A 1-dimensional hydrodynamic model for the cooling and heating of gas in dark matter halos from $z = 6$ to $z = 0$

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1 INTRODUCTION

Galaxy formation is known to be an outcome of gravitational processes (clustering) and gas-dynamical effects acting together. The success of N-body simulations lies in their ability to predict the detailed clustering properties of galaxies from initial conditions set by the ΛCDM cosmological model (Frenk et al. 1985, Davis et al. 1985, Colberg et al. 2000). However, radiative cooling and dissipation become important when attempting to predict the detailed physical properties of galaxies, such as their stellar masses and star formation rates. Semi-analytic models with radiative cooling, star formation and quenching by AGN and supernova feedback, provide a simple description of the relevant gas physics, but are unable to track the spatial distribution of the gas in any detail (Springel et al. 2005, Croton et al. 2006).

The advent of cosmological hydrodynamic simulations (Hernquist & Katz 1989, Springel 2000, Springel, Yoshida & White 2001, Murali et al. 2002, Vogelsberger et al. 2014) provided better means to treat gas dynamics more self-consistently. However, current simulations are very expensive or do not have the sufficient mass resolution to study the small scale physics in detail in very large cosmological volumes. In addition, “sub-grid” models are used for studying small-scale, unresolved physical processes like star formation and feedback (Kay et al. 2002, Di Matteo, Springel & Hernquist 2005, Puchwein, Sijacki & Springel 2008, Crain et al. 2015). Three dimensional hydrodynamic simulations have been able to track the spatial distribution of gas with reasonable accuracy. Compared to these, one-dimensional hydrodynamic models are computationally less expensive and often provide valuable insight into the physics of gas and its radial distribution, while retaining simplicity. In addition, it is possible to carry out controlled parameter studies to elucidate the physical inputs that are necessary to match different aspects of the observational data.

In the standard picture of spherical infall of gas (Gunn & Gott 1972, Bertschinger 1985) into the dark matter halos, the accreting IGM is heated to the halo virial tempera-
ture behind an expanding virial shock. The gas is first supported by pressure in quasi-hydrostatic equilibrium. It cools radiatively, gradually contracting and forming a cold disk in which star formation happens (Fall & Efstathiou 1980, Mo, Mao & White 1998). One-dimensional Lagrangian models of gas in dark matter halos, incorporating this standard picture, have been around for a long time. One of the earliest papers (Thomas 1988) studies multiphase gas in the intra-cluster medium with balance of cooling and inflow. However, only gas is evolved in this model and not dark matter. In contrast, Shapiro & Struck-Marcell (1985) included a collisionless component in 1D but a “pancake” collapse instead of spherical collapse.

Thoul & Weinberg (1995) introduced a spherical collapse of two-fluid system including gas and dark matter. They included radiative cooling and the effect of photoionizing background on the abundances of ionic species. In our model, we use a numerical scheme identical to Thoul & Weinberg (1995) but with different treatment of gravity due to dark matter, radiative cooling and heating that is expected from supernovae/AGNs. This combined model of dynamics and realistic energetics of gas has not been carried out before in a simple, Lagrangian set-up. In a one-dimensional model of dark matter and gas, Ch´eze, Teyssier & Alimi (1997) were the first to emphasize the need for the energy sources like star formation to explain core sizes and densities in X-ray clusters. The first treatment of gravity as a gradually deepening potential well, that we use in our model, was presented in Perrenod (1978). Knight & Ponman (1997) later added radiative cooling and concluded that in low-mass halos the baryon content is overpredicted in such a model. Thus the idea of a central source of heating became indispensable. Through a simple spherical 1D model which incorporates dark matter and gas, Birnboim & Dekel 2003 showed that the viral shock in smaller halos ($\lesssim 10^{12} M_{\odot}$), is not stable and in such cases gas is not heated to the virial temperatures and directly falls to the halo center (this is known as cold-mode accretion). However, these authors considered radiative cooling only and completely ignored heating due to feedback.

The viral shock is very stable at cluster scales and the hot gas is close to hydrostatic equilibrium (Oppenheimer 2018), except in the central parts where cooling time is usually shorter. The gas that cools out within the central region of the halo, is expected to form stars. The relation between the stellar mass and the host halo mass has been parameterized extensively. The mass function of dark matter halos, calibrated from large-scale simulations, is combined with the observed number density of galaxies as a function of their stellar mass using a technique called statistical abundance matching (Moster et al. 2010, Munshi et al. 2013). Recent observations also show that the black hole growth is intrinsically tied with the star formation rates (Martín-Navarro et al. 2018) and this co-evolution of the black hole and the galaxy is in accord with the observed tight relation between the black hole mass and the bulge mass for massive early type galaxies in the local Universe. However, the clearest evidence of black hole feedback comes from observations of synchrotron-emitting radio lobes that are co-spatial with the X-ray cavities found in galaxy clusters. The work done to inflate such huge bubbles, comparable to the energy required to prevent cooling flows (Fabian 1994, Cavagnolo et al. 2011) in clusters, can only come from accretion onto supermassive black holes.

Most of the observations of gas in dark matter halos are limited to X-ray in clusters and groups at reasonably low redshifts. The radial profiles of density, temperature and entropy have been extensively studied for low-redshift ($z = 0.05 − 0.2$) clusters (Vikhlinin et al. 2005, Pointecouteau, Arnaud & Pratt 2005, Vikhlinin et al. 2006, Kovtov & Vikhlinin 2006, Cavagnolo et al. 2009). The spatial distribution of the gas from X-rays can constrain the radial profiles from theoretical models of the ICM. The evidence of AGN jet and ICM interaction was studied in great detail in the ROSAT/Chandra observations of Perseus (Boehringer et al. 1993, Fabian et al. 2003) and other clusters (MacMara et al. 2000). Quenching of the cooling of hot gas at low-redshifts, has been seen in the observations of cavities that put important constraints on AGN outburst energy and mean jet power (Birzan et al. 2004, Cavagnolo et al. 2011). In groups, radial profiles of mass, density and temperature are traced out for relaxed systems (Gastaldello et al. 2007) and the differences between the gas content and the wide variety of luminosities within $r_{2500}$, compared to clusters, have been studied in Sun et al. 2009. The group observations imply lower gas content within $r_{2500}$ but similarity with clusters beyond $r_{2500}$.

