonthermal particles are accelerated via shocks (Sarazin 1999; Takizawa & Naito 2000; Totani & Kitayama 2000; Miniati et al. 2001; Ryu et al. 2003; Takizawa 2002; Inoue et al. 2005) and/or turbulence (Roland 1981; Schlickeiser et al. 1987; Blasi 2000; Ohno et al. 2002; Fujita et al. 2003; Brunetti et al. 2004), where the magnetic fields play a definitive role in most theoretical models of the particle acceleration. Heat conduction probably depends on the magnetic field configurations, because charged particles cannot freely move in the direction perpendicular to the field lines. As we wrote before, magnetic field seems to be a minor component in global ICM dynamics. However, it is possibly important on relatively small scales, where fluid instabilities such as Rayleigh-Taylor and Kelvin-Helmholtz ones might be suppressed by magnetic tension.

Cluster mergers have a significant impact on ICM magnetic field evolution. Turbulent motion excited by mergers would amplify the field strength via the dynamo mechanism. Moving substructures are expected to sweep the field lines and form the cold subclumps surrounded by the field lines (Vikhlinin et al. 2001). Asai et al. (2004, 2007) performed two- and three-dimensional magnetohydrodynamical (MHD) simulations of moving cold subclumps in the hot ICM in rather idealized situations, respectively, which confirmed that expectation. This field structure might be responsible for the suppression of heat conduction and fluid instabilities, which is essential for the maintenance of cold fronts.

A lot of numerical simulations about merging clusters have been done, most of which are N-body + hydrodynamical simulations (Roettiger et al. 1996; Takizawa 1999, 2000, 2006; Ricker & Sarazin 2001; Ritchie & Thomas 2002; Ascasibar & Markovich 2006; McCarthy et al. 2007; Springel & Farrar 2007). Although these simulations give us a great deal of understanding of the structures, evolution, and observational implications for merging clusters of galaxies, they make only a limited contribution to investigate the magnetic field structures. MHD simulations are essential in this regard. However, N-body + MHD simulations are rather rare, although Roettiger et al. (1999) did pioneering work. It is true that Lagrangian particle methods based on smoothed particle hydrodynamics (see Monaghan 1992) are extensively used in cosmological MHD simulations (Dolag et al. 1999, 2002). Considering that Eulerian mesh codes are essentially better at following the evolution of magnetic fields, simulations based on such codes are highly desirable. In this paper we present the results from N-body + MHD simulations of merging clusters of galaxies and investigate characteristic magnetic field structures during mergers and their implications.

1. INTRODUCTION

Clusters of galaxies have plenty of hot plasma as well as galaxies and dark matter (DM), which is called the intracluster medium (ICM). Some observational evidence indicates that the ICM is magnetized. For example, some clusters have diffuse nonthermal synchrotron radio emission that is called radio halos or relics, which shows that there are magnetic fields as well as relativistic electrons in the intracluster space (Giovannini et al. 1999; Kemper & Sarazin 2001). In addition, Faraday rotation measurements of polarized radio sources such as radio lobes and active galactic nuclei (AGNs) behind and/or in clusters indicate that the magnetic field structures are quite random (Clarke et al. 2001; Vogt & Ensslin 2003; Govoni et al. 2006).

Comparing the synchrotron radio flux with the hard X-ray one (or its upper limit) due to the inverse Compton scattering of cosmic microwave background (CMB) photons, we are able to estimate the volume-averaged magnetic field strength (or its lower limit). Typically, a strength of ~0.1 $\mu$G is obtained with this method (Fusco-Femiano et al. 1999, 2005), although those detections of nonthermal hard X-ray components are still controversial (Rossetti & Molendi 2004; Fusco-Femiano et al. 2007). On the other hand, somewhat higher values (~several $\mu$G) tend to be obtained with the Faraday rotation measure method (Clarke et al. 2001; Vogt & Ensslin 2003; Govoni et al. 2006), although these results depend on the detailed magnetic field structures that are not fully understood (Ensslin & Vogt 2003; Murgia et al. 2004).

Although the magnetic energy density is typically ~ a few percent of the thermal one in the intracluster space, it is believed that the magnetic fields play a crucial role in various aspects of the ICM. It is likely that nonthermal particles are accelerated via shocks (Sarazin 1999; Takizawa & Naito 2000; Totani & Kitayama 2000; Miniati et al. 2001; Ryu et al. 2003; Takizawa 2002; Inoue et al. 2005) and/or turbulence (Roland 1981; Schlickeiser et al. 1987; Blasi 2000; Ohno et al. 2002; Fujita et al. 2003; Brunetti et al. 2004), where the magnetic fields play a definitive role in most theoretical models of the particle acceleration. Heat conduction probably depends on the magnetic field configurations, because charged particles cannot freely move in the direction perpendicular to the field lines. As we wrote before, magnetic field seems to be a minor component in global ICM dynamics. However, it is possibly important on relatively small scales, where fluid instabilities such as Rayleigh-Taylor and Kelvin-Helmholtz ones might be suppressed by magnetic tension.

