1. INTRODUCTION

Numerical simulations of galaxy clusters have not yet succeeded in producing objects with properties identical to those of observed galaxy clusters. The most serious discrepancies with observations are in the cores (e.g., Nagai et al. 2007; Borgani & Kravtsov 2011; Skory et al. 2013). Within the central ∼100 kpc of clusters whose central cooling time is less than a Hubble time, radiative cooling of intracluster gas in the simulations tends to be too efficient, leading to the overproduction of young stars and excessive entropy levels in the core gas, as higher-entropy material flows inward to replace the gas that has condensed. The resulting cluster cores in such simulations consequently have temperature profiles that decline from ∼10 kpc outward, in stark disagreement with the observed temperature profiles of cool-core clusters, which rise in the ∼10−100 kpc range.

Feedback from a central active galactic nucleus (AGN) helps solve this problem because it slows the process of cooling. However, the core structures of clusters produced in cosmological simulations depend sensitively on the details of the implementation of feedback, the radiative cooling scheme, and numerical resolution. For example, Dubois et al. (2011) simulated the same cosmological galaxy cluster with two different mechanisms for feedback and two different cooling algorithms, one with metal-line cooling and one without. The core structures of the resulting clusters were quite diverse, with kinetic jet feedback and metal-free cooling producing the most realistic-looking cores. But discouragingly, allowing cooling via metal lines had the effect of supercharging AGN feedback and producing a cluster core with excessively high entropy.

Another oft-proposed solution to the core-structure problem in galaxy clusters is thermal conduction. It has been invoked many times to mitigate the prodigious mass flows predicted for cool-core clusters lacking a central heat source (e.g., Tucker & Rosner 1983; Bertschinger & Meiksin 1986; Narayan & Medvedev 2001) but cannot be in stable balance with cooling (e.g., Bregman & David 1988; Soker 2003; Guo et al. 2008) and does not completely compensate for radiative cooling in all cluster cores (Voigt et al. 2002; Zakamska & Narayan 2003). Nevertheless, the temperature and density profiles of many cool-core clusters are tantalizingly close to being in conductive balance, suggesting that AGN feedback is triggered only when inward thermal conduction, perhaps assisted by turbulent heat transport, fails to compensate for radiative losses from the core (e.g., Voit 2011).

The nebular line emission observed in many cool-core clusters also suggests that thermal conduction may be important. Ultraviolet light from young stars accounts for some of the line emission but cannot explain all of it, implying that another heat source is present (e.g., Johnstone et al. 1987; Voit & Donahue 1997; Werner et al. 2013). Observations of the optical and infrared line ratios are consistent with heat input by a population of suprathermal electrons, possibly entering the nebular gas from the ambient hot medium (Ferland et al. 2009). The emission-line studies are not conclusive on this point, but recent observations of ultraviolet emission from a filament in the Virgo Cluster have bolstered the case for conduction (Sparks et al. 2009, 2012). The C iv 1550 Å and He ii 1640 Å emission lines from this filament are consistent with a model in which the filament is surrounded by a conductive sheath that channels heat into the filament.
If conduction is indeed important in determining the structure of cluster cores and mediating the heat flow into cool gas clouds, then it ought to be included in cosmological simulations of galaxy clusters. Dolag et al. (2004) presented the first simulations of conductive intracluster media (ICMs) in cosmological galaxy clusters, using the Jubelgas et al. (2004) implementation of isotropic conduction in the GADGET smooth-particle hydrodynamics code. These simulations were notable in that they vividly demonstrated the inability of conduction to prevent a cooling catastrophe. Even at 1/3 of the full Spitzer rate, thermal conduction could not prevent a large cooling flow from developing in the cluster core. Conduction delayed the cooling catastrophe but did not diminish the magnitude of the eventual mass inflow. An updated set of cluster simulations using the same conduction algorithm was recently performed by Fabjan et al. (2011). Isotropic conductivity was set to 1/3 of the full Spitzer value for simulations including cooling and star formation, but without AGN feedback. As with the previous simulation set, there was not much difference between the global properties of the final states of clusters simulated with and without thermal conduction.

Conduction is also expected to smooth out temperature inhomogeneities in hot clusters (e.g., Markevitch et al. 2003). One of the clusters simulated by Dolag et al. (2004) was quite hot, with a peak gas temperature of $\sim 12$ keV, and conduction in that cluster produced a much more homogeneous ICM, with far less azimuthal temperature variation than its conduction-free counterpart. Producing simulated clusters with small-scale temperature variations similar to those of observed clusters continues to be somewhat tricky (see, e.g., Ventimiglia et al. 2012). In particular, some recent simulations of sloshing, magnetized cluster cores with anisotropic conduction suggest that conduction needs to be somewhat suppressed along the magnetic field lines in order to explain observations of cluster cores with cold fronts (ZuHone et al. 2013).

Implementations of anisotropic conduction directed along magnetic field lines are also of interest because the MHD instabilities that result may have important implications for the structure of cluster cores and the generation of turbulence in cluster outskirts (e.g., Parrish et al. 2009, 2012). Ruszkowski et al. (2011) recently simulated a cosmological cluster with anisotropic conduction using the FLASH adaptive mesh-refinement (AMR) code. Runs were performed with radiative cooling both on and off, but no AGN feedback was included. In this case, anisotropic conduction significantly reduced mass accretion into the central galaxy, even though the magnetic field geometry significantly suppressed radial heat conduction. However, the $\sim 31$ kpc/$h$ effective resolution was insufficient to adequately resolve core structure.

In this paper, we present a new set of cosmological cluster simulations performed with the AMR code Enzo, including isotropic Spitzer conduction. We simulate 10 clusters spanning a mass range of $2$–$8 \times 10^{14} M_\odot$. We also model a range of suppression factors in order to evaluate the effects of thermal conductivity on core structure. As with previous calculations, we do not include AGN feedback, and a cooling catastrophe therefore results, once again confirming that conduction alone cannot prevent strong cooling flows from developing in galaxy clusters. However, our focus here is on how conduction affects the radial profiles and homogeneity of gas density and temperature outside the central $\sim 20$ kpc.

