

Mass loss of galaxies due to an ultraviolet background

Takashi Okamoto,^{1*} Liang Gao^{1,2} and Tom Theuns^{1,3}

¹*Institute for Computational Cosmology, Department of Physics, Durham University, South Road, Durham DH1 3LE*

²*National Astronomical Observatories, Chinese Academy of Science, Beijing 100012, China*

³*Department of Physics, University of Antwerp, Campus Groenenborger, Groenenborgerlaan 171, B-2020 Antwerp, Belgium*

Accepted 2008 August 12. Received 2008 July 25; in original form 2008 May 30

ABSTRACT

We perform cosmological hydrodynamic simulations to determine to what extent galaxies lose their gas due to photoheating from an ionizing background. We find that the characteristic mass at which haloes on average have lost half of their baryons is $M_c \sim 6.5 \times 10^9 h^{-1} M_\odot$ at $z = 0$, which corresponds to a circular velocity of 25 km s^{-1} . This is significantly lower than the filtering mass obtained by the linear theory, which is often used in semi-analytical models of galaxy formation. We demonstrate it is the gas temperature at the virial radius which determines whether a halo can accrete gas. A simple model that follows the merger history of the dark matter progenitors, and where gas accretion is not allowed when this temperature is higher than the virial temperature of the halo, reproduces the results from the simulation remarkably well. This model can be applied to any reionization history, and is easy to incorporate in semi-analytical models.

Key words: methods: numerical – galaxies: evolution – galaxies: formation – cosmology: theory.

1 INTRODUCTION

An ultraviolet (UV) background produced by quasars and stars acts to quench star formation in small galaxies by photoheating their gas, which gets too hot to be confined in their shallow potential wells. This was originally pointed out by Doroshkevich, Zel'Dovich & Novikov (1967), and investigated in the context of the cold dark matter (CDM) model by Couchman & Rees (1986). Depending on the spectral shape of the ionizing background, the typical virial temperature below which haloes lose their gas is around 10^4 K . Photoheating associated with reionization therefore inhibits star formation in dwarf galaxies and affects the faint end of the galaxy luminosity function (e.g. Efstathiou 1992).

Since Rees (1986) argued that in dark matter haloes with circular velocities around 30 km s^{-1} gas can be confined in a stable fashion, with radiative cooling balancing photoheating, many attempts to quantify the effects of an ionizing background on galaxy formation have been made using semi-analytical calculations (e.g. Babul & Rees 1992; Efstathiou 1992; Shapiro, Giroux & Babul 1994; Nagashima, Gouda & Sugiura 1999; Bullock, Kravtsov & Weinberg 2000; Benson et al. 2002a,b; Somerville 2002), spherically symmetric simulations (Thoul & Weinberg 1996; Kitayama et al. 2000) and three-dimensional cosmological hydrodynamic simulations (Quinn, Katz & Efstathiou 1996; Navarro & Steinmetz 1997; Weinberg, Hernquist & Katz 1997; Gnedin 2000). The effects

of self-shielding have been investigated using three-dimensional radiative hydrodynamic simulations by Susa & Umemura (2004a,b).

However, there is still debate about the value of the characteristic mass, M_c , below which galaxies are strongly affected by photoionization. Gnedin (2000) argued that $M_c = M_F$, the *filtering mass* that corresponds to the scale over which baryonic perturbations are smoothed in linear perturbation theory (see Appendix B). Hoft et al. (2006) used simulations to argue that $M_c \ll M_F$, in particular at low redshift. They argued that M_c follows from considering the equilibrium temperature, T_{eq} , between photoheating and radiative cooling at a characteristic overdensity of $\Delta \simeq 1000$; M_c haloes have a virial temperature $T_{\text{vir}} \simeq T_{\text{eq}} (\Delta = 1000)$.

The relation $M_c = M_F$ proposed by Gnedin is often used in semi-analytic modelling for dwarf galaxy formation in order to describe the quenching of star formation (Bullock et al. 2000; Benson et al. 2002a,b; Somerville 2002). An important application of this relation is to estimate whether reionization can explain the apparent dearth of satellite galaxies in the Milky Way, as compared with the high abundance of dark matter subhaloes (Moore et al. 1999). Interestingly, Somerville (2002) and Nagashima & Okamoto (2006) argued that models underestimate the number of dwarf satellite galaxies in the Local Group if the effect is as strong as inferred from the filtering mass.

Given the uncertainty in M_c , we were motivated to perform high-resolution hydrodynamic simulations of galaxy formation in a Λ CDM universe with a time-evolving UV background, and to measure the evolution of the baryon fraction as function of halo mass. We compare our results to those of Gnedin and Hoft et al.,

*E-mail: takashi.okamoto@durham.ac.uk

Table 1. Numerical parameters of the performed runs. Simulations use periodic boxes with the indicated box size, ‘zoomed’ simulations evolve part of the simulation box at a higher resolution. m_{SPH} and m_{DM} denote the particle masses of SPH and (high-resolution) dark matter particles, respectively. ϵ is the comoving gravitational softening length in terms of the equivalent Plummer softening. The gravitational force obeys the exact r^{-2} law at $r > 2.8\epsilon$.

Simulation name	Box size (h^{-1} Mpc)	Zoom	N_{SPH}	m_{SPH} ($10^4 h^{-1} M_{\odot}$)	m_{DM} ($10^4 h^{-1} M_{\odot}$)	ϵ (h^{-1} kpc)
Reference	4	No	256^3	4.76	21.7	0.5
High resolution	4	Yes	256^3	0.596	2.71	0.25
Low resolution	8	No	256^3	38.1	174	1
Reference-lbox	8	Yes	256^3	4.76	21.7	0.5

and formulate an intuitive and a simple model that reproduces our simulations very well. The model can be incorporated easily into any semi-analytic model of galaxy formation.

This paper is organized as follows. The details of our simulations are described in Section 2. Section 3 contains our numerical results which focus on the baryon fraction of the simulated haloes, and the characteristic mass at which the baryon fraction is on average half the cosmic mean. We present a simple physical model that describes the evolution of the baryon fraction of individual haloes and compare it to the simulations in Section 4. Discussion and a summary follow in Sections 5 and 6.

2 SIMULATIONS

2.1 Background cosmology and radiative processes

The physics of how low-mass haloes fail to accrete gas, or lose it after the gas is being ionized, is largely independent of cosmology, but for definiteness we assume a geometrically flat, low-density CDM universe (Λ CDM) with cosmological parameters: $\Omega_0 = 0.25$, $\Omega_{\Lambda} \equiv \Lambda_0/(3 H_0^2) = 0.75$, $\Omega_b = 0.045$, $h \equiv H_0/(100 \text{ km s}^{-1} \text{ Mpc}^{-1}) = 0.73$ and $\sigma_8 = 0.9$. The symbols have their usual meaning.