Cluster mergers have a significant impact on ICM magnetic field evolution. Turbulent motion excited by mergers would amplify the field strength via the dynamo mechanism. Moving substructures are expected to sweep the field lines and form the cold subclumps surrounded by the field lines (Vikhlinin et al. 2001). Asai et al. (2004, 2007) performed two- and three-dimensional magnetohydrodynamical (MHD) simulations of moving cold subclumps in the hot ICM in rather idealized situations, respectively, which confirmed that expectation. This field structure might be responsible for the suppression of heat conduction and fluid instabilities, which is essential for the maintenance of cold fronts.

A lot of numerical simulations about merging clusters have been done, most of which are N-body + hydrodynamical simulations (Roettiger et al. 1996; Takizawa 1999, 2000, 2006; Ricker & Sarazin 2001; Ritchie & Thomas 2002; Ascasibar & Markovich 2006; McCarthy et al. 2007; Springel & Farrar 2007). Although these simulations give us a great deal of understanding of the structures, evolution, and observational implications for merging clusters of galaxies, they make only a limited contribution to investigate the magnetic field structures. MHD simulations are essential in this regard. However, N-body + MHD simulations are rather rare, although Roettiger et al. (1999) did pioneering work. It is true that Lagrangian particle methods based on smoothed particle hydrodynamics (see Monaghan 1992) are extensively used in cosmological MHD simulations (Dolag et al. 1999, 2002). Considering that Eulerian mesh codes are essentially better at following the evolution of magnetic fields, simulations based on such codes are highly desirable. In this paper we present the results from N-body + MHD simulations of merging clusters of galaxies and investigate characteristic magnetic field structures during mergers and their implications.
The rest of this paper is organized as follows. In § 2 we describe the adopted numerical methods and initial conditions for our simulations. In § 3 we present the results. In § 4 we summarize the results and discuss their implications.

2. THE SIMULATIONS

2.1. Numerical Methods

Numerical methods used here are basically similar to those in Takizawa (2006) except that ideal MHD equations are calculated for the ICM instead of ordinary hydrodynamical ones. In the present study we consider clusters of galaxies consisting of two components: collisionless particles corresponding to the galaxies and DM, and magnetized gas corresponding to the ICM. When calculating gravity, both components are considered, although the former dominates over the latter. Radiative cooling and heat conduction are not included. In calculating the MHD equations for the ICM, we use a simple linearized Riemann solver originally proposed by Brio & Wu (1988) and introduced into astrophysical MHD simulations by Ryu & Jones (1995). This Riemann solver is a similar and simpler version of Roe’s method (Roe 1981). Using the MUSCLE approach and a minmod total variation diminishing limiter (see Hirsch 1990), we obtain second-order accuracy without any numerical oscillations around discontinuities. To avoid negative pressure, we solve the equations for the total energy and entropy conservation simultaneously.

Gravitational forces are calculated by the particle-mesh method with the standard fast Fourier transform technique for the isolated symmetry, so that the DM particles would be in virial equilibrium. The coordinate system is taken in such a way that the center of mass is at rest of the origin. Two subclusters are initialized in the x-y plane, separated by a distance \(|(R_1 + R_2)^2 - b^2|^{1/2}\) in the x-direction and \(b\) in the y-direction, where \(R_1, R_2\) and \(b\) are the virial radii of each subcluster. It should be noted that the results will not be sensitive to the choice of this fraction unless the ICM dominates the DM in gravity.

The velocity distribution of the DM particles is assumed to be an isotropic Maxwellian. The radial profiles of the DM velocity dispersion are calculated from the Jeans equation with spherical symmetry, so that the DM particles would be in virial equilibrium in the cluster potential of the DM and ICM,

\[
\frac{d}{dr}\left(\rho_{\text{DM}}\sigma^2\right) = - \frac{GM(r)}{r^2} \rho_{\text{DM}},
\]

where \(\sigma^2\) and \(M(r)\) are the one-dimensional velocity dispersion at \(r\) and total mass inside \(r\), respectively, and \(G\) is the gravitational constant. We need boundary conditions to obtain \(\sigma^2(r)\) from differential equation (3), which is assumed to be

\[
\sigma^2(r_{\text{vir}}) = \frac{GM(r_{\text{vir}})}{3r_{\text{vir}}},
\]

where \(r_{\text{vir}}\) is the virial radius of the halo.

