A study of cool core resiliency and entropy mixing in simulations of galaxy cluster mergers

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ABSTRACT
We present results from a suite of binary merging cluster simulations. The hydrodynamical cluster simulations are performed employing a smoothed particle hydrodynamics (SPH) formulation in which gradient errors are strongly reduced by means of an integral approach. We consider adiabatic as well as radiative simulations, in which we include gas cooling, star formation and energy feedback from supernovae. We explore the effects of merging on the thermodynamic structure of the intracluster gas of the final merger remnant. In particular, we study how core entropy is generated during the merging and the stability properties of the initial cool-core profile against disruption. To this end, we consider a range of initial mass ratio and impact parameters.

Final entropy profiles of our adiabatic merging simulations are in good accord with previous findings (ZuHone 2011), with cool-cores being disrupted for all of the initial merging setups. For equal-mass off-axis mergers, we find that a significant contribution to the final primary core entropy is due to hydrodynamic instabilities generated by rotational motions, which are induced by tidal torques during the first pericenter passage. In radiative simulations, cool-cores are more resilient against heating processes; nonetheless, they are able to maintain their integrity only in the case of off-axis mergers with very unequal masses. We suggest that these results are robust against changes in the gas physical modeling, in particular to the inclusion of AGN thermal feedback.

Our findings support the view that the observed core cluster morphology emerges naturally in a merging cluster context, and conclude that the merging angular momentum is a key parameter in shaping the thermodynamical properties of the final merger remnant.

Key words: Hydrodynamics — methods: numerical — galaxies: clusters: general — galaxies: clusters: intracluster medium — X-rays: galaxies: clusters

1 INTRODUCTION
According to the hierarchical scenario, the formation of structure in the Universe proceeds under the action of gravity through merging and accretion of smaller structures. In this framework clusters of galaxies are the latest and most massive objects to be formed, with virial masses in the range $M \sim 10^{14} - 10^{15} M_\odot$ (Voit 2005).

During their formation process, the gas is heated by adiabatic compression and shock-heating to higher temperatures. At virial equilibrium, about ~ 90% of the baryons in a cluster will reside in the form of an hot, X-ray emitting intracluster medium (ICM) at temperatures $T \sim 10^7 - 10^8$ K.

Therefore, X-ray observations of the ICM provide X-ray maps with which to probe the spatial distribution of the cluster gas density, temperature and metallicities. Assuming hydrostatic equilibrium and spherical symmetry, these data can then be used to deduce the underlying dark matter (DM) distribution and to determine cluster virial masses.

An accurate determination of cluster masses it is necessary in order to exploit the usefulness of clusters as cosmological probes, since at any given epoch their number density is a sensitive function of the background cosmological model. This requires the clusters to be dynamically relaxed, since otherwise cluster mass estimates will be prone to uncertainties.

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However, there is a large variety of observations indicating that galaxy clusters can be broadly classified into two categories: relaxed and un-relaxed (see Buote 2002, for a review). The fraction of clusters exhibiting a disturbed morphology grows with redshift and at the present epoch can be even greater than ~50%, depending on the adopted criterion used to measure the amount of substructure present in the cluster (Buote 2002). It is then fundamental to study the physics of cluster merging, not only in order to assess the status of the cluster dynamical equilibrium, but because merging between substructures (or clusters themself) gives raise to a number of interesting physical processes (Sarazin 2002; Molnar 2016).

During the merging process, collisions between substructures drive shocks into the ICM, heating the gas and injecting turbulent motions. These X-ray shocks will boost X-ray luminosities, and leave a number of observational signatures in the ICM, such as contact discontinuities (or cold fronts) in the gas temperature, radio relics, relativistic electrons and other features (Markevitch & Vikhlinin 2007; Feretti et al. 2012).

Following the gas compression, mergers between clusters are also expected to drive star formation (Fujita et al. 1993; Roettiger et al. 1993; Roettiger et al. 1994), but with observations producing conflicting results. Some authors claim an increase in the star formation activity during mergers (Bekki et al. 2010; Stroe et al. 2017), while it is absent in other merging systems (Mansheim et al. 2017).

Moreover, major cluster mergers are the most energetic events since the Big-Bang, with energies $\gtrsim 10^{64}$ ergs. This renders these objects unique laboratories with which to study dark matter models. Because of the collisionless nature of dark matter (DM), the position of gas and DM centers will be offset during a merging process. By contrasting X-ray and weak lensing data, it is possible to derive upper limits on the cross section of self-interacting DM (Molnar et al. 2012; Zhang et al. 2018). These findings strongly suggest that the CC/NCC dichotomy can be naturally interpreted in terms of the cluster merging histories. In this framework, the population of NCC clusters originates as a consequence of major mergers that disrupt CC clusters. Conversely, the original core morphology is preserved for relaxed clusters that have not experienced a major merger recently.

This method has the advantage that it allows the detailed study of a single merging event. It also simplifies the interpretation of the simulation results because the initial conditions are kept under control and the simulation can be contrasted with a specific observation. This approach has been followed by many authors (Roettiger et al. 1993; Ricker & Sarazin 2001; Ritchie & Thomas 2002; Poole et al. 2006; McCarthy et al. 2007; Poole et al. 2008; Mitchell et al. 2009; Bonomet et al. 2017).

Specifically, idealized binary cluster mergers have been used to study the merging configuration of the ‘Bullet cluster’ (Springel & Farrar 2007; Mastroietro & Burkert 2008), as well as that of ‘El Gordo’ cluster (Zhang et al. 2012, 2018) and of other merging clusters (Machado & Lima Neto 2013; Molnar & Broadhurst 2018; Halbesma et al. 2019).

Simulated X-ray maps can be constructed to study cluster merging, for instance by assessing the degree of relaxation of a specific system (ZuHone et al. 2009). Moreover, the measured offset between X-ray and Sunyaev–Zel’dovich (SZ) maps allows the relative velocity of the two merging clusters to be determined (Molnar et al. 2013; Zhang et al. 2014). These limits in turn can be used to derive constraints on the assumed cosmological model.

Another important topic in which mergers of galaxy clusters play an important role is in the study of DM properties. As previously outlined, major mergers are very energetic events in which a self-interacting DM (SIDM) is expected to exhibit significant signatures. For this reason, merging simulations with a SIDM have been carried out by many authors (Robertson et al. 2017; Kim et al. 2017; ZuHone et al. 2019), the simulations being aimed at investigating the impact of a SIDM on gas and DM properties of the merged clusters.

Finally, numerical simulations of merging clusters have been widely used to study the origin of the observed central properties of the cluster gas. X-ray cluster surveys show that clusters can be divided into two categories according to ICM central properties (Cavagnolo et al. 2009; Johnson et al. 2009; Pratt et al. 2010; McDonald et al. 2013): cool-core (CC) and non-cool core (NCC) clusters. CC clusters are characterized by a peaked X-ray emission, very short cooling times (~10% of the Hubble time), central temperatures about $\sim 1/3$ of the virial ones and radial entropy profiles steeper in the core than those of NCCs (Cavagnolo et al. 2009).

These short cooling times should induce a run-away cooling process that is not observed; to balance radiative losses, some heating sources must be operating in the cluster cores. This is the so-called ‘cooling flow’ problem and various heating models have been proposed in the literature to offset cooling and regulate the cooling flow (see Soker 2014 and references cited therein).

To observationally define a CC cluster there are various criteria (Barnes et al. 2018), which depend on the available data. However, there is some consensus that CC clusters are correlated with a regular X-ray morphology (Chon et al. 2012), while this is not true for NCC clusters. The latter are often associated with a disturbed morphology (Pratt et al. 2010) and exhibit a much flatter radial entropy profile than CC clusters.

These findings strongly suggest that the CC/NCC dichotomy can be naturally interpreted in terms of the cluster merging histories. In this framework, the population of NCC clusters originates as a consequence of major mergers that disrupt CC clusters. Conversely, the original core morphology is preserved for relaxed clusters that have not experienced a major merger recently.

This scenario is an important issue for a better understanding of cluster formation and evolution, and N-body/hydro simulations of merging clusters have been widely employed (Ritchie & Thomas 2002; Gomez et al. 2004; McCarthy et al. 2007; Poole et al. 2008; Burns et al. 2009).
composed of gas and DM initially in equilibrium, we specify the radial DM density and gas entropy profiles. To define halo parameters, we use a ΛCDM cosmology, with \( \Omega_m = 0.3 \), \( H_0 = 70 \text{ km s}^{-1} \text{ Mpc}^{-1} \), and a baryon fraction of \( f_b = \Omega_b/\Omega_m = 0.162 \). For each of the gas and DM components, a particle realization of positions and velocities is then constructed, according to profiles computed under the assumption of hydrostatic equilibrium.

This procedure is used to construct both a primary and a secondary cluster, the virial mass of the two being related by the merging mass ratio. To initialize the merger simulation, the particle positions and velocities of the two halos are then shifted according to the initial orbital trajectory.

Our initial condition set up is analogous to that implemented by ZuHone (2011) in his adiabatic merger simulation study. In particular, we adopt the same range of collision parameters. In the present study, we have purposely chosen to adopt similar initial settings. This in order to compare with previous results on the effects of mergers on final CC properties, specifically when cooling is included in the simulations.

Moreover, the merger simulations of ZuHone (2011) were performed using an adaptive mesh-based Eulerian code. For adiabatic simulations, it is then interesting to compare the entropy profiles of the final merger remnants against the corresponding ones presented in ZuHone (2011). This because the two sets of simulations have been constructed by adopting very similar initial conditions, but the codes used to perform the simulations are based on two completely different numerical hydrodynamical schemes.

Our paper is organized as follows. In Section 2 we describe our hydrodynamical scheme, together with the method we use to initialize halos in equilibrium and the orbital properties. Section 3 is dedicated to the presentation of the results, in which we describe our findings from adiabatic and radiative simulations; a specific Section being dedicated to investigate the generation of entropy through mixing and shock-heating processes during the various merging phases. Finally, our main conclusions are summarized in Section 4.

2 SIMULATIONS

The simulations are performed by employing an entropy conserving SPH formulation (Price 2012). To estimate first-order SPH derivatives, the numerical scheme is improved by using an Integral Approximation (IA) accompanied by a matrix inversion, thus strongly reducing gradient errors in the momentum equation.

The IA scheme was originally proposed by García-Senz et al. (2012), and further tested in a variety of hydrodynamical tests (García-Senz et al. 2012, Rosewos 2015, V16). The results of the tests demonstrate that the new SPH formulation outperforms standard SPH and, in terms of accuracy, can be considered competitive with other numerical hydrodynamic schemes (V16). In particular, with respect standard SPH, it is found that the IA scheme greatly improves the numerical modeling of subsonic turbulence (V16, Valdarnini 2014). This aspect is particularly important for the simulations presented here, in which a significant amount of turbulence is expected to be injected into the ICM during cluster collisions (Schmidt et al. 2014).
We now outline the basic features of the hydrodynamical method — we refer the reader to García-Senz et al. (2012) and V16 for a comprehensive description of the IA method applied to SPH. In what follows, we will refer to the SPH scheme described here as integral SPH (ISPH).

2.1 Numerical method

In SPH, the fluid is described by a set of \( N \) particles with mass \( m_i \), velocity \( \mathbf{v}_i \), density \( \rho_i \) and specific entropy parameter \( A_i \).\footnote{We use the convention of using Latin indices to denote particles and Greek indices to denote the three spatial dimensions.} The latter is related to the particle pressure by \( P_i = A_i \rho_i^\gamma = (\gamma - 1) \rho_i u_i \), where \( \gamma = 5/3 \) and \( u_i \) is the thermal energy per unit mass \( u_i \).

At the particle position \( \mathbf{r}_i \), the SPH gas density \( \rho_i \) is given by the summation
\[
\rho_i = \sum_j m_j W(|\mathbf{r}_i - \mathbf{r}_j|, h_i) ,
\]
(1)

where the sum is over neighboring particles \( j \), and \( W(|\mathbf{r}_i - \mathbf{r}_j|, h_i) \) is a kernel with compact support. We define \( W_i(j, h_i) \equiv W(|\mathbf{r}_i - \mathbf{r}_j|, h_i) \). For the simulations presented here, we use the \( M_4 \) kernel (Price 2012), which is zero for \(|\mathbf{r}_i - \mathbf{r}_j| \geq 2h_i\).

In Equation (1) the sum is over a finite number of particles \( N_{nn} \), and the smoothing length \( h_i \) is implicitly defined by the equation
\[
\frac{4\pi(2h_i)^3 \rho_i}{3} = N_{nn} m_i ,
\]
(2)

which is solved iteratively for each particle by setting \( N_{nn} = 33 \).

The SPH momentum equation in the IA framework reads
\[
\frac{d\mathbf{v}_i}{dt} = - \sum_j m_j \left[ \frac{P_i}{\Omega_i \rho_i^2} A_{\alpha,ij}(h_i) + \frac{P_j}{\Omega_j \rho_j^2} \tilde{A}_{\alpha,ij}(h_j) \right] ,
\]
(3)

where \( \Omega_i \) is defined as
\[
\Omega_i = \left[ 1 - \frac{\partial h_i}{\partial \rho_i} \sum_k m_k \frac{\partial W_k(h_i)}{\partial h_i} \right] ,
\]
(4)

and the two terms \( A_{\alpha,ij}(h_i) \) and \( \tilde{A}_{\alpha,ij}(h_j) \) are the IA generalization to the SPH derivatives \( \nabla_i W_{ij}(h_i) \) and \( \nabla_j W_{ij}(h_j) \), respectively.

These IA terms are given by
\[
A_{\alpha,ij}(h_i) = \sum_{\beta,\gamma} C_{\alpha,\beta\gamma}(i) \Delta_{\beta\gamma} \nabla_i W_{ij}(h_i) ,
\]
\[
\tilde{A}_{\alpha,ij}(h_j) = \sum_{\beta,\gamma} C_{\alpha,\beta\gamma}(j) \Delta_{\beta\gamma} \nabla_j W_{ij}(h_j) .
\]
(5)

Here \( \Delta_{\beta\gamma} = (t^\beta - t^\gamma) \), and \( C_{\alpha,\beta\gamma}(i) \) are the elements of the matrix \( C = T^{-1} \) associated with the particle \( i \). The inverse of this matrix is a \( 3 \times 3 \) symmetric tensor \( T \), which for particle \( i \) takes the form
\[
T_{\alpha,\beta}(i) = \sum_j \frac{m_j}{\rho_j} \Delta_{\alpha,\beta} \nabla_j W_{ij}(h_i) .
\]
(6)

To properly handle shocks, the SPH momentum equation (3) must be generalized to include an artificial viscosity (AV) term:
\[
\frac{d\mathbf{v}_i}{dt} = - \sum_j m_j \Pi_{ij} A_{\alpha,ij} ,
\]
(7)

where \( \Pi_{ij} \) is the AV tensor and
\[
\tilde{A}_{\alpha,ij} = \frac{1}{2} \left[ A_{\alpha,ij}(h_i) + \tilde{A}_{\alpha,ij}(h_j) \right] .
\]
(8)

We adopt here the Riemann-based formulation proposed by Monaghan (1997) to write the AV tensor as:
\[
\Pi_{ij} = \frac{\alpha_{ij} \nu_{ij}^{AV} \mu_{ij}}{2 \rho_{ij}^2} f_{ij} ,
\]
(9)

where \( \nu_{ij}^{AV} = v_{ij} \cdot \mathbf{r}_{ij}/|\mathbf{r}_{ij}| \) if \( v_{ij} \cdot \mathbf{r}_{ij} < 0 \) but zero otherwise, \( v_{ij} = \mathbf{v}_i - \mathbf{v}_j \), \( \rho_{ij} \) is the arithmetic mean of the two densities and \( \alpha_{ij} = (\alpha_i + \alpha_j)/2 \) is the symmetrized AV parameter. The signal velocity \( v_{ij}^{AV} \) is estimated as
\[
v_{ij}^{AV} = c_i + c_j - 3\mu_{ij} ,
\]
(10)

with \( c_i \) being the sound velocity. The symmetrized AV limiter \( f_{ij} = (f_i + f_j)/2 \), where
\[
f_i = \frac{|\nabla \cdot \mathbf{v}_i|}{|\nabla \cdot \mathbf{v}| + |\nabla \times \mathbf{v}|} ,
\]
(11)

is introduced (Balsara 1995) to suppress AV when in presence of strong shear flows. The individual particle viscosity parameters \( \alpha_i(t) \) are allowed to evolve in time according to the Cullen & Dehnen (2010) scheme, which is found to significantly reduce AV away from shocks (Cullen & Dehnen 2010, V16: see equation (9) and following text). The \( \alpha_i \)'s can vary from a minimum value \( \alpha_{min} = 0.01 \) when shocks are absent, up to a maximum value \( \alpha_{max} = 1.5 \).