Our presentation of these conductive cluster simulations proceeds as follows. In Section 2, we describe the primary features of the Enzo code employed in this work and describe our recent implementation of isotropic Spitzer conduction, with the simulation setup given in Section 3. Section 4 discusses the qualitative morphological properties of the simulated clusters, while Section 5 compares their radial profiles. Section 6 focuses on how conduction affects the thermal homogeneity of the ICM, Section 7 looks at the star formation histories of the clusters, and Section 8 summarizes our results.

2. NUMERICAL METHODS

In this work, we use the open-source cosmological, AMR hydrodynamic $+$ N-body code Enzo\(^5\) (Bryan & Norman 1997a, 1997b; Norman & Bryan 1999; O'Shea et al. 2004, 2005). Below, we describe the primary features of Enzo used in this work, including newly added machinery for solving thermal conduction. All analysis is performed with the open-source simulation analysis toolkit, yt\(^6\) (Turk et al. 2011) using the Parallel-HOP halo finder (Skory et al. 2010).

2.1. The Enzo Code

The Enzo code couples an $N$-body particle-mesh solver (Efstathiou et al. 1985; Hockney & Eastwood 1988) used to follow the evolution of a collisionless dark matter component with an Eulerian AMR method for ideal gas dynamics by Berger & Colella (1989), which allows high dynamic range in gravitational physics and hydrodynamics in an expanding universe. This AMR method (referred to as structured AMR) utilizes an adaptive hierarchy of grid patches at varying levels of resolution. Each rectangular grid patch (referred to as a ”grid”) covers some region of space in its parent grid which requires higher resolution, and can itself become the parent grid to an even more highly resolved child grid. Enzo’s implementation of structured AMR places no fundamental restrictions on the number of grids at a given level of refinement, or on the number of levels of refinement. However, owing to limited computational resources, it is practical to institute a maximum level of refinement, $\ell_{\text{max}}$. Additionally, the Enzo AMR implementation allows arbitrary integer ratios of parent and child grid resolution, though in general for cosmological simulations (including the work described in this paper) a refinement ratio of two is used.

Multiple hydrodynamic methods are implemented in Enzo, the primary two being the piecewise parabolic method (PPM) of Colella & Woodward (1984) and the method from the ZEUS code (Stone & Norman 1992a, 1992b). PPM is more accurate than the ZEUS method, but is prone to instability in conditions with steep internal energy gradients, which are common in simulations of cosmological structure formation when radiative cooling is used. In these situations, PPM will occasionally predict negative internal energies, causing the simulation to crash. For this reason, we elect to use the ZEUS method for all the simulations presented here.

We use the metallicity-dependent radiative cooling method described in Smith et al. (2008) and Smith et al. (2011). This method solves the non-equilibrium chemistry and cooling for atomic H and He (Abel et al. 1997; Anninos et al. 1997) and calculates the cooling and heating from metals by interpolating from multi-dimensional tables created with the photo-ionization software, Cloudy\(^7\) (Ferland et al. 1998). Both the primordial

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\(^{5}\) http://enzo-project.org/

\(^{6}\) http://yt-project.org/

\(^{7}\) http://nublado.org/
and metal cooling also take into account heating from a time-dependent, spatially uniform metagalactic UV background (Haardt & Madau 2001).

To model the effects of star formation and feedback, we use a modified version of the algorithm presented by Cen & Ostriker (1992). A grid cell is capable of forming stars if the following criteria are met: the baryon overdensity is above some threshold (here 1000), the velocity divergence is negative (i.e., the flow is converging), the gas mass in the cell is greater than the Jeans mass, and the cooling time is less than the self-gravitational dynamical time. Because conduction can potentially balance radiative cooling, instead of the classical definition of the cooling time, \( \epsilon/\dot{\epsilon} \), we use a modified cooling time that includes the change in energy from conduction, defined as

\[
\dot{t}_{\text{cool}} = \frac{e}{\dot{\epsilon}_{\text{rad}} + \dot{\epsilon}_{\text{cond}}},
\]

(1)

where \( \dot{\epsilon}_{\text{rad}} \) is the cooling rate and \( \dot{\epsilon}_{\text{cond}} \) is the conduction rate. In the original implementation, if all of the above criteria are satisfied, then a star particle representing a large, coeval group of stars with the following mass is formed:

\[
m_* = f_* m_{\text{cell}} \frac{\Delta t}{t_{\text{dyn}}},
\]

(2)

where \( f_* \) is an efficiency parameter, \( m_{\text{cell}} \) is the baryon mass in the cell, \( t_{\text{dyn}} \) is the dynamical time of the combined baryon and dark matter fluid, and \( \Delta t \) is the timestep taken by the grid containing the cell in question. The factor of \( t_{\text{dyn}} \) is added to provide a connection between the physical conditions of the gas and the timescale over which star formation and feedback occur. When a star particle is created, feedback in the form of thermal energy, gas, and metals is returned to the grid at a rate given by

\[
\Delta m_{\text{sf}} = m_* \frac{\Delta t}{t_{\text{dyn}}} \frac{(t - t_*)}{t_{\text{dyn}}} e^{-(t-t_*)/t_{\text{dyn}}},
\]

(3)

where \( t_* \) is the creation time for the particle, so that \( \Delta m_{\text{sf}} \) rises linearly over one dynamical time, then falls off exponentially after that. This feature ensures that the feedback response unfolds over a dynamical time, even if star formation in the simulation is substantially instantaneous. We also use a distributed feedback model, in which feedback material is distributed evenly over a \( 3 \times 3 \times 3 \) cube of cells centered on the star particle’s location. This was shown by Smith et al. (2011) to transport hot, metal-enriched gas out of galaxies more effectively, and to minimize overcooking issues that are common to simulations of structure formation. This method was also shown by Skory et al. (2013) to provide a better match to observed cluster properties than injecting all feedback into a single grid cell.