The radial profiles of the ICM pressure are determined in a similar way so that the ICM would be in hydrostatic equilibrium within the cluster potential.

\[
\frac{dP}{dr} = - \frac{GM(r)}{r^2}\rho_{\text{gas}}.
\]

The boundary condition for equation (5) is

\[
P(r_{\text{vir}}) = \frac{1}{\beta_{\text{spec}}} \frac{GM(r_{\text{vir}})}{3r_{\text{vir}}} \rho_{\text{gas}}(r_{\text{vir}}),
\]

where \(\beta_{\text{spec}}\) is the specific energy ratio of the DM and gas at the virial radius.

In the case of the classical isothermal \(\beta\)-model (see Sarazin 1986), \(\beta_{\text{spec}}\) should be equal to the fitting parameter \(\beta\) in equation (2) to obtain isothermal temperature profile. Although our model is different than the isothermal \(\beta\)-model in its underlying DM distribution, we assume \(\beta_{\text{spec}} = \beta\) for simplicity. As a result, the temperature profile is not isothermal.

2.2. An Equilibrium Cluster Model

We consider mergers of two virialized subclusters with a Navarro-Frenk-White density profile (Navarro et al. 1997) in the \(\Lambda\)CDM universe \((\Omega_0 = 0.25, \lambda_0 = 0.75)\) for DM,

\[
\rho_{\text{DM}}(r) = \frac{\delta_c \rho_{\text{crit}}}{(r/r_y)(1 + r/r_y)^2},
\]

where \(\delta_c, \rho_{\text{crit}},\) and \(r_y\) are the characteristic (dimensionless) density, critical density, and scale radius, respectively. Given the cosmological parameters and the halo’s virial mass, we calculate these parameters following the method in the Appendix of Navarro et al. (1997).

The initial density profiles of the ICM are assumed to be those of a \(\beta\)-model,

\[
\rho_{\text{gas}}(r) = \rho_{\text{gas},0} \left[1 + \left(\frac{r}{r_c}\right)^3\right]^{-3\beta/2},
\]

where \(\rho_{\text{gas},0}\) and \(r_c\) are the central gas density and core radius, respectively. We assume that \(\beta = 0.6\), which is consistent with typical values obtained from the recent X-ray observations (Vikhlinin et al. 2006; Croston et al. 2008). On the other hand, it is not easy to obtain the observational relationship between \(r_c\) and \(r_y\) considering the difficulty of accurate determination of the mass profile. Following the discussion in Ricker & Sarazin (2001), we therefore chose \(r_c = r_y/2\) as a fiducial value. The gas mass fraction is set to be 0.1 inside of the virial radius of each subcluster. It should be noted that the results will not be sensitive to the choice of this fraction unless the ICM dominates the DM in gravity.

The velocity distribution of the DM particles is assumed to be an isotropic Maxwellian. The radial profiles of the DM velocity dispersion are calculated from the Jeans equation with spherical symmetry, so that the DM particles would be in virial equilibrium in the cluster potential of the DM and ICM,

\[
\frac{d}{dr}\left(\rho_{\text{DM}}\sigma^2\right) = - \frac{GM(r)}{r^2} \rho_{\text{DM}},
\]

where \(\sigma^2\) and \(M(r)\) are the one-dimensional velocity dispersion at \(r\) and total mass inside \(r\), respectively, and \(G\) is the gravitational constant. We need boundary conditions to obtain \(\sigma^2(r)\) from differential equation (3), which is assumed to be

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\sigma^2(r_{\text{vir}}) = \frac{GM(r_{\text{vir}})}{3r_{\text{vir}}},
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where \(r_{\text{vir}}\) is the virial radius of the halo.

The radial profiles of the ICM pressure are determined in a similar way so that the ICM would be in hydrostatic equilibrium within the cluster potential,

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\frac{dP}{dr} = - \frac{GM(r)}{r^2}\rho_{\text{gas}}.
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where \(\beta_{\text{spec}}\) is the specific energy ratio of the DM and gas at the virial radius.

In the case of the classical isothermal \(\beta\)-model (see Sarazin 1986), \(\beta_{\text{spec}}\) should be equal to the fitting parameter \(\beta\) in equation (2) to obtain isothermal temperature profile. Although our model is different than the isothermal \(\beta\)-model in its underlying DM distribution, we assume \(\beta_{\text{spec}} = \beta\) for simplicity. As a result, the temperature profile is not isothermal.