2.1.1 Dissipative terms

In SPH, the particle entropy \( A_i \) is generated at a rate
\[
\frac{dA_i}{dt} = \frac{\gamma - 1}{\rho_i^\gamma} (Q_{AV} + Q_{AC} - Q_R) ,
\]
(12)

where the \( Q_{AV} \) term refers to numerical viscosity effects (V16). \( Q_{AC} \) is an artificial conduction (AC) term and \( Q_R \) describes the effects of radiative cooling. The latter is defined as \( Q_R = \Lambda(\rho_i, T_i, Z_i)/\rho_i \), with \( \Lambda(\rho_i, T_i, Z_i) \) being the cooling function, \( T_i \) and \( Z_i \) the particle temperature and metallicity, respectively. Thus, \( Q_R \) is the cooling rate per unit mass.

The presence of the AC term is necessary in SPH simulations (Price 2008) for treating contact discontinuities, such as when studying the growth of Kelvin–Helmholtz instabilities. This term can be written as
\[
\left( \frac{du_i}{dt} \right)_{AC} = \sum_j \sum_\alpha m_j \nu_{ij}^{AC} \rho_{ij} \left[ \alpha_{ij}^C (u_i - u_j) \right] \Delta_{ij} \tilde{A}_{\alpha,ij}/|\mathbf{r}_{ij}| ,
\]
(13)

where \( \nu_{ij}^{AC} \) is the AC signal velocity, \( \alpha_{ij}^C \) is an AC parameter of order unity and \( \alpha_{ij}^C = (\alpha_{ij}^C + \alpha_{ij}^C)/2 \) its symmetrized value.

The form of the AC signal velocity depends on the problem under consideration (Price et al. 2018). An appropriate choice in the presence of self-gravity is found to be
\[
u_{ij}^{AC} = (|\mathbf{v}_i - \mathbf{v}_j| \cdot \mathbf{r}_{ij})/|\mathbf{r}_{ij}| ,
\]
(14)

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which has been checked in several test cases (Valdarnini 2012) and is zero for a self-gravitating system at equilibrium.

The time evolution of the AC parameter $\alpha_C^2$ is regulated by a source term which is proportional to the Laplacian of the particle thermal energy, and by a decaying term which quickly damps $\alpha_C^2$ away from discontinuities. A description of the AC numerical settings is given in Valdarnini (2012).

Finally it is worth noting that incorporating an AC term into the SPH thermal equation significantly improves the agreement with those produced using mesh codes. This difficulty is due to the Lagrangian nature of SPH, in which subgrid diffusion processes are missed. It is shown that introducing an AC term (Wadsley et al. 2008; Valdarnini 2012), the levels of entropies found in galaxy cluster cores are in much better agreement with those produced using mesh codes.

The $Q_R$ term is present in those runs for which radiative cooling is also included. For these simulations the gas physical modeling incorporates star formation and energy feedback from supernovae as well. For the numerical aspects of the cooling implementation we refer to Valdarnini (2006).

2.2 Initial condition setup

For a variety of initial conditions, we perform N-body/hydrodynamical ISPH simulations to study the collisions between galaxy clusters. Each cluster consists of an isolated spherical halo initially in equilibrium, composed of dark matter and gas particles. The initial conditions of our idealized binary cluster mergers are very similar, but not identical, to those of ZuHone (2011), hereafter Z11. We study collisions between a primary and a secondary cluster, with the primary mass always set to $M_{200} = 6 \times 10^{14} M_\odot \equiv M_1$. Here $M_{200}$ is the cluster mass within the radius $r_{200}$. We define $r_\Delta$ as the radius at which the cluster mean density is $\Delta$ times the cosmological critical density $\rho_c(z)$:

$$M_\Delta = \frac{4\pi}{3} \Delta \rho_c(z) r_\Delta^3.$$  

In the following, we assume $z = 0$ as the redshift at which $r_{200}$ is calculated. For the secondary, with cluster mass $M_{200} \equiv M_2$, we consider three different mass ratios $R = M_2/M_1 = 1.1, 1.3$ and $1.10$.

For each collision with mass ratio $R$, we consider three different impact parameters $b$: a head-on cluster collision with $b = 0$, and two off-axis mergers with $b/r_{200} = 0.3$ and $0.6$. Here $b$ is the impact parameter of the collision when the distance $d_{12}$ between the center of mass of the two clusters is $d_{12} = r_{200} + r_{200}$; see Figure 1 of Z11 for a geometric description of the collision set up. The procedure we use to assign initial separations and relative velocities between the two clusters is described in Section 2.2.3.

2.2.1 Dark matter halos

We assume spherical symmetry for the initial dark matter (DM) and gas mass distribution. For the DM density, we adopt an NFW profile (Navarro et al. 1997)

$$\rho_{DM}(r) = \frac{\rho_s}{r/r_s (1 + r/r_s)^2}, \quad 0 \leq r \leq r_{200},$$

where $c_{200} = r_{200}/r_s$ is the concentration parameter. To avoid a divergent total mass, outside $r_{200}$ the DM density profile is suppressed exponentially (Kazantzidis et al. 2004) up to a final radius $r_{\max} = \xi r_{200}:

$$\rho_{DM}(r) = \rho_{DM}(r_{200}) (r/r_{200})^\delta \exp \left(-\frac{r - r_{200}}{r_{\text{decay}}}ight), \quad r_{200} < r < r_{\max},$$

where $r_{\text{decay}} = \eta r_{200}$ is the truncation radius, and the parameter $\delta$ is set by requiring the first derivative of the DM density profile to be continuous at $r = r_{200}$

$$\delta = -\frac{r_s + 3r_{200}}{r_s + r_{200}} + \frac{r_{200}}{r_{\text{decay}}}. $$

For the runs presented here, we set the truncation parameters to the values $(\xi, \eta) = (2, 0.2)$; this choice will be motivated in Section 2.3. For a cluster of given mass, the density profile is then specified by the parameter $c_{200}$. For our two test clusters with $M_{200} > 10^{14} M_\odot$ we set the value of $c_{200}$ using the CLASH $c - M$ relation of Groener et al. (2016)

$$c_{200} \approx 3.66(M_{200}/M_\odot)^{0.32},$$

where $M_\star = 8 \times 10^{14} M_\odot h^{-1}$. For the cluster C3 with $M_{200} < 10^{14} M_\odot$ we set $c_{200} = 7.03$. This value is obtained by using the following $c_{200} - M_{500}$ relation for galaxy groups (Gastaldello et al. 2007; Sun et al. 2008)

$$c_{500} = 3.96(M_{500}/10^{14} M_\odot)^{0.226},$$

and solving numerically the halo profile to obtain $c_{200}$. Table 1 lists several initial parameters of the three idealized clusters we use to construct our simulation suite.

A numerical realization of the DM density profile is then constructed by first evaluating the enclosed DM mass $M_{DM}(r)$ within the radius $r$, which is normalized so that it is equal to $(1 - f_b)M_{200}$ at $r_{200}$. We subsequently invert $q(r) = M_{DM}(r)/M_{DM}(r_{\max}) = y$, where $y$ is a uniform random number in the interval $[0,1]$, to obtain the radial particle coordinate $r$. Finally, Cartesian coordinates are assigned to the particle by randomly orienting the particle position vector $\mathbf{r}$.

Kazantzidis et al. (2004) showed that, for exponentially truncated NFW halos, particle velocities are accurately determined if their energies are drawn from the system distribution function $f(E)$. For spherical symmetric systems

$$\rho_v(r) = 4\pi \int_0^{\Psi(r)} f(E) \sqrt{2[\Psi(r) - \frac{E}{\Psi}]} \, dE,$$

where $\Psi(r) = -\Phi(r) = -(\Phi_h + \Phi_g)$ is the relative gravitational potential and $E = \Psi - v^2/2$ is the relative energy. Here the subscripts $h$ and $g$ denote the DM and gas components, respectively.

Equation (21) can be inverted (Binney & Tremaine 1987) to give

$$f(E) = \frac{1}{\sqrt{8\pi}} \left[ \int_0^E \frac{d^2 \rho_v}{d\Psi^2} \frac{d\Psi}{\sqrt{\Psi - E}} + \frac{1}{\sqrt{E}} \frac{d\rho_v}{d\Psi} \Bigg|_{\Psi=E} \right]. $$

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The boundary term on the rhs of the equation is zero for any sensible choice of \( \Psi(r) \) and \( \rho(r) \) [Kazantzidis et al. 2004]. The second order derivative \( d^2\rho_s/d\Psi^2 \) can be expressed as

\[
\frac{d^2 \rho}{d\Psi^2} = \left( \frac{r^2}{GM} \right)^2 \left[ \frac{d^2 \rho}{dr^2} + \frac{d \rho}{dr} \left( \frac{2}{r} - \frac{4\pi \rho r^2}{M} \right) \right], \tag{23}
\]

which has the advantage of avoiding numerical differentiation in the integral [22], since \( \rho \) is known analytically.

The function \( f(\mathcal{E}) \) is then evaluated numerically and its values tabulated over a grid of energies. For a given energy \( \mathcal{E} \) the value of \( f \) is obtained by interpolation. For a particle at position \( r \) with energy \( \mathcal{E} \in [0, \Psi(r)] \), we randomly draw pairs \((\mathcal{E}, \Psi)\) and use an acceptance-rejection method [Kuijken & Dubinski 1994; Zemp et al. 2008; Drakos et al. 2013] to obtain the particle speed \( v = \sqrt{2\Psi(r) - \mathcal{E}} \). As for the particle position, the direction of the velocity vector is randomly oriented.

These prescriptions for generating DM particle distributions at equilibrium are widely used by many authors in numerical simulations of merging cluster galaxies, for which initially DM halos are described by an exponentially truncated NFW profile. However, a major drawback of the method introduced by Kazantzidis et al. [2004] is that the second derivative \( d^2\rho_s/d\Psi^2 \) is discontinuous at \( r = r_{200} \). This implies that the behavior of the function \( f(\mathcal{E}) \) can become inconsistent for certain values of \( r_{\text{decay}} \) (Zemp et al. 2008; Drakos et al. 2017), thus compromising halo stability over cosmological timescales. Zemp et al. [2008] recommend the choice \( \eta = r_{\text{decay}}/r_{200} = 0.3 \); here we show in Section 2.2.2 that by setting \( \eta = 0.2 \) one can obtain sufficiently stable halos.

### 2.2.2 Baryonic halos

We assume hydrostatic equilibrium to construct the cluster gas initial conditions. Following [21], we initialize gas density and temperature profiles by specifying analytically the initial cluster entropy profile. Physically, this would be best-represented by giving the specific physical entropy per particle in the gas, \( s(r) \). However, both observations of the gas in clusters and previous simulations have instead utilized a related entropy parameter [written as \( S(r) \) or \( K(r) \)]. To be consistent and allow easier comparison to observations and previous simulations, we will adopt this convention. Thus, in the following we refer to the “gas entropy” as this entropy parameter \( S \equiv k_B T/n_e^{2/3} \), where \( k_B T \) is the gas temperature in keV and \( n_e \) the electron number density. From the Sackur-Tetrode equation [Landau & Lifshitz 1981], it is easily shown that \( s = (3/2)k_B \ln(S) + \text{const.} \)

where the constant is not important here. Thus, it is straightforward to convert between \( s \) and \( S \). However, one should be aware that differences in the entropy are exaggerated by \( S \), since it is only its logarithm that enters into the physical entropy.

CC clusters are observationally characterized by dense, compact cores with cooling times shorter than \( H_0^{-1} \). A key feature of CC clusters is that of having a level of central entropy below a threshold value \( \approx 60 \text{ keV cm}^2 \).

We then adopt for the gas entropy profile an observationally motivated functional form which consists of a power law behavior and an entropy floor value:

\[
S(r)/S_{500} = S_0 + S_1 \left( \frac{r}{r_{500}} \right)^\alpha, \tag{24}
\]

where \( S_0 \) and \( S_1 \) are related entropy parameter \([\text{written as } S_0 \text{ or } S_1]\). Zemp et al. (2008) to obtain the particle speed \( v_{500} \)

\[
\begin{align*}
\frac{d^2 \rho}{d\Psi^2} &= \left( \frac{r^2}{GM} \right)^2 \left[ \frac{d^2 \rho}{dr^2} + \frac{d \rho}{dr} \left( \frac{2}{r} - \frac{4\pi \rho r^2}{M} \right) \right], \tag{23}
\end{align*}
\]

The numerical integration of these equations is found more manageable if one integrates the temperature instead of pressure. The latter can be expressed as

\[
P = \frac{\rho g k_B T}{\mu m_p} = K T^{5/3}, \tag{27}
\]

where \( K(r) = S(r)(\mu/\mu_e)^{2/3}/(\mu m_p)^{5/3} \) and for the mean molecular weights we assume \( \mu = 0.59 \), \( \mu_e = 1.14 \). Equations 28a and 28b now read

\[
\begin{align*}
\frac{dT}{d\log r} &= -\frac{2\mu m_p}{5k_B} \frac{GM_{\text{tot}}(\lesssim r)}{r} + \frac{3}{5} \frac{d\log K}{d\log r}, \tag{28a}
\end{align*}
\]

\[
\begin{align*}
\frac{d\log M_g}{d\log r} &= \frac{4\pi r^3 \rho_g}{M_g(\lesssim r)}. \tag{28b}
\end{align*}
\]

To integrate these equations it is necessary to specify two boundary conditions. Our first condition is that \( r(M_g = 0) = 0 \), whilst the second requires that the halo gas mass at \( r = r_{500} \) yields a gas mass fraction \( f_g = M_g/M_{\text{tot}} \) given by the measured relation [Sun et al. 2009]

\[
f_g = 0.035 h^{-3/2}(M_{500} h/7 \times 10^{12} M_\odot)^{0.135}. \tag{29}
\]
Table 2. The simulation ID is defined according to the mass ratio $R$ and the impact parameter $b$ of the collision, second column the corresponding ID of the Z11 runs. The simulations start at $t_s = 0$, the time $\tau_{\text{num}}$ is the simulation time at which the center of mass position and velocities of the two halos are closest to the corresponding Z11 initial conditions. $\varepsilon_x$ and $\varepsilon_v$ are the corresponding relative errors, $\tau_{\text{K}}$ is the analytic solution given by the Kepler’s problem.

