TWO-TEMPERATURE INTRA ClUSTER MEDIUM IN MERGING CLUSTERS OF GALAXIES

MOTOKAZU TAKIZAWA

Department of Astronomy, Faculty of Science, Kyoto University, Sakyo-ku, Kyoto 606-8502, Japan; takizawa@kusastro.kyoto-u.ac.jp

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ABSTRACT

We investigate the evolution of the intracluster medium during a cluster merger, explicitly considering the relaxation process between the ions and electrons using N-body and hydrodynamical simulations. When two subclusters collide, a bow shock is formed between the centers of the two substructures and propagates in both directions along the collision axis. The shock primarily heats the ions because the kinetic energy of an ion entering the shock is larger than that of an electron by the ratio of their masses. In the postshock region, the energy is transported from the ions to the electrons via Coulomb coupling. However, since the energy-exchange timescale depends on both the gas density and the temperature, the distribution of the electron temperature becomes more complex than that of the plasma mean temperature, especially in the expanding phase. After the collision of two subclusters, gas outflow occurs not only along the collision axis but also in its perpendicular direction. The gas originally located in the central parts of the subclusters moves in both the parallel and perpendicular directions. Since the equilibrium timescale of the gas along these directions is relatively short, the temperature difference between ions and electrons is larger in the directions tilted at angles of $\pm 45^\circ$ with respect to the collision axis. The electron temperature could be significantly lower than the plasma mean temperature, by at most $\sim 50\%$. The significance of our results for the interpretation of X-ray observations is briefly discussed.

Subject headings: galaxies: clusters: general — hydrodynamics — intergalactic medium — plasmas — shock waves — X-rays: galaxies

1. INTRODUCTION

Clusters of galaxies (CG) contain collisionless particles, galaxies, dark matter, and a diffuse gas component. The gas component is called the intracluster medium (ICM). The ICM is a plasma with temperatures of about $10^8$ K, thus emitting X-rays mainly through thermal bremsstrahlung of electrons (Sarazin 1988). Since the ICM is optically thin for X-rays and spreads over the entire cluster, it can provide good information regarding the physical properties of the cluster.

According to the hierarchical clustering scenario, it is believed that CGs are formed through subcluster mergers and/or absorption of smaller galaxy groups. The rate of mergers is dependent on cosmological parameters, especially on the density parameter, $\Omega_0$ (Richstone, Loeb, & Turner 1992). However, when this is applied to more realistic situations, such as in dealing with observational results (e.g., West & Bothun 1990; Rhee, van Harlem, & Katgert 1991; Jones & Forman 1992; Mohr, Fabricant, & Geller 1993; Bird 1994; Serna & Gerbal 1996; Gurzadyan & Mazure 1998; Solanes, Salvador-Solé, & González-Casado 1999) or cosmological numerical simulations (Evrard et al. 1993; Jing et al. 1995; Crone, Evrard, & Richstone 1996), uncertainties in the dynamical properties of CGs during mergers seriously affects the estimation of $\Omega_0$ (Nakamura, Hattori, & Mineshige 1995).

In studying the evolution of CGs, plenty of numerical simulations have been performed so far using N-body and hydrodynamical codes. Perrenod (1978) was the first to calculate the evolution of a spherically symmetric CG with a standard mesh hydrodynamical code. Recently, Takizawa & Mineshige (1998) calculated the spherical CG and investigated the effect of preheating on ICM density profiles. To investigate the formation and evolution in more realistic situations, many simulations starting from cosmological initial density perturbations have been performed (e.g., Evrard 1990; Katz & White 1993; Metzler & Evrard 1994; Navarro, Frenk, & White 1995; Bryan & Norman 1998; Eke, Navarro, & Frenk 1998; Yoshikawa, Itoh, & Suto 1998). Schindler & Müller (1993) find that characteristic temperature structures occur in merging clusters through shock heating and adiabatic compression and expansion. There are also simulations starting from rather idealized initial conditions designed to study the mergers in detail. Roettiger, Loken, & Burns (1997) study general behaviors of head-on collisions with a high-resolution finite-difference code. The observational consequences are discussed in more detail in Roettiger, Burns, & Loken (1996). Ishizaka (1997) found that the specific energy ratio of ICM and galaxies, $\beta_{\text{spec}}$, is a good indicator of merging clusters and that it can be used to determine the phase of the merger. Off-center collisions have been investigated in particular by Ricker (1998) and Takizawa (1999).

In recent years, rather complex electron-temperature structures have been revealed in merging CGs through X-ray observations (e.g., Fujita et al. 1996; Honda et al. 1996; Churazov et al. 1999; Donnelly et al. 1998, 1999; Davis & White 1998; Markevitch, Sarazin, & Vikhlinin 1999). In addition, Hanami et al. (1999) find that the energy of the iron K$\alpha$ line is higher than expected from the temperature in the ionization equilibrium plasma in SC 1329–313 in the Shapley concentration. These results imply that there exists strong bulk-flow motion in the ICM of merging clusters. In addition to $\beta_{\text{spec}}$ temperature maps of the ICM also contain important information on mergers. Some numerical simulations have tried to explain these observational results; Roettiger, Burns, & Pinkney (1995) tried modeling A2256, and Ishizaka & Mineshige (1996) Coma. A more detailed “anatomy” is presented for A754 by Roettiger, Stone, & Mushotzky (1998).