2.3. Initial Conditions

For typical runs, the DM masses within the virial radius of the larger and smaller subclusters are \(5.0 \times 10^{14} \, M_\odot\) and \(1.25 \times 10^{14} \, M_\odot\), respectively. Thus, the mass ratio is 4 : 1. It is useful to introduce an angular momentum parameter \(\lambda \equiv J/E^{1/2}/(GM^{3/2})\) in order to characterize off-center collisions, where \(J, E,\) and \(M\) are the angular momentum of the two subclusters around the centers of mass, the binding energy between the two, and the total mass, respectively. Please note that \(\lambda\) is the ratio between the actual angular velocity and the angular velocity needed to provide rotational support (see Binney & Tremaine 1987).

Given each subcluster’s mass and an angular momentum parameter of the system, we estimate “typical” initial conditions for cluster mergers in a similar way to that of Takizawa (1999). How to estimate these conditions is described in detail in the Appendix, which is a natural generalization of the method for head-on collisions (see Takizawa 1999, § 2) to cases including off-center collisions. The coordinate system is taken in such a way that the center of mass is at rest at the origin. Two subclusters are initialized in the x-y plane, separated by a distance \(|(R_1 + R_2)^2 - b^2|^{1/2}\) in the x-direction and \(b\) in the y-direction, where \(R_1, R_2\) and \(b\) are the virial radii of each subcluster and the
impact parameter, respectively. The initial relative velocity is directed along the $x$-axis. The centers of the larger and smaller subclusters were initially located at the sides of $x < 0$ and $x > 0$, respectively.

At present, we have only limited information about magnetic field configurations in the intracluster space. Roughly speaking, however, they have random structures whose power spectrum can be approximated by a power law on a smaller scale, and the mean field strength is an increasing function of the gas density. Taking these features into account, we make the initial magnetic field as follows. First, we generate a random Gaussian vector potential in wavenumber space with a zero mean and single power-law spectrum, $A(k) \propto k^{-\mu}$. We assume $\mu = 5/3$. Then, the generated potential is transformed into real space by using a three-dimensional fast Fourier transform and multiplying by a factor of $\rho_{\text{gas}}^{2/3}$, which is expected assuming a uniform spherical collapse and flux freezing. A random and divergence-free magnetic field is calculated via $B = \nabla \times A$. The normalization of the magnetic field is determined so that the total magnetic energy is 1% of the thermal one. Strictly speaking, each subcluster in the initial conditions is not in hydrostatic equilibrium with magnetic field. However, it does not matter in our simulations, because the magnetic energy is much less than the thermal one as we wrote above.

The parameters for each model are summarized in Table 1, where $M_i$, $r_{s,i}$, $c_i$, $N_i$ are the total mass, scale radius, concentration parameter, and number of $N$-body particles for the $i$th subcluster, respectively, and $N_{\text{grd}}$ is the total number of grid points. Run ST, which is a simple head-on collision case with a mass ratio of 4:1, is regarded as a standard run. Run OC is a case of an off-center merger with $\lambda = 0.05$. Run MM is a relatively minor merger case with a mass ratio of 8:1. Run LR is a lower resolution run to investigate how the results depend on numerical resolution.

### 3. RESULTS

The left panels of Figure 1 show snapshots of the density (contours) and temperature (colors) on the $x$-$y$ plane, which is perpendicular to the collision axis, at $t = 1.11, 1.56, 2.00, 2.89$, and $3.78$ Gyr of run ST. The right panels are the same but for the magnetic field strength. At $t = 1.11$ Gyr, two subclusters are approaching each other, and the temperature between two density peaks is higher than elsewhere. However, magnetic field strength in this region is lower than that in the density peaks. This is simply because the initial mean magnetic field strength increases with the density. Although the magnetic field there is slightly amplified by the adiabatic compression, this does not compensate for the trend of the initial magnetic field. At $t = 1.56$ and 2.0 Gyr, slightly amplified magnetic field is seen just behind the shock. However, two kinds of structure appear more notably. One is a relatively strong magnetic field region associated with a contact discontinuity between the ICM that originated from the smaller and larger subclusters, which will be recognized as a cold front with strong magnetic field in the observational point of view. As a result, a cool region surrounded by the magnetic field appears. The other is

| Run | $M_1/M_2$ ($10^{14}$ $M_\odot$) | $r_{s,1}/r_{s,2}$ (kpc) | $c_1/c_2$ | $\lambda$ | $N_1/N_2$ | $N_{\text{grd}}$ |
|-----|---------------------------------|--------------------------|------------|----------|------------|--------------|
| ST  | 5.0/1.25                        | 256/139                  | 6.12/7.08  | 0.0      | 1,677,721/419,431 (256) |
| OC  | 5.0/1.25                        | 256/139                  | 6.12/7.08  | 0.05     | 1,677,721/419,431 (256) |
| MM  | 5.0/0.625                       | 256/104                  | 6.12/7.56  | 0.0      | 1,864,135/23,3017 (256) |
| LR  | 5.0/1.25                        | 256/139                  | 6.12/7.08  | 0.0      | 209,715/52,429 (128) |