| ID | ID Z11 | $\tau_{\text{num}}$[Gyr] | $\tau_{\text{K}}$[Gyr] | $\varepsilon_x$ | $\varepsilon_v$ |
|----|--------|--------------------------|------------------------|----------------|----------------|
| R01600 | S1 | 3.81 | 3.73 | 1.06 $\times 10^{-2}$ | 0.115 |
| R01603 | S2 | 3.94 | 4.17 | 2.3 $\times 10^{-2}$ | 0.09 |
| R01606 | S3 | 4.81 | 5.10 | 4.1 $\times 10^{-2}$ | 0.13 |
| R03000 | S4 | 2.81 | 2.73 | 1.4 $\times 10^{-2}$ | 0.09 |
| R03003 | S5 | 2.87 | 2.99 | 1.2 $\times 10^{-2}$ | 0.08 |
| R03006 | S6 | 3.31 | 3.41 | 2.8 $\times 10^{-2}$ | 0.11 |
| R10600 | S7 | 2.32 | 2.32 | 3.7 $\times 10^{-2}$ | 0.06 |
| R10603 | S8 | 2.44 | 2.53 | 1.03 $\times 10^{-2}$ | 0.08 |
| R10606 | S9 | 2.75 | 2.87 | 2.4 $\times 10^{-2}$ | 0.09 |

To construct our gas density and temperature profiles we proceed as follows. For a given set of entropy parameters $(S_0, S_1, \alpha)$, we initially choose an arbitrary value of $P(r = 0) \equiv P_0$ which is used to compute $\rho_0$ and $T_0$. Equations (28a) and (28b) are then numerically integrated up to $r = r_{200}$, after which the value of $f_g$ is contrasted with that of $f_{\text{gas}}$. We iterate the whole procedure in order to bracket the value of $P_0$ until the quantity $|f_g - f_{\text{gas}}|/f_{\text{gas}}$ is below a certain threshold value ($\lesssim 1\%$). When this condition is satisfied and $P_0$ is a root value, we propagate the solution outward to $r_{\text{max}} = 2r_{200}$. This normalization procedure implies a gas mass fraction at $r = r_{200}$ that may differ from the cosmic value $f_b$, and therefore that may not be entirely consistent with the normalization adopted in Section 2.2.1 for the DM component, the latter making use of the cosmic gas fraction $f_0$ to set the halo DM mass to $(1 - f_0)M_{200}$ at $r_{200}$. However, we have verified that for all the considered halos the difference between the two gas fractions at $r = r_{200}$ is very small (say $\lesssim 1\%$).

The radius up to which the gas density profiles is continued beyond $r_{200}$ coincides here with the DM truncation radius. The choice of the gas truncation radius requires some care. Setting this radius to $r_{200}$ would imply that gas particles at the halo edge will begin to flow outward, owing to the absence of an external pressure. This effect implies a steepening of the gas density profile and a mass leak, which can persist even in the absence of an external pressure. This effect implies a steepening of the gas density profile and a mass leak, which can persist even in the absence of an external pressure.

To avoid this mass-leaking issue, an approach commonly employed (Turner et al. 1995; Ricker & Sarazin 2001; Poole et al. 2006; McCarthy et al. 2002; Donnert 2014) consists of surrounding the gaseous halo with a low-density, dynamically negligible, confining medium. However, this procedure comes at the cost of adding a large number of SPH particles to the simulations. We choose here a different approach by extending the gaseous halo beyond $r_{200}$ and up to a maximum radius $r_{\text{max}} = 2r_{200}$.

The shell between $r_{200}$ and $r_{\text{max}} = 2r_{200}$ then acts as a buffer zone that is able to keep the gas particles within $r_{200}$ confined. Clearly the edge particles at $r \simeq 2r_{200}$ will begin to flow outward, leading to a steepening of the density profile. This steepening can be considered unimportant as long as collisions between clusters occur on timescales much shorter than that necessary to modify the initial gas density profile at radii $r \lesssim r_{200}$. It will be seen in Section 2.2.1 that these conditions are always verified for the cluster collisions we consider.

We assume $\alpha = 0.93$ in Equation (28) and we report in Table 1 the coefficients $S_0$ and $S_1$, together with $S_{500}$, for each of the three test clusters. These coefficients are the best-fit values of the entropy profile given by Ghirardini et al. (2019) for their CC subsample of X-COP clusters (their Table 3, CC entry), but increased by about a factor of $\approx 30\%$. It has been found necessary to introduce this offset in the coefficients to obtain a physical meaningful solution to Equations (28a) and (28b) up to $r_{\text{max}}$. This is because, using the original coefficients, it is not possible to propagate the numerical solution beyond $r_{200}$, the radial profile of $P(r)$ being characterized by a very steep decline with radius. As a result, at radii slightly beyond $r_{200}$, the pressure becomes numerically consistent with zero.

By setting $r \simeq 10^{-2}r_{500}$, we obtain entropy core values for our test clusters of $S/S_{500} \approx 3 \times 10^{-2}$ from Table 1. For cluster C1 this gives $S(r = 10^{-2}r_{500}) \approx 28\,\text{keV cm}^2$, still below the threshold value for CC clusters (Cavagnolo et al. 2004). These core values at $r \simeq 10^{-2}r_{500}$ are in the lower portion of the observed range of entropies for CC clusters; see Figure 6 of Ghirardini et al. (2019).

Note that for cluster C3, the coefficient $S_1$ has been further increased by $\approx 70\%$, with respect the best-fit value of 1.35 given by Ghirardini et al. (2019). This is because all of the CC clusters of the X-COP sample used by the authors have $M_{200} \geq 6 \times 10^{14} M_\odot$, whereas here cluster C3 has $M_{200} = 6 \times 10^{13} M_\odot$. From Figure 10 of Sun et al. (2009) it can be seen that a core entropy of $S/S_{500} \approx 3 \times 10^{-2}$ at $r \simeq 10^{-2}r_{500}$ is well below the observed range of entropy values reported by Sun et al. (2009) for their group sample. Finally, for cluster C1 (C3) we have at $r = 10\,\text{kpc}$ an the entropy core value of $\approx 25\,\text{keV cm}^2$ ($\approx 12\,\text{keV cm}^2$), which is about $20\%$ higher than the entropy value found at the same radius for the corresponding test cluster of Z11.

Once the initial gas density and temperature profiles are computed for the range $0 \leq r \leq r_{\text{max}}$, we use them to construct the initial particle configuration. To this end, we store the profiles on a very fine grid. This in order to obtain $\rho(r)$ and $T(r)$ at a generic radius $r$ using grid values.

In SPH, a non-trivial issue is the realization of the gas particle distribution which must reproduce the required density profile $\rho(r)$. From Equation (1), one can see that in SPH the density at a given particle position depends on the masses and relative positions of nearby particles. This in turn implies that any specific realization of particle positions which satisfies the prescribed profile $\rho(r)$, must be found by solving simultaneously for all the particle positions. This is a difficult task and several methods have been devised to solve the problem.

A relatively simple approach, which starts from a uniform distribution and solves Equation (1) by varying particle masses, cannot be applied in SPH. This is because numerical instabilities are found to arise (Monaghan & Price 2004) when large mass contrasts are present between particles. Similarly, a random realization of particle positions generated from the specified $\rho(r)$, as we did in Section 2.2.1, to construct the DM particle positions, cannot be used here.
The noise induced by Poisson sampling implies the development of large fluctuations which in turn quickly perturb the initial equilibrium profile.

These difficulties have led many authors to develop alternative methods to solve the problem of properly generating initial conditions in SPH. These methods can be summarized as follows: lattice stretching (Herant 1994; Rosswog & Price 2002), viscous damping (Wang & White 2007; Price & Monaghan 2007; Pakmor et al. 2012; Price et al. 2013), relaxation (Zurek & Benz 1986; Nagasawa et al. 1988), space partition based either on tessellation (Pakmor et al. 2012; Raskin & Owen 2016; Reinhardt & Stadel 2017) or weighted Voronoi tessellations (Diehl et al. 2013; Vela et al. 2013; Arth et al. 2019).

Initially, we implemented the method of Diehl et al. (2013) to setup our SPH initial conditions. The algorithm is based on a Voronoi tessellation in which particles are moved iteratively toward a relaxed configuration. However, during the merging simulations the halo profiles constructed according to this procedure were found to deviate from the initial equilibrium solution. In several cases this happened on a timescale shorter than that occurring between the start of the simulation and the cluster collision.

We interpret this behavior as a direct consequence of the entropy profile we use to construct our initial conditions. For CC clusters, the average entropy profile implies an equilibrium solution with a very steep density profile toward the center of mass of the two clusters being at the origin. The radial transformation is such that the new gas position is obtained by a radial transformation applied to the coordinates of a uniform glass distribution of points. This is a minimal noise configuration, and it is generated by applying to an initial Poisson distribution of points a reversed gravitational acceleration together with a damping force (Wang & White 2007). Particle positions are advanced until a low energy state is reached.

The radial transformation is such that the new gas positions are consistent with the desired initial gas mass profile. This profile is now much more stable than that obtained from a random realization (Section 2.2.3), but its stability properties are not yet sufficient to consistently satisfy the initial condition set up required by the merging runs studied here.

To further improve the stability properties of the gaseous halo, we add to the momentum equation a time-dependent friction term (Price & Monaghan 2007; Pakmor et al. 2012)

\[
\frac{dv}{dt} = \frac{v - V_{em}}{\tau_{damp}} + \left( \frac{dv}{dt} \right)_{SPH},
\]

where the second term on the rhs is given by Equation (3). \(V_{em}\) is the halo center of mass velocity, and \(\tau_{damp}\) is a time-dependent damping time scale. This is written as

\[
\tau_{damp} = \frac{\tau_{dyn}}{\alpha_{damp}},
\]

where \(\tau_{dyn} = 1/\sqrt{4\pi G \rho_0}\) is the local dynamical timescale, \(\rho_0(x)\) is the particle DM density, and \(\alpha_{damp}\) is a friction parameter which controls the strength of the friction. The DM density \(\rho_0\) must be calculated at run time according to an SPH prescription, but one can introduce a DM to gas density ratio: \(\beta_{dyn} = \rho_g/\rho_0\). The halo stability can then be exploited to avoid the calculation of the DM density by using \(\beta_{dyn} \rho_0\) in place of \(\rho_0\) in Equation (31). As it will be seen from the plots of Section 2.2.3 a conservative value for \(\beta_{dyn}\) is obtained by setting \(\beta_{dyn} \simeq 50\). It has been found that this choice also has an impact on the value of \(\alpha_{damp}\). In principle \(\tau_{damp}\) should be a small fraction of \(\tau_{dyn}\), but with the adopted value of \(\beta_{dyn}\) very stable halos are already obtained when \(\alpha_{damp} \simeq 3\).

Our merging simulations are then performed by using the generalized momentum equation (30) as the simulations start, and switching off (\(\alpha_{damp} = 0\)) the friction parameter at a simulation time which depends on the initial merging kinematics (Section 2.2.3). Initially, we set gas particle velocities to zero and temperatures are assigned by interpolating grid values. These are calculated from the numerical solution and the interpolation is done according to the radial particle coordinates.

Finally, we set the mass of DM and gas particles according to the scaling \(m_d \simeq 8 \times 10^8 M_\odot (M_{200}/2 \times 10^{14} M_\odot)\) and \(m_g = f_{sm} m_d/(1 - f_0)\) \(\simeq 0.16 m_d\), respectively. These mass assignments are consistent with previous findings (Valdarnini 2014), in which ICM profiles extracted from a set of hydrodynamical cluster simulations were found numerically converged when similar settings were adopted for the particle masses. However, as outlined before, in SPH numerical instabilities can arise in the presence of very different gas particle masses. This implies that the simulation numerical resolution is enforced by the smallest mass of the binary system. We thus write

\[
m_d \simeq 8 \times 10^8 M_\odot (M_{200}^{(2)}/2 \times 10^{14} M_\odot).
\]

For cluster C1, the total number of gas particles ranges then from \(N_g^{(1)} \simeq 2 \times 10^7\) when the collision mass ratio is \(R = 1 : 1\), up to \(N_g^{(1)} \simeq 2 \times 10^8\) when \(R = 1 : 10\). For cluster C3 one has \(N_g^{(2)} \simeq 1.27 \times 10^5\). The gravitational softening parameters of the particles are set according to the scaling \(\epsilon_i = 15.8 \cdot (m_i/6.2 \times 10^8 M_\odot)^{1/3}\). Additionally, in some test cases we run high-resolution (HR) simulations in which the particle masses are scaled down by a factor 4, with respect the reference value given by equation (32).

### 2.2.3 Initial merger kinematics

To construct the orbits of our merging simulations, we choose a Cartesian system of coordinates \(x, y, z\), with the center of mass of the two clusters being at the origin. The orbits are initialized in the \((x, y)\) plane at \(z = 0\), with \(\{d_0, V_0\}\) being the initial separation and relative velocity vectors, respectively. Thus, the initial \((x, y)\) coordinates of the two cluster center of mass read \(-d_0\cos 0/1 + R)\) and \((d_0\cos 0, d_0\sin 0) R/(1 + R)\). Similarly, the velocity components are given by \(-V_0\sin 0/(1 + R)\) and \((V_0\cos 0, V_0\sin 0) R/(1 + R)\).
As already outlined, our collision parameter space is the same as in Z11. However, there is here a significant difference in the initial condition setup of the halos. For the reasons discussed in Sections 2.2.1 and 2.2.2 of the initial DM and gas mass profiles are continued beyond $r_{200}$ and extended up to $2r_{200}$. This implies that, unlike in Z11, the relative initial separation cannot be set here to the sum of the two virial radii, but to twice its value: $d_{in} = 2(r_{200} + r_{200})$. As discussed in Section 2.2.3 this is to avoid a significant overlap at the start of the simulation between the mass profiles of the two halos, which in turn would soon put the profiles out of equilibrium.

Our initial condition vectors $\{d^{in}, V^{in}\}$ at $t = 0$ must then be chosen such that, at some later simulation time $t$, the orbit of the binary cluster system produces the initial conditions of Z11. These consists of a separation $d_{in}/2$ between the two cluster centers of mass, with a collision impact parameter $b$ and relative infall velocity $V \simeq 1.1/\sqrt{GM_{200}/r_{200}}$. The latter value is justified by cosmological simulations (Vitvitska et al. 2002).

In order to realize these settings, we adopt a procedure similar to that described in Poole et al. (2004). For a specified set of initial conditions taken from Z11, we first approximate the two clusters as point-like and accordingly assign positions and velocities to the two points. We tag this orbital status as occurring at the time $t_f$, i.e. the start of the simulations. We now numerically solve Kepler’s problem by seeking the time $t_i < t_f$ such that the separation between the two points is $d_{in}$. The orbital positions and velocities at $t_i$ then complete the solution vectors $\{d^{in}, V^{in}\}$.