The radiative processes that we take into account are Compton cooling off the microwave background, thermal Bremsstrahlung cooling, and line cooling and photoionization heating from Hydrogen and Helium in the presence of an imposed ionizing background as computed by Haardt & Madau (2001). Although our cooling routine handles cooling and photoheating for many species, we assume primordial gas throughout this paper for simplicity. The routine was developed for a different project and will be described elsewhere.¹ Briefly, it employs tabulated reaction rates evaluated with CLOUDY (Ferland et al. 1998), assuming ionization equilibrium. We assume H and He I reionization occurs at $z = 9$, and He II reionization at $z = 3.5$.

We increase the thermal energy by 2 eV per atom both during H and He II reionization to mimic radiative transfer and non-equilibrium effects (Abel & Haehnelt 1999). This brings the evolution of T_0 and γ_{EOS} that characterize the low-density temperature–density relation, $T(\rho) = T_0 (\rho/\langle\rho\rangle)^{\gamma_{\text{EOS}}-1}$ closer to the measurements of Schaye et al. (1999).² The extra energy is distributed over redshifts with a half-Gaussian distribution with a dispersion $\sigma_z = 0.0001$ from $z = 9$ for Hydrogen reionization and with a Gaussian distribution centred on $z = 3.5$ and with $\sigma_z = 0.5$

for He II reionization. While our numerical results depend on the assumed ionizing background, our model, described in Section 4, can be applied to any assumed thermal history.

2.2 Runs

We use a modified version of the PM-TREESPH code GADGET2 (Springel 2005), the successor of the TREESPH code GADGET (Springel, Yoshida & White 2001). Hydrodynamics is treated with smoothed particle hydrodynamics (SPH; Lucy 1977; Gingold & Monaghan 1977), employing a ‘conservative entropy’ formulation that manifestly conserves energy and entropy (Springel & Hernquist 2002).

We have performed a series of runs which differ in numerical resolution and box size (Table 1). We also use a run with a different imposed ionizing background. The ‘reference’ simulation uses a periodic cubic box of comoving size $L = 4 h^{-1}$ Mpc, and equal numbers of $N = 256^3$ dark matter and SPH particles. The ‘high-’ and ‘low-resolution’ simulations are used to gauge numerical convergence. The ‘high-resolution’ simulation used zoomed initial conditions, where a region of size $L_h = 2 h^{-1}$ Mpc inside the ‘reference’ box is simulated with eight times better mass resolution. The region outside the high-resolution region is populated with dark matter particles only, at coarse resolution. Linear waves common to both simulations have identical phase and amplitude. The ‘low-resolution’ simulation uses an $8 h^{-1}$ Mpc comoving cube, with uniform mass resolution, eight times coarser than the ‘reference’ simulation. Simulation ‘reference-lbox’ has identical resolution to ‘reference’, using zoomed initial conditions in a larger box.

We convert SPH particles into collisionless (‘star’) particles in high-density regions (those with gas overdensity $\Delta \equiv \rho/\langle\rho\rangle > 1000$ and physical hydrogen number density $n_H > 5 \text{ cm}^{-3}$) in a time given by the local dynamical time, $t_{\text{dyn}} = (4\pi G\rho)^{-1/2}$. This decreases the computational cost of the simulations significantly,³ and we have verified it does not change the results as compared to simulations without star formation.

The evolution of gas in the temperature–density plane is shown in Fig. 1. Before reionization, low-density gas lies along the adiabat $T \propto \rho^{\gamma-1}$, with $\gamma = 5/3$, except where shocked, whereas cooling introduces a sharp maximum temperature of $T \sim 10^4$ K at higher density. Bigger simulation boxes will contain more massive haloes which do contain shocked gas above 10^4 K. Heat input during reionization makes the gas nearly isothermal at $T \sim 10^4$ K. At later times, gas is above a minimum temperature $T_{\text{min}}(\rho, z)$, set by the

¹ We would like to thank our colleagues J. Schaye, C. Dalla Vecchia and R. Wiersma for allowing us to use these rates.

² Throughout this paper, ρ and $\langle\rho\rangle$ denote physical gas density and the mean baryon density in the universe, respectively, unless otherwise stated.

³ We do not include stellar feedback, since we concentrate here on radiative feedback only.

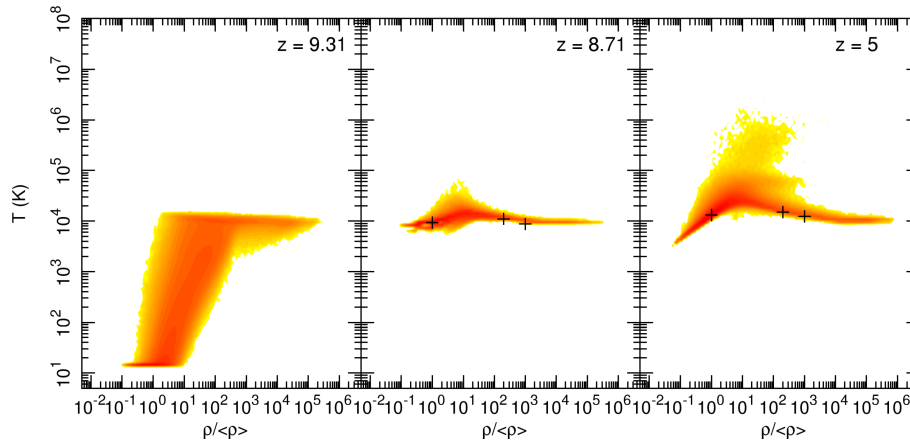


Figure 1. Distribution of the particles in the density–temperature phase diagram for the ‘reference’ simulation, colour coded according to the fraction of particles at given ρ and T , at redshifts $z = 9.31$ (before reionization), $z = 8.71$ (just after reionization) and $z = 5$, as indicated in the panels. Adiabatic cooling, impulsive heating during reionization, photoheating, radiative cooling and shocks from structure formation determine the distribution and its evolution. Plus signs indicate temperatures obtained assuming gas evolves at constant overdensity. This reproduces $T(z)$ well for the indicated overdensities of $\Delta = 1, 200$ and 1000 .

balance between photoheating and adiabatic cooling at low density, and photoheating and radiative cooling at higher density (Hui & Gnedin 1997; Theuns et al. 1998). Shocked gas forms the plume above $T_{\min}(\rho, z = 5)$ seen in Fig. 1.