In practice, Enzo simulations often implement an additional restriction, allowing the creation of a star particle only if \( m_* \) is above some minimum mass. This prevents the formation of a large number of low-mass star particles whose presence can significantly slow down a simulation. However, the use of this algorithm in concert with conduction must be done with care. As we discuss in Section 2.2, explicit modeling of thermal conduction can require timescales that are significantly shorter than the hydrodynamic Courant condition, which has the effect of considerably reducing the value of \( m_* \). If the value of \( m_* \) is only slightly less than the minimum particle size, this condition will likely be satisfied in a relatively short time as condensation proceeds, causing the gas density to increase and its dynamical time to decrease. When the conduction algorithm is active, the resulting time delay can allow cold, dense gas to persist while in thermal contact with the hot ICM. Since radiative cooling is proportional to the square of the density, heat conducted into high-density gas can be easily radiated away. This artificial time delay therefore produces a spurious heat sink in the ICM that can potentially boost the rates of cooling and star formation. In order to avoid this unphysical behavior, we remove the factor of \( \Delta t/t_{\text{dyn}} \) from Equation (2) and adopt a constant value of 10 Myr for \( t_{\text{dyn}} \) for use in Equation (3). We have tested the effects of these changes by comparing two simulations run without conduction, using both the original star-formation implementation and this modification, and find the stellar masses at any given time differ by less than 1%.

2.2. Conduction Algorithm

We implement the equations of isotropic heat conduction in a manner similar to that of Jubelgas et al. (2004), where the heat flux, \( j \), for a temperature field, \( T \), is given by

\[
j = -\kappa \nabla T,
\]

(4)

where \( \kappa \) is the conductivity coefficient. In an ionized plasma, heat transport is mediated by Coulomb interactions between electrons. In this scenario, referred to as Spitzer conduction (Spitzer 1962), the conductivity coefficient is given by

\[
\kappa_{sp} = 1.31n_e \lambda_e k \left( \frac{k_B T_e}{m_e} \right)^{1/2},
\]

(5)

where \( m_e, n_e, \lambda_e, \) and \( T_e \) are the electron mass, number density, mean free path, and temperature, and \( k_B \) is the Boltzmann constant. The electron mean free path is

\[
\lambda_e = \frac{3^{3/2}(k_B T_e)^2}{4\pi^{1/2} e n_e \ln \Lambda},
\]

(6)

where \( e \) is the electron charge and \( \ln \Lambda \) is the Coulomb logarithm, defined to be the ratio of the maximum and minimum impact parameters over which Coulomb collisions yield significant momentum change in the interacting particles, which are electrons in this case. The Coulomb logarithm for electron–electron collisions is

\[
\ln \Lambda = 23.5 - \ln \left( n_e^{1/2} T_e^{-5/4} \right) - \left( 10^{-5} + \ln (T_e) - 2 \right)^{1/2}
\]

(7)

(Huba 2011, p. 34). Because of its relative insensitivity to \( n_e \) and \( T_e \), we follow Jubelgas et al. (2004) and Sarazin (1988) and adopt a constant value of \( \ln \Lambda = 37.8 \), corresponding to \( n_e = 1 \) cm\(^{-3} \) and \( T_e = 10^8 \) K. For values more appropriate to the ICM (\( n_e = 10^{-3} \) cm\(^{-3} \), \( T_e = 10^7 \) K), \( \ln \Lambda = 43.6 \). At a constant density, \( \ln \Lambda \) varies by \( \sim 10\%-20\% \) over the range of temperatures relevant here. The Spitzer conductivity then reduces to

\[
\kappa_{sp} = 4.9 \times 10^{-7} T^{3/2} \text{ erg cm}^{-1} \text{ s}^{-1} \text{ K}^{-1}.
\]

(8)

Since \( \ln \Lambda \) increases with \( T \) and \( \kappa_{sp} \propto 1/\ln \Lambda \), including a direct calculation of \( \ln \Lambda \) would serve to somewhat soften dependence of \( \kappa_{sp} \) on \( T \). In low-density plasmas with large temperature gradients, the characteristic length scale of the temperature gradient, \( \ell_T = T/|\nabla T| \), can be smaller than the electron mean
free path, at which point Equation (5) no longer applies. Instead, the maximum allowable heat flux is described by a saturation term (Cowie & McKee 1977), given as

$$j_{\text{sat}} \simeq 0.4n_e k_B T \left( \frac{2k T}{\pi m_e} \right)^{1/2}.$$  

(9)

To smoothly connect the saturated and unsaturated regimes, we use an effective conductivity (Sarazin 1988) given by

$$\kappa_{\text{eff}} = \frac{\kappa_{sp}}{1 + 4.2 \alpha_e / \ell_T}.$$  

(10)

The rate of change of the internal energy, $u$, due to conduction is then

$$\frac{du}{dt} = - \frac{1}{\rho} \nabla \cdot j.$$  

(11)

Our numerical implementation closely follows that of Parrish & Stone (2005). We use an explicit, first-order, forward time, centered space algorithm, which is the most straightforward to implement in an AMR code. Equations (8)–(11) are evaluated on a per-level basis and is taken to be the minimum of all such values on a given level, ensuring that Equation (11) is solved by summing the heat fluxes from all grid cell faces, and calculating the electron density and temperature on the grid cell face as the arithmetic mean of the cell and its neighbor sharing that face. The time step stability criterion for an explicit solution of the conduction equation is

$$dt < 0.5 \frac{\Delta x^2}{\alpha},$$  

(12)

where $\Delta x$ is the grid cell size and $\alpha$ is the thermal diffusivity, defined as

$$\alpha = \frac{\kappa}{n_e k_B}.$$  

(13)

Equation (12) has the potential to be considerably more constraining than the hydrodynamical Courant condition, which is proportional to $T^{-1/2} \Delta x$, and so the conduction timestep in conditions typical of the ICM is often much shorter than the hydrodynamic timestep. The conduction timestep is calculated on a per-level basis and is taken to be the minimum of all such values on a given level, ensuring that Equation (11) is applied in a stable manner for all grid cells. Enzo pads each AMR grid patch with three rows of ghost zones from neighboring grids. This allows the conduction routine to take three sub-cycled time steps for every hydrodynamic step, and thus allows the minimum grid timestep to be a factor of three larger than Equation (12). After three steps, temperature information from the outermost ghost zone has propagated to the edge of the active region of the grid, so performing additional conduction cycles would yield inconsistent solutions with neighboring grids. By default, Enzo rebuilds the adaptive mesh hierarchy for a given refinement level after each timestep, which is computationally expensive. However, because conduction does not explicitly change the density field (the value of which is the only quantity that determines mesh refinement), we modify this behavior such that the hierarchy is only rebuilt after an amount of time has passed equivalent to what the minimum timestep would be if conduction were not enabled.