To account for tidal distortions we first run a DM only merging simulation, using as initial conditions the solution vectors $\{d^{in}, V^{in}\}(t = 0)$ previously determined. During the simulations we denote as $\{\mathbf{X}_{cl}, \mathbf{V}_{cl}\}(t)$ the center of mass position and velocities of the two halos. These vectors are contrasted with the Z11 initial conditions $\{\mathbf{X}_{cl}, \mathbf{V}_{cl}\}(\tau_K)$, which the binary system must reproduce at the simulation time $t_{sim} = t_f - t_i \equiv \tau_K$. To quantify the deviations between the specified set of initial conditions and the numerical solution we define the following norms

$$
\left\{
\begin{array}{l}
\varepsilon_{x}^{(d)}(t) = \left\| \mathbf{X}_{cl}(t) - \mathbf{X}_{cl}^{(2)}(\tau_K) \right\| \\
\varepsilon_{v}^{(d)}(t) = \left\| \mathbf{V}_{cl}(t) - \mathbf{V}_{cl}^{(2)}(\tau_K) \right\|
\end{array}
\right.
$$

with $d = 1, 2$ being the halo index.

We define as position error $\varepsilon_{x}(t)$ the maximum of the two error norms: $\varepsilon_{x}(t) = \text{MAX}(\varepsilon_{x}^{(1)}(t), \varepsilon_{x}^{(2)}(t))$, the velocity error $\varepsilon_{v}(t)$ being similarly defined. These errors are computed and saved at run times $t_m = \tau_K + m\Delta t$, centered around $\tau_K$. We set the grid spacing to $\Delta t = 1/16$ Gyr and $m$ is an integer ranging between $-20$ and $20$. Finally, the simulation time at which our constructed set of error values has a minimum, is identified as the simulation time $t \equiv 0$ in the corresponding merging run of Z11. We label as $\tau_{num}$ this solution time obtained numerically. Table 2 lists the values of $\tau_{num}$ and $\tau_K$ for each of our merging runs, together with the notation we use to label the different simulations. Note that the difference $\tau_{num} - \tau_K$ is smallest for head-on collisions, while it is largest ($\simeq 0.3$ Gyr) for the $R = 1.1$ off-axis merger with $b = 0.6$. The simulation time $t_s$ here is then related to that of Z11 by the relation:

$$
t^2 = t_s - \tau_{num}.
$$

Our merging simulations are performed up to a simulation time $t_s = \tau_{num} + t_{fin}$, where $t_{fin} = 10$ Gyr. We analyze simulation results when $t^2 \geq 0$, at epochs spaced by 1 Gyr. To ease comparisons between our results and those of Z11 hereafter we will always use the simulation time $t^2$, which we abbreviate as $t$.

The damping factor $\alpha_{\text{damp}}$ in Equation 25 is set to zero when $t \geq 0$, but with some exceptions (see later). This guarantees that at $t = 0$ our halo entropy profiles possess the correct radial behavior. Our merging simulation suite is constructed by performing both adiabatic and radiative simulations. For adiabatic runs we consider all of the collision parameter space, consisting of nine different merging simulations. For radiative simulations, we run only a limited number of mergers because of the high computational cost of the simulations.

### 2.3 Stability tests

As already discussed in Section 2.2.1 the stability of spherically symmetric DM halos with an exponentially truncated NFW profile depends critically of how initial particle velocities are assigned. According to Kazantzidis et al. (2004), the long term halo evolution is significantly affected if the particle velocities are initialized using the local Maxwellian approximation. By contrast, much more stable halos are obtained when the initial particle energies are consistently extracted from the equilibrium distribution function $f(\mathbf{E})$. However, the choice of the halo truncation parameters $\xi = r_{max}/r_{200}$ and $\eta = r_{decay}/r_{200}$ is not entirely arbitrary. In particular, an overly sharp truncation ($\eta \leq 0.1$) can lead to instabilities in the halo evolution (Zemp et al. 2008; Drakos et al. 2017). The solution is increase the truncation radius ($\eta \simeq 0.3$), truncating the atmosphere more smoothly, as already done in some merging runs (Zhang et al. 2014).

To validate our choice of the truncation parameters ($\xi, \eta$) we studied the evolution over cosmological timescales of three isolated DM halos. All of the halos have $M_{200} = 6 \times 10^{14} M_{\odot}$, but their initial density profiles have different truncation parameters ($\xi, \eta$). The three pairs of values we consider are ($\xi, \eta$) = (3, 0.3), (1.4, 0.1) and (2, 0.2). We refer to the corresponding halo realizations as DMa, DMb and DMc, respectively. Initial particle position and velocities are initialized according to the procedures described in Section 2.2.1 and we use Equation 26 to set the DM particle mass to $m_{p} \simeq 2.4 \times 10^{5} M_{\odot}$. The number of DM particles $N_{p}$ then ranges from $N_{p} \simeq 3 \times 10^{5}$ for the DMb halo, up to $N_{p} \simeq 3.8 \times 10^{5}$ in the case of the DMa halo.

Figure 4 shows the time evolution of the density and velocity dispersion profiles for the three halos at four different time slices: $t = 0$, 1, 5 and 10 Gyr. For the range of initial conditions we consider, the time span $t_{\text{fin}}$ occurring between the start of the simulation and the direct hit between the primary and secondary cluster cores (see later) ranges from $\simeq 5$ Gyr, in the case of head-on collisions, up to $\simeq 8$ Gyr for the R10606 initial merging configuration. The different epochs displayed in Figure 4 have been chosen with the criterion of covering the whole range of time spans $t_{\text{fin}}$.

As it can be seen, the best stability properties are exhibited by the DMa halo. For this halo, both density and velocity dispersion profiles are quite stable up to $\simeq 10$ Gyr.
The halos are initially in equilibrium, with the particle distributions realized according to the procedures described in Figure 1. All of the halos have $M_{200} = 6 \times 10^{14} M_{\odot}$, but their initial conditions differ in the choice of the truncation parameters $(\xi, \eta)$ (see text). Density is in units of the cosmological critical density and velocity dispersion in units of $\sigma_{200} = \sqrt{GM_{200}/r_{200}} \equiv \sigma_v$. Different line styles refer to different epochs, as indicated in the bottom left panel, where the time is in Gyr. (Note that the times decrease from top to bottom in the legend in this and subsequent figures.)

This is in accord with previous findings (Zemp et al. 2008; Drakos et al. 2017), and confirms that setting $\eta = 0.3$ is the safest choice when stability is an issue. However, this choice of $r_{\text{decay}}$ requires the continuation of the DM halo beyond $r_{200}$ and up to $r_{\text{max}} \simeq 3r_{200}$, if one wants to avoid an abrupt truncation in the density profile. This choice of the truncation parameters $(\xi, \eta)$ in turn implies that the simulations will begin with a significant overlap between the two DM halos, if the initial separation between the two clusters is chosen to be $d_{\text{in}} = r_{200}^2 + r_{200}^2$, as in $Z11$. These initial settings then put the initial mass profiles of the two halos out of equilibrium, and it is not clear what the impact of these initial conditions is on ICM properties during cluster mergers (McCarthy et al. 2007).

We choose here to put the initial separation between the two cluster center of mass to the value $d_{\text{in}} = \xi (r_{200}^2 + r_{200}^2)$, this choice of the initial setup being clearly advantageous because the merging runs are then performed with halos initially at equilibrium. To avoid very large values ($\gtrsim 10$ Gyr) of the simulation time $\tau_{\text{sim}}$, when the two clusters orbital parameters best approximate the $Z11$ initial conditions, we decided here to use a value of $\xi$ smaller than that of the DMa halo ($\xi = 3$).

Figure 1 shows that the stability properties of the DMc halo, having $(\xi, \eta) = (2, 0.2)$, are much better than those of the DMb halo with $(\xi, \eta) = (1.4, 0.1)$. There is some evolution in the density profile beyond $r_{200}$, but the velocity dispersion profile is much more stable than those of the DMb halo. To set up initial conditions for our DM halos, we thus adopt as truncation parameters the pair of values $(\xi, \eta) = (2, 0.2)$.

This choice is motivated by the criterion of having the DM halos be as stable as possible, but without extending them very far beyond $r_{200}$. It must be stressed that this choice does not necessarily imply that final gas profiles are significantly affected by using the pair $(\xi, \eta) = (1.4, 0.1)$. For instance, Figure 2 of $Z11$ shows little evolution in the gas profiles of an isolated halo, although the DM component is initialized by setting $\eta = 0.1$.

We now investigate the stability properties of halos which contain both DM and gas. As for the DM only tests, we always set the halo mass at $r_{200}$ to $M_{200} = 6 \times 10^{14} M_{\odot}$. We setup the gas density and temperature profiles according to the procedure described in Section 2.2.2, the entropy profile parameters being those of cluster C1. All of the halos have then the same physical parameters and analytical profiles. The total DM and gas halo masses at $r_{\text{max}} = 2r_{200}$ are then $M_{DM} \simeq 6.6 \times 10^{14} M_{\odot}$ and $M_{\text{gas}} \simeq 9.2 \times 10^{13} M_{\odot}$, respectively. Accordingly, from Equation (32) the number of DM (gas) particles is $N_{\text{p}} \simeq 2.7 \times 10^5 \ (2 \times 10^5)$. 

![Figure 1.](image-url)
We initially consider two different particle realizations of the initial gas density profile. The first halo (RN) has gas particle positions drawn from a uniform random distribution. This is the simplest approach to realize the desired density profile, but for the reasons discussed in Section 2.2.2, the stability of its gas profiles can be considered very poor. Thus, we use the profile evolution of this halo realization as a benchmark, against which to assess the stability properties of other procedures. For the second halo (GL) gas positions are obtained by transforming the radial coordinates of a glass-like configuration of points. The transformation is consistently done by numerically solving for the radial coordinate of each particle that satisfies the requested mass profile.

For the two halo realizations, Figure 2 shows the time evolution of the gas density, temperature, and entropy profiles. The temperature is in units of $T_{200}$, the mass-weighted temperature within $r_{200}$, and entropy in units of $S_{500}$. Additionally, we also show the radial profiles of the dynamical time in gas and the cooling time (defined in equation 35 below).

As expected, the plots clearly show the very poor stability properties of the RN halo. On the other hand, there is some improvement when using glass-like initial conditions. The entropy profile $S_{GL}(r)$ exhibits a better stability, with deviations from the initial reference profile systematically smaller than in the RN case. At $t = 1$ Gyr there are small deviations in the very inner region ($r \lesssim 0.02 r_{500}$), and at $t = 10$ Gyr the profile $S_{GL}(r)$ is similar to that of $S_{RN}(r)$ at $t = 5$ Gyr.

These results demonstrate that in order to improve the profile stability of our SPH particle realization, one must resort to more sophisticated methods. As outlined in Section 2.2.2, the use of a relaxation method (Diehl et al. 2013) was found to improve the profile stability, but not in a very significant way with respect the GL run. Motivated by previous findings (Price & Monaghan 2002, Pakmor et al. 2012), in order to keep the initial configuration in equilibrium we then add a time-dependent damping force to the SPH momentum equation. The procedure and the parameter settings are described in the previous Section.

For this test case, which we label as FD, Figure 3 shows the time evolution of the different gas profiles. The meaning of the different panels and lines being the same of Figure 2. The profile evolution clearly indicates that the damping method is very good in maintaining the stability of the initial SPH particle realization, and in turn the gas profiles. We have verified that this behavior holds for clusters C2 and C3 as well.

Accordingly, we implement this setup procedure to construct stable gas profiles. The initial particle positions are extracted from a uniform glass-like distribution, as for the GL halo. The hydrodynamic SPH force equation is then generalized in Equation (43) to incorporate a friction term. The latter is present from the start of the simulation ($t_{hit} = 0$) up to the time when the binary system has reached the optimal configuration aimed at reproducing the initial conditions ($t_{hit} = \tau_{num}$). After this epoch ($t = t_{hit} - \tau_{num} \geq 0$) the friction term is switched off ($\sigma_{damp} = 0$) in the momentum equation. With these settings, we can consistently compare our simulation results with those of (Z13) having realized the same cluster orbital and gas profile initial conditions.

However, it must be stressed that in some mergers the clusters will come in contact having higher core entropies than those initially specified. We define the time span $t_{hit}$ as that occurring between the start of the simulation and when the two clusters cores collide or interact strongly. For our head-on mergers ($b = 0$), we find that this is well approximated by the epoch when the distance between the two cluster centers of mass is smaller than $r_{h}$. However, for our offset mergers with $b = 0.3$ or 0.6, we find that the secondary core passes by the primary core without being significantly affected during the first pericentric passage. After the secondary reaches the apocenter, it falls more directly into the primary core. Thus, this second encounter is nearly head-on, and we therefore apply the same definition as for $b = 0$ to this second encounter. Empirically, we find that this timescale does approximate the time when the secondary core is significantly affected. For example, for the R10606 merging run one has $t_{hit} \approx 8$ Gyr.

In general, we find that there is some small evolution in the inner ($r \lesssim 0.1 r_{500}$) level of initial ($t = 0$) entropy of the primary cluster when $t_{hit} \approx 4$ Gyr. This shows that long term stability in the initial profiles is not always achieved, even after the application of a friction term to the motion of the SPH particles. To assess the impact of this behavior on the final ($t = 10$ Gyr) entropy profile of the merged clusters, we performed some of our merging simulations with the friction term still active up to $t = t_{hit}$. These runs will be discussed in detail later, unless otherwise stated in the following, we will discuss merging simulations in which the friction term $\sigma_{damp}$ is set to zero when $t \geq 0$.

Finally, we also show in Figure 3 the time evolution of the different gaseous halo profiles when the SPH entropy equation (12) incorporates radiative cooling. Following (Z11) we adapt the bremsstrahlung cooling time approximation:

$$\tau_c \approx 28.7 \text{ Gyr} \left( \frac{S}{100 \text{ keV cm}^{-2}} \right)^{1/2} \left( \frac{n_e}{10^{-3} \text{ cm}^{-3}} \right)^{-2/3}.$$  \hspace{1cm} (35)

The test runs with radiative cooling are indicated as CR in the panels. As for the FD runs, the initial conditions are the same as for the GL halo, but here the damping term is absent in the momentum equation (30). These settings allow us to assess the impact of radiative cooling on the thermal evolution of an isolated halo initially in equilibrium. The cooling time profiles in Figures 2 and 3 show that the condition $\tau_c \gg \tau_{dyn}$ is always satisfied at all radii, thus suggesting that radiative processes are not very important dynamically (e.g., motions induced by cooling will be very subsonic).

The results indicate that in the halo inner regions ($r \lesssim 0.1 r_{500}$), radiative losses become significant on time scales $t_{rad} \gtrsim 5$ Gyr, in accordance with the range of cooling times $\tau_c(r)$ displayed by the CR halo in the bottom left panel. From Table 2 one can see that the condition $t_{rad} \gg \tau_{num}$ is not always satisfied, this in turn implies that for some merging runs with cooling the entropy profile at $t_{hit} = \tau_{num} (t = 0)$ will not satisfy the prescribed initial conditions.

To construct the initial setup for the merging simulations with cooling, we then proceed as follows. The simulations are performed up to $t_{hit} = \tau_{num}$ as in the adiabatic case, with the friction term present and, in particular, the cooling term $Q_R$ in equation (12) switched off. This guarantees...
that both adiabatic and radiative simulations will start at $t = 0$ with the same profiles. After this epoch the radiative merging runs are performed with the $Q_R$ term now present in Equation (12) and the damping term switched off.

Note that the previous discussion about the appropriate level of core entropy when $t_{\text{hit}}$ is large (say $\gtrsim 5$ Gyr), is not relevant here. This is because the time evolution of the CR halo profiles demonstrate that core heating due to numerical effects is subdominant with respect radiative losses.