**Fig. 1** — Left: Snapshots of the density (contours) and temperature (colors) distributions on the $z = 0$ surface at $t = 1.11, 1.56, 2.00, 2.89$, and $3.78$ Gyr of run ST. Right: Same as the left panels, but for the magnetic field strength distribution.
an ordered magnetic field in the $x$-direction just behind the subcluster. This is because flow behind the subcluster is converging on the collision axis. As a result, the magnetic field is collected and amplified because of compression. At $t = 2.89$ Gyr, the bow shock has gone outside of the panel. The field structure associated with the contact discontinuity and ordered field along the collision axis are still present. In addition, eddylike field structures start to appear above and below the ordered field along the collision axis in the $x$-$y$ plane. These structures grow as Kelvin-Helmholtz instabilities develop and become clearer at $t = 3.78$ Gyr.

The initial conditions of run ST have an axial symmetric structure. Thus, the above-mentioned features could change significantly in off-center mergers. Figure 2 is the same as Figure 1, but for run OC. Again, the prominent magnetic structures are associated with a contact discontinuity rather than a bow shock. Because of the asymmetry, the magnetic field at the contact discontinuity is stronger on the side closer to the larger cluster core (or bottom in each panel) at $t = 2.00$ and 2.89 Gyr. A cool region surrounded by the magnetic field certainly appears, but the field structure on the side farther from the larger cluster core (or top in each panel) becomes less clear at $t = 3.78$ Gyr. As in the case of
run ST, an ordered and relatively strong magnetic field structure appears behind the moving substructure. However, this is not just behind the clump but shifted toward the larger cluster side. Eddylike magnetic fields generated by the Kelvin-Helmholtz instabilities are clearer on the side farther from the larger cluster. Cluster mergers cause various characteristic magnetic field configurations as we have seen above. Unfortunately, it is very difficult to observe the magnetic field structure itself directly at present. The observation of the Faraday rotation measure is an indirect, but quite useful, way to obtain information about the magnetic field structure. The rotation measure (RM) is given by (see Rybicki & Lightman 1979)

$$\text{RM rad m}^{-2} = 812 \int n_n [\text{cm}^{-3}] B_{||} [\mu \text{G}] dl,$$

where $n_n$ and $B_{||}$ are the electron number density and line-of-sight component of the magnetic field, respectively. Figure 3 presents snapshots of the RM map for run ST. Left and right panels show the maps seen from the z- and x-axes, respectively. Please
remember that the collision axis is along the $x$-axis. When the line of sight is perpendicular to the collision axis, RM maps give us relatively rich information. Cool regions with high magnetic field are seen as those with high absolute values of RM. On the other hand, the structure is relatively featureless when we observe the system along the collision axis. The absolute values of RM tend to be higher in the right panels, which is not surprising because the density distribution is elongated along the collision axis and because there is the ordered magnetic field in the same direction behind the subclump.

Figure 4 shows profiles along the collision axis at $t = 2.00$ Gyr of run ST for (a) pressure, (b) electron number density, (c) temperature, (d) the $x$-component of the velocity, (e) the $x$-component of the magnetic field, and (f) the absolute value of the magnetic field perpendicular to the axis. There is a bow shock at $x \approx -1.5$, where clear discontinuities are seen in the profiles of pressure, density, temperature, and the $x$-component of the velocity. A somewhat blunt contact discontinuity is at $x \approx -1$, where density and temperature have a jump, but pressure and $x$-velocity are smooth. The magnetic field perpendicular to the axis also increases across the bow shock. The tangential magnetic field is amplified between the shock and the contact discontinuity. As a result, the cool region is wrapped by the magnetized layers. This characteristic structure probably has an influence on transport processes such as heat conduction, which is discussed further below.