This implementation models the case of isotropic Spitzer conduction, where heat flows unimpeded along temperature gradients. If magnetic fields are present and strong enough that the electron gyroradius is small compared to the physical scales of interest (which is likely to be true in the ICM), then heat is restricted to flow primarily along magnetic field lines. Therefore, the heat flux becomes the dot product of the temperature gradient with the magnetic field direction (i.e., $j = -\kappa \mathbf{b} \cdot \nabla T$, where $\mathbf{b}$ is the unit vector pointing in the direction of the magnetic field). Diffusion perpendicular to magnetic field lines is generally considered to be negligible. We do not consider magnetic fields in this work, but instead approximate the suppression of conduction by magnetic fields by adding a suppression factor, $f_{sp}$, varying from 0 to 1, to the heat flux calculation. In reality, the strength and orientation of magnetic fields in galaxy clusters is poorly understood. Hence, the degree to which conduction is suppressed from its maximum efficiency is not known. The extreme limit of tangled magnetic fields is equivalent to isotropic conduction with $f_{sp} = 1/3$. A number of works have shown that the presence of magnetic fields alongside conduction can create a variety of instabilities that greatly influence the effective level of isotropic heat transport (e.g., Parrish & Stone 2005, 2007; Parrish et al. 2009; Ruszkowski & Oh 2010; McCourt et al. 2011, 2012). For this reason, we simulate a large range of values of $f_{sp}$, from strongly suppressed ($f_{sp} = 0.01$) to fully unsuppressed ($f_{sp} = 1$).

3. SIMULATION SETUP

The simulations described in this work are initialized at $z = 99$ assuming the Wilkinson Microwave Anisotropy Probe Year 7 “best fit” cosmological model (Larson et al. 2011; Komatsu et al. 2011): $\Omega_m = 0.268$, $\Omega_b = 0.0441$, $\Omega_{\text{CDM}} = 0.2239$, $\Omega_\Lambda = 0.732$, $h = 0.704$ (in units of $100$ km s$^{-1}$ Mpc$^{-1}$), $\sigma_8 = 0.809$, and using an Eisenstein & Hu power spectrum (Eisenstein & Hu 1999) with a spectral index of $n = 0.96$. We use a single cosmological realization with a box size of 128 Mpc/h (comoving) and a simulation resolution of 256$^3$ root grid cells and dark matter particles. From this realization, we select the 10 most massive clusters and resimulate each individually, allowing AMR only in a region that minimally encloses the initial positions of all particles that end up in the halo in question at $z = 0$. We allow a maximum of 5 levels of refinement, each by a factor of two, refining on baryon and dark matter overdensities of eight. This gives us a maximum comoving spatial resolution of 15.6 kpc/h. We set the Enzo parameter $\text{MinimumMassForRefinementLevelExponent}$ to $-0.2$ for both dark matter and baryon-based refinement, resulting in slightly super-Lagrangian refinement behavior. For each halo in our sample, we perform a control simulation with $f_{sp} = 0$ and conductive simulations with $f_{sp} = 0.01, 0.1, 0.3,$ and 1, evolving each to $z = 0$. All of the simulations employ the star formation/feedback and radiative cooling methods described above in Section 2.1. Table 1 lists the masses and equivalent temperatures for all 10 clusters in the sample.

4. QUALITATIVE MORPHOLOGY

Perhaps the most surprising finding to emerge from our simulations is the insensitivity of the qualitative morphological features of the ICM to thermal conduction, even in clusters with typical gas temperatures $\sim 5$ keV. Figures 1–3 show gas density, temperature, and X-ray surface brightness projections for three clusters at the minimum and maximum conduction efficiency. We choose these clusters in order to capture a wide range of potential influence from conduction. Halo 0 is the most massive cluster in the sample (with a mass of $8 \times 10^{14} M_\odot$), Halo 2 is a highly unrelaxed cluster, and Halo 5 is relatively relaxed and is
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Figure 1. Projections of a 500 kpc thick slab centered on the most massive cluster at $z = 0$ ($M_{200} = 8.0 \times 10^{14} M_\odot$, $T_{200} = 5.6$ keV) for the simulations with $f_{sp} = 0$ (left) and 1. Right: the top and middle rows show average density and temperature, weighted by X-ray emissivity in the energy range of 0.5–2 keV. The bottom row shows the X-ray surface brightness in the same energy range. (A color version of this figure is available in the online journal.)

Table 1

| ID | $M_{200}$ ($10^{14} M_\odot$) | $T_{200}$ (keV) |
|----|-----------------------------|-----------------|
| 0  | 8.04                        | 5.56            |
| 1  | 8.26                        | 5.66            |
| 2  | 5.60                        | 4.37            |
| 3  | 5.56                        | 4.35            |
| 4  | 3.76                        | 3.35            |
| 5  | 4.64                        | 3.85            |
| 6  | 3.85                        | 3.40            |
| 7  | 3.52                        | 3.20            |
| 8  | 1.91                        | 2.13            |
| 9  | 3.41                        | 3.14            |

Notes. Masses and equivalent temperatures (as given by Equation (15)) for the galaxy clusters in the sample. Note that the clusters are ordered by the masses computed by a halo finder in an original exploratory simulation.

close to the average mass for the sample ($M_{200} \simeq 4.9 \times 10^{14} M_\odot$). In order to highlight differences between values of $f_{sp}$, we limit the projected line of sight to 0.5 Mpc. The morphological similarity between these conductive and non-conductive clusters starkly contrasts with the obvious differences found by Dolag et al. (2004) for a $\sim 12$ keV cluster. In their work, it is immediately obvious which simulation included conduction.

There are, however, some subtle but noticeable morphological differences between our conductive and non-conductive

Figure 2. Projections of a 500 kpc thick slab centered on the third most massive cluster at $z = 0$ ($M_{200} = 5.6 \times 10^{14} M_\odot$, $T_{200} = 4.4$ keV) for the simulations with $f_{sp} = 0$ (left) and 1 (right). (A color version of this figure is available in the online journal.)