The equilibrium temperature $T_{\text{eq}}(\rho, z)$ where photoheating balances radiative cooling is indicated by crosses for overdensities $\Delta = 200$ and 1000 , whereas we also include adiabatic cooling for the cross at $\Delta = 1$. As expected these points fall on the $T_{\min}(\rho, z)$ curves both for $z = 8.71$ (just after reionization) and the later redshift $z = 5$.

3 ANALYSIS OF SIMULATIONS

3.1 Halo identification

We identify haloes and the baryons associated with them as follows. First, we identify overdense regions using the friends-of-friends (FoF) algorithm (Davis et al. 1985) on the dark matter particles only, choosing the linking according to Okamoto & Nagashima (2003). This picks-out overdensities according to the spherical top-hat model appropriate for the assumed cosmological parameters (Eke, Cole & Frenk 1996).

We refine the location of the centre of each given halo, by determining the centre of mass of all the particles within a sphere of given radius R , and progressively shrinking R (Lacey & Cole 1994). The location of the density peak is identified as the centre of the halo. Finally, we determine the virial radius of the halo, R_{vir} , such that the mean dark matter density within R_{vir} equals the virial density of dark matter component.

Finally, SPH and star particles within R_{vir} from the halo centre are assigned to that halo. We restrict the analysis below to haloes with at least 100 dark matter particles, and discard haloes contaminated by boundary particles in the zoomed simulations.

3.2 Baryon fraction of simulated haloes

Before reionization, the baryon fraction $f_b \equiv M_b/M$ of simulated haloes scatters around the cosmic mean, $\langle f_b \rangle \equiv \Omega_b/\Omega_0$, as expected (Fig 2; the drop towards lower masses is a resolution effect). After reionization, f_b drops sharply below the characteristic mass $M_c(z)$;

haloes with $M \leq M_c(z)$ have too shallow potential wells to hold on to their photoheated gas. The drop of f_b from the cosmic mean to just of few per cent occurs over approximately a decade in total mass. The characteristic mass M_c where $f_b = \langle f_b \rangle / 2$ is the same for ‘reference’ and ‘high-resolution’ simulations demonstrating numerical convergence.⁴

Some haloes have $f_b \gg \langle f_b \rangle$. By tracking their merger histories we find that most of these were subhaloes (i.e. part of a larger halo) at earlier times, which have been tidally stripped from their hosts leaving baryon-rich cores. Ludlow et al. (2008) found that many apparently isolated low-mass haloes are subhaloes ejected by multiple-body interactions. This phenomenon of tidal stripping may have occurred for haloes with $f_b \leq \langle f_b \rangle$ as well; therefore, the baryon fraction of such low-mass haloes may be artificially high.

To reduce the contamination from tidally stripped haloes, we define *isolated* haloes as those that lie outside $6R_{\text{vir}}$ of more massive haloes. The baryon fraction of isolated haloes shows much less scatter as function of halo mass (Fig. 2, bottom panels), and we restrict our analysis to isolated haloes.⁵

We use the following function proposed by Gnedin (2000) to describe how the baryon fraction depends on mass and redshift:

$$f_b(M, z) = \langle f_b \rangle \left\{ 1 + (2^{\alpha/3} - 1) \left[\frac{M}{M_c(z)} \right]^{-\alpha} \right\}^{-3/\alpha} \quad (1)$$

in terms of the cosmic mean baryon fraction, $\langle f_b \rangle$. This function has two fitting parameters: α and $M_c(z)$. The baryon fraction of haloes with $M \gg M_c$ equals the cosmic mean while that of haloes with $M \ll M_c$ goes to zero $\propto (M/M_c)^3$. The parameter α controls how rapidly f_b drops for low-mass haloes; a value of $\alpha = 2$ fits the simulations well (Fig. 2), and is also consistent with the results of Hoefl et al. (2006). Haloes with mass equal to the characteristic mass, $M = M_c$ have lost half their baryons. As expected, the value of M_c for isolated haloes is slightly higher than for the full sample

⁴ We stop the ‘high-resolution’ simulation at $z = 5$ because of its small volume: the number of haloes with $M > M_c$ becomes too small to determine M_c . We stop the ‘reference’ simulation at $z = 2.09$ for the same reason.

⁵ This decreases the number of haloes from 3618 to 1185 for the ‘reference’ simulation at $z = 2.09$.

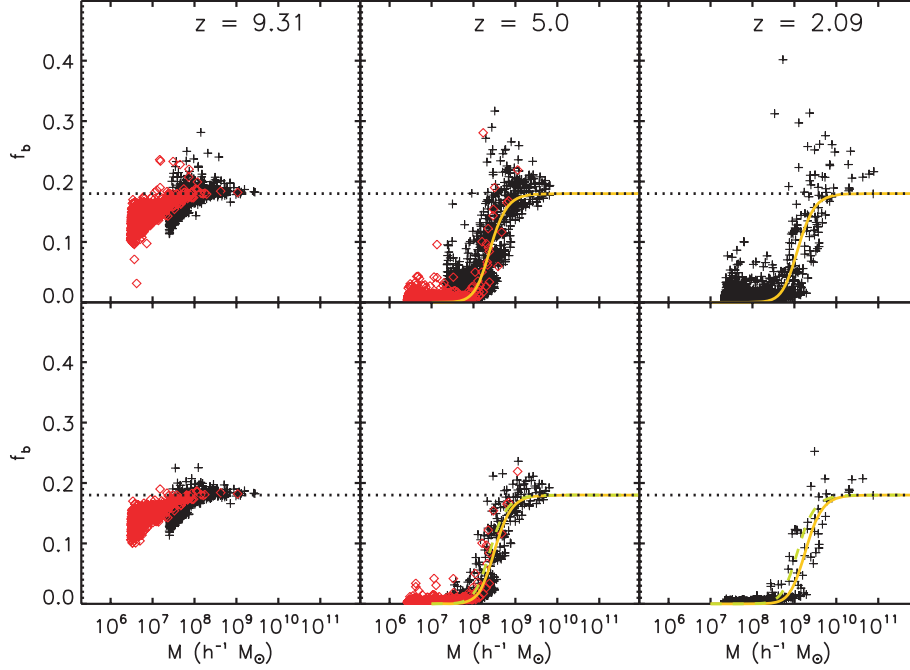


Figure 2. Fraction of baryonic mass as function of total mass, M , for simulated haloes at three redshifts z for all haloes (top row) and ‘isolated haloes’ (bottom row). The cosmological baryon fraction $\Omega_b/\Omega_0 \approx 0.18$ is indicated as the dotted horizontal line. Plus signs and diamonds indicate haloes from the ‘reference’ and ‘high-resolution’ simulations, respectively. Their comparison shows that the drop in M_b/M below $M \sim 10^8 M_\odot$ before reionization ($z = 9$ in these simulations) is an artefact of numerical resolution; at higher masses or after reionization, both simulations give similar values, demonstrating numerical convergence. The baryon fraction M_b/M falls from the cosmic mean to just a few per cent over about a decade in mass below a critical value M_c which depends on z . Halo-to-halo scatter is much reduced when considering only isolated haloes (compare top panels with bottom panels), which removes tidally stripped haloes. Full lines indicate fits using equation (1) to the ‘reference’ model; these fit the simulated haloes very well. The dashed lines in the bottom panel repeat the fits to the top row; the critical mass in isolated haloes is slightly higher than for all haloes.