3 RESULTS

In this Section, we present our main results from the simulations we performed. We first discuss results from adiabatic simulations and subsequently those obtained from the cooling runs. Our findings are qualitatively discussed in light of the impact on the final entropy profiles of the different merging processes we consider.

3.1 Adiabatic runs

For the nine merging simulations we show in Figure 4 the final gas density and temperature profiles of the resulting
merging clusters. The plots are depicted at $t = 10$ Gyr, an elapsed time since the start of the collision which should be sufficiently large to guarantee a relaxed state for all of the considered merging configurations.

The radial profiles are calculated for each radial bin by spherical averaging the extensible physical quantities. These are total electron number for the average electron density, total thermal energy divided by $3/2$ of the total particle number for the gas temperature, and total physical entropy (adding up the specific physical entropy per particle for all the particles, dividing by the total number of particles, and converting to the average entropy parameter $S$ as discussed as discussed at the start of Section 2.2.2) for the entropy parameter.

As in Z11, for each physical quantity we have subdivided the plots by showing in each panel of Figure 4 profiles extracted from merging runs with the same mass ratio but different impact parameters. This layout is common also to the other Figures, and allows a better comparison with previous findings.

The radial behavior of the final density profiles depicted in Figure 4 exhibit the common feature of a flattened density ($n_e \simeq 2 \times 10^{-3} \text{ cm}^{-3}$) at cluster radii $r \lesssim 300$ kpc. This flattening is in sharp contrast with the initial density profiles, which are constructed so as to reproduce that of cooling flow clusters. The initial profiles steadily increase toward the cluster centers and have much higher central densities ($n_e \simeq 2 \times 10^{-2} \text{ cm}^{-3}$).

Similarly, the final temperature profiles no longer show the initial inversion and steadily increase toward the cluster centers. There is a wide range of central temperature values, from $\simeq 5$ keV up to $\simeq 15$ keV, depending on the mass ratio $R$ and impact parameter $b$ of the merging simulation.

These findings strongly suggest that the initial cool-core cluster configurations do not survive the impact on the gas of the processes that occur during the collisions. This issue
is central to the paper and will be addressed later, when discussing the final entropy profiles.

It is instructive to compare the profiles of Figure 3 with the corresponding ones shown by Z11 (Figure 15 and 16 of his paper). There are strong similarities, but also interesting differences. In general, both density and temperature profiles have the same radial behavior as the corresponding profile of Z11. In particular, for a specific mass ratio $R$, the hierarchy of the profiles at any given radius $r$ is always reproduced. This is reassuring because it validates our setup procedure and the code we are using.

Nonetheless, when contrasted against Z11 values, the central temperatures are found smaller by a factor lying in the range $\simeq 20 - 30\%$. Differences in the final profiles of thermodynamic quantities between our simulations and those of Z11 can be attributed to a number of reasons. To be specific, the largest impact will be caused by differences in the setup of the initial cluster kinematic and physical parameters, and by the different numerical hydrodynamical schemes used to perform the simulations. The latter can be significant, and in order to pin down its impact it is necessary to reduce as much as possible the effects of differing initial conditions.

To this end, we use as reference the merger with $R = 1:1$ and $b = 0$. This merging configuration has the advantage of having a very short collision time ($t_{\text{hit}} \simeq 2$ Gyr), so that differences between our initial orbital settings and those of Z11 are minimized. In what follows, we will refer to this simulation in brief as $R1b$.

For a better comparison of our merging simulations with those of Z11 we perform a head-on merging simulation with mass ratio $R = 1:1$ and initial conditions constructed as follows: The DM halo of each of the two clusters has a mass of $M_{200} = 6 \times 10^{14} M_{\odot}$, equal to that of cluster C1 in Table II but we set the cluster radius to $r_{200} \simeq 1.55$ Mpc. This value is that reported in Table I of Z11 for his cluster C1, and its a bit smaller ($\simeq 10\%$) than our corresponding value ($r_{200} \simeq 1.76$ Mpc). Note that the concentration parameter is the same for the two clusters ($c = 4.5$). It must be stressed that this small difference in the cluster radii has a significant impact when comparing final results, such as entropy profiles. This is because small differences in the reference radius $r_{500}$ or $r_{200}$ (Equation 24 or Equation 1 of Z11) induce differences in the initial gas entropy at the same physical radius, which in turn imply much larger differences in the final entropy profiles.

To construct our DM density profile we truncate the central region $r_{\text{dec}} = 0.1 r_{200}$ (as in Z11). These radii are smaller than those adopted in our initial conditions: $(\xi, \eta) = (0.1, 0.0)$.

For the core radii $R_{\text{core}}$, we adopt $r_{\text{core}} = r_{500} - 0.25 r_{200}$, as suggested by Z11.

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(2.0, 2.0), but for this merging configuration the collision time is very short and the considerations of Section 2.2.3 about DM stability can be considered secondary. Moreover, as in Z11 we initialize our cluster center of mass positions with a relative initial separation of \( d_{\text{rot}} = (r_{200}^1 + r_{200}^2) \).

Finally, our initial gas profiles are constructed according to the procedures described in Section 2.2.2, but truncating the profiles at \( r = r_{200} \) and using Equation 1 of Z11 with the same parameters \( S_0 \) and \( S_1 \) to specify the initial entropy profile \( S(r) \). In the following, we refer to this simulation as ZuH and we will use it as our reference run against which to contrast our simulation results with those of Z11. The initial physical settings and kinematics of the ZuH simulation are now identical to those of simulation S1 in Z11, so that differences between the final thermodynamic profiles of the two runs can be entirely attributed to the different numerical schemes used to perform the simulations.

In the left panels of Figure 4 the solid blue lines indicate the density and temperature profiles of this simulation extracted at \( t = 10 \) Gyr. The difference between these profiles and the corresponding ones of the R180 run (solid red lines) can then be interpreted as originating from the different settings in the initial conditions between the two simulations.

A visual inspection shows that the difference in the ZuH density profile and its Z11 counterpart S1 is minimal. Both of the profiles have the same central density \( (n_e \approx 10^{-2} \text{ cm}^{-3} \) at \( r = 10 \) kpc) and a knee at the same radius \( r \approx 300 \) kpc. Similarly, the temperature profiles are also in accord. The left bottom panel of Figure 4 shows a central temperature of \( \approx 13 \) keV for the ZuH run, whereas in Z11 the central temperature of the S1 run is \( T \approx 14 \) keV. These agreements strongly suggest the validity of the hydrodynamic code used here to carry out the simulations. We postpone further discussion of this topic to later when we address the radial behavior of the final entropy profiles.

The final profiles of the R180 run can also be contrasted with the corresponding ZuH profiles in order to assess the impact of different initial conditions and collision parameters on the final merged cluster. In particular, Figure 4 shows that the ZuH temperature profile is in good accord with the profile of its parent simulation, whilst it can be seen that in the inner cluster region \( r \lesssim 300 \) kpc the density profile \( n_e(r) \) is higher than that of R180 by about a factor \( \sim 2 \) and has a steeper decline with radius at \( r \gtrsim 500 \) kpc.

This difference in the final density radial behavior is a consequence of two distinct effects. In the ZuH simulation, the initial entropy profile is the same as that of Z11 and from Figure 15 (left panel) of Z11 it is easily seen that this leads to a much steeper initial density profile than that of the R180 run. This initial difference is not destroyed during the merging phases and still has an impact on the density profiles at \( t = 10 \) Gyr. On the other hand, at large radii the initial ZuH density profile is truncated at \( r_{\text{max}} = r_{200} \), half the value of the R180 simulation. As already discussed in Section 2.2.2, this implies a significant leakage of gas particles in the cluster outer regions during the merger. Thus, the final gas density at large cluster radii will be smaller than in the R180 run.

These findings demonstrate that final differences between the gas cluster profiles of our simulations and those of Z11 can be entirely interpreted in terms of the adopted initial entropy profile. This will be confirmed later when studying the radial behavior of the entropy profile.

However, the approach used here to initialize cluster dark matter particle orbits differs in several ways from that of Z11. As can be seen in Figures 5 and 6, this has an impact on the final DM velocity dispersion \( \sigma_{\text{DM}} \) in several runs. For ease of comparison with the gas temperature \( T_{\text{gas}} \) and previous findings (Z11), we introduce the DM temperature \( T_{\text{DM}} \):

\[
 k_B T_{\text{DM}} = \mu m_p \sigma_{\text{DM}}^2/3. \tag{36}
\]

Final profiles of the DM to gas temperature ratio \( \kappa(r) = T_{\text{DM}}/T_{\text{gas}} \) are shown at \( t = 10 \) Gyr in Figure 6 for the different merging simulations. In accord with Z11 (Figure 19), the ratio \( \kappa(r) \) is of order unity \( (\approx 0.9) \) at all cluster scales. The only exception is in the innermost cluster regions \( (\lesssim 300 \) kpc) where the \( \kappa(r) \)'s tend to zero. This is expected, since baryons in the core will raise their entropy through mixing processes with post-shocked high-entropy material.

However, Figure 6 shows that there is a significant difference between the initial ratio \( \kappa(r) \) of the head-on merger with \( R = 1 : 1 \) and the others. In fact, for the R180 merger run the initial \( \kappa \) is systematically higher by \( \approx 20\% \) compared to the other simulations. This is in sharp contrast with the corresponding profile in Figure 19 of Z11, which does not exhibit such a feature and whose behavior is in line with the others simulations.

We argue that this difference can be interpreted as originating from the adopted initial conditions. At variance with Z11, we initially set the center of mass of our clusters separated by a distance \( d_{12} = 2(r_{200}^1 + r_{200}^2) = d_{\text{vir}} \). The merging simulation time \( t = 0 \) is then defined when \( d_{12} = d_{\text{vir}}/2 \). This procedure then implies that at \( t = 0 \), the two clusters have already had time to interact. For the gas component, the impact on the initial profiles of this interaction is negligible (Figure 4), but for the DM halos one expects some amount of heating and an increase in the DM velocity dispersion. The strength of this effect will be weaker as the mass ratio \( R \) gets higher.

Figure 6 shows the radial profiles \( T_{\text{DM}}(r) \), corresponding to the ratios depicted in Figure 5. The left panel \( (R = 1 : 1) \) shows that the initial profile \( T_{\text{DM}}(r) \) (solid black line) is a bit higher \((\gtrsim 20\%)\) than the initial profiles displayed in the other two panels \( (R = 1 : 3 \) and \( R = 1 : 10) \), thus confirming the previous reasoning. In fact, this effect is significant only when \( R = 1 : 1 \).

To demonstrate the correctness of this interpretation, we ran an additional merger simulations. As with the R01600 run, we study a head-on merger with both cluster masses being \( M_{200} = 6 \times 10^{14} M_\odot \). At variance with the initial condition setup described in Section 2.2.1, here we truncate the DM halos at a cut-off radius \( r_{\text{max}} = 1.4 r_{200} \). We then perform the simulation and study the final \( \kappa(r) \) and \( T_{\text{DM}}(r) \) profiles. If this heating effect depends on the cut-off radius \( r_{\text{max}} \), then at any given radius the final profiles of this simulation should approach the profiles of simulations with lower mass ratios. The profiles are shown (solid magenta lines) in Figures 7 and 8 and confirm these expectations.

Finally, a comparison with Figure 18 of Z11 shows that the final profile \( T_{\text{DM}}(r) \) of the ZuH run (solid blue line, Figure 6) is in accord with the corresponding profile of simulation S1. Moreover, at \( r \approx 200 \) kpc the ZuH profile has a
peak value ($\simeq 7.2$ keV) that is about 15% higher than the peak of the $R01b00$ run at the same location. This offset between the two runs in the peak of the final DM velocity dispersion is interpreted as originating from the differences in the adopted initial conditions. In particular, for simulation $R01b00$ the two clusters have initial radii of $r_{200} \simeq 1.76$ Mpc, whilst initially $r_{200} \simeq 1.55$ Mpc for the ZuH simulation.

We show in Figure 7 the mean radial ($V_r$) and circular ($V_c$) gas velocities profiles at $t = 10$ Gyr. The latter is defined at the cluster radius $r$ as $V_c(r) = \sqrt{\bar{v}_\phi^2(r) + \bar{v}_\theta^2(r)}$, where $\bar{v}_\phi$ and $\bar{v}_\theta$ are the mean azimuthal and polar velocities, respectively. The profiles of Figure 7 can be contrasted with the corresponding profiles in Figures 21 and 22 of [Z11]. All of them exhibit a radial behavior which is in accord with their [Z11] counterparts, with the only exception being the head-on $R = 1:3$ merger (simulation S4 of [Z11]). The final velocity profiles of this merged cluster are significantly different from those of simulation S4; in particular, the mean radial velocity is not close to zero. Values of $V_r \simeq 50$ km/sec persist up to $r \simeq 800$ kpc. Similarly, the circular velocity $V_c$ is as high as $V_c \simeq 400$ km/sec within $r \lesssim 50$ kpc. These values suggest that for this merger a fully relaxed status has not yet been achieved at $t = 10$ Gyr. To verify this possibility we have continued the simulation until $t = 11$ Gyr. The velocity profiles corresponding to this epoch are shown in the middle panels of Figure 7 as solid black lines, and they clearly show lower velocities.

Figure 8 shows the final entropy profiles of the merged clusters. Note that the astrophysical entropy parameter $S \equiv kT/n\sigma_{DM}^2/3$ is not the physical entropy and is not an extensive, additive quantity. Thus, in averaging $S$ over spherical shells, $S$ was converted into the physical entropy (which is proportional to $\ln S$), and this was averaged over the spher-
Figure 7. Mean gas radial ($V_r$, top panels) and circular ($V_c$, bottom panels) cluster velocities profiles at $t = 10$ Gyr for mergers without gas cooling. The meaning of the symbols is the same as in Figure 4. The solid black lines in the middle panels refer to the profiles of the $R = 1:3$ and $b = 0$ run evaluated at $t = 11$ Gyr.

Figure 8. Final entropy profiles for the nine non-radiative merging runs at $t = 10$ Gyr (red lines). The panel presentation and line styles are identical to Figure 4, solid black lines represent the initial entropy profile of the primary.

The presence of an entropy core is common to all of the profiles, with its level and extent depending on the mass ratio and impact parameter of the simulation. A comparison with the corresponding Figure 24 of Z11 shows a substantial agreement in the radial behavior of the profiles, with differences in the central levels of core entropy which can be reconciled in light of the previous discussions.

In particular, at any specified radius and for a given mass ratio, the hierarchy of the entropy profiles as a function of the impact parameters is strictly reproduced. Following Z11, differences in the various levels of entropy profiles can be interpreted in terms of the different amounts of entropy mixing taking place during the mergers.

The core entropy of the primary increases during the merger owing to the mixing of low- with high-entropy gas. This high-entropy gas is made available by the secondary as it falls through the ICM of the primary and is ram-pressure stripped. The gas of the secondary is then efficiently mixed...
with that of the primary through the development of Kelvin–Helmholtz instabilities. This scenario has been confirmed by various authors in several merging simulations (Takizawa 2002, Mitchell et al. 2009, ZuHone 2011). Following this line of argument, the level of core heating of the primary should depend on the impact parameter $b$ of the simulation. The higher the impact parameter, the lower is the amount of mixing. This follows because in off-center collisions, the amount of gas stripped from the secondary depends on the ram pressure it encounters, and in turn on the initial mass ratio and angular momentum of the merger. For a given mass ratio the quantity of stripped material, which is available in the inner regions of the primary to raise core entropy through mixing, is then expected to depend sensitively on the orbit traced by the secondary.