Figure 3. Projections of a 500 kpc thick slab centered on the sixth most massive halo at $z = 0$ ($M_{200} = 4.6 \times 10^{14} M_\odot$, $T_{200} = 3.9$ keV) for the simulations with $f_{sp} = 0$ (left) and 1 (right). (A color version of this figure is available in the online journal.)
clusters. The overall density structure is slightly smoother for $f_{sp} = 1$ in all three clusters, and the temperature field is visibly smoother in Halo 5. In all three clusters shown, the core appears to be marginally less dense with conduction on, and is also markedly hotter in Halos 2 and 5. Nevertheless, while it is mostly apparent that conduction is present in our cluster sample with $f_{sp} = 1$, its effects are not nearly as striking as in Dolag et al. (2004). This difference is quite possibly due to the strong temperature dependence of Spitzer conductivity, suggesting that it may be possible to estimate the typical conductivity of the ICM by analyzing the dependence of azimuthal temperature homogeneity on cluster temperature (or lack thereof) in a large sample of galaxy clusters.

5. RADIAL PROFILES

Systematic conductivity-dependent differences among our simulated clusters are hard to see in individual cases, but are more apparent in comparisons of average properties of sample sets with differing conductivity. Therefore, in order to understand the systematic effects of increasing the conduction efficiency, we aggregate the radial profiles from all the clusters in our simulated sample for each value of $f_{sp}$, as shown in Figure 4. To minimize artifacts that result from the rebinning of profile data, we perform the initial profiling in physical distance units to calculate the virial radius for each cluster (which, for convenience, we take to be $r_{200}$ with respect to the critical density). We then create a second set of profiles in units of $r/r_{200}$ that are combined to make the aggregate profiles presented here. We refer to these as “averaged profiles” and employ this method for Figures 4 and 5 and all subsequent figures showing aggregate properties.

5.1. Cluster Interiors

The presence of conduction leads to higher gas temperatures both inside and outside of the conduction-free cluster’s characteristic temperature peak at $r/r_{200} \sim 0.07$. Within $r/r_{200} \sim 0.07$, gas temperatures in conductive clusters are greater than those in the non-conductive control simulation, with values of $f_{sp} \gtrsim 0.1$ leading to nearly isothermal cores and central temperatures $\sim 40\%$ greater than those in non-conductive cluster simulations. The temperature profiles of our conductive clusters are in reasonably good agreement with those of Ruszkowski et al. (2011) for their cluster with isotropic conduction and $f_{sp} = 1/3$, although the core of their cluster is less isothermal than ours, with cooler temperatures near the center. This could be because the simulations of Ruszkowski et al. (2011)
did not include a prescription for star formation and feedback in cluster galaxies, which can increase the core temperature both by consuming the cold gas and through thermal and mech-anical feedback. Agreement with the temperature profiles of the simulated clusters from Dolag et al. (2004) is not nearly as good. Those clusters show a significant reduction in both the peak and central temperatures with conduction present, which we do not observe. The theoretical temperature profiles of McCourt et al. (2013) for the $10^{14.5}M_\odot$ clusters show an elevation in the core temperatures in conductive clusters in rough agreement with our results. We also see a marginal inward movement of the location of the temperature peak as they predict, although not nearly to the same degree. The temperature peak for our clusters is also at a smaller radius to begin with.

Despite its ability to maintain approximate isothermality in the cluster cores, conduction at even maximal efficiency is unable to avert a cooling catastrophe at the very center of the halo. This cooling catastrophe and the resulting condensation and star-formation activity produce sharp peaks at $r/r_{200} \lesssim 0.03$ in all the gas density profiles in Figure 4, as well as a slight temperature increase at the same location in the more highly conductive clusters. The elevated central temperatures at the centers of conductive clusters reduce the central gas density relative to the control clusters by $\sim 20\%$–$30\%$. Material that would have fallen into the center is displaced to larger radii, as can be seen by the elevated density at $0.04 \lesssim r/r_{200} \lesssim 0.6$ in the more conductive runs. Gas density is enhanced primarily in the range $0.04 \lesssim r/r_{200} \lesssim 0.2$, but the enhancement appears to saturate at $\sim 20\%$ for $f_{\text{sp}} = 0.33$. For $f_{\text{sp}} = 1$, the density is actually lower than for $f_{\text{sp}} = 0.33$ in this range, but is then higher for $0.2 \lesssim r/r_{200} \lesssim 0.6$. It is unclear what causes the outward transport (or prevention of inflow) to stall just inside $0.2\ r_{200}$, but it seems that it can be overcome for some value of $f_{\text{sp}}$ above 0.33, likely closer to 1.

Further evidence of conduction-driven inflation of the core can be seen in Figure 5, where we plot averaged profiles of $M/r$, including the individual contributions from dark matter, stars, and gas. Despite the lack of perfect monotonicity in the density and temperature profiles, the gaseous component of the potential decreases monotonically with increasing $f_{\text{sp}}$, indicating that the clusters are indeed responding to the presence of conduction by puffing up their cores and redistributing gas to larger radii. As is the case for most simulated galaxy clusters, the stellar component is extremely centrally concentrated, dominating the gravitational potential inside $0.04\ r_{200}$. The stellar and dark matter components of the potential are largely unaffected by conduction.

In the right panel of Figure 4, we plot normalized entropy profiles, where the normalization term, $K_{200}$ (Voit 2005; Voit et al. 2005), is given by

$$K_{200} = k_B T_{200} n_e^{-2/3},$$

where $T_{200}$ is

$$k_B T_{200} = \frac{GM_{200}\mu m_p}{2r_{200}},$$

and $n_e$ is the average electron number density within $r_{200}$, assuming a fully ionized plasma of primordial composition. The decrease in density and increase in temperature in the cluster cores with conduction yield entropy profiles that are enhanced by $40\%$–$70\%$, but still show the steep decline in the very center that is typical of simulated clusters. This enhancement decreases out to $r/r_{200} \sim 0.07$, where the mean entropy values are approximately equivalent for all values of $f_{\text{sp}}$. The entropy profiles for the clusters with conduction then dip below the control sample by roughly $10\%$ out to $r/r_{200} \sim 0.2$. In the range $0.2 \lesssim r/r_{200} \lesssim 0.6$, the combined temperature and density enhancements appear to be in perfect balance, producing nearly identical entropy profiles across the entire cluster sample with a very small variance.