As mentioned above, the temperature structure of the ICM provides us with very important information about
the phase of the mergers. However, until now we could only find the electron temperature through X-ray observations. In all the previous studies the discussion is based on the assumption that the observed electron temperature \((T_e)\) is the same as the plasma mean temperature \((\bar{T})\). However, we believe that this assumption is problematic. The shock primarily heats the ions because the kinetic energy of an ion entering the shock is larger than that of an electron by the ratio of their masses. In addition, the equilibrium timescale between electrons and ions \((t_{eq})\) in a typical CG becomes \(10^{8} \sim 10^{9}\) yr if Coulomb coupling is considered as the relaxation process (see Spitzer 1962):

\[
t_{eq} = 2.0 \times 10^8 \times \left(\frac{\ln \Lambda}{40}\right)^{-1} \left(\frac{n_i}{10^{-3} \text{ cm}^{-3}}\right)^{-1} \left(\frac{T_e}{10^8 \text{ K}}\right)^{3/2},
\]

where \(n_i\) is the ion density, \(T_e\) is the electron temperature, and \(\ln \Lambda\) is the Coulomb logarithm. Indeed, some CGs have larger electron-temperature gradients than those expected in numerical simulations (Markevitch 1996). A two-temperature model of the ICM is one solution to this discrepancy (Fox & Loeb 1997; Chieze, Alimi, & Teyssier 1998; Ettori & Fabian 1998; Takizawa 1998). In cluster mergers, it is believed that shock heating plays an important role in the evolution of the ICM. Furthermore, according to previous studies of mergers, the timescale for surviving complex temperature structures is an order of \(10^9\) yr, which is comparable to \(t_{eq}\). Therefore, the difference between \(T_e\) and \(\bar{T}\) should be properly considered in a study of the evolution of \(T_e\) structure in mergers.

To investigate the evolution of two-temperature ICM in cluster mergers, we perform \(N\)-body and hydrodynamical simulations considering the relaxation process between ions and electrons. For the hydrodynamical part, we choose the smoothed-particle hydrodynamics (SPH) method. Although SPH codes are not better than mesh codes for problems with shocks, their Lagrangian nature is suitable for dealing with the electron-ion coupling simply.

The rest of this paper is organized as follows. In \(\S 2\) we estimate analytically some physical quantities relevant to cluster mergers. In \(\S 3\) we describe the adopted numerical methods and initial conditions for our simulations. In \(\S 4\) we present the results. In \(\S 5\) we summarize the results and discuss their implications.

2. BASIC CONSIDERATIONS OF EQUILIBRIUM TIMES

In this section, we consider the scaling laws of CGs and the initial conditions of cluster mergers. We assume that the power spectrum of the cosmological density fluctuation field, \(P(k)\), obeys a power law, \(P(k) \propto k^n\). According to the spherical model for nonlinear collapse (see Peebles 1980; Padmanabhan 1993; etc.), the virial radius \((r)\), the density \((\rho)\), and the virial temperature \((T)\) of the dark halo obey the scaling laws

\[
r \propto M^{(n+5)/6},
\]

\[
\rho \propto M^{-(n+3)/2},
\]

\[
T \propto M^{1-(n)/6},
\]

where \(M\) is the total mass of the halo. The effective index of \(P(k)\) of the cold dark matter (CDM) is \(-1\) to \(-2\) in the galactic and CG scale. We consider the case of \(n = -2\). Then the above relations are

\[
r \propto M^{1/2},
\]

\[
\rho \propto M^{-1/2},
\]

\[
T \propto M^{1/2}.
\]

Using these relations and the gas fraction, \(f_g\), we obtain the scaling laws of the equilibrium timescale between ions and electrons, \(t_{eq} \propto (f_g \rho)^{-1/3} t^{3/2}\)

\[
t_{eq} \propto f_g^{-1} M^{5/4}.
\]

Thus, the more massive the CG is, the longer the \(t_{eq}\) if \(f_g\) is constant. However, the observational results for the luminosity-temperature (LT) relation suggest that \(f_g\) depends on \(M\) (Mushotzky 1994; Edge & Stewart 1991; David et al. 1993; Fukazawa 1997; Mushotzky & Scharf 1997; Markevitch 1998; Arnaud & Evrard 1999). Since \(L \propto (f_g \rho)^2 T^{1/2} R^3\), \(f_g \propto M^{3/8}\) if we assume \(L \propto T^3\). Then we obtain

\[
t_{eq} \propto M^{7/8}.
\]

Thus \(t_{eq}\) is also longer in more massive CGs in this case, although the dependence is slightly weaker than in the case where \(f_g\) is constant.