A cluster of higher redshift ($z = 1.46$) has been seen in the X-rays combining with Spitzer infrared (IR) observations (Hilton et al. 2010) in which the luminosity is less than the $L_X − T$ expectation. This is also the second highest redshift X-ray selected cluster, following a cluster at $z = 1.62$ (Tanaka, Finoguenov & Ueda 2010). Recently the redshift evolution of galaxy clusters is traced for SZ selected clusters using the X-ray data (Marriage et al. 2011, McDonald et al. 2014, McDonald et al. 2013, Nurgaliev et al. 2017, McDonald et al. 2017). It hints at the invariant cooling properties and X-ray morphology of clusters till $z \approx 2$. These latter SZ selected observations imply that massive clusters are remarkably similar out to $z \approx 2$. We present comparisons of the radial gas density profiles of such SZ selected halos with that of the cluster in our model. Additionally, there have been efforts to quantify the baryon content in clusters and groups by tracking the free electrons (Schaan et al. 2016).

In this paper, we propose a very simple yet useful 1-D model for the evolution of the gas in dark matter halos with only a few adjustable parameters. The model is tuned to match the observed relation between stellar mass and halo mass. We consider isolated halos of a wide range of masses and disregard effects of mergers as well as processes affecting satellite galaxies, such as ram pressure stripping and strangulation. We account for the cosmological accretion using average mass accretion histories of dark matter halos derived from cosmological simulations. From our models, we get the following important perspectives on the state of baryons within the halos: (1) the radial profiles of gas density and temperature in massive clusters which are dominated by hot-mode accretion, (2) total feedback energy required to approximately regulate cooling and heating in halos of different masses (3) how the baryon fraction evolves when we constrain the global energetics using the relation of stellar mass and halo mass. The models of the kind presented in this paper may be easy to implement in semi-analytic models of galaxy formation that do not follow gas-dynamical
processes in detail. This model may also provide an useful middle-ground between the idealized simulations of isolated clusters and groups that usually do not evolve dark matter halo and the cosmological hydrodynamic simulations that often lack sufficient resolution to study an individual object.

The paper is organized as follows. In section 2 we describe the set-up for our numerical experiments including the initial conditions, in section 3 we present all our results for four different halos with a wide range of masses under adiabatic and non-adiabatic conditions and finally in section 4 we discuss the implications of our results and caveats in detail.

2 PHYSICAL SET-UP

In this paper we evolve four halos namely $M_{200, z=0} = 5 \times 10^{14} M_\odot (M_{14}), 5 \times 10^{13} M_\odot (M_{13}), 5 \times 10^{12} M_\odot (M_{12}), 5 \times 10^{11} M_\odot (M_{11})$ from $z = 6$ to $z = 0$. The dark matter halo does not evolve dynamically. We use a parametric form for dark matter density profile and vary the parameters as functions of time/redshift. The gas, however, is evolved with the code time-step determined by the hydro CFL conditions appropriate for this setup. At all times the halos have an inner radius at $r_{\text{min}} = 1$ kpc. We use reflective boundary condition at $r_{\text{min}}$, such that the velocity of that gas shell is zero. For runs with radiative cooling, we include a careful treatment of the gas shells that move too close to this innermost shell (described in section 2.3.1).

2.1 Evolution of dark matter halos

We do not evolve the equation of motion for the dark matter shells and simply approximate the increase in the gravitational potential due to the growth of the halo using the fitted mass accretion histories proposed by van den Bosch (2002). These authors use extended Press-Schechter formalism (Press & Schechter 1974) to show that the average mass accretion history follows a simple universal form for a wide range of halos and cosmologies. They then follow the most massive progenitor of the halos at the current redshift, back in time, in the Millennium simulation and compare their analytical formulae with the mass accretion histories derived for halos in the Millennium simulation. In this paper we adopt the recipe that calculates the analytical formulae. The mass accretion history or $M_{200}(t)$ is evaluated using Eq. (A1) in the APPENDIX A of van den Bosch (2002) in which we use the parameters ($z_f$ and $\nu$) as given by the same recipe in its following equations.

Figure 1 shows the halo mass as function of redshift in our halos. In order to evolve the halo masses with time, we use the following cosmological parameters, $\Omega_m = 0.75$, $\Omega_m = 0.25$ and $\sigma_8 = 0.9$. The concentration parameter ($c$) is modelled using the analytical formula proposed by Zhao et al. (2009), which takes into account that the halo concentration is tightly correlated with the time at which the main progenitor of a halo gains $4\%$ of its final mass. It can be written in the following form:

$$ c = 4 \left(1 + \left(\frac{t}{3.75 t_{0.04}}\right)^{8.4} \right)^{1/8}, $$

where $t_{0.04}$ is the time when the main progenitor gained $4\%$ of its final mass. We use a gravitational acceleration due to $M_{200}(t)$ of the following form (Navarro, Frenk & White 1996),

$$ g(r, t) = -\frac{d\Phi_{\text{NFW}}}{dr} \left(\frac{\ln(1 + r/r_s(t))}{r^2} - \frac{1}{r(r_s(t) + r)}\right), $$

where

$$ F(c) = \ln(1 + c) - \frac{c}{1 + c}. $$

Thus the entire information of the dark matter is encapsulated in the expression of the gravitational acceleration that becomes a function of time (or equivalently, redshift). The inherent assumption is that within each hydrodynamic step for evolving the gas, the dark matter comes to equilibrium instantaneously.

2.2 Initial conditions for gas

We start our runs at $z = 6$ where we initialize the gas within the approximate virial radius, $r_{200}$ (the mean density within which is 200 times the critical density of the universe), to be in hydrostatic equilibrium confined by the gravity due to the dark matter halo. Outside $r_{200}$ we use a parametric density profile. This is the density profile that we set in the outskirts of the halo so that the mass accreted at each time can preserve the universal baryon fraction for the adiabatic case. This parametric profile may not describe gas density...
profile outside the halo accurately but serves as a reservoir of gas that falls onto the halo over the entire range of redshifts. We will describe these two profiles (inner and outer) in the next two subsections (section 2.2.1 and 2.2.2). The general density profile can be described as follows:

\[
\rho(r) = \rho_{\text{HSE}}(r) \quad r \leq r_{200},
\]

\[
\rho(r) = \rho_0(r) \quad r > r_{200},
\]

where \(\rho_{\text{HSE}}(r)\) is what we initialize for the gas density inside \(r_{200}\) and \(\rho_0(r)\) is the density profile that we set outside \(r_{200}\). The matching condition for the density profile inside and outside initially at \(z = 6\), is

\[
\rho_{\text{HSE}}(r_{200}) = \rho_0(r_{200}).
\]