of haloes (by 27 per cent at $z = 5$ and 50 per cent at $z = 2.09$ for the ‘reference’ simulation; see Fig. 2).

The characteristic mass increases from $M_c(z) \approx 10^7 h^{-1} M_\odot$ just after reionization to $M_c \approx 6.49 \times 10^9 h^{-1} M_\odot$ at $z = 0$ (Fig. 3). The ‘high-resolution’ simulation has a smaller value of M_c by 20 per cent as compared to the ‘reference’ simulation, at $z = 5$. This may simply reflect the small number of massive haloes in the ‘high-resolution’ simulation (see the lower middle panel of Fig. 2). Haloes with fewer than 10^3 dark matter particles may suffer from two-body heating due to massive dark particles (Steinmetz & White 1997); they are shown as open symbols. Values of M_c for these less well-resolved haloes are only slightly higher than of haloes resolved with more particles. Overall though simulations that differ in numerical resolution and/or box size give similar results, demonstrating numerical convergence. This gives confidence in the value of $M_c(z = 0) \approx 6.48 \times 10^9 h^{-1} M_\odot$ from the ‘low-resolution’ simulation. The amazing agreement with the value of $6.5 \times 10^9 h^{-1} M_\odot$ obtained by Hoefl et al. (2006) for simulations of voids may be slightly fortuitous.

We convert halo mass M to virial temperature T_{vir} using

$$T_{\text{vir}} = \frac{1}{2} \frac{\mu m_p}{k_B} V_c^2, \quad (2)$$

where

$$V_c = \left(\frac{GM}{R_{\text{vir}}} \right)^{1/2} \quad (3)$$

is the circular velocity of the halo at the virial radius. While we compute the mean molecular weight μ self-consistently in our simulations; here, we use $\mu = 0.59$ to compute T_{vir} for simplicity. In terms of the overdensity Δ_{vir} within the virial radius R_{vir} , T_{vir}

depends on halo mass and redshift as

$$T_{\text{vir}}(M, z) = \frac{1}{2} \frac{\mu m_p}{k_B} \left(\frac{\Delta_{\text{vir}} \Omega_0}{2} \right)^{1/3} (1+z)(GMH_0)^{2/3}. \quad (4)$$

The value of T_{vir} at the critical mass $M_c(z)$ rises rapidly to ≈ 5000 K following reionization, keeps increasing to $\sim 10^4$ K at $z = 2$, then increases much more slowly to $\approx 2.2 \times 10^4$ K ($V_c \approx 25 \text{ km s}^{-1}$) at $z = 0$ (Fig. 3, bottom panel). Our $z = 0$ value is much lower than the temperature below which galaxy formation is assumed to be strongly suppressed in semi-analytical models of galaxy formation (e.g. $\sim 10^5$ K in Benson et al. 2002b).

4 MODELLING THE SUPPRESSION OF THE BARYON FRACTION BY THE PHOTOIONIZING BACKGROUND

A simple model of gas accretion compares the halo’s virial temperature T_{vir} with the gas temperature T : if $T > T_{\text{vir}}$, there is no accretion. However, the temperature T of the intergalactic medium (IGM) depends on its density ρ (Fig. 1): at which density should we evaluate T ? Gnedin’s filtering mass suggests to use a density close to the mean (see Appendix B) whereas Hoefl et al. (2006) used $\rho = 10^3 \langle \rho \rangle$ appropriate for significantly overdense gas inside haloes. However, if one considers the density profile of a halo, gas at the edge of the halo (at R_{vir}) has a local density of order of one-third of the virial gas density, $\rho_{\text{vir}}/3$. This suggests that gas will be prevented from accreting when $T_{\text{eq}}(\rho_{\text{vir}}/3) > T_{\text{vir}}$. This temperature,