In accord with this scenario, the third panel of Figure 8 shows an increase in the central level of final entropy as the impact parameter decreases. However, this behavior is clearly seen for the mergers with mass ratio 1:10 (third panel) but is progressively less pronounced as the mass ratio $R$ becomes higher. In fact, for the 1:1 mass ratio case the dependency of the final core entropy level on the impact parameter $b$ is reversed, i.e. the first panel of Figure 8 shows that simulation R01606 has an higher level of central entropy than R01500.

The likely origin for this difference with the results from the $R = 1:10$ merger is that in the equal-mass mergers, a significant amount of core heating is provided dynamically by the secondary during the final merging with the primary. This effect is almost completely absent in the $R = 1:10$ cases, in which the mass of the secondary is small with respect to that of the primary, and for off-axis mergers the secondary is totally stripped by instabilities before coalescing with the primary.

For a better understanding of this scenario, in Section 3.3 we present a thorough discussion of how entropy is generated during the merging process. For several merging runs, we investigate in detail the time evolution of entropy and other related quantities, in order to demonstrate how the final entropy profile of the merged clusters depends critically on the mass ratio and angular momentum of the collision.

### 3.2 Stability issues

As a convergence test, we compared the final entropy profiles for several merger runs with simulations in which we varied the numerical resolution and/or the adopted initial conditions.

For two equal-mass mergers ($b = 0$ and $b = 0.6$), in Figure 9 (left panel) we show the final entropy profiles together with those extracted from the corresponding higher-resolution runs (HR, black lines). These simulations were performed by adopting the same initial conditions as the baseline runs, but with the particle masses reduced by a factor $\sim 4$. The plots show a radial behavior of HR profiles which is in excellent agreement with the corresponding standard resolution profiles, a result which leads us to conclude that the simulations presented here are numerically converged.

Similarly, the entropy profile (blue line) of the ZuH run is contrasted with its parent simulation S1 (open circles, the points are taken from Figure 24 of [Z11]). There is a significant agreement between the two profiles, the only exception being the outermost point ($r \approx 1800$ kpc) for which the entropy of the ZuH simulation is higher than that of S1. This result is interpreted in light of the steepening of the ZuH density profile at large radii (Figure 1). As already outlined, this outer behavior follows from the adopted initial conditions and the lack of an external buffer surrounding the SPH particles.

This strict agreement between the entropy profiles of two independent simulations is very significant and it has a number of implications. Firstly, for a given merging configuration, it definitively shows that the only parameter which determines the thermodynamic structure of the final merged cluster is the initial entropy profile. The other direct consequence is that the numerical scheme used here produces, for the same initial conditions, a final entropy profile which is identical to that obtained by Z11 using the adaptive mesh refinement (AMR) code FLASH. This is a non trivial issue, and consistency between hydrodynamical test cases performed using Lagrangian SPH schemes and mesh-based codes has been the subject of many investigations.

Specifically, Agertz et al. (2007) found that the standard formulation of SPH (SSPH, Price 2012) fails to reproduce the results of several hydrodynamic test cases, when contrasted against those obtained from Eulerian mesh based codes. In particular, non-radiative SSPH simulation of galaxy clusters exhibit entropy profiles with a power-law behavior. This is in sharp contrast with the constant entropy cores produced in Eulerian mesh simulations (Mitchell et al. 2004). These discrepancies are due, in part, to the intrinsic difficulty SSPH has in modeling density gradients around contact discontinuities, which in turn implies that there is a surface tension effect that inhibits the growth of fluid instabilities (Agertz et al. 2007).

To address these problems, several solutions have been proposed (Hopkins 2015, and references cited therein). In particular a possible solution is to add a dissipative term to the SPH thermal equation, with the purpose of smoothing the thermal energy at fluid interfaces (Price 2008, Wadsley et al. 2008). The presence of this AC term has the effect of smoothing entropy transitions at contact discontinuities, thus enforcing pressure continuity and in turn removing the artificial surface tension effect that suppresses the growth of the instabilities at fluid interfaces.

The SPH scheme employed here is based on this AC formulation, but the adopted signal velocity (equation 14) is different from that originally proposed by Price (2008) and it is better suited when gravity is present (Wadsley et al. 2008, Valdarnini 2012). Other formulations of SPH aimed at solving these issues are the SPH scheme proposed by Read & Hayfield (2012, SPHIS), which is based on the use of a high order dissipation switch, and the density-independent scheme of Saitoh & Makino (2016, DISPH). To validate these numerical schemes, it is important to assess the degree of consistency between the level and radial extent of the core entropies produced by these codes in cluster simulations. To this end, radial entropy profiles extracted from galaxy cluster simulations can be contrasted with the corresponding ones obtained from their AMR counterparts.

On this issue, the results reported in the literature show the absence of a general agreement between the various final entropy profiles. From DISPH simulations of galaxy clus-
Figure 9. For several merging runs, final entropy profiles from Figure 8 (red lines) are contrasted against the corresponding profiles extracted from merging simulations with different numerical parameters. Left panel: black lines labeled HR show the entropy profiles of high-resolution simulations, performed by using a number of particles about four times higher than in the baseline runs. The solid blue line is the entropy profile extracted from a merging run with $R = 1:1$ and $b = 0$, but with its initial entropy profile identical to that in Z11. Open circles are taken from the entropy profile of the corresponding simulation S1, as given in Figure 24 of Z11. Right panel: blue lines ($\alpha_d$) refer to merging runs in which the friction term is switched off at a simulation time $t > 0$ which depends on the impact parameter $b$. For the merging simulation $R = 1:10$ and $b = 0.6$, the black dots ($\alpha_d + HR$) are from a high-resolution run that also had the damping term switched off.

Sembolini et al. (2016) carried out a systematic comparison between the final entropy profiles extracted from a suite of simulations of an individual cluster. Their simulation set is constructed by using different codes. Their results showed that a flat inner entropy profile, such as that obtained using the AMR code RAMSES (Teyssier 2002), is similarly formed in cluster simulations produced using SPH variants which are based on some form of artificial dissipation. In particular, both the improved SPH code of Beck et al. (2016) and the SPHS scheme (Read & Hayfield 2012) give entropy profiles in accord with mesh-based results. The AC implementation of the former SPH scheme is very similar to the one employed here, thus reinforcing the consistency between our ZuH entropy profile and that of the corresponding S1 run of Z11.

It must be stressed that the core entropy level and size of the core are mainly regulated by the maximum value $\alpha_{C,MAX}$ of the AC particle parameter $\alpha_C$. For the simulations presented, here we set $\alpha_{C,MAX} = 1.5$, this upper limit being derived from the consistency of self-gravity tests with mesh results (Valdarnini 2012). This limiting value is also in accord with the DISPH runs of Saitoh & Makino (2016), who concluded that a core entropy is established when $\alpha_{C,MAX} \gtrsim 1$.

Power et al. (2014) criticized the AC formulation of SPH; based on the results from Wadsley et al. (2008), they suggest that the AC scheme may not always achieve numerical convergence. However, the HR profiles of Figure 9 are fully converged and do not support this view. We argue that it is the adopted method to estimate gradients using a matrix inversion that is more relevant in this context. As demonstrated in V16, our scheme is seen to exhibit excellent convergence properties.

Finally, it must be emphasized the strict agreement between the final entropy profile of our ZuH test run with the corresponding S1 profile of Z11 does not imply that the produced core entropy levels are correct. It just demonstrates that the two codes consistently obtain the same results, when adopting the same initial conditions. It remains unclear which is the correct core entropy level in these sort of simulations, with SSPH lacking of any mixing process and Eulerian codes having the tendency to overestimate mixing effects because of numerical diffusion (Springel 2010).

In the right panel of Figure 9 we analyze the consistency of our setup procedure for several merger configurations. Specifically, the friction parameter $\alpha_{d,\text{damp}}$ introduced in Section 2.2.2 is switched off when the two cluster orbits have reached the initial conditions of Z11 ($t = 0$). This friction term is introduced to maintain a stable realization of...
the CC entropy profile before the occurrence of the cluster collision. However, in merging simulations with large angular momentum, the time interval between $t = 0$ and the direct collision can be large ($\gtrsim 5$ Gyr). For these mergers, a certain amount of numerical heating can modify the core of the original entropy profile before the collision (see Figure 2), thereby increasing the final level of core entropy.

In order to assess the impact of this effect on the final entropy profiles, we ran three additional merger simulations. Among the simulation suite, we have chosen three merging configurations with the criterion of having the highest angular momentum. These are the two merging simulations with mass ratio $R = 1 : 10$ and impact parameters $b = 0.6$ and $b = 0.3$, and the simulation with $R = 1 : 3$ and $b = 0.6$. For these simulations, the damping parameter $\alpha_{damp}$ is switched off at a simulation time $t_{hit} > 0$ and not at $t = 0$. This procedure guarantees that the initial entropy profile of the primary cluster maintains its form for a certain period of time after $t = 0$. Henceforth, we will generically refer to these simulations as $\alpha_{d}$.

The choice of the time $t_{hit}$ is a compromise between the need to avoid the possible numerical heating of the primary cluster core, and at the same time, to not damp significantly the primary core’s heating due to entropy mixing driven by dynamical interactions with the secondary. As outlined in Section 2.3, we define $t_{hit}$ approximately as the epoch when, after the first pericenter passage, the distance between the two clusters center of mass becomes smaller than $\simeq r_{200}$.

Our estimates give $t_{hit} \simeq 9$ Gyr for the merging simulation R10606, and $t_{hit} \simeq 5$ Gyr in the case of the R10603 and R03606 merging runs. Clearly, a higher value of $t_{hit}$ leads to a longer time required for the damping term to keep the entropy profile stable.

The first value of $t_{hit}$ corresponds to the maximum required period of damping, whereas in the case of off-axis mergers $t_{hit} \simeq 5$ Gyr constitutes an approximate lower limit to $t_{hit}$. This choice of different off-axis merger cases allows us to assess the impact of numerical heating on the entropy profile of the final merger remnant.

By comparing the final entropy profiles of these simulations ($\alpha_{d}$, blue lines) in the right panel of Figure 9 with the corresponding ones in Figure 3 we see that a certain amount of numerical heating is always present. All of the $\alpha_{d}$ profiles have core entropy levels systematically smaller than their standard counterpart. For instance, in the case of R03606 the final central entropy of the $\alpha_{d}$ run is smaller by about $\simeq 30\%$. For the other two merging simulations the difference is even higher, being almost a factor of two in the case of R10603. Note that the core entropy levels of the $\alpha_{d}$ runs are now in better agreement with the corresponding ones displayed in Figure 24 of [11].

Moreover, to demonstrate that the final profiles are numerically converged, for the simulation R10606, we also run a high resolution simulation ($\alpha_{d}$+HR, black dots). In Figure 8, it can be seen that the entropy profiles of the two simulations are almost coincident, thus confirming that our simulations are not affected by insufficient resolution.

These results demonstrate that our setup procedures as described in Section 2.3 are not entirely free of relaxation effects, with some amount of numerical heating being present in the final entropy profiles of the merged clusters. The result of the $\alpha_{d}$ run with $R = 1 : 3$ and $b = 0.6$ suggests that this effect leads to an overestimate of the final core entropy level by $\simeq 30\%$. For the merging run R10606 the increase is similar ($\simeq 50\%$), whereas for the merging simulation R10603 the difference in the central entropy values is larger, $\sim 100$ keV cm$^{2}$.

Thus, we conclude that with the setup procedure adopted here, in merger simulations with high angular momentum and a 1 :10 mass ratio there is the tendency to overestimate final entropy in cluster cores by about $\simeq 30\%$. This relaxation effect can be compensated for by switching off the damping term at a later time $t_{\alpha_{d}} > 0$; however, the correct implementation of this correction requires a careful choice of $t_{\alpha_{d}}$ to avoid over-correction leading to artificially lower entropy levels.

### 3.3 How is central entropy generated during mergers?

In this Section, we examine the time evolution of entropy and other related quantities for several merger simulations. The analysis is aimed at investigating the origin and amount of final central entropy as a function of the initial merging parameters.

To this end, we construct radial profiles of gas density and entropy at given times. The radial profile of a given quantity is calculated for each radial bin by averaging grid values computed on spherical shells, defined by a set of $(\theta, \phi) = 40 \times 40$ grid points uniformly spaced in $\cos \theta$ and $\phi$. Unlike in Section 3.1, the origin of the shells is centered at the gas density peak of the primary, tracked at run time. For equal-mass mergers we adopt the convention of defining as primary the cluster on the left in the $(x, y)$ plane of the collision.

We first discuss the head-on, 1 : 1 mass ratio case. This is the most energetic event and its dynamics are relatively simple. In the left panel of Figure 10 we show the time evolution of the gas density radial profile of the primary at different epochs. Similarly, the evolution of the entropy radial profile is shown in the right panel. Additionally, we also show in the left panel the density profile of the secondary; this quantity has been evaluated in the frame of the primary.

The gas profiles of each cluster member are constructed separately by culling from the set of SPH particles the corresponding gas particles. These are tagged at the start of the simulation according to their membership.

We have chosen to evaluate the profiles at five different times: $t = 1.25$, 1.50, 1.54, 1.70, and 2 Gyr. These are centered around $t_{\rho} \sim 1.6$ Gyr, the epoch at which the distance between the centers-of-mass of the two colliding cores attains its first minimum. In other mergers with non-zero impact parameters, this epoch is identified as the first pericenter passage and we denote it as $t_{p}$.

From the time evolution of the profiles of Figure 10 we conclude that most of the final core entropy is generated at the time $t_{\rho}$ of the first core collision. From the right panel it can be seen that the level of central entropy is already at $S(0) \sim 300$ keV cm$^{2}$ at $t \simeq 2$ Gyr. This entropy level is about $\sim 70\%$ of its final value, as can be inferred from the right panel of Figure 10.

This scenario is also confirmed by the radial behavior of the temperature profiles displayed in Figure 11. The two panels show the temperature profiles of both primary and
Figure 10. For the head-on 1:1 mass ratio case, we show radial profiles of gas density (entropy) in the left (right) panels. Different line styles refer to different epochs, and the time is in Gyr. In the left panel, red lines indicate the gas density profiles of the primary (I), whilst blue and green lines those of the secondary (II). At each epoch the profiles of both clusters are evaluated in a frame centered on the peak of the gas density of the primary. For the secondary blue (green) lines are used to indicate profiles extracted at times before (after) $t \approx 1.6$ Gyr, an epoch which is approximately identified as that when the distance between the two cluster centers-of-mass reaches its first minimum. Right panels shows the entropy profiles of the primary at different epochs. The solid black line refers to the initial profile.

Figure 11. For the same merger as in Figure 10, we show here the evolution of the gas temperature profiles, red lines are for the primary and green lines refer to the secondary. For the sake of clarity, in the left panels we show profiles evaluated at $t < 1.6$ Gyr, when the two clusters are approaching each other, and in the right panels, the profiles are evaluated at $t > 1.6$ Gyr, after the first pericenter passage. As in Figure 10, at each epoch radial profiles are evaluated in the primary frame, with its origin being defined as the gas density peak of the primary tracked at run time.
secondary at different times. As in Figure 10, the secondary profiles are evaluated in the primary frame. To avoid overcrowding, we have divided the profiles in two categories: before (left panels) and after (right panels) the epoch \( t_p = 1.6 \) Gyr.