We assume that initially the gas is isentropic throughout; this assumption breaks down inside the halo when the accretion shock is formed. With this density profile for all the gas shells we calculate pressure \((p)\), temperature \((T)\) and energy density per unit mass \((u)\) in the following way:

\[
p(r) = K_\rho \rho(r)^\gamma,
\]

\[
T(r) = \frac{p(r)/\mu m_p}{n(r)k_B},
\]

\[
u(r) = \frac{\rho(r)}{(\gamma - 1) \rho(r)},
\]

where \(\gamma = 5/3\), \(k_B\) is the Boltzmann constant, \(n(r) = \rho(r)/\mu m_p\) is the number density, \(K_\rho = K_{fr}k_B/\mu m_p (\mu_c m_p) \gamma^{-1}\) is an entropy index in which \(f_T = 1.16 \times 10^7\) is the conversion factor of temperature from keV to kelvin \((K)\). We have used the mean molecular weights, \(\mu = 0.62\) and \(\mu_c = 1.17\). We also use \(K = C \left(\frac{M_{200}}{M_0}\right)^{\frac{3}{2}}\) keV cm\(^2\) s\(^{-1}\), which is an initial self-similar entropy profile and \(C = 2\) is a normalization constant that we fix using the halo \(M_{200}\) that has mass \(M_0\) at \(z = 6\). This scaling is motivated by the X-ray entropy \(K = T_{\text{keV}}/n_e^{2/3}\), popularly used in cluster literature. We have the same \(K\) for the entire gas initially, but the temperature of the IGM gas outside \(r_{200}\) is much lower than the virial temperature because the density set by \(\rho_0\) profile (see Figure 4) is also low. Therefore, the gas falling onto the halo is supersonic and forms a virial shock close to \(r_{200}\).

We have tested that the evolution of the gas is independent of the temperature profile outside \(r_{200}\). All the shells (both inside and outside) have an initial Hubble-flow velocity \(v_i = H(z = 6)r\) as given in section 3.2 of Thoul & Weinberg (1995) for the case of self-similar collapse of the shocked accretion of a collisional, non-radiative gas.

2.2.1 Density inside \(r_{200}: \rho_{\text{HSE}}\)

We calculate \(\rho_{\text{HSE}}\) assuming hydrostatic equilibrium inside \(r_{200}\). Therefore, we solve the following equations,

\[
\frac{dp}{dr} = -\rho g,
\]

\[
p = nk_B T,
\]

\[
K = \frac{T_{\text{keV}}}{n_e^{\gamma^{-1}}},
\]

where \(T_{\text{keV}} = T/fr\), \(g\) is calculated from Eq. 2 at the time \(t(z = 6)\) and all the other symbols are as described above. In order to solve these equations we need to assume a boundary condition for gas density at \(r_{200}\). Since \(r_{200} \propto M_{200}^{1/3}\), the outer density for all the halos can be assumed to be identical. We make sure here that the parameters \(K\) (as discussed in the previous section) and density at \(r_{200}\), are such that the total amount of gas inside the halo at \(z = 6\) is close to the universal baryon fraction (within around 10\% - 15\%, Figure 2). It is worth noting that this is only the initial boundary condition and the density will vary according to the hydrodynamic evolution for times \(t > t_{z=6}\). The red curves in Figure 4 shows the initial number density and temperature for \(M_{14}\) and \(M_{11}\).

2.2.2 Density outside \(r_{200}: \rho_0\)

For the initial outer (> \(r_{200}\)) gas density profile, we use a broken power-law. If density at \(r_{200}\) is denoted as \(\rho_{200}\) and the density at \(2r_{200}\) as \(\rho_{2r_{200}}\), the general form of the density power-law is chosen to be

\[
\rho_0(r) = \rho_{200} \left(\frac{r}{r_{200}}\right)^{\alpha} \quad r_{200} < r < 2r_{200},
\]

\[
\rho_0(r) = \rho_{2r_{200}} \left(\frac{r}{2r_{200}}\right)^{-\alpha} \quad r > 2r_{200},
\]

where a different \(\alpha\) is obtained for different halos (5th column in Table 1).

Note that, \(\rho_{200} = \rho_{\text{HSE}}(r_{200})\) by Eq. 5 and \(\rho_{2r_{200}} = \rho_0(2r_{200})\) for a given halo. The motivation of a power-law in the outskirts is primarily the simplicity of the form. In order to fix \(\alpha\) for a given halo, we carry out a large number of test runs under adiabatic condition for that halo, for a range of \(\alpha\). Such runs are computationally inexpensive and hence very fast. This way we find the \(\alpha\) corresponding to which baryon fraction follows the universal value (≈ 0.17) particularly at late times. Thus we obtain Figure 2.

As is clear from the values of \(\alpha\) in the Table 1, for \(M_{14}\) the index is maximum (\(\alpha = 0\)). For smaller halos, the values of \(\alpha\) decreases slightly, making outer densities falling very slowly at very large radii. We also scan a range of values for the most appropriate extreme outer trajectory (> \(2r_{200}\)) initially, to make sure that by current time there are enough number of shells to fall around the virial radius. Note that the broken power-law has been obtained by trial and error. Also, the gas that we distribute outside the halo, solely serves as a reservoir from which the halo can accrete slowly with time. This outer density profile is not observationally constrained.

2.2.3 Number of shells

For all the four halos we tested the adiabatic runs with different number of shells. We set up the representative runs shown in this paper with the number of shells stated in Table 1, such that the evolution of the gas fraction (as discussed in section 2.2.2) remains approximately unchanged and the number of shells is not so high that the computation becomes expensive.

2.3 Evolution of gas

We solve the continuity equation, momentum equation and energy equation for concentric gas shells and use the ideal
The baryon fraction reflects a transient as the halo adjusts to a self-similar Gyr it remains at the universal value. The initial dip in the baryon fraction (0 times for adiabatic evolution). After around 3 Gyr the deviation from universal baryon fraction shows a maximum.

Figure 2. The gas fractions (see section 2.2.2) within r200 at all times for adiabatic evolution. After around 3 Gyr the deviation from universal baryon fraction (0.17) is less than 10% and after 5 Gyr it remains at the universal value. The initial dip in the baryon fraction reflects a transient as the halo adjusts to a self-similar profile.

The gas equation of state to relate pressure, density and temperature at each time-step. We use the numerical scheme implemented by Thoul & Weinberg (1995) (section 2.3), which is a standard, second-order accurate (both in space and time), Lagrangian finite-difference scheme, to solve the following hydrodynamic equations for Lagrangian shells:

\[
\frac{dm}{dt} = 4\pi r^2 p dr,
\]

\[
\frac{dv}{dt} = -4\pi r^2 \frac{dp}{dm} - g(r, t),
\]

\[
\frac{du}{dt} = \frac{p dp}{p^2 dt} + \frac{\Gamma_b - \Lambda_c}{\rho},
\]

\[
p = (\gamma - 1)\rho u,
\]

where \(p, \rho, u\) have usual meanings as described in section 2.2, \(v\) denotes velocity of the shells, \(dm\) is the mass in each shell, \(\Gamma_b\) is the net energy injection rate density into the gas from a central source of heating (discussed in detail in section 2.3.1), \(\Lambda_c = n_n n_c \Lambda(n, T)\) is the net radiative cooling rate and \(\Lambda(n, T)\) is the cooling function, the form of which we will specify in section 2.3.1. For adiabatic runs in section 3.1, the cooling and heating term in which the latter is simply the net radiative cooling rate (Gaspari, Ruszkowski & Oh 2013, Prasad, Sharma & Babul 2017). However, more relevant is the total energy input is the minimum of the cooling times of all the shells and update the internal energy \(u\) with the last term in Eqn 15. This way the main time-step is not too short. We combine the cooling and heating terms in which the latter is simply zero for pure radiative cooling runs. We use a first-order explicit (semi-implicit) method to update the specific energy of the gas shells in each cooling step if heating (cooling) dominates (Sharma, Parrish & Quataert 2010). Cooling is turned on only inside \(r_{200}\) at each time, although the shock radius may be larger than \(r_{200}\) at late times. However, we are interested in the energetics within the inner regions of the halo.