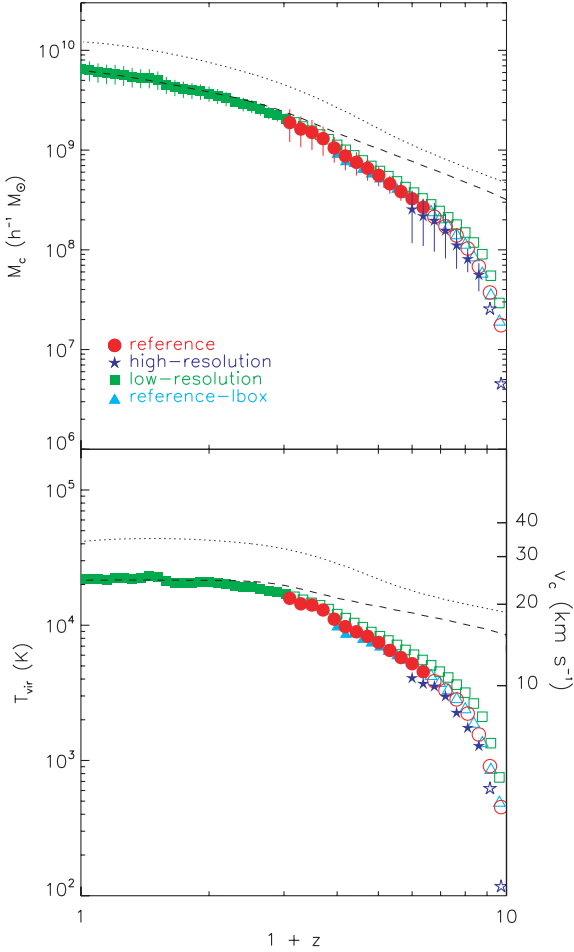


Figure 3. Top panel: evolution of the characteristic mass M_c with redshift z . Different symbols refer to simulations with the same assumed ionizing background but different resolution and/or size of the simulation box (‘reference’: circles; ‘high resolution’: stars; ‘low resolution’: squares; ‘reference-lbox’: triangles). When haloes of $M = M_c$ contain fewer than 10^3 dark matter particles, the simulation is shown with open symbols. Overlap of the different filled symbols demonstrates numerical convergence. M_c rises from $\sim 10^7 h^{-1} M_\odot$ just after reionization to $\sim 6.49 \times 10^9 h^{-1} M_\odot$ at $z = 0$. Lines indicate the virial masses corresponding to virial temperatures of $T_{\text{vir}} = T_{\text{eq}}(\Delta = 1/3\Delta_{\text{vir}})$ (dotted line) and $T_{\text{eq}}(\Delta = 1000)$ (dashed line); the redshift dependence of these lines is quite different from that of M_c , demonstrating a single temperature cut is not sufficient to compute the characteristic mass. It is necessary to follow the merger history of the haloes. Lower panel: corresponding evolution of the virial temperature (left-hand scale) and circular velocity (right-hand scale). M_c has been converted to T_{vir} using equation (4). The characteristic circular velocity $V_c(M_c)$ rises rapidly to $\approx 10 \text{ km s}^{-1}$ following reionization, then continues to increase more gradually reaching $\approx 20 \text{ km s}^{-1}$ at $z = 2$, and remains roughly constant thereafter. The equilibrium temperatures at two overdensities ($\Delta = 1/3\Delta_{\text{vir}}$: dotted line; $\Delta = 1000$: dashed line) are also shown; they are nearly constant, whereas the virial temperature of haloes with mass M_c increases rapidly especially at early times.

$T_{\text{eq}}(\rho_{\text{vir}}/3)$, is uniquely defined as a function of redshift⁶ from cooling and heating functions used in the simulations.

Simply identifying M_c from the condition $T_{\text{eq}}(\rho_{\text{vir}}/3) = T_{\text{vir}}$ does not work well, overestimating $M_c(z = 0)$ by a factor of ~ 2 as

compared to the simulations (dotted line in Fig. 3). In addition, the redshift dependence of $M_c(z)$ is quite different as well. Clearly, this model is too simple, as it ignores the baryon fraction of the progenitors of a given halo, as well as the photoevaporation of gas from haloes during reionization. For the same reasons, the agreement at $z = 0$ between a model by Hoefl et al. (2006) which compared the equilibrium temperature at high density ($\Delta = 1000$) with T_{vir} (dashed line in Fig. 3) and the simulations is probably a coincidence. Next, we describe an improved model that takes the *merger history* of haloes into account, as well as modelling photoevaporation and the decrease in accretion rate if gas is too hot.

Before reionization, we assume each halo has the cosmic baryon fraction, $\langle f_b \rangle$. We begin to construct merger trees of FoF haloes with more than 32 dark matter particles once the Universe is reionized. We model photoevaporation during reionization by computing the equilibrium temperature at large overdensity and comparing it to the halo’s virial temperature: when $T_{\text{eq}}(\Delta_{\text{evp}}) > T_{\text{vir}}$ the baryon fraction decreases $\propto \exp(-t/t_{\text{evp}})$, where the evaporation time-scale $t_{\text{evp}} = R_{\text{vir}}/c_s(\Delta_{\text{evp}})$. The time t is counted either starting from reionization, or from the previous merger (whichever is later), R_{vir} is the virial radius and the sound speed c_s is evaluated at $T_{\text{eq}}(\rho)$ is nearly independent of ρ for high densities, the exact value of the overdensity is not important as long as $\Delta_{\text{evp}} \gg \Delta_{\text{vir}}$; for reference, we use $\Delta_{\text{evp}} = 10^6$.

When two haloes merge, we add the baryon masses of the progenitors, and we also allow the halo to ‘accrete’ baryons, but only when the temperature of the accreting gas is lower than the virial temperature: $T_{\text{eq}}(\rho_{\text{vir}}/3) \leq T_{\text{vir}}$. Haloes that form after reionization at the resolution limit of the FoF are assumed to have the minimum baryon fraction,⁷ f_b^{min} . This completes our model ‘A’; it reproduces $f_b(M)$ as measured from the simulations reasonable well (Fig. 4, left-hand panel), with a small tendency to underpredict f_b for haloes with $f_b \sim 0.04$ in the simulations. The appendix describes a more involved model ‘B’ that does better for such haloes (Fig. 4, right-hand panel).

In detail, model A works as follows. Before reionization, each halo has the maximum baryon mass given by the cosmic mean (f_b):

$$M_b = M_b^{\text{max}} \equiv \frac{\langle f_b \rangle}{1 - \langle f_b \rangle} M_{\text{DM}}. \quad (5)$$

The total mass of the halo is simply the sum of its baryonic and dark matter masses: $M = M_b + M_{\text{DM}}$. After reionization, the baryon mass of a halo that forms from the merger of several progenitors is

$$M_b = M'_b + M_{\text{acc}} \\ M'_b = \sum^{\text{prog}} \exp\left(-\frac{\delta t}{t_{\text{evp}}}\right) M_b, \quad (6)$$

the sum of baryon masses of its progenitors, M'_b , and the accreted baryons, M_{acc} . The decrease in baryon fraction $\exp(-\delta t/t_{\text{evp}})$ is due to photoevaporation, with δt the elapsed time since the previous merger. For a progenitor with $T_{\text{eq}}(\Delta_{\text{evp}}) < T_{\text{vir}}$, we set $\delta t/t_{\text{evp}} = 0$ since evaporation does not take place in such a halo. The mass available for accretion is given by

$$M_{\text{acc}} = \max(M_b^{\text{max}} - M'_b, 0). \quad (7)$$

⁷ We assume the baryon density in a halo is at least the mean baryon density in the universe. Thus, $f_b^{\text{min}} \equiv \langle \rho \rangle / (\rho_{\text{vir}}^{\text{DM}} + \langle \rho \rangle) < 10^{-2}$. This gives a good estimate of the baryon fraction of low-mass haloes with $M \ll M_c$.

⁶ Note that the equilibrium temperature also depends on gas metallicity.