Next, let us consider the merger of two subclusters with masses \(M_1\) and \(M_2\), respectively. The collision velocity \(V\) is estimated as follows. From the dynamical energy conservation law, we obtain

\[
\frac{GM_1 M_2}{2R} = \frac{GM_1 M_2}{r_1 + r_2} + \frac{1}{2} M V^2,
\]

where \(G\) is the gravitational constant, \(r_1\) and \(r_2\) are the virial radii of each subcluster, \(R\) is the maximum expansion radius of the whole system, and \(M = M_1 M_2/(M_1 + M_2)\) is the reduced mass. Since the maximum expansion radius is twice the virial radius, using equation (5) we obtain

\[
R = 2 \left(\frac{M_1 + M_2}{M_1}\right) \frac{1}{r_1}. \tag{11}
\]

The sound speed of subcluster 1, \(c_1\), is

\[
c_1^2 = \gamma (\gamma - 1) \frac{GM_1}{2r_1}, \tag{13}
\]

where \(\gamma\) is the ratio of specific heat. Thus, the Mach number, \(\mathcal{M}_1 = V/c_1\), is

\[
\mathcal{M}_1^2 = \frac{4(1 + \alpha)}{\gamma (\gamma - 1)} \left[ \frac{1}{1 + \alpha^{1/2}} - \frac{1}{4(1 + \alpha^{1/2})} \right]. \tag{14}
\]

From the scaling relation given in equation (5), the Mach number for subcluster 2, \(\mathcal{M}_2\), is

\[
\mathcal{M}_2^2 = \alpha^{-1/2} \mathcal{M}_1^2. \tag{15}
\]

Figure 1 shows \(\mathcal{M}_1^2\) (solid line) and \(\mathcal{M}_2^2\) (dotted line) with respect to \(\alpha\) when \(\gamma = 5/2\). For major mergers (\(\alpha > 0.1\),
these values are greater than $\sim 2$. Thus, in a CDM universe, supersonic collisions are quite natural in major mergers.

The Mach number gives the maximum temperature difference between ions and electrons in the postshock regions. The specific inertial energies of ions and electrons are, in the limit of no mutual energy exchange,

$$\frac{k_B T_i}{\mu m_p} \sim c^2 + V^2,$$

$$\frac{k_B T_e}{\mu m_p} \sim c^2,$$

where $\mu$ is the mean molecular weight, $m_p$ is the proton mass, $k_B$ is the Boltzmann constant, and $c$ is the sound speed in the preshock regions. Thus, the difference between $T_e$ and $T \sim (T_i + T_e)/2$ is

$$\frac{T_e}{T} \sim \frac{1}{1 + \alpha^2/2}.$$  

From Figure 1, we expect that $T_e/T$ is less than about 0.6. We will confirm this later by numerical simulations.

3. THE SIMULATIONS

In the present study, we consider a CG consisting of two components: collisionless particles corresponding to galaxies and dark matter (DM), and two-temperature 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 our simulations.

3.1. The Numerical Method

To solve the hydrodynamical equations for the gas component, we used the smoothed-particle hydrodynamics (SPH) method (see Monaghan 1992). Although more accurate finite-difference methods handle shocks better than SPH methods, the fully Lagrangian nature of SPH is convenient for calculating the evolution of two-temperature plasmas, since the coupling between ions and electrons can be treated more simply in a Lagrangian view, as follows.

As the standard SPH codes for one-temperature fluid, we solve the continuity equation, the momentum equation, and the thermal energy equation with artificial viscosity to treat

the shocks. In addition to these equations, we solve one more equation for the normalized electron temperature, $\bar{T}_e \equiv T_e/T$, where $T_e$ is the electron temperature and $T$ is the plasma mean temperature. We assume that artificial viscous heating is effective only for ions and that only the Coulomb coupling is considered in the relaxation process. Then the Lagrangian time evolution of $\bar{T}_e$ is (see Appendix),

$$\frac{d\bar{T}_e}{dt} = \frac{T_i - \bar{T}_e}{t_{eq}} - \frac{\bar{T}_e}{u} Q_{vis} / u,$$  

where $\bar{T}_i \equiv T_i/T$ is the normalized ion temperature, $Q_{vis}$ is the artificial viscous heating per unit mass, and $u$ is the thermal energy per unit mass. For the numerical integration of equation (19), we should take a time step that is sufficiently shorter than both the Courant and viscous timescale as usual SPH codes (see Monaghan 1992), and also $t_{eq}$. In our simulations, however, the latter is much shorter than the former in the central high-density regions. Thus, the integration with a time step relevant to $t_{eq}$ takes a huge amount of computational time. In addition, since the two-temperature nature is not important in such regions, the time-step control mentioned above is not a good choice. To integrate equation (19) with the Courant and viscous time-step control, we use the results of Fox & Loeb (1997) as follows. Let us integrate equation (19) from $t = t_0$ to $t = t_0 + \Delta t$. First, we integrate equation (19), regarding the second term on the right-hand side as being negligibly small. In this case, neglecting the small change of the Coulomb logarithm in $t_{eq}$, we can integrate analytically (Fox & Loeb 1997):

$$\frac{\Delta t}{t_{2s}} = \frac{n_i}{n_i + n_e} \left[ F(\bar{T}_e(t_0 + \Delta t)) - F(\bar{T}_e(t_0)) \right],$$  

where

$$F(x) = \ln \left( \frac{1 + \sqrt{x}}{1 - \sqrt{x}} \right) - 2\sqrt{x} \left( 1 + \frac{x}{3} \right),$$  

and $t_{2s} \equiv t_{eq}(t_0)\bar{T}_e(t_0)^{-3/2}$. Next, we add the contribution from the second term to $\bar{T}_e(t_0 + \Delta t)$ using the second-order Runge-Kutta method. If the time step is chosen correctly, the displacement of the position of each SPH particle is sufficiently smaller than the spatial resolution and $Q_{vis}/u$ is sufficiently smaller than unity. Thus, the above method can follow the evolution of $\bar{T}_e$ efficiently and with reasonable accuracy.