Voit et al. (2005) find that for a sample of non-radiative, non-conducting clusters simulated with Enzo, the entropy profiles in the range $0.2 \lesssim r/r_{200} \lesssim 1$ are best fit by the power law

$$K(r)/K_{200} = 1.51(r/r_{200})^{1.24},$$

which we overplot on Figure 4 with a black dashed line. This power law matches our sample well in normalization, but has a slightly steeper slope. We find that our sample is best matched by

$$K(r)/K_{200} = 1.50(r/r_{200})^{1.09}.$$  

For comparison, we also plot the predicted entropy profile for a pure cooling model (Voit & Bryan 2001; Voit et al. 2002) of a 5 keV cluster. This model assumes gas and dark matter density distributions that follow an Navarro–Frenk–White profile (Navarro et al. 1997) with concentration $c = 6$. The gas is initially in hydrostatic equilibrium and allowed to cool for a Hubble time. The pure cooling model agrees quite well with our clusters in the range $0.2 \lesssim r/r_{200} \lesssim 1$, and especially well in the range $0.2 \lesssim r/r_{200} \lesssim 0.5$. Outside of $0.5\ r_{200}$, the slope of the pure-cooling model is slightly too steep and is closer in slope to the Voit et al. (2005) fit.

5.2. Cluster Exteriors

The outskirts of galaxy clusters, while having lower temperatures and hence lower conductivities, are not subject to intermittent heat injection from stellar feedback, and thus offer an intriguing laboratory for studying the effects of conduction. We identify two distinct regions in the cluster outskirts where conduction appears to have influence, at $\sim r_{200}$ and at $\sim 3r_{200}$. At $r/r_{200} \lesssim 0.6$, the density excess seen in the cluster interiors turns into a deficit for all values of conduction simultaneously, as can be seen in Figure 4. From this point out to well past the virial radius, there exists a perfectly monotonic trend of lower gas densities for higher values of $f_{\text{sp}}$. For the maximum value of $f_{\text{sp}}$, the average gas density at the virial radius is $10\%$ lower than that of the clusters simulated without conduction. In fact, the density is measurably lower for all values of $f_{\text{sp}} \geq 0.1$ out to a few virial radii. The temperature just inside $r_{200}$ is marginally higher, while the temperature just outside $r_{200}$ is reduced.

Conduction transports heat outward in these regions because the gravitational potential there produces a declining gas temperature gradient. This heat transfer causes the entropy of the outer gas to increase. However, because gas in the cluster outskirts is not pressure-confined like the gas in the core, it is free to expand outward and decrease in both density and temperature, while its temperature gradient remains determined by the gravitational potential. This happens because the timescale for conduction in the outer regions is substantially greater than the sound-crossing time. Consequently, outward conduction of heat causes the outer gas to expand without much change in the temperature gradient. Beyond $r_{200}$, however, we see a slight steepening in the temperature profile as predicted for conductive clusters by McCourt et al. (2013), presumably because inflation...
of the ICM due to conduction pushes gas near the virial radius farther from the cluster center without adding much thermal energy. The fact that the entropy profiles of conductive clusters are lower than those of non-conductive clusters beyond the virial radius supports this interpretation.

Further out, at $\sim 3r_{200}$, there is another systematic decrease in both density and temperature. Interestingly, as was pointed out by Skillman et al. (2008), this is the typical location for a galaxy cluster’s accretion shocks, which are responsible for heating gas up to the virial temperature. This raises a critical question: How can conduction affect a large halo’s accretion shocks?

To test this question in a more controlled environment than a cosmological simulation, we performed a pair of one-dimensional simulations designed to mimic the conditions of an accretion shock around a galaxy cluster. Following the shock properties characterized by Skillman et al. (2008), we initiated a Mach 100 standing shockwave with preshock conditions of $n = 10^{-4} \text{ cm}^{-3}$ and $T = 10^4 \text{ K}$, comparable to gas that is falling directly onto a galaxy cluster (i.e., through spherical accretion of the surrounding intergalactic medium, rather than being accreted via filaments). We ran simulations with $f_{sp} = 0$ and 1, with profiles shown in Figure 6. Almost immediately, a conduction front in the $f_{sp} = 1$ simulation races ahead of the original shock. The conduction front, referred to as the foreshock (Tsai et al. 2002; Fadeyev et al. 2002; Tsai et al. 2007), continues to advance with ever-decreasing speed and settles into a near-steady state within 50–100 Myr, roughly 30 kpc upstream of its initial location. The final position of the foreshock can be related to the conduction timescale (Tsai et al. 2007; Parrish et al. 2009) given by

$$t_{\text{cond}} = \frac{L^2}{\alpha},$$

where $L$ is a length scale and $\alpha$ is the thermal diffusivity (Equation (13)). The average velocity of the conduction front over the length, $L$, is then inversely proportional to $L$. One can then calculate the value of $L$ at which the average velocity of the conduction front is equal to the upstream velocity in the frame of the original shock, here $\sim 1100 \text{ km s}^{-1}$. Using the average of the upstream and downstream temperature and density to calculate $\alpha$, we arrive at a value of $L \sim 28 \text{ kpc}$, which agrees very well with Figure 7. This behavior has been seen in previous works, simulating solar coronae and stellar winds.

Figure 6. Profiles of a one-dimensional, Mach 100 standing shockwave after 200 Myr for $f_{sp} = 0$ and 1 (thick black and blue lines, respectively). Note, the curves for $f_{sp} = 0$ also denote the initial configuration of the shock, which is stationary in the case where conduction is negligible. Thin blue lines denote the shock with $f_{sp} = 1$ at 10, 20, 50, and 100 Myr. The lower middle panel shows the ratio of the gas velocity to the local sound speed, while the lower right panel shows the shock Mach number based on the temperature jump and calculated with the shock finding method of Skillman et al. (2008). The red lines in the lower right panel correspond to the Mach number calculated based on the density jump for the simulation with $f_{sp} = 1$.

(A color version of this figure is available in the online journal.)

Figure 7. Location of the foreshock and true shock with respect to the initial shock position as a function of time for the one-dimensional shock simulation with $f_{sp} = 1$ shown in Figure 6. Positive displacement refers to motion upstream of the original shock.

(A color version of this figure is available in the online journal.)
The transfer of heat upstream causes a slight reduction in the pressure at the location of the original shock front, pushing the true shock downstream. As seen in Figure 7, the motion of the true shock proceeds much more slowly, requiring roughly 1 Gyr to reach an equilibrium position about 15 kpc downstream.