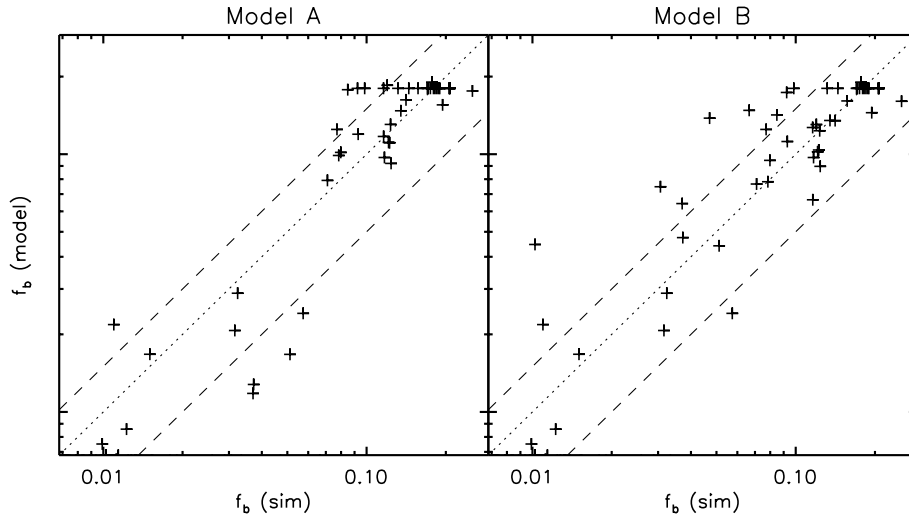


Figure 4. Baryon fraction of isolated haloes calculated by our models compared to the values from the SPH simulation for isolated haloes with more than 100 dark matter particles and with $f_b > f_b^{\min}$ picked from the ‘reference’ simulation at $z = 2.09$. The simple model A (left-hand panel) typically reproduces the baryon fraction obtained from the simulation within 50 per cent, whereas model B (right-hand panel) improves the agreement especially for haloes having low-baryon fraction, $f_b < \langle f_b \rangle / 2$. The dashed lines indicate $|f_{b, \text{model}} - f_{b, \text{sim}}| / f_{b, \text{sim}} = 0.5$.

When the temperature of the accreting gas, $T_{\text{acc}} \equiv T_{\text{eq}}(\rho_{\text{vir}}/3)$, is greater than the virial temperature of the halo, T_{vir} , the halo does not accrete any gas: $M_{\text{acc}} = 0$.

We can now predict the baryon fraction for all haloes selected from a simulation at a given redshift, and fit equation (1) to the distribution of $f_b(M, z)$; this yields the characteristic mass $M_c(z)$. This function is compared to that obtained from the simulations directly in Fig. 5; model A gets the simulated value to better than 20 per cent whereas model B is almost indistinguishable from the simulation. We also show model B without taking photoevaporation into account; this leads to an underestimate of M_c by more than a factor of 2 around $z \approx 6$, illustrating the importance of including this effect.

5 DISCUSSION

Our numerical results are in good agreement with those of Hoesft et al. (2006) who used the same simulation code (GADGET2) but a different implementation of photoionization and cooling (Haardt & Madau (1996) UV background with reionization at $z = 6$). The characteristic mass, M_c , at which haloes lose half of their baryons is much less than the value found by Gnedin (2000): we find that reionization is not effective at suppressing gas accretion in haloes of the filtering mass M_F (Appendix B), only haloes with $M \leq M_c \ll M_F$ are strongly affected. Our simulations suggest M_c is mainly determined by the temperature of the accreting gas at a density of one-third of the virial density, $\Delta \simeq 1/3\Delta_{\text{vir}}$, whilst Gnedin (2000) argued it is the temperature of the IGM at the mean density,⁸ T_0 . The values of T_0 in his and our simulations are in fact quite similar. A much stronger effect found by Gnedin (2000) was probably due to insufficient resolution (N. Gnedin, private communication).

It is, however, important to note that radiative transfer may lead to spectral hardening and may preferentially heat gas in overdense regions. To mimic this effect, we ran a simulation where we artificially increased the photoheating rate by a factor of 10, for gas at

overdensity $\Delta > 10$:

$$\mathcal{H}_{\text{mod}}(\Delta, z) = \begin{cases} \mathcal{H}(\Delta, z) & \Delta < 10 \\ 10\mathcal{H}(\Delta, z) & \Delta \geq 10. \end{cases} \quad (8)$$

$\mathcal{H}(\Delta, z)$ is the original rate at overdensity Δ and redshift z , and $\mathcal{H}_{\text{mod}}(\Delta, z)$ is the modified rate. This arbitrary modification heats gas at $\Delta = 1/3\Delta_{\text{vir}}$ to much higher temperatures yet does not affect T_0 (Fig. 6). Since T_0 is not affected, neither is the filtering mass. However since the accreting gas is hotter, our model predicts that M_c will be increased.

The volume-averaged temperature $\langle T \rangle_V$ in the ‘reference’ simulation with the modified photoheating is shown in the upper panel of Fig. 6 and is, as expected, almost identical to that in the original ‘reference’ simulation. Consequently, the filtering mass obtained with the modified photoheating is indistinguishable from that in the ‘reference’ simulation even if we compute M_F from $\langle T \rangle_V$ instead of T_0 (lower panel). The characteristic mass is, however, significantly higher, as predicted by our model.

This numerical experiment provides three important insights.

- (i) The mean IGM temperature is *not important* for the evolution of the baryon fraction of the haloes. Hence, any agreement between the characteristic and filtering mass must be a coincidence.
- (ii) The characteristic mass is sensitive to the temperature at the overdensity around $\Delta = 1/3\Delta_{\text{vir}}$, and therefore spectral hardening from radiative transfer effects can, in principle, increase the characteristic mass significantly.
- (iii) Our model works well for any reionization histories as long as the gas temperature as a function of density and redshift is known.

The modified heating rate used above probably significantly overestimates the effect of spectral hardening. Although radiative transfer effects may be important during reionization, their effect will lessen after the gas is ionized. Therefore, we believe our results with the original heating rate are robust, at least for lower redshifts ($z \leq 6$).

Interestingly, Rasera & Teyssier (2006) found a good agreement between their simulations and their model on the baryon budget, which was based on the filtering mass. The reason for this

⁸ Gnedin used the volume-averaged temperature $\langle T \rangle_V$ instead of T_0 .