Because minor mergers are more frequent in actual situations, it is interesting how the results change in minor mergers. Figure 5 is the same as Figure 4 but for run MM, where the mass ratio is 1:8. The Mach number of the bow shock becomes lower, which is clearly seen in the profiles of pressure, density, temperature, and the $x$-component of the velocity. As a result, amplification of the magnetic field component perpendicular to the collision axis is amplified especially between a bow shock and contact discontinuity. As a result, a cool region wrapped by field lines appears. A relatively ordered field structure along the collision axis appears just behind the moving substructure. Eddylike field structures are also generated by Kelvin-Helmholtz instabilities. Although the detailed structures change slightly, similar features are seen in off-center mergers. RM maps have some information about the magnetic field structure. The above-mentioned characteristic structures are partly recognized in the RM maps when we observe the merging systems in the direction perpendicular to the collision axis. On the other hand, RM maps observed nearly along the collision axis are less informative in this respect. Typical absolute values of the RM become higher when the system is observed along the collision axis because of both the density distribution elongated toward the axis and the ordered magnetic field along the same direction. In minor mergers, amplification of the magnetic field component perpendicular to the collision axis becomes less prominent. The results about magnetic field along.

4. SUMMARY AND DISCUSSION

We perform $N$-body and MHD simulations of merging clusters of galaxies. We find that cluster mergers cause various characteristic magnetic field structures because of strong bulk flow motion. The magnetic field component perpendicular to the collision axis is amplified especially between a bow shock and contact discontinuity. As a result, a cool region wrapped by field lines appears. A relatively ordered field structure along the collision axis appears just behind the moving substructure. Eddylike field structures are also generated by Kelvin-Helmholtz instabilities. Although the detailed structures change slightly, similar features are seen in off-center mergers. RM maps have some information about the magnetic field structure. The above-mentioned characteristic structures are partly recognized in the RM maps when we observe the merging systems in the direction perpendicular to the collision axis. On the other hand, RM maps observed nearly along the collision axis are less informative in this respect. Typical absolute values of the RM become higher when the system is observed along the collision axis because of both the density distribution elongated toward the axis and the ordered magnetic field along the same direction. In minor mergers, amplification of the magnetic field component perpendicular to the collision axis becomes less prominent. The results about magnetic field along.
the collision axis do not change very much. Although numerical resolution effects have a significant impact on the small-scale structures, such as the widths of a bow shock and contact discontinuity and detailed small-scale fluctuations of the magnetic field, our results about overall global structures are reasonably reliable.

Some observational and theoretical studies suggest that heat conduction in the ICM is suppressed from the Spitzer value. For example, a sharp temperature jump across a cold front in A2142 requires that the heat conductivity is reduced by a factor of between 250 and 2500 at least in the direction across the front (Ettori & Fabian 2000). The spatial temperature variations in the central region of A754 also suggest that the conductivity is at least an order of magnitude lower than the Spitzer value (Markevitch et al. 2003). It is well known that fine-tuning of the heat conductivity is necessary in order to reproduce the observed nature of cool cores assuming that radiative cooling in the cores balances with the conductive heating from the outer part of the clusters (Zakamska & Narayan 2003). Although the detailed process involved there is still unclear, magnetic fields likely play an important role. Heat conduction in the direction across the magnetic field lines is likely suppressed, because electrons cannot travel freely in that direction. Using two-dimensional MHD simulations with anisotropic heat conduction, Asai et al. (2004) show that a moving subclump naturally forms the magnetic field structure along the contact discontinuity and that the temperature jump there is maintained if the heat conduction perpendicular to the field line is sufficiently suppressed. Similar results are obtained in three-dimensional cases (Asai et al. 2007). Basically, our simulations confirmed their results about magnetic field configurations in more realistic situations, although the heat conduction is not included.

We show that Faraday RM maps are useful to obtain information about characteristic magnetic field structures formed by cluster mergers. This could be another probe of the internal dynamics of clusters. However, there is a serious problem in this method at present. We need polarized radio sources such as radio lobes and/or AGNs in and/or behind the clusters of galaxies to do that. In other words, we are able to obtain RM maps in the limited regions where we have suitable polarized sources by chance. Naturally, these do not always correspond to the regions that we are interested in. However, this difficulty could be overcome if we have suitable polarized sources that cover a cluster entirely. One possible solution is to use the CMB as the polarized source (Ohno et al. 2003), although this is still challenging in the present status of CMB observations. There are a lot of observations that are ongoing and planned for measuring the CMB polarization. We hope that future observations of the CMB polarizations will enable us to make RM maps that cover clusters entirely, which would give us important clues to understand the internal dynamics as well as the magnetic field structures in clusters of galaxies.