In the runs with only radiative cooling, a large number of shells move to the central region and the time-step reduces to very small values for all the runs with cooling (including the ones with feedback). In order to get rid of this difficulty, whenever the shells have temperature less than \(2 \times 10^4\) K (for \(M_{14}\) and \(M_{13}\)) and \(1.8 \times 10^4\) K (for \(M_{12}\) and \(M_{11}\)) and they have reached \(1.05 t_{min}\) (1.05 kpc), we put them \(t_{min} = 1\) kpc and set their velocity to zero and no longer evolve them. This means the smallest allowed time-step will correspond to a trajectory coming as close as 0.05 kpc to the innermost one at \(t_{min} = 1\) kpc. The exact value of the smallest time-step is typically 3 – 4 orders of magnitude less than \(t_{z=6}\). Also, it is worth noting that cooling time-step within each hydrodynamic time-step is allowed to be even smaller due to subcycling. For lower mass halos the threshold for freezing is slightly less because the virial temperature is not as high as in higher mass halos. Otherwise a large number of shells go into the frozen zone and the halo evolution becomes numerically unstable. Thoul & Weinberg 1995 use a very similar technique to freeze shells but instead of a criterion on temperature they use the ratio \(t_{cool}/t_{dyn} = c_f\) to determine the freezing condition \((c_f \approx 0.01)\). On the other hand, Forcada-Miro & White 1997 use a criterion on temperature to merge overlying shells with the inner one whenever the trajectories come too close due to rapid cooling.

In the cases with pure radiative cooling, our freezing condition ensures that the amount of cold gas mass in the central regions of all our halos are not extremely high, or in other words, only the shells within central 30 – 40 kpc (within a factor of 1.5, depending on the halo mass) may cool without getting virialized, in the first 1 – 2 Gyrs. At later times gas usually goes through a shock and cools out relatively slowly for all the halos except the smallest one which shows a stable shock slightly later. Reducing freezing temperature below what we use, will enhance the run time and will not make significant difference to the total amount of cold gas \((\approx 10^4\) K). It is worth noting that in our runs with both cooling and feedback heating, which is the most realistic scenario, very few shells (typically < 10) are frozen. Additionally, in order to verify our freezing conditions quantitatively, we compare the time-averaged hot gas content in the SAMS and the cluster in our model and get a close match (discussed in section 3.3.3).

For runs with feedback heating, we use an idealized Bondi-Hoyle-Lyttleton (Bondi & Hoyle 1944, Hoyle & Lyttleton 1939) accretion rate to compute the heating rate. This is an idealized treatment of heating and Bondi accretion rate is not necessarily the best estimate for black hole accretion rate (Gaspari, Ruszkowski & Oh 2013, Prasad, Sharma & Babul 2017). However, more relevant is the total energy in-
The following equations and parameters are used to implement heating,

\[
\dot{M}_{\text{HBL}} = \frac{4\pi G^2 M_{\text{BH}} \rho}{(c_s^2 + v^2)^{\frac{3}{2}}},
\quad \dot{M}_{\text{Edd}} = \frac{4\pi G M_{\text{BH}} m_p}{\epsilon_r \sigma T C},
\quad \dot{E}_{\text{feed}} = \epsilon_f \min(\dot{M}_{\text{HBL}}, \dot{M}_{\text{Edd}}) c^2,
\]

where \(\epsilon_r = 0.1\) is the radiative efficiency and \(\epsilon_f\) is the feedback efficiency which is adjusted to obtain the baryonic properties of the halos. The total energy that is injected within an injection radius \(R\) (which we take as a parameter) in time \(dt\) is \(E_{\text{feed}} dt\). It is useful to note that this energy is limited by the Eddington limit. Now within \(R\) there are a few shells so the energy should be distributed between them. A simplest model is assumed in which the energy injection rate is \(E_{\text{feed}} / (\frac{4}{3} \pi R^3)\) in each shell within \(R\). For calculating BHL rate we use the density \((\rho)\) and velocity \((v)\) of the first active shell and trajectory (i.e non-frozen) respectively.

In some hydrodynamic simulations a pre-factor, \(\alpha_{\text{inf}} > 1\) is used to enhance \(\dot{M}_{\text{HBL}}\). This is because the Bondi radius, at which the density will be higher (hence \(\dot{M}_{\text{HBL}}\) will be higher), is much smaller (see Eq. 11 in Prasad, Sharma & Babul 2017) and usually not resolved in large scale simulations. However, we do not use a pre-factor here and instead only adjust \(\epsilon_f\) to evaluate the feedback energy that can be useful to prevent cooling flow. Here we are more concerned with the total feedback energy required rather than a specific feedback mechanism.

The heating runs are computationally more expensive and to maintain some reasonable distance between shells (so that the time-step is not less than 5 – 6 orders of magnitude lower than \(t_{\text{cool}}\)), we set a cut-off cooling temperature for different halos below which we do not cool the shells (second last column in Table 1). Specifically, when a feedback event creates a low-density bubble/cavity in the central region of the halo, shells start cooling out and compressing around it and this reduces the time-step significantly. We control the time-step by a cut-off cooling temperature. This technique of controlling rapid cooling around cavities, is similar to how we control extremely rapid cooling at the innermost trajectory with a threshold. Also, for different halos we have used different values of heating radius \((R)\) and efficiency \((\epsilon_f)\) to match the observed stellar mass versus halo mass relation (values given in Table 1) as discussed in the results (section 3.3).

3 RESULTS

A summary of all the different initial parameters tested for different halos is given in Table 1.