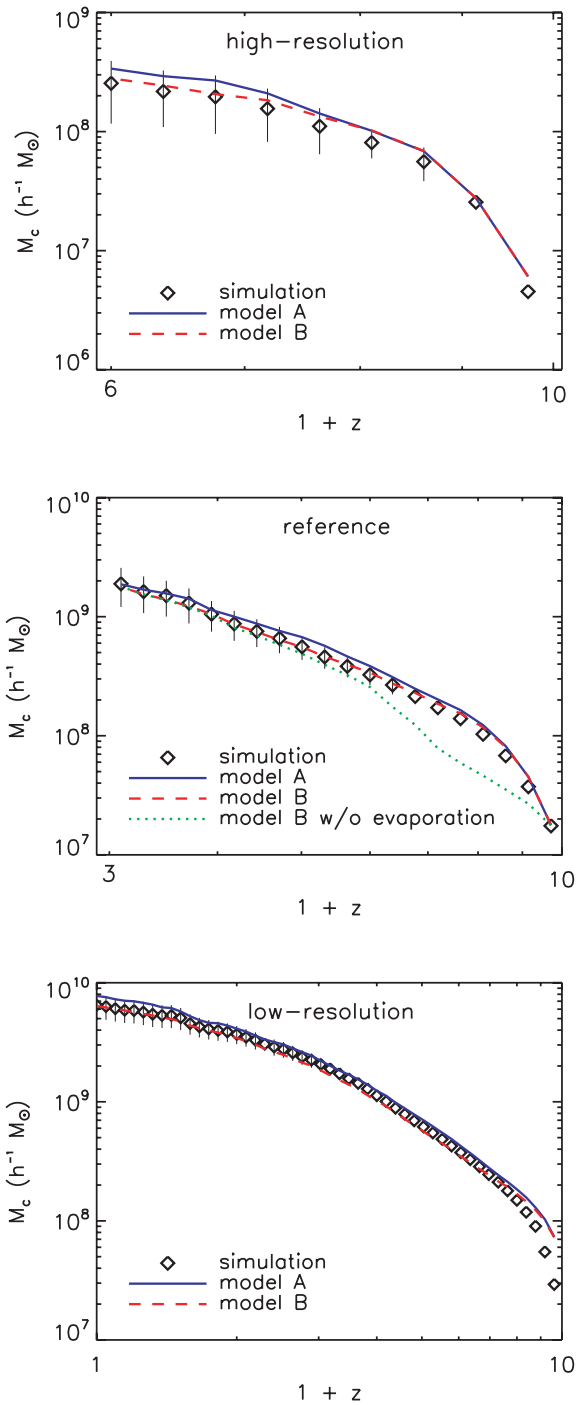


Figure 5. Evolution of the characteristic mass, M_c , below which haloes lose half of their baryons, comparing numerical simulations (symbols) with the simple models A and B described in the text. These models use the dark matter merger trees with simple recipes when haloes lose and accrete gas. Both models do well, with $M_c(z)$ from model B nearly indistinguishable from the full simulation. The middle panel also illustrates model B but neglecting photoevaporation during reionization: this leads to a significant underestimate in M_c .

agreement is probably that the minimum halo mass above which baryons are able to fall into the dark matter potential well was controlled by the numerical resolution rather than the filtering mass as they discussed. Although their highest resolution simu-

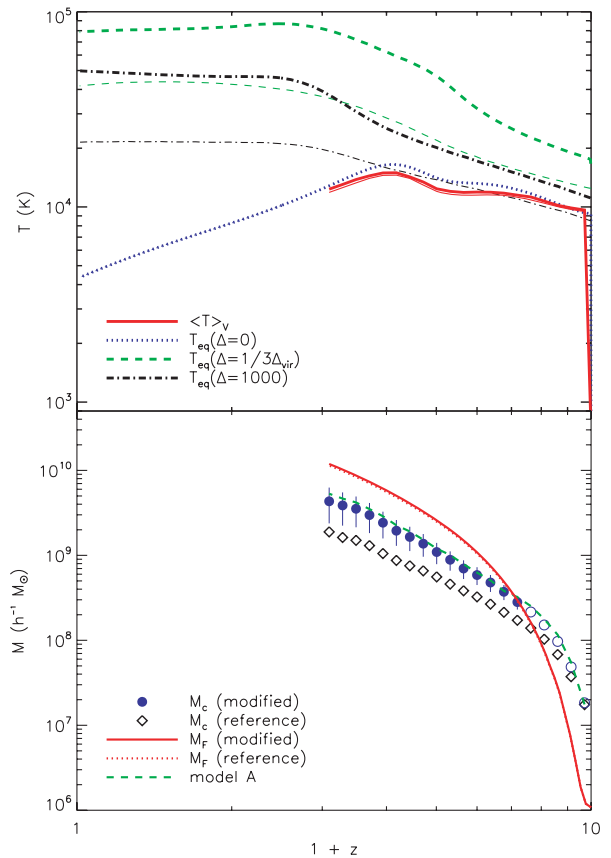


Figure 6. Upper panel: temperature evolution of gas at several overdensities indicated in the legend, for simulations with the Haardt & Madau (2001) photoheating rate (thin lines), and an experiment with increased photoheating at high density (thick lines). Lower panel: filtering mass, M_F , computed from $\langle T \rangle_v$ and characteristic mass, M_c , for the original photoheating rate (dotted line and circles) and modified photoheating rate (full line and diamonds). (Open circles indicate that haloes of total mass M_c do not contain a sufficient number of dark matter particles for reliable measurements.) Modifying the photoheating rate does not change M_F , but does increase M_c significantly. Model A with the modified heating rate (dashed line) reproduces this increase well. (Model B is not shown; it gives nearly identical results to model A).

lations well-resolved the haloes whose mass is below the filtering mass, these simulations were stopped just after the reionization where the discrepancy between the filtering mass and the characteristic mass is small (see Appendix B). We expect their model can be greatly improved by using our model of gas accretion and evaporation.

6 CONCLUSIONS

We have performed high-resolution cosmological hydrodynamical simulations with an imposed, spatially uniform but time-dependent UV background to investigate to what extent haloes lose baryons due to photoevaporation. Our numerical results are consistent with the simulations of void regions by Hoefft et al. (2006), but inconsistent with those found earlier by Gnedin (2000). The effect of a photoionizing background on galaxy formation is much weaker than previously assumed in semi-analytic models (e.g. Bullock et al. 2000; Benson et al. 2002a,b; Somerville 2002) based

on the filtering mass approach proposed by Gnedin (2000) (see Appendix B).

We have presented a novel model which uses the merger tree of each halo, and very simple physical arguments about gas retention and accretion, to predict how much the baryonic mass is in haloes of given total mass. The model is in excellent agreement with the simulations. Gas accretion is modelled very accurately if one simply assumes that gas at the outskirts of the halo, at an overdensity of a third of the virial gas density, cannot accrete if it is hotter than the halo's virial temperature. A slightly more involved model does even better in predicting baryon fractions as function of halo mass and merger history, as compared to simulations. These models allow us to predict the baryon fraction of individual haloes and hence the characteristic mass $M_c(z)$, at which haloes lose half of their baryons, for any assumed reionization history, by computing the predicted baryon fraction as a function of mass and redshift. The predicted $M_c(z)$ evolution is in very good agreement with the full simulations.