Gravitational forces are calculated by the Barnes-Hut tree algorithm (Barnes & Hut 1986) and softened using a Plummer potential profile $\Phi \propto (r^2 + \epsilon^2)^{-1}$, where $\epsilon$ is the softening parameter. We set $\epsilon$ be 1/10 of the initial core radius of the smaller subcluster in the simulation. Tree structure is also used to search for nearest neighbors in SPH calculations (Hernquist & Katz 1989). In the standard one-dimensional shock-tube test, our code nicely reproduced the analytic solution. The shock and contact discontinuity are broadened over a range of about 3 times the smoothing length. During the calculations described below, the total energy is conserved within $\sim 0.4\%/(1000 \text{ steps})$.

3.2. Models and Initial Conditions

We consider mergers of two virialized subclusters of galaxies with masses $M_1$ and $M_2$. The initial configuration of each subcluster is described as follows. The spatial distribu-
tion of the DM in each subcluster is represented by the King distribution,

$$\rho_{\text{DM}}(r) = \begin{cases} 
\rho_0 \left[ 1 + (r/r_c)^2 \right]^{-3/2}, & r \leq r_{\text{out}}, \\
0, & r > r_{\text{out}}, 
\end{cases} \quad (22)$$

where $\rho_0$ is the central density, $r_c$ is the core radius, and $r_{\text{out}}$ is the radius of the cluster. We set $r_{\text{out}} = 5r_c$. This is nearly equal to the virial radius of the King distribution. The value of $\rho_0$ is determined by the total mass of the DM ($M_{\text{DM}}$),

$$\rho_0 = \frac{M_{\text{DM}}}{4\pi r_c^3} \left( -\frac{t_{\text{out}}}{\sqrt{t_{\text{out}}^2 + 1}} + \ln \frac{t_{\text{out}}}{\sqrt{t_{\text{out}}^2 + 1}} + 1 \right)^{-1}, \quad (23)$$

where $t_{\text{out}} = r_{\text{out}}/r_c$. We assume the velocity distribution of the DM particles to be an isotropic Maxwellian. Then, from the second moment of the collisionless Boltzmann equation, the one-dimensional velocity dispersion at $r$, $\sigma^2_1(r)$, is (e.g., Binney & Tremaine 1987)

$$\sigma^2_1(r) = \frac{1}{\rho_{\text{DM}}(r)} \int_0^{r_{\text{out}}} \rho_{\text{DM}}(r') \frac{GM_r}{r'^2} dr', \quad (24)$$

where $M_r$ is the mass inside $r$. We assume that the initial ICM temperature is isothermal and equal to the virial temperature at $r_{\text{out}}$. The ICM is initialized in hydrostatic equilibrium within a cluster potential of the DM and the ICM itself.

We set the initial conditions as follows. Two subclusters are initially placed at rest with a separation between the center of each subcluster, $r_{\text{sep}}$. The coordinate system is taken in such a way that the center of mass is at rest in the origin. Two representative cases are examined in the present study: a collision between two clusters of equal mass (run A) and an absorption of a smaller cluster by a larger one (run B). The scaling law between the two subclusters and $r_{\text{sep}}$ is determined as in the case of $P(k) \propto k^{-2}$ in §2. In run A, the two subclusters have the same masses: $M_1 = M_2 = 0.5 \times 10^{15} M_\odot$. We set the core radii to be $r_c = 0.2$ Mpc. Each subcluster consists of 5000 collisionless particles and 5000 SPH particles. The total gas mass fraction is 10% for both clusters. In run B, the mass ratio is $M_1 : M_2 = 4 : 1$. The larger cluster has the same mass as that in run A, and the particle numbers of the smaller cluster are one-fourth those of the larger one. The parameters in our calculations are summarized in Table 1.

4. RESULTS

4.1. Time Evolution of the Total Features during the Mergers

Figure 2 shows the time evolution of various physical quantities representing the total features of the system for run A: from top to bottom, the total luminosity ($L$), the emissivity-weighted mean electron temperature ($T_{\text{e,ew}}$), the square root of the velocity dispersion parallel to the collision axis ($\sigma_\parallel$), and the specific energy ratio of DM and the ICM ($\beta_{\text{spec}}$).

When calculating $L$, we assume thermal bremsstrahlung of the optically thin plasma (Rybicki & Lightman 1979) and neglect the line emissions.

Around the collision, $L$, $T_{\text{e,ew}}$, and $\sigma_\parallel$ all rise, but on different timescales. Since the rise of $T_{\text{e,ew}}$ is quicker than that of $\sigma_\parallel$, $\beta_{\text{spec}}$ has two peaks in the time evolution. The minimum of $\beta_{\text{spec}}$ is coincident with the maximum of $T_{\text{e,ew}}$.