3.1 Adiabatic runs

The adiabatic runs reflect the effect of the growth of the dark matter halos on the inflow of gas, and an almost self-similar evolution of the gas density profiles. A temperature gradient develops over time within the halo, with a characteristic non-isothermal profile.
A simple model for gas in CGM and ICM

Figure 4. The evolution of number density (left) and temperature (right) for runs $M_{14}$ and $M_{11}$. The red lines show the initial conditions as described in section 2.2. The blue, green and black lines show the profiles at later times. The IGM gas outside the virial radius has very low temperatures at later times because of adiabatic expansion and the region close to the virial radius has the maximum temperature, which is consistent with observations.

Table 1. A summary of the halo models and different initial conditions, cooling and heating parameters

| Model | Current halo mass ($M_{200} (z = 0)$) ($M_\odot$) | Number of shells | Initial outer shell radius (kpc) | $\alpha$ | Shells frozen at (K) in runs with cooling | $\epsilon_f$ | $R$ (kpc) | Cooling stopped at (K) in runs with heating | Seed black hole mass ($M_\odot$) |
|-------|---------------------------------|-----------------|-------------------------------|--------|-----------------------------------|-------|--------|---------------------------------|-----------------|
| $M_{14}$ | $5 \times 10^{14}$ | 180 | 901.0 | 0.0 | $2 \times 10^4$ | 0.3 | 27.0 | $5 \times 10^4$ | $1 \times 10^8$ |
| $M_{13}$ | $5 \times 10^{13}$ | 150 | 501.0 | 0.63 | $2 \times 10^4$ | 0.08 | 18.0 | $4 \times 10^4$ | $1 \times 10^8$ |
| $M_{12}$ | $5 \times 10^{12}$ | 135 | 271.0 | 0.24 | $1.8 \times 10^4$ | 0.03 | 18.0 | $3 \times 10^4$ | $1 \times 10^5$ |
| $M_{11}$ | $5 \times 10^{11}$ | 120 | 136.0 | 0.36 | $1.8 \times 10^4$ | 0.005 | 12.0 | $2 \times 10^4$ | $1 \times 10^5$ |

Notes: All the models are tested in adiabatic and/or non-adiabatic runs. The inner most radius is fixed at 1.0 kpc for all runs. The initial outer shell radius is selected in the adiabatic runs such that almost all the shells join the halo by the end of the evolution.

The lower panel of Figure 3 shows the evolution of trajectories of all the halos we consider. The shock radius moves out with time and the gas gets shock-heated and settles in almost parallel layers after it joins the halo, indicating that the shells are close to hydrostatic equilibrium. For the adiabatic cases, this is the typical nature of evolution of the gas shells for all the halos. At later times the shock radii of all the halos are slightly greater than the corresponding $r_{200}(t)$.

Figure 4 shows the time evolution of the number densities and temperatures of two representative small and large halos for different times. The most massive halo, corresponding to cluster, contains a large amount of hot gas which is shown in the temperature profiles of $M_{14}$ where the peak
Temperature goes up to a few keV. The smallest halos show a lower overall temperature which is consistent with a self-similar evolution. The shock discontinuity and the corresponding jump in the halo temperature are distinctly seen in all our halos.

Figure 6 shows the profiles of gas density to dark matter density ratio at different times. The radial density profile of gas in hydrostatic equilibrium that we impose initially at $z = 6$ is shown in red. At later times, however, inside the halo (within $r_{200}$), the gas begins to follow the dark matter gradually (except at very low radii where the gas profile starts forming small cores). This shows that the halo is slowly accreting from the gas reservoir that we have initially outside $r_{200}$. We do not impose any condition on the gas fraction outside the halo in this model and verifying how that varies far away from the halo, is beyond the scope of this current formalism. Figure 5 shows the density gradient as a function of radius in which the approximate self-similar evolution is also quite evident. The important aspect to note from these two figures is that, the evolution at later times is not dependent on the initial distribution of gas within the halo and the redistribution of gas from the outer region happens over time.

Now we move to more realistic runs, starting from runs with cooling but no heating, and then moving on to runs with both cooling and feedback heating.

### 3.2 Runs with only cooling

The middle row of Figure 3 shows all the halos in which we turn on radiative cooling as described in section 2.3.1. In the smallest halo, which has a current mass of $5 \times 10^{11} M_\odot$, the virial shock is initially unstable and the gas cools out and falls directly to the center. This is the cold mode accretion, prominent in halos smaller than $M_{12}$. Birnboim & Dekel (2003) shows that the critical mass below which cold mode accretion dominates is approximately $3 \times 10^{11} M_\odot$, irrespective of the redshift. We see that in $M_{11}$, by the time...
the shock becomes stable (around $6 - 7$ Gyr), the mass of the halo is close to $10^{11}$ $M_\odot$, which is slightly less than the critical mass discussed before. For bigger halos, the virial shock is stable at all times and all the halo masses are $\gtrsim 10^{11}$ $M_\odot$ almost from the beginning (see the blue lines in the panels of the middle row of 3).

For all our runs with only cooling, we define gas to be cold if the temperature is $\lesssim 2 \times 10^4$ K. In Figure 7 we compare the total cold gas within 40 kpc (solid line) with pure radiative cooling and the total amount of gas within $r_{200}$ (dashed line) in all the halos. It is important to note that not all these cold shells are frozen. Only the shells that reach $1.05$ kpc are frozen while the rest of the shells within 40 kpc are dynamically active (see section 2.3.1 where we describe freezing shells for numerical reasons). In the biggest halo there is a huge reservoir of hot gas and only the gas in the central region of the halo cools out and falls to the center. In $M_{11}$ a large amount of gas cools out easily as the virial temperature is lower than that in the bigger halos and cooling is more efficient. The mass accretion rate in bigger halos are significantly higher and hence there is an inflow of a larger amount of shock-heated gas continuously into the halo that may keep re-heating the gas inside.

3.3 Runs with cooling & feedback

We inject realistic amount of energy as feedback into our halos to prevent cooling flows. Feedback is triggered by the accretion of gas across the inner radius $r_{\text{min}}$. We use an idealized prescription for the accretion rate of gas. The seed black hole masses are given in the last column of Table 1. We are more interested in an estimate of the total energy budget for each halo and not in the specific feedback mode (AGN or supernova). However, we can predict the dominant source of feedback from the average feedback power in this model.

3.3.1 Constraining feedback parameters from stellar mass-halo mass relation

As an estimate of the stellar mass, we use the cold gas mass within 40 kpc. Again it is important to note here that all the shells within 40 kpc are not frozen and most are dynamically active. For these cases with feedback, we have slightly different threshold cooling temperatures for different halos to control rapid cooling around cavities (as discussed in the last paragraph of section 2.3.1), below which radiative cooling is turned off. Accordingly, we define cold gas for $M_{14}$ to be below $5 \times 10^4$ K, for $M_{13}$ below $4 \times 10^5$ K, for $M_{12}$ below $3 \times 10^5$ K and for $M_{11}$ below $2 \times 10^5$ K. This should be a good estimate of stellar mass, averaged over several dynamical times, since cooling gas ultimately forms stars. Once we have an estimate of the stellar mass, we can compare our simulations to the abundance matching results.