The left panel of Figure 11 shows how the temperature profile of the primary increases progressively, as the secondary approaches the primary and the gas is shock heated. This increase is first characterized by a peak in the primary outskirts, which moves inward and gets wider at late times \((t \rightarrow t_p)\). The temperature of the secondary increases too, in fact at \( t = 1.54 \) Gyr both the two core temperatures approach the same level. These findings consistently support the view that for head-on mergers most of the final core entropy is generated during the first core collision, with the remainder of the high-entropy gas being accreted later during the final phases of the merging.

A different scenario of entropy generation emerges when analyzing mergers with an initial angular momentum (AM). We have chosen to discuss first the case of the off-axis merger \( \alpha \), with a mass ratio and \( \alpha = 0.6 \). Among the equal mass mergers this has the highest AM, so its study is particularly interesting in order to analyze how entropy is generated during these collisions.

The two panels of Figure 12 depict density and entropy profiles as done in Figure 10, but for the R01506 merging run. In particular, the first pericenter passage occurs at \( t_p \sim 2 \) Gyr, and we show profiles extracted at five different times centered around this epoch.

From the entropy profiles displayed in the right panel of Figure 12 one can recognize that, unlike the head-on case, most of the core entropy of the primary is generated well after the pericenter passage, between \( t \sim 3 \) and \( t \sim 5 \) Gyr. This last epoch corresponds to when the secondary has passed the apocenter and is falling back onto the primary. Note that in the left panel of Figure 12 the density profile of the secondary at \( t = 5 \) Gyr extends closer to the core of the primary than at \( t = 3 \) Gyr.

At \( t \sim 3 \) Gyr, core heating of the primary proceeds as previously discussed for the head-on merger, with the secondary coalescing with the primary and the low-entropy gas in the core being mixed with the high-entropy gas generated during the core collision. This suggests that we can decompose the generation of entropy in the central regions of the primary into two distinct phases: a first one when the primary has a grazing encounter with the secondary, and a second phase when the secondary finally collapses onto the primary. We now investigate how entropy generation proceeds during the first phase, when the secondary first approaches the primary.

To this end, we first show in the two panels of Figure 13 the time evolution of the primary and secondary temperature profiles. Their time variations exhibit a behavior in line with that seen with the corresponding profiles of Figure 10 with a significant increase in the gas temperature of the primary at late epochs.

However, the right panel of Figure 12 shows a jump in the entropy of the primary between \( t = 2 \) and \( t = 3 \) Gyr. We argue that this entropy increase is a consequence of a transfer of AM between the secondary and the primary as the secondary is getting closer. This in turn is due to tidal torques that become significant as the two clusters reach their closest approach along their orbits.

This transfer of AM leads to an increase of the gas circular velocity of the primary and, subsequently, to the development of instabilities and entropy mixing. To better quantify this point, we show in Figure 14 the radial profiles of the mean circular velocity \( \bar{V}_c(r) \) (left panel) and artificial viscosity parameter \( \alpha(r) \) (right panel) for the primary cluster. The latter is constructed by radial averaging the AM parameters \( \alpha \) in a manner similar to that adopted to calculate the other radial profiles. The mean velocity \( \bar{V}_c(r) \) is evaluated by subtracting the center-of-mass velocity of the primary from the gas velocities. The profiles have been extracted at the following epochs: \( t = 2, t = 2.5, t = 2.75, t = 3 \) and \( t = 5 \) Gyr.

The time evolution of the mean circular velocity profiles reveals several important features. In particular, there is a progressive increase in the amplitude of the profiles as \( t \rightarrow 3 \) Gyr. All of the profiles at different epochs have the tendency to reach their peak values around \( \sim 100 \)–\( 200 \) kpc, these are of the order of \( \sim 600 \) km s\(^{-1}\) at \( t \sim 3 \) Gyr.

This behavior is shared by the corresponding artificial viscosity profiles \( \alpha(r) \), which at epochs \( t \gtrsim 3 \) Gyr exhibit peak values of \( \sim 0.12 \)–\( 0.15 \) in the same radial range. This clearly shows evidence of significant gradients in the flow velocity at \( \sim 100 \)–\( 300 \) kpc, induced by the strong increase in the gas rotational motions.

These motions will generate local instabilities in the medium which in turn will lead to the development of turbulence, thus driving the diffusion of entropy. This is indirectly confirmed by the radial behavior of the AC profiles \( \alpha^C(r) \), which are constructed from the AC parameters \( \alpha^C \) in the same way as we did for the AV profiles \( \alpha(r) \). The qualitative behavior of the profiles \( \alpha^C(r) \) mirrors very closely that of the corresponding \( \alpha(r) \), thus showing the presence of diffusive processes associated with the appearance of rotational motions.

We argue that the increase in core entropy seen between \( t = 2 \) and \( t = 3 \) Gyr is then a consequence of the rotational gas motions, induced by the passage of the secondary at the pericenter. However, to validate this picture, the local diffusion time scale must be of the same order or lower than the estimated time span (\( \sim 1 \) Gyr) over which entropy undergoes its changes. To further elucidate this issue, we now try to assess the diffusion time scale associated with the development of Kelvin-Helmholtz instabilities (KHI).

The generation of KHI leads to the development of turbulent motions and to the formation of eddies at different spatial scales. The size of these eddies is expected to be of the order of \( \sim 100 \)–\( 300 \) kpc [Takahawa 2002; Subramanian et al. 2006]. These estimates are also in accord with length scales found in simulations aimed at studying turbulent properties of the ICM [Vazza et al. 2012; Valdarnini 2013, see in particular Figure 9 of the first authors].

From Eq. 8 of Vazza et al. (2012), we estimate the coefficient for turbulent diffusion as given by

\[
D_{\text{turb}} \approx 0.1 \lambda \sigma \approx 3 \times 10^{30} \left( \frac{\lambda}{100 \text{kpc}} \right) \left( \frac{\sigma}{1000 \text{ km/sec}} \right) \text{cm}^2 \text{s}^{-1},
\]

\( \lambda \) and \( \sigma \) represent the core scale and the velocity dispersion of the ICM, respectively.
and the corresponding diffusion time scale by

$$\tau_D \simeq R^2 / D_{\text{turb}} \sim 3 \left( \frac{R}{100 \text{ kpc}} \right)^2 \frac{1}{D_{30}} \text{Gyr}, \quad (38)$$

where $\sigma$ is the gas velocity dispersion, $D_{30} \equiv D / (10^{30} \text{ cm}^2 \text{s})$, and $R$ is the considered scale.

Between $t = 2$ Gyr and $t = 3$ Gyr the gas velocity dispersion of the primary is found to change drastically from $\sim 800$ km/sec down to $\sim 200$ km/sec within radii $r \leq 100$ kpc. By assuming an upper limit of $\lambda \sim 300$ kpc for KHI eddies, in the given time interval we then obtain values of $D_{30}$ lying in the range $7 \lesssim D_{30} \lesssim 1.8$. This correspondingly gives $0.4 \lesssim \tau_D / (R/100\text{ kpc})^2 \text{Gyr} \lesssim 1$ between $t = 2$ Gyr and $t = 3$ Gyr, showing the consistency of the diffusion time scale with the increase in central entropy seen between $t = 2$ and $t = 3$ Gyr at $r \leq 100$ kpc.

Figure 12. Identical to Figure 11 but for the $R = 1 : 1$ off-center merger with $b = 0.6$. The first pericenter passage occurs approximately at $t_p \approx 2$ Gyr.

Figure 13. As in Figure 11 for the merging case of Figure 12 ($R = 1 : 1, b = 0.6$), we show here the time evolution of the gas temperature profiles for both the primary and secondary cluster.
To further elucidate this critical point, we show in Figure 14 the entropy profiles of the primary at different epochs. Unlike in the right panel of Figure 12 we show the radial entropy profiles extracted from a wider range of time frames; moreover, for each radial bin we also evaluate the entropy dispersion $\sigma_S(r)$. As noted in the text discussing Figure 8 above, the average $S$ over a spherical shell is determined by converting the entropy parameter $S$ into a physical entropy, averaging this, and then converting the mean physical entropy back to $S(r)$. Similarly, the dispersion of the physical entropy is added and subtracted from the mean physical entropy, and these two values are converted back to $S$ to give the upper and lower limits of the error bars in Figure 15.

Within each panel of Figure 15 we show the entropy profiles at two distinct epochs; for the sake of better understanding the profile referring to the latest epoch is shifted upward by one order of magnitude. The time evolution of these entropy profiles clearly illustrates a progressive increase in entropy at outer radii, with a subsequent propagation toward the inner regions. If we now adopt the reasonable assumption that the entropy dispersion $\sigma_S(r)$ can be taken as a metric for assessing the amount of mixing present at that radius, $\sigma_S(r)$ (and presumably mixing) are initially very small at radii $r \lesssim 500$ kpc. The entropy dispersion becomes progressively wider in the outer regions of the primary as the secondary approaches ($t \rightarrow t_p \approx 2$ Gyr), but with an amplitude which is still negligible within $r \lesssim 100$ kpc at $t = 2$ Gyr. At $t = 3$ Gyr, there is a widening of $\sigma_S(r)$, which is significant at radii $r \gtrsim 100$ kpc. However, it is only when $t \gtrsim 5$ Gyr that the entropy dispersion is approximately constant across all the cluster and the gas has now a higher degree of mixing. To summarize, we conclude that for the considered merging configuration, part of the final core entropy owes its origin to rotational motions induced by tidal torques occurring during the collision.

We now discuss the two off-axis mergers for the 1 : 10 mass ratio case. Figure 16 is the analogue of Figure 10 but for the $R = 1$ : 10 and $b = 0.3$ ($t_p \sim 1.6$ Gyr) merging, and Figure 17 for $b = 0.6$ ($t_p \sim 2$ Gyr). In both cases the level of central entropy at late times is around $\sim 100$ keV cm$^2$, while in the left panel of the Figures the time evolution of the secondary density profiles shows a negligible interaction with the primary’s core. In fact, for the $b = 0.6$ merging the impact of the secondary on the density profile of the primary at $t = 4$ Gyr can be considered completely absent.

As already discussed in Section 3.2 in merging simulations with 1 : 10 mass ratio and AM, numerical heating effects can modify the level of core entropy before the two clusters merge together. For instance, one can assume that for the $b = 0.6$ run, all of the core entropy increase at $t = 4$ Gyr is due to this effect. This level can be contrasted with that expected in an isolated halo when the damping term is absent. From the GL run in the right panel of Figure 4 we obtain an entropy value of $S/S_{\text{iso}} \approx 0.15$ at $r/r_{500} = 0.1$ at $t = 5$ Gyr. This translates into $S \sim 130$ at $\sim 115$ kpc, where we have taken $S_{500}$ and $r_{500}$ from Table 1. This implies that at the same epoch, for the $b = 0.3$ merging, at most $\sim 30\%$ of the entropy core level is due to merging effects.

To illustrate the impact of this effect, in Figure 18 we depict density and entropy profiles for the $R = 1$ : 10 and

\[\text{Figure 14. Average radial profiles of the mean circular velocity } V_c(r) \text{ (left panel) and artificial viscosity parameter } \alpha(r) \text{ (right panel) are shown at different epochs for the primary cluster of the merging run of Figure 12. The center-of-mass velocity of the primary has been subtracted from the mean gas velocities before evaluating the velocity profiles. The } \alpha(r) \text{ profiles have been extracted from the AV parameters } \alpha_\alpha \text{ of the gas particles of the primary cluster, in the same way in which the radial profiles of other gas properties (e.g., temperature) have been derived. Time is in Gyr.}\]
Figure 15. For the equal-mass merger with $b = 0.6$ of Figure 8, we show here the entropy profiles of the primary at different epochs. At a given time, we compute the entropy dispersion (the error bars) by averaging over angular grid values of the shell for each radial bin (see text). As in Figure 8, the solid black lines refer to the initial profile. For the sake of clarity, in each panel the entropy profile referring to the later epoch has been shifted upward by a factor 10.

$b = 0.3$ merger at various epochs, but extracted from the corresponding $\alpha_d$ simulation presented in Section 3.2. For this simulation, the damping parameter $\alpha_{damp}$ was switched off at a simulation time $t_{hit} = 5$ Gyr in order to stably maintain the initial entropy profile until $t = t_{hit}$. The density profiles are time-centered around $t = 7$ Gyr, an epoch at which the secondary begins to coalesce with the primary.

The final entropy profiles depicted in the right panel of Figure 15 show a core entropy level of $\sim 100$ keV cm$^2$ at $t = 10$ Gyr. This level is in accord with previous findings and demonstrates that in mergers with low mass ratios and AM, the bulk of core heating occurs at late stages. A similar level of core entropy at $t = 10$ Gyr is obtained from the $\alpha_d$ simulation with $b = 0.6$ (Figure 9). To summarize, our findings indicate that heating of the core in off-axis mergers depends critically on the initial merging mass ratio as well as on the AM of the system. For equal-mass mergers and high AM, a significant contribution to the central entropy level is sourced by instabilities generated by tidal torques, as the secondary first reaches its pericenter. On the contrary, for unequal-mass mergers with $1:10$ mass ratio, the secondary is progressively disrupted along its orbit by ram pressure and the development
Figure 16. The same as Figure 10 but for the $R = 1 : 10$ off-center merger with $b = 0.3$. The first pericenter passage is estimated to occur at $t_p \approx 1.6$ Gyr.

Figure 17. The same as in Figure 10 but for the off-center case with $b = 0.6$. Here the first pericenter passage is around $t_p \approx 2$ Gyr.

of hydrodynamical instabilities, and core heating becomes significant only during the late merging phases.

3.4 Radiative runs

We now investigate the heating of gas cores and the survival of CCs in a more realistic set of merging simulations. In this section, we present results extracted from simulations where the physics of the gas includes cooling, star formation, and energy feedback following supernova explosions.

The computational cost of these simulations is much higher than that of their adiabatic counterparts. Because of cooling, the development of short cooling times and large central densities during the simulations requires very small timesteps. For this reason, we refrain from resimulating all of the merging cases previously discussed and perform radia-
Figure 18. The same as for the merging case shown in Figure 16 but for the \( \alpha \) simulation of Section 3.2 (see text), in which the parameter \( \alpha_{damp} \) is switched off at a simulation time \( t_{hit} = 5 \) Gyr. The first pericenter passage is between \( t = 1 \) and \( t = 2 \) Gyr, and after it the secondary falls back onto the primary shortly after \( t = 7 \) Gyr.

Figure 19. The final entropy profiles of several equal-mass radiative merger simulations are displayed in the left panel of Figure 19. We first consider the two merging cases with impact parameter \( b = 0 \) and \( b = 6 \), previously investigated in Section 3.3. For the sake of comparison their adiabatic counterparts are also shown (blue lines).

In the head-on case we expect the core entropy to undergo a very steep increase because of the strong shock following the collision of the cores, with a subsequent decrease due to radiative cooling. The difference at \( t = 10 \) Gyr between the core entropy of the adiabatic simulation and the radiative one (solid lines, left panel of Figure 19) is \( \Delta S(0) \approx 200 \) keV cm\(^2\).