Similar results are also seen in run B (Fig. 3). However, the amplitudes of the time fluctuations are smaller than in run A.

Since the emissivity is proportional to the square of the gas density, the emissivity-weighted mean temperature (EWMT) more closely resembles the temperature in the central high-density regions, where $t_{\text{eq}}$ is $\sim 10^8$ yr or less. Thus, the difference in EWMTs between electrons and ions is practically negligible. Indeed, this difference becomes $\sim 1\%$ even at maximum in both runs. Therefore, our results agree qualitatively with those of the previous one-temperature simulations, such as that of Ishizaka (1997).

| Parameter | Run A | Run B |
|-----------|-------|-------|
| $M_1/M_2$ ($10^{15} M_\odot$) | 0.5/0.5 | 0.5/0.125 |
| $r_{\text{sep}}$ (Mpc) | 0.2/0.2 | 0.2/0.1 |
| $k_b T_1/k_b T_2$ (keV) | 4.78/4.78 | 4.78/2.39 |
| $r_{\text{sep}}$ (Mpc) | 4.0 | 3.3 |
| $t_0$ | 0.1 | 0.1 |
| $\epsilon$ (Mpc) | 0.02 | 0.01 |
| $N_1/N_2$ (SPH) | 5000/5000 | 5000/1250 |
| $N_1/N_2$ (DM) | 5000/5000 | 5000/1250 |

TABLE 1

MODEL PARAMETERS

| Parameter | Run A | Run B |
|-----------|-------|-------|
| $M_1/M_2$ ($10^{15} M_\odot$) | 0.5/0.5 | 0.5/0.125 |

| parameter | Run A | Run B |
|-----------|-------|-------|
| $M_1/M_2$ ($10^{15} M_\odot$) | 0.5/0.5 | 0.5/0.125 |

| parameter | Run A | Run B |
|-----------|-------|-------|
| $M_1/M_2$ ($10^{15} M_\odot$) | 0.5/0.5 | 0.5/0.125 |

| parameter | Run A | Run B |
|-----------|-------|-------|
| $M_1/M_2$ ($10^{15} M_\odot$) | 0.5/0.5 | 0.5/0.125 |
4.2. Time Evolution of the Temperature Distribution of Run A

As seen in § 4.1, the two-temperature nature of the ICM does not manifest itself in the EWMT of the CG. This is not the case, however, when we can resolve the spatial distribution of the ICM temperature. Let us first examine the case of run A.

First of all, let us examine the $T$ map and the gas velocity field and discuss the dynamical evolution of run A. Figure 4 shows snapshots of the X-ray surface brightness (contours) and emissivity-weighted $T$ (colors) distribution seen from the direction perpendicular to the collision axis. X-ray surface brightness contours are equally spaced on a logarithmic scale and separated by a factor of 7.4. The blue, green, yellow, and red colors correspond to 10, 15, maximum of and the minimum of respectively. The times are listed above each panel; $t = 4.1, 4.3, 4.6$, and 4.85 Gyr correspond to the first maximum of $\beta_{\text{spec}}$, the minimum of $\beta_{\text{spec}}$, the second maximum of $\beta_{\text{spec}}$, and the minimum of $T_{\text{ew}}$, respectively. The times are listed above each panel; $t = 4.1, 4.3, 4.6$, and 4.85 Gyr correspond to the first maximum of $\beta_{\text{spec}}$, the minimum of $\beta_{\text{spec}}$, the second maximum of $\beta_{\text{spec}}$, and the minimum of $T_{\text{ew}}$, respectively. Figure 5 shows the evolution of the gas velocity field, seen from the same direction at the same times as in Figure 4. The longest vector corresponds to the maximum velocity listed below each panel.

When two subclusters just contact each other at $t = 3.5$ Gyr (Fig. 4, top left panel), $T$ rises slightly at the interface between the two subclusters due to the adiabatic compression, but we can clearly distinguish two individual clusters through the X-ray image. Then the two approach each other, and we see double peaks in the X-ray emissivity profile of “one cluster” at $t = 4.1$ Gyr (Fig. 4, top middle panel). Just between the peaks, $T$ rises to $\sim 20$ keV due to a shock. The shock front is perpendicular to the collision axis. However, since $T$ still remains at nearly the initial value ($\sim 5$ keV) around the peaks of the X-ray image, the EWMT does not increase so much that $\beta_{\text{spec}}$ becomes rather large (see Fig. 2). At the most contracting phase ($t = 4.3$ Gyr; Fig. 4, top right panel), two peaks merge into one peak in the X-ray image and the X-ray image elongates in the directions perpendicular to the collision axis. Since the high-temperature ($\sim 20$ keV) region is located around the X-ray peak, the EWMT becomes maximum (see Fig. 2). Then the cluster expands and two shocks propagate in opposite directions along the collision axis ($t = 4.6$ Gyr; Fig. 4, bottom left panel). We can clearly see two “lens-shaped” high-temperature regions associated with the shocks in the $T$ structure. We can also see the cooler region ($\sim 5$ keV) spread in the directions perpendicular to the collision axis. These characteristic temperature structures reflect the fact that the gas outflow occurs not only along the collision axis but also in perpendicular directions (see Fig. 4). In these directions the gas is effectively cooled by the adiabatic expansion and emits more X-rays than the hottest regions. Thus, the EWMT decreases, although high-temperature regions still exist. At the most expanding phase ($t = 4.85$ Gyr; Fig. 4, bottom middle panel), the X-ray image elongates along the collision axis. The shocks reach a very low density region. We see an X-shaped region with a slightly high temperature (indicated by a light blue color). This structure is also due to adiabatic expansion associated with the gas outflows. Then the system contracts again and settles down to the spherical structure ($t = 7$ Gyr; Fig. 4, bottom right panel). Owing to the subsonic accretion of the gas along the collision axis, slightly high temperature regions are created on the both sides of the collision axis (see Fig. 5).