The lower right panel of Figure 6 shows the Mach numbers determined using the shock-finding algorithm described in Skillman et al. (2008). Using this method, we calculate the Mach number separately using the temperature (blue lines) and density (red lines) jumps. Since conduction has led to a physical separation of the two jumps, these two methods will give different results. The true shock, defined as the position of the velocity discontinuity and hence where the kinetic energy is converted to thermal energy, has a Mach number of just under 3 as calculated from the density jump. The foreshock has a temperature jump corresponding to a Mach number of roughly 70. While the foreshock is not the true shock, measuring its equivalent Mach number is important since the primary role of an accretion shock is to heat infalling gas to the virial temperature of the halo.

These results are independent of resolution. The simulations shown in Figure 6 are for a grid 128 cells across, and we observe nearly identical behavior down to a resolution of 16 cells. For this configuration, the results are strongly dependent on $f_{\text{sp}}$. At $f_{\text{sp}} = 0.67$, the distance between the separated shocks is approximately half of that at $f_{\text{sp}} = 1$ and the Mach number of the foreshock is only reduced to just under 90. However, the Mach number of the true shock, based on its density jump, is reduced to nearly the same value as the case of $f_{\text{sp}} = 1$. It appears that if conduction is strong enough to separate the foreshock form the true shock, the resulting Mach number of the true shock, unlike the foreshock, is relatively insensitive to the value of $f_{\text{sp}}$. For $f_{\text{sp}} \lesssim 0.4$, the results are indistinguishable from those without conduction. However, we find that it is possible to produce a significant shockwave alteration for lower values of $f_{\text{sp}}$ simply by lowering the initial Mach number and increasing the preshock temperature (resulting in gas in a thermodynamic regime comparable to gas that is being accreted from cosmological filaments). As long as the postshock gas is able to reach temperatures in the range of $10^7$ K, where the Spitzer conductivity becomes considerable, conduction is capable of bifurcating the shock front.

Figure 8 shows the average Mach number calculated from temperature jumps as a function of gas density in our simulations of the most massive halo for all shocks in the subvolume in which refinement is allowed. Conduction reduces the Mach numbers of the strongest shocks, which occur preferentially in the lowest density gas, by roughly 10% from $f_{\text{sp}} = 0$ to $f_{\text{sp}} = 1$. We therefore conclude that the density, temperature, and entropy deficits observed beyond the virial radius in our conductive clusters are indeed due to shock alterations similar to those seen in our idealized, one-dimensional simulations of conductive shock fronts. However, as was seen in the one-dimensional shock tube simulations, identification of the true shock is not as straightforward when conduction is present. Therefore, we caution the reader that the precise Mach numbers of the accretion shocks may be somewhat different.

However, the physics of the actual accretion shocks around real galaxy clusters is undoubtedly more complex. In particular, it is important to note that the electron mean free path in that gas is several times greater than the intershock distance in our one-dimensional simulations, and furthermore that real accretion shocks are likely to be collisionless and magnetically mediated. Nonetheless, the general qualitative point these simulations illustrate is interesting: heating of preshock gas by a hot electron precursor has the potential to alter the expected relationships between the sizes and locations of the density and temperature jumps in accretion and merger shocks. The displacement of the foreshock and true shock is small (10s of kpc) compared to the typical location of accretion shocks around galaxy clusters ($r \sim 10$ Mpc), but observation of such an effect for merger shocks may be possible.

Progress in understanding the effects of conduction on accretion shocks will require modeling the gas as a fully ionized plasma (Zel’dovich 1957; Shafranov 1957), which is beyond the scope of this work. The high Mach number of an accretion shock, combined with a postshock temperature high enough for significant heat flux, create conditions similar to a supercritical radiative shock, as described by Lowrie & Rauenzahn (2007). When simulated with a two-fluid approach, the combination of preheating the preshock medium via conduction in the electrons, electron–ion coupling, and compression of the ion fluid in the postshock region can produce a small region, known as a Zel’dovich spike, where the ion temperature is slightly higher than the equilibrium postshock temperature (Lowrie & Rauenzahn 2007; Lowrie & Edwards 2008; Masser et al. 2011). Our single-fluid simulations, despite their limitations, produce shocks quite similar to the non-equilibrium results of Lowrie & Edwards (2008), where a diffusion term proportional to $T^{7/2}$ (similar to the $T^{5/2}$ of Spitzer conductivity) is used. Nevertheless, a more detailed study of the characteristics of accretion shocks employing a two-fluid MHD treatment along with physically motivated conduction and cooling rates seems warranted.
6. TEMPERATURE HOMOGENEITY

Conduction strong enough to alter the temperature gradients in cluster cores should also be effective at smoothing out small-scale thermal variations in the ICM. However, those effects turn out to be rather subtle, as shown in Figure 9, which plots the normalized variance of the temperature field as a function of radius, averaged over all clusters in the sample that have the same level of conductivity. It reveals an extremely weak trend of greater temperature homogeneity with increasing \( f_{sp} \), but at all radii the difference in homogeneity among cluster sets with different values of \( f_{sp} \) is less than the cluster-to-cluster variation. Furthermore, this trend reverses beyond the virial radius, at \( 2 \lesssim r/r_{200} \lesssim 3 \), albeit at a nearly marginal level.

The homogenizing effects of conduction are even harder to see in projection, which perhaps is not surprising given the scarcity of obvious conduction-dependent morphological differences in Figures 1–3. Those differences are further diluted when projected over a 4 Mpc line of sight through each cluster, as in Figures 10–12. These latter figures show mean temperature weighted by the X-ray emission in the 0.5–2 keV energy band, and X-ray emission is calculated by interpolating from density, temperature, and metallicity-dependent emissivity tables computed with the Cloudy code. Since X-ray emission is proportional to \( n_e^2 T_e^{3/2} \), this weighting should highlight clumpiness and differences in temperature. However, as stated previously, these results contrast considerably with the cluster map comparison of Dolag et al. (2004), which shows a significantly hotter cluster in which conduction should be much more efficient.