The off-axis merger (dot lines) exhibits a similar difference in final core entropies. In Figure 20 we show the time evolution of density and entropy profiles for this run, as we did in Figure 12 for the adiabatic case. A comparison between the two sets of entropy profiles shows a much more modest increase in entropy for the cooling run.

In fact, at \( t = 5 \) Gyr the level of core entropy is about \( S_{cr}(0) \approx 100 \) keV cm\(^2\), while in the adiabatic run it is about...
Figure 19. Final entropy profiles extracted from adiabatic (AD: blue lines) and cooling simulations (CR: red lines). The left panel shows some runs with the mass ratio $R = 1:1$ and the right panels refers to merging simulations with mass ratio $R = 1:10$. In the left panel, thick lines indicate an head-on merging with cluster masses $M_1 = M_2 = 6 \times 10^{13} M_\odot$.

Figure 20. As in Figure 12, we show here the gas density and entropy radial profiles for the $R = 1:1$ off-center merger with $b = 0.6$. At variance with the run of Figure 12, the simulation here incorporates radiative cooling.
$S_{\text{ad}}(0) \simeq 800 \text{ keV cm}^2$. The difference is subsequently reduced soon after $t = 5 \text{ Gyr}$, as the secondary finally merges with the primary and the process raises the central entropy up to $S_{\text{ad}}(0) \simeq 500 \text{ keV cm}^2$ at $t = 10 \text{ Gyr}$.

This is in accord with Poole et al. (2008). From their merging simulations the authors argue that to re-establish a CC it is necessary for the remnant to be relaxed for a significant fraction of the cooling time $\tau$. For the considered merger, we estimate $n_s(0) \sim 10^{-2} \text{ cm}^{-3}$ at $t = 5 \text{ Gyr}$ and $\tau_s(0) \sim 6 \left( S(0)/100 \text{ keV cm}^2 \right)^{1/2} \text{ Gyr} \sim 6 \text{ Gyr}$, from Equation (4). This time scale is larger than that set by the final collision of the secondary, which raises again the central entropy and in turn $\tau_s$ to above $\sim 10 \text{ Gyr}$.

Finally, note that in the left panel of Figure 20 the density profile of the secondary is significantly reduced in the cluster inner regions, when contrasted with the adiabatic case at $t = 5 \text{ Gyr}$. Similarly, in the same regions the core gas density of the primary is higher by a factor $\sim 10$. These differences follow because of radiative cooling, with the primary developing a steeper profile than the secondary.

In the right panel of Figure 19 we show the entropy profiles of two unequal mass mergers with the mass ratio $R = 1 : 10$. The dynamics of the head-on case mirrors that of the $R = 1 : 1$ merging, but to a lesser extent because of the reduced mass of the secondary. There is a smaller increase in entropy as the two core collide, with radiative losses subsequently reducing entropy down to $S_{\text{ad}}(0) \simeq 100 \text{ keV cm}^2$ at $t = 10 \text{ Gyr}$. In contrast to the adiabatic case, this level of core entropy is a factor $\sim 2$ smaller.

In the most off-center ($b = 0.6$) merger, the impact of the secondary on the core is negligible. As in the adiabatic case, low-mass subclumps with large AM are progressively disrupted by instabilities before being able to significantly shock-heat the primary’s core. However, in the adiabatic run there is a certain amount ($\Delta S(0) \simeq 100 \text{ keV cm}^2$) of core heating taking place during the late phases ($t \sim 9 - 10 \text{ Gyr}$) of the merger. This small entropy jump is now absent because of radiative losses, thus allowing the CC to survive.

Additionally, we ran another head-on radiative merging simulation with the mass ratio $1 : 1$. In contrast to the merging case previously discussed, for the two clusters we now adopt a halo mass of $M_{\text{halo}} = 6 \times 10^{13} \text{ M}_\odot$. The initial condition set-up being the same as described in Section 2.2.

For this simulation we show in the right panel of Figure 19 the final entropy profile of the merger remnant (solid thick line), together with the profile of its adiabatic counterpart. The entropy profile of the radiative run exhibits a near power-law behavior and, unlike the head-on $R = 1 : 10$ merging case, the CC has been rapidly re-established. This happens because now the mass of the primary is an order of magnitude smaller, thus the collision with the secondary (at $t \sim 2 \text{ Gyr}$) is able to raise the central entropy only up to $S(0) \simeq 80 \text{ keV cm}^2$. This is a factor $\sim 2$ smaller than in the the head-on $R = 1 : 10$ merger, so that the cooling time is $\sim 5 \text{ Gyr}$ and the CC is soon restored.

These findings demonstrate that radiative cooling dominates the final ICM core properties, with physical processes governed by timescales much shorter than those set by diffusion (Biffi & Valdarnini 2015).

To summarize, the results of this Section demonstrate that the final level of core entropy, and thus the resiliency of CCs to disruption, depends critically on the merging mass ratio and initial orbit. Specifically, CCs are destroyed in head-on high-mass mergers, but can survive low-mass mergers or off-axis low mass ratio mergers. This suggests that the merging AM is a key parameter which determines the final remnant core entropy. In merging systems with high AM, a CC is re-established after the final collision as long as the cooling time is shorter than the core free-fall time. This condition depends on the merging mass ratio as well.

Our results are, partially, in agreement with Hahn et al. (2017). The authors argue that AM is a fundamental quantity to determine whether a CC can survive a cluster merger or not. They found CC disruption to occur in major mergers with low AM, but it is absent if the AM is high. This is in contrast with our findings, for which the disruption of CC in major mergers (mass ratio higher than $1 : 5$) of massive clusters (i.e. with a primary mass $M_1 \gtrsim 6 \times 10^{14} \text{ M}_\odot$) occurs regardless of whether the merging is head-on or off-axis. Note also that the CC is not disrupted, and the AM becomes unimportant, when the primary mass is small ($M_2 \lesssim 10^{14} \text{ M}_\odot$).

Finally, our findings are in contrast with those of Z11. For the same set of merging initial conditions, that paper finds final levels of core entropy high enough to erase CCs, regardless of the considered merging case. This is at variance with the results of this Section, in which the merging simulations now include gas cooling. This shows that a realistic physical modeling of the simulations is crucial to address the issue of CC survival in merging clusters.

4 CONCLUSIONS

In this paper we have presented results extracted from a suite of idealized binary cluster merger simulations, realized using an N-body/hydro code which employs an improved SPH scheme. Each merging cluster simulation was performed by constructing two isolated gas+DM halos in equilibrium; initial positions and velocities of the halos are then assigned according to the specific orbital trajectory. We purposely adopted the same range of initial mass ratios and impact parameters as in a previous paper (Z11), so as to consistently compare our results with previous findings.

Our simulations are aimed at investigating how the heating of the gas core leads to an increase of its entropy and to the disruption of the original CC profile during cluster mergers. In order to assess the impact of different physical processes on the survival of CC systems, we consider adiabatic as well as radiative merging simulations. The latter incorporate cooling, star formation and supernovae feedback. Our main conclusions can be summarized as follows.

i) For adiabatic simulations and head-on mergers, the dominant source of entropy is the shocking of the gas at the epoch of the first collision, with core heating being later driven by mixing processes.

ii) In the case of off-axis mergers, core heating depends critically on the initial merging mass ratio, as well as on the angular momentum (AM) of the system. For equal-mass mergers part of the final core entropy owes its origin to the transfer of AM, induced by tidal torques occurring during the first encounter of the two clusters along their orbit. The corresponding increase in the primary circular velocity at this epoch (see Section 3.3) generates instabilities in the
cluster inner regions, and in turn an increase in its core entropy.

iii) For mergings with low mass ratios the previous effect is negligible or absent and the bulk of core heating occurs at late stages, when the secondary accretes onto the primary. The final increase in core entropy can be modest because, before the final merging with the primary, most of the secondary mass has been stripped along its orbit by instabilities and ram pressure.

iv) In general, our results from adiabatic simulations are in accord with previous findings (Z11). The initial CC profiles do not survive the various merger cases we considered, and because of the different physical processes occurring during cluster merging, high-entropy gas is always present in the cluster core after the merger.

v) From a numerical point of view, it worth noting the good hydrodynamical behavior of the ISPH code presented here. For a specific run (see Figure 9), we find the final entropy profile in good accord with the corresponding one shown in (Z11). This is a non trivial result since it demonstrates how the ISPH scheme, based on a Lagrangian formulation, can be considered competitive with Eulerian-based AMR codes in terms of hydrodynamical performance.

The scenario outlined above changes in several respects when cooling is incorporated in the simulations. The most important differences are:

i) The increase in core entropy during cluster merging is now counteracted in part by radiative cooling, thus leading to lower levels of final entropy in the cluster inner regions than in the corresponding adiabatic case. CCs are found to survive if the merger is only able to raise the central entropy to \( S(0) \lesssim 50 \text{ keV cm}^2 \). This implies that the cooling time is shorter than the Hubble time and the CC is restored.

ii) For high mass mergers, CCs are destroyed in major merger, but are resilient to off-axis mergers with low mass ratios. This suggests that the survival of CCs depends both on the initial mass ratio and AM, and is thus characterized by a two-parameter dependency. This can be considered as the most important result of this paper.

iii) Finally, this dependence on AM tends to disappear as one considers low-mass cluster mergers. We ran a head-on merger with a primary mass an order of magnitude smaller \((M_{200} = 6 \times 10^{13} \text{ M}_{\odot})\) than in the baseline simulations, and the final core entropy is found to be low and cooling-dominated.

Overall, these findings support the observational evidence (Pratt et al. 2013; Chon et al. 2012) of a correlation between the CC/NCC core morphology and cluster mergers (but see Barnes et al. 2018, for a different viewpoint). The results presented here are also in broad agreement with previous works (Ritchie & Thomas 2002; Burns et al. 2008; Poole et al. 2008; Hahn et al. 2017), aimed at investigating the impact of merging clusters on core properties.

In particular, Hahn et al. (2017) argued that CCs can survive major mergers with large AM. While our merging simulations also clearly indicate a significant role of AM in determining the status of the final core remnant, for the considered merging case we find that CCs are destroyed in high mass major mergers even in the case large AM. We suggest that this discrepancy is not significant and is of statistical origin. Our results are obtained from binary cluster mergers realized in isolation with specific initial conditions, whereas the Hahn et al. (2017) sample is comprised of only ten clusters extracted from a previous cosmological simulation.

A possibility which has been left open by our study of merging clusters occurs when the merger is between a NCC and a CC cluster. In such a case, if the CC of the secondary is able to survive the merger process, it might settle in the center of the primary, leading to a transition of the primary from a NCC to a CC state. However, we argue that this scenario is unlikely to occur.

This is justified by the chosen initial condition set up of our merger simulations, in which both the primary and the secondary are initialized as CC clusters. Our previous findings indicate that the CC of the secondary does not survive ram pressure and shock heating as it enters the atmosphere of the primary, regardless of the initial mass ratio and AM of the binary system. This result is valid for a primary CC cluster, but we expect it to be little affected by the level of core entropy of the primary.

The validity of our simulation results depends both on the numerical resolution of our simulations as well as on the adopted physical modeling of the gas. For several merging runs, in Section 3.2 we contrasted the final entropy profiles against parent merging simulations performed using a higher resolution. The stability of the corresponding profiles is shown in Figure 9 indicating that our simulations can be considered free of resolution effects and are numerically converged.

In our simulations the physical modeling of the ICM is based on a number of simplifying assumptions. In particular, the most relevant is the absence of a subgrid model for the energy injection from AGNs, which can offset radiative cooling in cluster cores. For the purpose of the present investigation, it is then important to assess the impact of AGN feedback on the results presented in Section 3.3. Rasia et al. (2015) argued that CC thermal properties are affected by AGN feedback, and that its absence renders CCs more resilient against late-time mergers. This is in contrast with the findings of Hahn et al. (2017), for whom CC stability is not affected by incorporating AGN feedback, regardless of the adopted feedback parameters. Hahn et al. (2017) suggest that this discrepancy is directly connected to the hydrodynamic codes used in the two sets of simulations. The authors performed their simulations using an Eulerian AMR code, while Rasia et al. (2015) employed an improved SPH scheme (see Sembolini et al. 2016). Hahn et al. (2017) argued that the treatment of thermal diffusion in the two codes is critical in determining thermal properties of the simulated CCs.

This topic has been discussed at length in Section 3.2 and, as mentioned in point v) above, there is a full consistency between the final entropy profile of our \( R = 1:1 \) b = 0 merging simulation with the parent one of Z11. This strongly suggests that the artificial diffusion parameters of our hydrodynamical scheme are correctly calibrated. We thus assume that the inclusion of AGN feedback in radiative merging simulations will mirror the thermal behavior of ICM seen by Hahn et al. (2017), and in turn should be of limited impact on the findings of Section 3.3.

However, the entropy profile of the head-on low-mass radiative merger depicted in the left panel of Figure 19 clearly exhibits an overcooling behavior in its inner regions. Incorporating thermal AGN feedback in this simulation will avoid...
runaway cooling and will bring the final level of core entropy to higher values. Without a dedicated simulation including AGN feedback, it is difficult to assess the fate of the CC at the end of this merger.

From the parameters of cluster C3 ($M_{200} = 6 \times 10^{13} M_\odot$) given in Table 1 and the entropy profile (21), we estimate an initial entropy value of $\sim 12 \text{ keV cm}^2$ at $r = 10 \text{ kpc}$. Assuming that AGN feedback during the collision will maintain the core entropy of the primary at approximately this level, we require that the gain in core entropy be limited to $\lesssim 50 - 60 \text{ keV cm}^2$ during the collision to keep the integrity of the CC. In this case, we expect the final central entropy level of the merger remnant to be below the CC threshold ($\simeq 80 \text{ keV cm}^2$) previously given.

An upper limit to the level of core entropy achieved by the primary during the collision can be inferred by looking at the adiabatic merging simulation. From Figure 19 we obtain a final core entropy $\simeq 100 \text{ keV cm}^2$, but in the radiative simulations cooling effects will reduce this level to lower values. We thus conclude that for this specific merging case, the inclusion of thermal AGN feedback will change the ICM thermal state of the core. The final level of core entropy is likely to be close to the threshold above which the CC will be destroyed.

To summarize, our findings support the scenario in which the observed CC/NCC dichotomy is driven by cluster mergers. The difference in the disruption histories of the CCs, between adiabatic and radiative merging simulations, demonstrates that a realistic modeling of ICM physics is crucial in order to properly investigate the behavior of core morphology during the merging phase. We argue that the inclusion in our simulations of AGN thermal feedback is unlikely to impact most of our findings, at least in merging simulations in which the final cluster remnant has a virial mass $M_{200} \gtrsim 6 \times 10^{14} M_\odot$.

However, it must be stressed that in our high mass merging simulations, CCs survive only in mergers with low mass ratios and high AM. This leaves open the problem if such a result is consistent with the observed fraction of CC/NCC clusters at the present epoch (Barnes et al. 2018).

This issue can only be addressed in a cosmological framework, in which the evolution of simulated clusters can be followed self-consistently in a cosmological volume. Our idealized merging simulations are performed in isolation, so that environment effects are absent. In a cosmological simulation, these effects are automatically taken into account, and one expects merging environments to be affected.

On the other hand, our results indicate that the majority of core heating occurs when the secondary enters the innermost regions of the primary cluster. We thus suggest that environmental effects will be of limited impact on our findings.

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DATA AVAILABILITY

The data underlying this article will be shared on reasonable request to the corresponding author.

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Cool core resiliency in galaxy cluster mergers

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