Note that there can be complex temperature structures in the ICM when X-ray images do not have definite substructures. It is certain that image elongation occurs in these case, but it is severely affected by the viewing angles. Thus, the temperature map is a more suitable diagnostic for merging clusters.

Next, let us examine the $T$ map and discuss the difference between $T$ and $T_e$. Figure 6 shows the same conditions as Figure 4, but for $T_e$. Although in the contracting phase ($t \leq 4.3$ Gyr) $T_e$ maps are similar to the $T$ maps, in the expanding phase ($t > 4.3$ Gyr) the two maps are rather different. To see the difference between $T$ and $T_e$ in detail, we make snapshots of the distribution of $T_e = T_e/T$. Figure 7 shows the $T_e$ maps of run A. The red, yellow, green, and blue colors correspond to $T_e \sim 0.1, 0.3, 0.5$, and 0.7, respectively.

In the contracting phase ($t \leq 4.3$ Gyr), the difference between $T$ and $T_e$ is rather small. At $t = 4.1$ Gyr, $T_e$ is slightly lower than $T$ at the shock. Overall, there are no significant two-temperature regions. Then, at $t = 4.3$ Gyr, split two-temperature regions appear due to the propagation of shocks. Since $t_{\text{eq}}$ is longer in the outer parts where density is low, two-temperature regions are more spread out than in the central parts. The two-temperature nature is less important in this phase than in the later phase, however, because the shocks are nearly standing shocks and $t_{\text{eq}}$ is shorter in the central parts.

On the other hand, in the expansion phase ($4.3 \leq t \leq 4.85$ Gyr), the distributions of $T$ and $T_e$ are clearly qualitatively different. From Figure 7 we find two common features at $t = 4.6$ and 4.85 Gyr. One feature is that $T_e$ is significantly lower (by $\sim 50\%$) than $T$ in the postshock high-$T$ regions. These regions are spread out on a $\sim 0.5$ Mpc scale behind
Fig. 4.—Snapshots of X-ray surface brightness (contours) overlaid with emissivity-weighted $T$ (colors) distribution seen from the direction perpendicular to the collision axis for run A. X-ray surface brightness contours are equally spaced on a logarithmic scale and separated by a factor of 7.4. The blue, green, yellow, and red colors correspond to $k_B T \sim 5, 10, 15, \text{ and } 20$ keV, respectively. The times are listed above each panel.
the shocks. The other feature is that the difference between $T$ and $T_e$ is larger in the directions tilted at angles of $\pm 45^\circ$ with respect to the collision axis. This is also due to the gas outflows, both along the collision axis and in its perpendicular directions. The gas originally located in the central part, the $t_{eq}$ of which is relatively short, moves outward in these directions. Note that $t_{eq}$ is hardly changed by the adiabatic compression and expansion when $\gamma = 5/3$. Thus, the difference between $T$ and $T_e$ is less in these directions. As a result, we see four peaks in the $T_e$ map at $t = 4.6$ Gyr, and the X-shaped high-temperature region becomes less contrasted in the $T_e$ map at $t = 4.85$ Gyr. Note that high-temperature regions around $\sim 20$ keV cannot be seen in the $T_e$ maps of the expanding phase. Therefore, it is probable that we underestimate the collision velocity by $\sim 30\%$ if we estimate it from the hottest regions in the $T_e$ map in the expanding phase. The two-temperature nature still remains at $t = 7$ Gyr in the outer parts, as in the case of spherical models (Fox & Loeb 1997; Takizawa 1998).

4.3. Time Evolution of the Temperature Distribution of Run B

First, we describe the dynamical evolution of run B through the $T$ map and the gas velocity field. Figure 8 shows the same conditions as Figure 4, but for run B. The temperature color scale is adjusted for run B. The blue, green, yellow, and red colors correspond to $k_B T \sim 4, 7, 10,$ and 13 keV, respectively. The times are listed above each panel; $t = 3.9, 4.1, 4.3,$ and 4.75 Gyr correspond to the first maximum of $\beta_{spec}$, the minimum of $\beta_{spec}$, the second maximum of $\beta_{spec}$, and the minimum of $T_{spec}$, respectively. Figure 9 shows the same conditions as Figure 5, but for run B.