In actual clusters, the magnetic field has random components on small scales. It is very difficult to treat such scales by numerical simulations that concern the global structures. Clearly, observed RM maps are influenced by the small-scale magnetic field fluctuations, which are not considered in RM maps calculated from our results. As a result, the coherent length of the magnetic fields tends to be effectively overestimated, which means absolute values of RM from the simulations are also overestimated. However, it is probable that global spatial patterns of RM maps are nearly independent of such small-scale fluctuations. In addition, it is also probably robust that higher RMs are expected when the merging systems are observed along the collision axis.

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APPENDIX

AN ESTIMATION METHOD OF THE INITIAL CONDITIONS OF Mergers

The estimation method introduced here is a natural extension of that for head-on collisions in Takizawa (1999) to cases including off-center collisions. We consider a merger of two subclusters with masses $M_1$ and $M_2$ ($M_1 \geq M_2$). Let us consider a region that contains both subclusters. It expands following the cosmological expansion at first. However, the expansion is decelerated and then the region will collapse if the mean density is higher than the critical one. We assume that they are separated by a distance of $2r_{1a}$ at the maximum expansion epoch. In the case of head-on collisions, it is obvious that they are at rest with respect to their center of mass at that time. In the case of off-center mergers, on the other hand, they have the relative tangential velocity $v_0$ at that time. Taking into account the conservation of energy and angular momentum between the maximum expansion epoch and just before the merger, we obtain following equations,

$$- \frac{GM_1 M_2}{2r_{1a}} + \frac{1}{2} M_1 v_0^2 = \frac{GM_2 M}{R_1 + R_2} + \frac{1}{2} M v^2,$$

$$2M_0 v_0 r_{1a} = M v b,$$

where $R_1$, $R_2$, $v$, and $b$ are the virial radius for each subcluster, the initial collision velocity, and the initial impact parameter, respectively, and $M \equiv M_1 M_2/(M_1 + M_2)$ is the reduced mass of the system. It is natural that the virial radius has a correlation with the mass. According to a spherical collapse model, the virial mass is proportional to the square of the virial radius (see Peebles 1980). Thus, we assume the following relation,

$$R_2 = \left( \frac{M_2}{M_1} \right)^{1/3} R_1.$$
The spherical collapse model also tells us that the turnaround radius is twice the virial radius. Thus, we also assume

\[ r_{ta} = 2 \left( \frac{M_1 + M_2}{M_1} \right)^{1/3} R_1. \]  

(A4)

Using equations (A1), (A2), (A3), and (A4), we eliminate \( v_b \) and obtain \( v \) as follows,

\[ v^2 = \frac{2GM_1}{R_1} (1 + \alpha) \left[ \frac{1}{1 + \alpha^{1/3}} - \frac{1}{4(1 + \alpha)^{1/3}} \right] \left[ 1 - \frac{1}{16(1 + \alpha)^{2/3}} \left( \frac{b}{R_1} \right)^2 \right]^{-1}, \]

where

\[ \alpha \equiv M_2/M_1. \]  

(A6)

It is useful to introduce an angular momentum parameter \( \lambda \equiv \sqrt{E}/(GM^{3/2}) \) in order to characterize off-center mergers, where \( J, E \), and \( M \) are the angular momentum of the two subclusters around the centers of mass, the binding energy between the two, and the total mass, respectively. Please note that \( \lambda \) is the ratio between the actual angular velocity and the angular velocity needed to provide rotational support (see Binney & Tremaine 1987).

From equations (A1) and (A2), we obtain

\[ \lambda = \frac{v_b}{(GM_1 R_1)^{1/2}} \alpha^{3/2} \left[ \frac{1}{1 + \alpha^{1/3}} - \frac{R_1 v^2}{2GM_1 (1 + \alpha)} \right]^{1/2}. \]  

(A7)

Given \( M_1, R_1, \alpha, \) and \( \lambda \), therefore, we are able to calculate \( v \) and \( b \) from equations (A5) and (A7).