The change in temperature homogeneity due to conduction is quite small in our simulated cluster sample, but important insights into the physics of the ICM could be gained if the effective value of \( f_{sp} \) could be measured observationally. To evaluate this possibility, we created X-ray-weighted temperature...
maps of the central 300 kpc for each of the clusters in the sample, masking out the pixels within 40 kpc of the cluster centers to remove features that would be considered part of the central galaxy. We then quantified the amount of azimuthal temperature structure by dividing each temperature map into azimuthal bins and calculating both the mean temperature in each bin and the temperature variance among all azimuthal bins in the map. We performed this calculation multiple times for each map while rotating the azimuthal bins, and took the maximum variance calculated as the value for that map. Finally, we averaged the values together for all clusters in the sample and performed the entire exercise over a range in the total number of azimuthal bins, from two to nine. Figure 13 plots the maximum variance as a function of the number of bins for each value of $f_{\text{sp}}$. We find that, in general, the maximum variance is lower for the simulations with conduction, but only by approximately 10%. There does not appear to be any sort of monotonic trend with increasing $f_{\text{sp}}$. We repeated this experiment, varying the inner and outer radius for the temperature maps, but were unable to find conditions that produce a better trend than can be seen in Figure 13. Thus, we conclude that the efficiency of conduction is difficult to determine solely from the observable temperature homogeneity of the ICM, at least for clusters of temperature $\lesssim 6$ keV.

7. STAR FORMATION

Figure 14 shows the average difference in stellar mass between the clusters with conduction and those without as a function of time. The difference in stellar mass at $z = 0$ for all levels of conduction is less than 5%. Surprisingly, the clusters with higher levels of conduction form more stars in our calculations (see also Dolag et al. 2004). However, the shaded regions in Figure 14 give an indication of just how much variation there is between clusters. As we have shown, even the highest level of conduction is unable to prevent a cooling catastrophe in the center of a cluster. Therefore, one should not expect a large difference in the amount of star formation.
Given the coarseness with which the star-forming regions are resolved, it is possible that the enhancement in star formation with increasing $f_{\text{sp}}$ is numerical and not physical. We propose two potential numerical explanations. First, the short timesteps required for the stability of the conduction algorithm may not provide enough time for a single star particle to sufficiently heat up a grid cell and quench star formation in a given cell in the following timestep. Second, conduction may transport thermal energy too quickly out of regions heated by recent star formation. This would allow a star-forming region to recool and form additional stars too rapidly.

Because conduction is able to create a nearly isothermal core for $f_{\text{sp}} \gtrsim 0.1$, yet the final stellar mass from the clusters with $f_{\text{sp}} = 0.1$ is consistent with no change, we find it reasonable to conclude that the isothermality of the core in a conductive cluster has no influence on the star formation rate Enzo calculates for the central galaxy. However, because we do not resolve the interface between the ICM and the interstellar medium (ISM), we cannot definitively state the effect of conduction on star-forming gas in a cosmological simulation.

8. CONCLUSIONS

We have performed cosmological simulations of 10 galaxy clusters using isotropic thermal conduction with five values of the conductive suppression factor in order to study the effects of conduction on galaxy cluster cores and the ICM. By studying the aggregate properties of the clusters in our sample, we find that the presence of conduction even at its maximum possible efficiency induces changes to the density and temperature structure on the order of only 20%–30%. For efficiency inducing changes to the density and temperature that the presence of conduction even at its maximum possible effect of conduction is instantly recognizable. The cluster shown in Dolag et al. (2004) is significantly more massive than our most massive cluster, so it is possible that a hotter ICM, with a higher thermal conductivity, is made more homogeneous, suggesting that the temperature dependence of temperature inhomogeneity in a large cluster sample could help reveal the typical conductivity of the ICM. We attempted to find a means of distinguishing the level of conduction observationally by measuring the variance in our projected temperature maps, but without success.

Conduction appears to have very little influence on the star formation rate within our simulated clusters. When determining whether a grid cell should form a star particle, we include the energy change from conduction in the calculation of the cooling time, but this seems to have very little influence. This is likely because star-forming grid cells are mostly surrounded by cells that are also quite cool. Somewhat surprisingly, we observe a marginal increase in the total stellar mass with increasing conduction, such that the sample with $f_{\text{sp}} = 1$ shows an enhancement in star formation rate of $5\%$. The fact that conduction cannot suppress star formation is directly related to its inability to prevent the cooling catastrophe in the very center of the cluster. However, the reasons for the slight increase in star formation may be more numerical than physical. Further progress on understanding the effects of thermal conduction on star formation in cluster cores will require properly resolving the interface between ISM and the ICM, which is presently impractical in cosmological galaxy cluster simulations.

Finally, we caution the reader to not forget the various simplifying assumptions that were employed for this work. The simulations presented here ignored the effects of magnetic fields, whose primary role in this context is to restrict heat to flow along magnetic field lines. This will preserve temperature gradients that are perpendicular to magnetic fields, especially when the field is regularly structured. Thus, the degree to which conduction can isotropize the temperature structure of the ICM may be even lower than our results indicate. Additionally, the physics of conduction occurring over shock fronts is more complex than the treatment in this work. Our simulations only...
consider a single fluid while, in reality, the electron and ion temperatures will be quite different. Additionally, the electron mean free path in the ICM is larger than both the typical grid cell size in the simulation, making the assumption of the ICM as a somewhat questionable fluid, and the separation observed between the foreshock and the true shock. Even so, our results suggest that a more detailed investigation of the influence of conduction on the outskirts of galaxy clusters is not without merit.

We are very grateful to the anonymous referees for their insightful comments. This work was supported by NASA through grants NNX09AD80G and NNX12AC98G, and by the NSF through AST grant 0908819. The simulations presented here were performed and analyzed on the NICS Kraken and Nautilus supercomputing resources under XSEDE allocations TG-AST090040 and TG-AST120009. We thank Greg Bryan, Gus Evrard, Eric Hallman, Megan Donahue, Andrey Kravtsov, Jack Burns, Matthew Turk, and Stephen Skory for helpful discussions during the course of preparing this manuscript. S.W.S. has been supported by a DOE Computational Science Graduate Fellowship under grant number DE-FG02-97ER25308. B.W.O. was supported in part by the MSU Institute for Cyber-Enabled Research—Theory and Applications, ed. T. Plewa, T. Linde, & G. Weirs (New York: Springer), 341

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Their commitment to open science has helped make this work possible.

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