In Figure 9, the underlying dashed red ($z = 0$) and blue ($z = 1$) lines show the stellar mass and halo mass correlation deduced from observational data and large scale simulations (Moster et al. 2010). On top of those we put circles whose sizes correspond to the halo masses. For each of the four halos, we take the cold mass in the inner region (40 kpc) as a function of time and interpolate this to find the cold mass at a specific redshift. Thus red circles and blue circles denote two different redshifts as do the dashed red and blue lines ($z = 0$ and $z = 1$ respectively). The circles, which are entirely opaque, are derived from the cases with pure radiative cooling. This gives an idea of how much energy needs to be injected to match with the dashed lines. Accordingly, we try out a range of values for our two feedback parameters $\epsilon_f$ and $R$ (refer to section 2.3.1 and Table 1) so that all the

![Figure 8](https://example.com/figure8.png)

Figure 8. The trajectories of the fiducial feedback runs of $M_{11}$ and $M_{14}$ in the central region. Half of the current $r_{200}$ is marked by the blue line. The entire radial extent is upto $r_{200}(z = 0)$. The smallest halo (left) shows a brief rapid accretion phase initially, which grows the black hole significantly and generates a huge feedback event. It is evacuated for most times thereafter while the biggest halo retains enough gas upto a comparable scaled radius.
circles are as close to the dashed lines as possible. Thus we select four runs for our four halos respectively with specific $\epsilon_f$ and $R$. These are denoted by transparent red and blue circles of same sizes as those of the purely radiative cooling cases. The opaque circles show that in the absence of a central source of heating, there is an excess cold gas formation. We regulate the feedback energy injected by $\epsilon_f$, and the region in which this energy is distributed, by $R$, to obtain the transparent circles which are as close to the dashed lines as possible. Thus the dashed lines constrain our values of $\epsilon_f$ and $R$ and hence the total feedback energy.

These four selected runs, for which we adjust the abundance matching results, are considered the fiducial ones and the trajectories are shown in the upper panel of Figure 3. The smallest halo is completely evacuated by an early feedback event which can also be seen in the left panel of Figure 8. There is a large inward flow of gas at the very beginning which triggered the growth of black hole to a great extent. The figure shows the inner $r_{200}(z = 0)$ of the halo and $0.5r_{200}(z = 0)$ is marked by the blue lines. The gas shells that are thrown out, are very slowly accreted back into the halo and there are no significant bursts after that till the current time. For $M_{12}$, there is one significant burst almost at the beginning of the evolution at $z = 6$, just like in $M_{11}$. For $M_{13}$ and $M_{14}$ there are two significant bursts over the entire timespan and their evolutions are very similar. The right panel of Figure 8 shows the inner $r_{200}(z = 0)$ of the biggest halo, $M_{14}$, and the dynamics of the inner shells over time. This halo retains enough gas within the halo almost all the time.

### 3.3.2 Baryon fraction evolution

We consider the fiducial runs for all the halos and we consider the ratio of the mass of the gas within $r_{200}$ and the halo mass, which we define as the gas fraction or the baryon fraction. Recall that, we choose the initial density profile outside the halo such that the gas fraction follows the universal value after an initial transient (Figure 2). Our fiducial feedback runs show the deviations from the universal value (Bregman et al. 2017, Anderson & Bregman 2010) and this explains why baryons are often observed to be missing from smaller halos. Figure 10 shows the baryon fraction evolution for the fiducial runs. For the smallest halo, there is a large accretion event at the beginning which grows the black hole and generates a significant feedback event subsequently which easily throws away most of the gas out of the halo. These gas shells take a long time to enter the halo and move...
back towards the bottom of the potential well. For most of the time, the halo lacks more than 50% of the gas. For $M_{12}$, the halo is bigger than $M_{11}$ and the gas is not thrown far out of $r_{200}$. Consequently, around 7 Gyr, we find the baryon fraction rising sharply. Thereafter small cycles of cooling and heating follow. In $M_{13}$ and $M_{14}$, the cycles are progressively more prominent. Particularly for $M_{14}$, only 50% of baryons are missing for a brief period around 5 Gyr and the halo maintains its gas content continuously afterwards (Ettori 2003, Bregman 2007, Nicastro et al. 2005).

It is usually predicted from observations and simulations that most of the missing baryons are available in the cosmic filaments (Nicastro, Mathur & Elvis 2008, Hill et al. 2016) and recent observations (de Graaff et al. 2017) corroborate to such predictions. Figure 8 (discussed before) also clearly shows how the smallest halo is completely evacuated for most of the time while in the right panel, $M_{14}$ retains a large amount of gas in the central region.

Schaan et al. 2016 discuss the baryon content of clusters and groups by combining data from Atacama Cosmology Telescope and “Constant Mass” CMASS galaxy sample from the Baryon Oscillation Spectroscopic survey to measure kinetic SZ over the redshift range 0.4 – 0.7. They consider the proportionality constant between the kSZ signal obtained and the expected kSZ signal to be a proxy for the average baryon content. Note that these observations have several uncertainties and the proxy of the baryon fraction is only proportional to the free electron fraction $f_{\text{free}}$. In our model, $M_{14}$ and $M_{13}$ which have masses in the range of galaxy clusters and groups at current redshift, show slightly more baryon content than what is predicted by them, around $z = 0.5$, in the central regions. In our halos, the baryon fraction in the outskirts of the halo, roughly follows the universal value around that time, if we consider NFW dark matter profile. At smaller radii where cavity is blown out by the feedback, the average baryon fraction falls to around half the universal value.

3.3.3 Radial profiles of clusters

We test our simple model by comparing the radial profiles of $M_{14}$ with observations available from Chandra clusters and recent Sunyaev-Zel’dovich-selected clusters (McDonald et al. 2017).

Figure 11 (left panel) shows the time-averaged density profiles for different cases, within a redshift range $z = 1.2 - 1.9$. Firstly the red line shows the density profile for pure adiabatic evolution. In the outskirts of the halo, this line coincides with the hot gas profile of the SAM (cyan dashed line) which assumes that isothermal gas with $p \propto r^{-2}$ resides in the halos. The match at large radii is expected with the same mass accretion history. However, in the central region, the SAM isothermal profile (that subtracts the gas cooling out) does not accurately trace out the observed gas distribution. The case with pure cooling flow (in yellow) shows a very large density in the central region, as expected. However, if we compute the time-averaged (between $z = 1.2$ and $z = 1.9$) hot gas content in SAM (using same mass accretion history and cooling prescription) and our model with pure radiative cooling (that is, time-averaged after subtracting the mass of gas below $2 \times 10^4$ K at each time), we see that the amount of hot gas is $6.84 \times 10^{12}$ $M_\odot$ and $6.06 \times 10^{12}$ $M_\odot$ respectively. The orange line shows the fiducial case with radiative cooling and feedback heating. It almost coincides with the average density profile obtained from observations of 8 clusters (of different masses and redshifts, within $z = 1.2 - 1.9$), except for tiny wiggles (when the latter is shifted inward by 0.3 Mpc to match the shock location; note that some of the 8 clusters can be bigger/smaller than $M_{14}$). The fiducial run also delineates that some amount of gas cools and falls to the center (frozen), with a peak density at 0.001 Mpc.