REFERENCES

Asai, N., Fukuda, N., & Matsumoto, R. 2004, ApJ, 606, L105
———. 2007, ApJ, 663, 816
Ascasibar, Y., & Markevitch, M. 2006, ApJ, 650, 102
Binney, J., & Tremaine, S. 1987, Galactic Dynamics (Princeton: Princeton Univ. Press)
Blasi, P. 2000, ApJ, 532, L9
Brio, M., & Wu, C. C. 1988, J. Comput. Phys., 75, 400
Brunetti, G., Blasi, P., Cassano, R., & Gabrieli, S. 2004, MNRAS, 350, 1174
Clarke, T., E., Kronberg, P. P., & Bohringer, H. 2001, ApJ, 547, L111
Creston, J. H., et al. 2008, A&A, 487, 431
Dolag, K., Bartelmann, M., & Lesch, H. 1999, A&A, 348, 351
———. 2002, A&A, 387, 383
Ensling, T. A., & Vogt, C. 2003, A&A, 401, 835
Ettori, S., & Fabian, A. C. 2000, MNRAS, 317, L57
Fujita, Y., Takizawa, M., & Sarazin, C. L. 2003, ApJ, 584, 190
Fusco-Femiano, R., Dal Fiume, D., Feretti, L., Giovannini, G., Grandi, P., Matt, G., Moleoni, S., & Santangelo, A. 1999, ApJ, 513, L21
Fusco-Femiano, R., Landi, R., & Orlandini, M. 2005, ApJ, 624, L69
———. 2007, ApJ, 654, L9
Giovannini, G., Tordi, M., & Feretti, L. 1999, NewA, 4, 141
Govoni, F., Murgia, M., Feretti, L., Giovannini, G., Dolag, K., & Taylor, G. B. 2006, A&A, 460, 425
Hirsch, C. 1990, Numerical Computation of Internal and External Flows (New York: John Wiley & Sons)
Hockney, R. W., & Eastwood, J. W. 1988, Computer Simulation Using Particles (London: IOP)
Inoue, S., Aharonian, F. A., & Sugiyama, N. 2005, ApJ, 628, L91
Kempner, J. C., & Sarazin, C. L. 2001, ApJ, 548, 639
Navarro, J. F., Frenk, C. S., & White, S. D. M. 1997, ApJ, 490, 493
Markevitch, M., et al. 2003, ApJ, 586, L19
McCarthy, I. G., et al. 2007, MNRAS, 376, 497
Mentati, F., Jones, T. W., Kang, H., & Ryu, D. 2001, ApJ, 562, 233
Monaghan, J. J. 1992, ARA&A, 30, 543
Murgia, M., Govoni, F., Feretti, L., Giovannini, G., Dallacasa, D., Fanti, R., Taylor, G. B., & Dolag, K. 2004, A&A, 424, 429
Ohno, H., Takada, M., Dolag, K., Bartelmann, M., & Sugiyama, N. 2003, ApJ, 584, 599
Ohno, H., Takizawa, M., & Shibata, S. 2002, ApJ, 577, 658
Peebles, P. J. E. 1980, The Large-Scale Structure of the Universe (Princeton: Princeton Univ. Press)
Ricker, P. M., & Sarazin, C. L. 2001, ApJ, 561, 621
Ritchie, B. W., & Thomas, P. A. 2002, MNRAS, 329, 675
Roc, P. L. 1981, J. Comput. Phys., 43, 357
Roettiger, K., Burns, J. O., & Loken, C. 1996, ApJ, 473, 651
Roettiger, K., Stone, J. M., & Burns, J. O. 1999, ApJ, 518, 594
Roland, J. 1981, A&A, 93, 407
Rossetti, M., & Molendi, S. 2004, A&A, 414, L41
Rybicki, G. B., & Lightman, A. P. 1979, Radiative Process in Astrophysics (New York: John Wiley & Sons)
Ryu, D., & Jones, T. W. 1996, ApJ, 442, 228
Ryu, D., Kang, H., Hallman, E., & Jones, T. W. 2003, ApJ, 593, 599
Ryu, D., Ostriker, J. P., Kang, H., & Cen, R. 1993, ApJ, 414, 1
Sarazin, C. L. 1986, Rev. Mod. Phys., 58, 1
———. 1999, ApJ, 520, 529
Schlickeiser, R., Sievers, A., & Thiemann, H. 1987, A&A, 182, 21
Springel, V., & Farrar, G. R. 2007, MNRAS, 380, 911
Takizawa, M. 1999, ApJ, 520, 514
———. 2000, ApJ, 532, 183
———. 2002, PASJ, 54, 363
———. 2006, PASJ, 58, 925
Takizawa, M., & Naito, T. 2000, ApJ, 535, 586
Takizawa, M., & Kitayama, T. 2000, ApJ, 545, 572
Vikhlinin, A., Kravtsov, A., Forman, W., Jones, C., Markevitch, M., Murray, S. S., & Van Speybroeck, L. 2006, ApJ, 640, 691
Vikhlinin, A., Markevitch, M., & Murray, S. S. 2001, ApJ, 549, L47
Vogt, C., & Ensslin, T. A. 2003, A&A, 412, 373
Wada, K., & Norman, C., A. 2001, ApJ, 547, 172
Zakamska, N., & Narayan, R. 2003, ApJ, 582, 162