When the two subclusters approach each other, a bow shock with an arc shape is formed just between them. Thus, the arc-shaped hot region is seen between the two peaks of the X-ray image at $t = 3.9$ Gyr (Fig. 8, top middle panel). This region clearly separates the component of the larger subcluster ($\sim 5$ keV) from the component of the smaller subcluster ($\sim 2.5$ keV). Then the two peaks merge into one triangular image in the X-ray ($t = 4.1$ Gyr; Fig. 8, top right panel). The hottest temperature region ($\sim 13$ keV) is located in the peak of the X-ray image, and the hot region associated with the bow shock is seen elongated slantingly backward with respect to the motion of the smaller subclusters. Then the gas expands and the two shocks propagate outward along the collision axis. The one in front of the motion of the smaller subcluster (at the left of the figure) is arc-shaped, and the other is rather flat ($t = 4.3$ Gyr; Fig. 8, bottom left panel). At the most expanding phase ($t = 4.75$
Fig. 6. Same as Fig. 4, but for $T_2$ instead of $T$ for run A.
Fig. 7. Snapshots of the distribution of normalized electron temperature, $\tilde{T}_e = T_e / T_i$, for run A. The red, yellow, green, and blue colors correspond to $\tilde{T}_e \sim 0.1, 0.3, 0.5$, and 0.7, respectively.
Fig. 8. Same as Fig. 4, but for run B. Temperature color scale is adjusted for run B. The blue, green, yellow, and red colors correspond to $k_B T \sim 4$, 7, 10, and 13 keV, respectively.
FIG. 9.—Same as Fig. 5, but for run B

Gyr; Fig. 8, bottom middle panel) the X-ray image elongates along the collision axis and the shocks reach very low density regions. The elongation of the image is more significant at the left of the figure. Then the system contracts again and settles down to a spherical structure (\( t = 7 \) Gyr; Fig. 8, bottom right panel). Accretion flows are seen only from the front side with respect to the motion of the smaller subcluster.

Next, we describe the two-temperature nature of run B through the \( T_e \) map and the \( T_i \) map. Figure 10 shows the same conditions as Figure 8, but for \( T_e \), and Figure 11 shows the same conditions as Figure 7, but for run B. As in the case of run A, in the contracting phase (\( t \leq 4.1 \) Gyr) the \( T_e \) distribution is similar to that of \( T_i \), although \( T_e \) is slightly lower than \( T_i \) around the shock. On the other hand, in the expanding phase (\( t > 4.1 \) Gyr) a substantial deviation emerges, especially around the hot regions associated with the shocks. At \( t = 4.3 \) Gyr in the \( T_i \) map, the two hot regions have almost the same temperature (\( \sim 13 \) keV), whereas in the \( T_e \) map, the former (with respect to the motion of the smaller subcluster) hot region becomes less than \( \sim 10 \) keV. On the other hand, in the backward hot region, the \( \sim 13 \) keV component still remains. In this region \( t_{eq} \) is rather short because the gas there originates from the smaller subcluster (see the scaling relation given in eq. [8]). At \( t = 4.75 \) Gyr, as in the case of run A, \( T_i \) is low in the directions tilted at angles of \( \pm 45^\circ \) with respect to the collision axis, although this tendency is less significant than in run A.

5. SUMMARY AND DISCUSSION

We investigate the evolution of the ICM during mergers, considering the relaxation process between the ions and electrons. From the simple analytical estimation, we find that in the CDM universe the equilibrium timescale between ions and electrons is longer in more massive CGs even if the gas fraction is dependent on the CG mass in such a way that the X-ray luminosity is proportional to the cube of the temperature. We estimate the collision velocity and show that supersonic collisions are quite natural in a CDM universe in major mergers and that the electron temperature can be less than half the plasma mean temperature in the postshock regions. Temperatures estimated from X-ray observations can be significantly lower than ion temperature.

We carry out numerical simulations of the mergers by \( N \)-body and hydrodynamical simulations, incorporating the relaxation process between the ions and electrons. To solve the evolution of the normalized electron temperature, we adopt the results of Fox & Loeb (1997). The difference between the emissivity-weighted mean temperatures of ions
FIG. 10. Same as Fig. 6, but for run B.
Fig. 11 — Same as Fig. 7, but for run B
and electrons is practically negligible during the merger; the discrepancy is \( \sim 1\% \) even at maximum. On the other hand, the spatial distribution of the electron temperature differs significantly from that of the plasma mean temperature, especially in low-density regions in the expanding phase. In this phase, the electron temperature is at most \( \sim 50\% \) lower than the plasma mean temperature in the postshock hot regions. In addition, the temperature difference between them is more enhanced in directions tilted at angles of \( \pm 45^\circ \) with respect to the collision axis. When the two subclusters have different masses, the hot region located at the former position (with respect to the direction of the motion of the smaller subcluster) has a lower electron temperature than that located backward in the expanding phase.

Recently, Markevitch et al. (1999) estimated the subcluster collision velocity of Cygnus A using the electron-temperature map obtained by ASCA. From this temperature map and the ROSAT PSPC image, it is likely that Cygnus A is in just such a contracting phase as run A at \( t = 4.1 \) Gyr. The two-temperature nature is almost negligible in this case. However, collision velocities are likely to be underestimated by \( \sim 30\% \) if we use the electron temperature of the hottest regions of merging clusters in the expanding phase, such as in SC 1329–313 (Hanami et al. 1999).