The right panel of Figure 11 shows the comparison of the time-averaged temperature profiles for $M_{14}$ within $z = 0.05 - 0.15$ with those of the clusters at redshifts $z = 0.1, 0.14, 0.11$. The radial profile of Abel 1446 at $z = 0.1$ coincides with those of $M_{14}$ within the radial range of a few times 0.01 Mpc to a few times 0.1 Mpc. However, the profile corresponding to the adiabatic run, has lower temperatures as one moves towards the center. The fiducial heating run of $M_{14}$ has a temperature profile that varies around that of Abel 1446. The SAM temperature profile falls almost on top of Abel 1446. The two clusters in the middle of this redshift range, have around twice the temperature of Abel 1446 and $M_{14}$. The fiducial case also shows drastic temperature changes in the central core because of the regions around the cavity where several shells are cooling out rapidly and enhancing the density. In multidimensions, the temperature at this location will be shell averaged and hence will be higher because of the presence of multiphase (both cold and hot) gas in the same shell (Prasad, Sharma & Babul 2015, Choudhury & Sharma 2016, Fielding et al. 2017). The physics of the local cooling rather than the monolithic collapse of the entire shell, is missing in our 1D runs. The cavity temperature is as high as 10 keV. These jumps in the temperature will be smoothed out for a multidimensional run. But our temperature profiles approximately reproduce the observa-
tions particularly in the outskirts of the halo. Hot gas profiles of density and temperature could be better constrained with our future multidimensional simulations.

3.3.4 Qualitative changes in the equation of state

One of the useful ways to understand how the density of the halo gas is related to the corresponding temperature and pressure is to understand the equation of state of the medium. The assumption of an EoS is helpful to construct the gas profiles once the global SZ signal or the total X-ray luminosity is obtained. Hence, EoS is a relevant feature of gas evolution.

We try to see the qualitative changes in the equation of state of the gas in the ICM ($M_{14}$), specifically how $\Gamma (\equiv d \ln p / d \ln \rho)$ changes for the gas that has fallen into the dark matter halos. Figure 12 (left column) shows the profiles of best fitted $\Gamma$ for the pure adiabatic case in $M_{14}$. The initial index is 5/3, which is clearly seen by the slope of the red line (at $z = 6$) and that of the lines outside the shock. As the gas gets shock-heated, the index falls to around 1.0 in the adiabatic case and around 1.1 in the non-adiabatic case. The trends associated with different times, fall almost along these dashed lines except in the cooling zone.

Figure 12 (right column) shows qualitatively the flattening of the trend at higher densities for non-adiabatic evolution. The pressure-density relation grows flatter near the central cores of dark matter halos. Simple theoretical model of ICM, that matched with hydrodynamic simulations, has found the polytropic index to be 1.15 (Ostriker, Bode & Babul 2005) except inside the rapid cooling zone, while X-ray observations have found the index to be around 1.2 (Solanes et al. 2005). It is worth noting that recently Fender, Nagai & McDonald 2017 describe the ICM at $z = 0$ by a broken power-law for pressure-density relation from X-ray observations of SZ selected clusters. The cooling break is seen around $p_{500c}$, where the index falls around $\gtrsim 0.1$ (varying systematically with redshift) inside the rapid cooling zone. The consideration of non-thermal pressure support ($p = p_{\text{thermal}} + p_{\text{non-thermal}}$) will modify the index in this rapid cooling zone as well.

3.3.5 Growth of black holes

In our model we constrain the cold gas formation within 40 kpc at $z = 1$ and $z = 0$. Thus from a range of efficiencies and heating radii, we select the ones that provide us with the desired relation of stellar mass and halo mass at $z = 0$ and 1. However, the assumption of spherical Bondi accretion (Eddington limited) gives only a crude estimate of the accretion rate. With this accretion rate, we can also get an estimate of the growth of black hole, starting from a seed. It must be noted that the multidimensional version of our model, with jet inflated cavities along a specific direction, will give a more accurate picture and a smoother growth of black hole.

In our model, the small halos $M_{11}$ and $M_{12}$ show a brief growth of black hole at around the Eddington rate, quite early in time. Thereafter, they do not have any significant growth. In Willott et al. (2010), the authors present the data for quasars around $z = 6$ from Canada-France High-z Quasar Survey (CFHQS). They find that most quasars
are accreting at close to the Eddington rate, which gives an exponential growth. Earlier works, like Yu & Tremaine 2002, emphasize that many of the massive black holes grow mostly in the bright QSO-phases.

Figure 13 shows the black hole mass (solid lines) and the cold gas mass within 40 kpc for different times in the current day group and cluster of our model. This shows a similar average evolution for group and cluster. Additionally, this plot also reflects that most of the stellar mass accumulates at very early times ($\lesssim 4$ Gyr) and to some extent at very late times, which is in agreement with Martín-Navarro et al. 2018. For clusters with observed cavities, Rafferty et al. 2006 have shown that in general, the black hole grows at a rate which is roughly 3 orders of magnitude less than the star formation rate at very low redshifts. Our black holes grow at a higher average rate because of the unrealistic 1D assumption.

It is interesting to note that for the biggest halo ($M_{14}$) which grows to cluster scale by the current time, the black hole mass goes up to $\lesssim 10^{10} M_\odot$, which is quite high compared to what is observed in clusters. This can be partly an artefact of the fact that other sources of heating like thermal conduction, falling galaxies, turbulence, dark matter subhalos etc. are not taken into account. Moreover, Hlavacek-Larrondo et al. (2012) show the possibility of ultramassive black holes in brightest cluster galaxies with masses around a few times $10^{10} M_\odot$. For halo masses of around $\approx 10^{12} M_\odot$, black hole masses greater than $\approx 10^9 M_\odot$ have not been detected by current observations (Läsker et al. 2016). In our model, the final black hole mass for $M_{12}$ (with current mass $5 \times 10^{12} M_\odot$) is also $\gtrsim 10^9 M_\odot$.