Although we consider only head-on collisions for simplicity in the present calculations, off-center collisions should be investigated to model more realistic situations. Ricker (1998) find that spiral bow shocks occur in off-center mergers, which probably produce rather complex electron-temperature distributions. We will investigate this issue in a future work (Takizawa 1999).

We consider only the classical Coulomb coupling for the relaxation process between ions and electrons. It is possible, however, that in the ICM more efficient relaxation processes could be effective (McKee & Cowie 1977; Pistinner, Levinson, & Eichler 1996). In this case the equilibrium time-scale could be shorter than the value given by equation (1). Therefore, the temperature difference between ions and electrons could be less than our results. If a magnetic field exists in the ICM, it is possible that electrons are also significantly heated by shocks associated with MHD instabilities. In this case the temperature difference between ions and electrons could also be less.

We neglect the heating process from the galaxies to the ICM. If a substantial amount of thermalized hot gas is injected into the ICM from the galaxies, the temperature difference could be less. It is possible that mergers of CGs activate star formation in the member galaxies through galaxy-galaxy interaction, galaxy-ICM interaction, etc. (Caldwell et al. 1993). In this case the effect would not be negligible, but the effect on the ICM temperature distribution is rather sensitive to the detailed modeling of the star formation activity (Fujita et al. 1999).

For simplicity, we adopt the SPH method for the hydrodynamical part of the present simulations to treat the coupling between ions and electrons. However, the resolution of the shock in SPH is not better than in high-resolution finite-difference codes such as TVD, PPM, CIP, etc., because artificial viscosity is used to treat the shocks in SPH. Furthermore, bulk viscosity is included in our code and possibly affects the electron temperature slightly. Although we believe that our simulations can follow the qualitative behavior, it would be useful to simulate the two-temperature ICM during mergers with finite-difference codes.

APPENDIX

DERIVATION OF THE EQUATION FOR THE NORMALIZED ELECTRON TEMPERATURE

In the SPH calculations, shocks are treated through artificial viscosity. We assume that artificial viscous heating is effective only for ions and that only Coulomb coupling is considered in the relaxation process. Then the Lagrangian time evolution of the electron temperature \( T_e \) and the mean temperature \( \bar{T} = (n_e T_e + n_i T_i)/(n_e + n_i) \) are

\[
\frac{dT_e}{dt} = (\gamma - 1) \frac{T_e}{n} \frac{dn}{dt} + \frac{T_i - T_e}{T_{\text{eq}}} , \tag{A1}
\]

\[
\frac{d\bar{T}}{dt} = (\gamma - 1) \frac{T}{n} \frac{dn}{dt} + 2 \frac{\mu m_p}{3 k_B} Q_{\text{vis}} , \tag{A2}
\]

where \( T_i \) is the ion temperature, \( n \) is the gas density, \( \gamma = 5/3 \) is the ratio of specific heats, and \( Q_{\text{vis}} \) is the artificial viscous heating per unit mass. We use the artificial viscosity in the form described in Monaghan (1992), which is the one most commonly used. Then the explicit expression of \( Q_{\text{vis,i}} \) for the \( i \)th SPH particle is

\[
Q_{\text{vis,i}} = \frac{1}{2} \sum_j m_j \Pi_{ij} (v_i - v_j) \cdot \nabla_i W_{ij} , \tag{A3}
\]

where \( \Pi_{ij} \) is given by

\[
\Pi_{ij} = \begin{cases} \frac{-\alpha \tilde{e}_{ij} \mu_{ij} + \beta \mu^2_{ij}}{\bar{\rho}_{ij}}, & \text{for } (v_i - v_j) \cdot (r_i - r_j) < 0 , \\ 0, & \text{for } (v_i - v_j) \cdot (r_i - r_j) > 0 , \end{cases} \tag{A4}
\]

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and
\[ \mu_{ij} = \frac{h_i (v_i - v_j) \cdot (r_i - r_j)}{(r_i - r_j)^3 + \eta^2}. \] (A5)

In the above expressions, \( m_i, v_i, \) and \( r_i \) are the mass, velocity, and position of the \( j \)th SPH particle, respectively, and \( c_{ij}, \rho_i, \) and \( h_i \) are the averages of the sound speed, density, and smoothing length of particles \( i \) and \( j \), respectively. We set \( \alpha = 1, \beta = 2, \) and \( \eta^2 = 0.01 \), which are typical values in SPH calculations. \( V_i W_{ij} \) is the gradient of the \( j \)th particle's kernel function at \( r_i \).

Introducing the temperature normalized by \( T \), i.e., \( T_e \equiv (T / T) \) and \( T_i \equiv (T_i / T) \), we find
\[ \frac{dT_e}{dt} = \frac{\bar{r}_e - r_e}{\bar{r}_e} - \frac{Q_{vis}}{u} \] (A6),

where \( u = 3k_B T / (2m_p) \) is the thermal energy per unit mass.

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