3.3.6 Estimated sources of feedback

It is imperative to note that the injected feedback power is more important than the feedback mode – whether it is black
Figure 13. The black hole mass as computed from the Eddington limited mass accretion rate in our idealized spherical Bondi model (solid lines) and the cold mass within 40 kpc (which is a proxy for an average stellar mass in our model; dashed lines) as functions of time. The discrete growth of black hole mass will be smoothed out in a multidimensional model with feedback inflated hot cavities in specific directions.

Table 3. Comparison of time-averaged power for different halos

| Halo (Current mass in $M_\odot$) | Time-averaged feedback power (erg s$^{-1}$) | Time-averaged SN power (erg s$^{-1}$) | Final black hole mass ($M_\odot$) |
|----------------------------------|------------------------------------------|-----------------------------------|---------------------------------|
| $M_{14}(5 \times 10^{14})$       | $1.2 \times 10^{46}$                     | $6.2 \times 10^{43}$              | $8.9 \times 10^9$               |
| $M_{13}(5 \times 10^{13})$       | $6.2 \times 10^{44}$                     | $7.8 \times 10^{42}$              | $1.8 \times 10^9$               |
| $M_{12}(5 \times 10^{12})$       | $6.8 \times 10^{44}$                     | $5.5 \times 10^{42}$              | $5.0 \times 10^9$               |
| $M_{11}(5 \times 10^{11})$       | $4.8 \times 10^{43}$                     | $1.7 \times 10^{41}$              | $2.1 \times 10^9$               |

Notes: The time-averaged supernova power is estimated with a Kroupa IMF (using Eq. 2 of Kroupa 2001) in which the cold gas mass within 40 kpc is used as a proxy for the stellar mass.
the importance of central energy source, by adding realistic amount of feedback proportional to the spherical accretion rate. This enables us to compare the temporally varying state variables of the medium as well as the time-integrated quantities with those of existing models and observations. Additionally, we can predict whether the feedback is provided by AGN jets or if supernova feedback is sufficient.

The thermodynamics of the gas in clusters and groups are interesting to study, particularly with the advent of future galaxy redshift surveys and the measurement of the Sunyaev-Zel’dovich signal (Lim et al. 2017, Park, Alvarez & Bond 2018, Schaaf et al. 2016) combined with the existing X-ray observations. In order to test a large number of ICM models and reasonable parametric profiles to match new observations, simple semi-analytical models (Flender, Nagai & McDonald 2017) are preferred over large scale hydrodynamic simulations. However, our 1D model, with hydrodynamic evolution of baryons, is computationally less expensive and provide a reasonable middle ground. Moreover, this model can be easily extended to multidimensional simulations of gas in halos without directly evolving dark matter particles. Thus great simplification can be achieved without entirely losing the physics of cosmological accretion of baryons and dark matter.

We obtain the following interesting results from our 1D model:

- We see the cosmological infall of gas, starting from a small initial halo, undergoing virial shock at radius slightly greater than $r_{200}$, which is fixed by the mass accretion history. Using a simple profile for the gas reservoir outside the halo, we can ascertain that the average baryon fraction within the halo, in absence of cooling and feedback, is close to the universal value. This provides an attractive setup for multi-dimensional simulations in which the outer boundary is much further out than the virial radius.

- For pure radiative cooling, cold mode accretion may dominate for very small halos ($< 10^{11} \, M_{\odot}$). This halo mass is a few times smaller than in Birnboim & Dekel (2003). So our model incorporates both hot and cold mode accretion in appropriate conditions. In the smallest halo, the gas cools out and loses pressure support fast in the absence of central heating and this often raises the baryon fraction beyond the universal value. The biggest halo retains enough hot gas even in the absence of feedback. This causes it to maintain the universal baryon fraction as the outskirts are hot and pressure-supported while only the gas in the central region cools and falls to the center. The cosmological accretion rate in the cluster-scale halos is high and a large amount of shock-heated gas continuously joins the halo.

- We tune our feedback (modelled as Bondi accretion) parameters to obtain a stellar mass-halo mass relation consistent with abundance matching (Moster et al. 2010). We use the total cold gas mass within the central 40 kpc as a proxy for the stellar mass. These fiducial runs including cooling and feedback provide realistic estimates of the energy budget in each halo. These fiducial runs show some interesting characteristics:

  - The baryon fraction evolution of all the halos show signatures commonly predicted; e.g., smaller halos have majority of baryons missing due to the ejection by feedback (de Graaff et al. 2017) and the biggest halos maintain the baryons by intermittent cooling and heating cycles (Prasad, Sharma & Babul 2015, Bregman 2007, Anderson & Bregman 2010). In the former case, the cycles are delayed as it takes a long time for the gas to be recycled, while on cluster scales, the feedback-heated gas remains inside the halo and moves to the center quite fast.

  - The time-averaged density profile for our cluster-scale halos match very well with Chandra clusters and recent Sunyaev-Zel’dovich-selected clusters (McDonald et al. 2017). (Figure 11). This is the most interesting result of our model because a very simple 1D picture of cluster over a range of redshift can reproduce the estimated average density profiles of clusters quite accurately. Note that, we are only concerned with the average number density profiles in this work and the feedback efficiencies, in our model, can be tweaked to reproduce density profiles of CC (cool-core) and NCC clusters (also discussed in McDonald et al. 2017). We will revert to this in our future multidimensional version of the model. The interplay of local thermal instability and gravity, which gives the important parameter $t_{\text{cool}}/t_{\text{ff}}$, can also be modeled only in more than one dimensions (Choudhury & Sharma 2016).

  - A qualitative flattening of the equation of state (or broken power-law as discussed in Flender, Nagai & McDonald 2017) is seen in all the halos that include cooling and heating. Inside the virial shock, $\Gamma$ (where $p \propto \rho^\Gamma$) is around 1.1 while around the rapid cooling zone it falls down to around $\approx 0.4 - 0.6$.

  - We use a crude estimate of the black hole mass from the idealized Bondi accretion rate (Eddington limited). We see that in groups and clusters the black holes, with a reasonably high seed mass at $z = 6$, grow significantly till only around $z = 2$.

  - The supernova power estimated from Kroupa IMF and our simple estimate of the star formation rate, suggests that AGN feedback is relevant for all the halos and absolutely necessary for clusters. The average feedback power ($\approx 10^{46}$) in $M_{14}$ is typically observed for many cooling clusters with X-ray cavities and radio lobes.

Lastly, there are predictable artefacts due to the limitation caused by a simple 1D model in this work, particularly the lack of multifluid gas in a single shell, which cools and collapses entirely. This causes some of our fiducial (with central heating) radial profiles to have drastic density/temperature which will be smoothed out in an identical, multidimensional model. In the latter case, we will also need to incorporate more realistic star formation and feedback mechanisms. However, the general properties of this simple model are remarkable and closely match observations and other existing models.

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