Our simulations and models both neglect self-shielding (Susa & Umemura 2004a), and radiative transfer and non-equilibrium effects (Abel & Haehnelt 1999). We have performed a numerical experiment in which we artificially increase the photoheating rate at higher density, to estimate how our results might be changed by these effects. The modified heating rate does not change the mean IGM temperature and hence also the filtering mass remains the same, yet M_c determined from the simulations increases significantly. Our models correctly predict this increase, demonstrating that the key physical quantity regulating the baryon fraction in haloes is neither the mean temperature of the IGM favoured by Gnedin (2000), nor the temperature of gas of $\Delta \simeq 1000$ favoured by Hoeft et al. (2006), but the temperature of the gas hanging around haloes ($\Delta \simeq 1/3\Delta_{\text{vir}}$).

Our models are easy to implement in any semi-analytic model of galaxy formation, provide one can compute the equilibrium temperature at $\Delta = 1/3\Delta_{\text{vir}}$ as a function of redshift for the assumed UV background.

ACKNOWLEDGMENTS

We are grateful to Volker Springel for providing us with the GADGET2 code and to Joop Schaye, Claudio Dalla Vecchia and Rob Wiersma for providing us with tabulated radiative cooling and heating rates. We thank Nick Gnedin and Masahiro Nagashima for their careful reading. We acknowledge support from a STFC rolling grant. The simulations were carried out at the Cosmology Machine at the ICC, Durham.

REFERENCES

- Abel T., Haehnelt M. G., 1999, ApJ, 520, L13
 Babul A., Rees M. J., 1992, MNRAS, 255, 346
 Benson A. J., Frenk C. S., Lacey C. G., Baugh C. M., Cole S., 2002a, MNRAS, 333, 177
 Benson A. J., Lacey C. G., Baugh C. M., Cole S., Frenk C. S., 2002b, MNRAS, 333, 156
 Bullock J. S., Kravtsov A. V., Weinberg D. H., 2000, ApJ, 539, 517
 Couchman H. M. P., Rees M. J., 1986, MNRAS, 221, 53
 Davis M., Efstathiou G., Frenk C. S., White S. D. M., 1985, ApJ, 292, 371
 Doroshkevich A. G., Zel'Dovich Y. B., Novikov I. D., 1967, Soviet Astronomy, 11, 233

- Efstathiou G., 1992, MNRAS, 256, 43P
 Eke V. R., Cole S., Frenk C. S., 1996, MNRAS, 282, 263
 Ferland G. J., Korista K. T., Verner D. A., Ferguson J. W., Kingdon J. B., Verner E. M., 1998, PASP, 110, 761
 Gingold R. A., Monaghan J. J., 1977, MNRAS, 181, 375
 Gnedin N. Y., 2000, ApJ, 542, 535
 Gnedin N. Y., Hui L., 1998, MNRAS, 296, 44
 Haardt F., Madau P., 1996, ApJ, 461, 20
 Haardt F., Madau P., 2001, in Neumann D. M., Tran J. T. V., eds, Clusters of Galaxies and the High Redshift Universe Observed in X-rays Modelling the UV/X-ray cosmic background with CUBA
 Hoeft M., Yepes G., Gottlöber S., Springel V., 2006, MNRAS, 371, 401
 Hui L., Gnedin N. Y., 1997, MNRAS, 292, 27
 Kitayama T., Tajiri Y., Umemura M., Susa H., Ikeuchi S., 2000, MNRAS, 315, L1
 Lacey C., Cole S., 1994, MNRAS, 271, 676
 Lucy L. B., 1977, AJ, 82, 1013
 Ludlow A. D., Navarro J. F., Springel V., Jenkins A., Frenk C. S., Helmi A., 2008, preprint (arXiv:801.1127)
 Moore B., Ghigna S., Governato F., Lake G., Quinn T., Stadel J., Tozzi P., 1999, ApJ, 524, L19
 Nagashima M., Okamoto T., 2006, ApJ, 643, 863
 Nagashima M., Gouda N., Sugiura N., 1999, MNRAS, 305, 449
 Navarro J. F., Steinmetz M., 1997, ApJ, 478, 13
 Okamoto T., Nagashima M., 2003, ApJ, 587, 500
 Quinn T., Katz N., Efstathiou G., 1996, MNRAS, 278, L49
 Rasera Y., Teyssier R., 2006, A&A, 445, 1
 Rees M. J., 1986, MNRAS, 218, 25P
 Schaye J., Theuns T., Leonard A., Efstathiou G., 1999, MNRAS, 310, 57
 Shapiro P. R., Giroux M. L., Babul A., 1994, ApJ, 427, 25
 Somerville R. S., 2002, ApJ, 572, L23
 Springel V., 2005, MNRAS, 364, 1105
 Springel V., Hernquist L., 2002, MNRAS, 333, 649
 Springel V., Yoshida N., White S. D. M., 2001, New Astron., 6, 79
 Steinmetz M., White S. D. M., 1997, MNRAS, 288, 545
 Susa H., Umemura M., 2004a, ApJ, 600, 1
 Susa H., Umemura M., 2004b, ApJ, 610, L5
 Theuns T., Leonard A., Efstathiou G., Pearce F. R., Thomas P. A., 1998, MNRAS, 301, 478
 Thoul A. A., Weinberg D. H., 1996, ApJ, 465, 608
 Weinberg D. H., Hernquist L., Katz N., 1997, ApJ, 477, 8

APPENDIX A: AN IMPROVED MODEL FOR GAS ACCRETION ON TO HALOES

Once density of accreting gas, ρ_{acc} , is specified, its temperature can be computed as $T_{\text{acc}} = T_{\text{eq}}(\rho_{\text{acc}})$. In model A, we assumed that the density of gas accreting on to a halo is one-third of the virial gas density, i.e. $\rho_{\text{acc}} = 1/3\rho_{\text{vir}}(z)$. If the temperature at this density, $T_{\text{eq}}(\rho_{\text{vir}}/3)$, is higher than the halo's virial temperature, T_{vir} , the gas cannot accrete. If merging haloes have an unusually low-baryon fraction, it may be that a larger fraction of baryons is available for accretion. We attempt to include this idea in our model 'B', by writing the density of accreting gas as

$$\rho_{\text{acc}} = \beta \frac{M_{\text{acc}}}{M_{\text{b}}^{\text{max}}} \rho_{\text{vir}}, \quad (\text{A1})$$

where β is a free parameter of the order of unity. A halo with baryon poor progenitors has higher $M_{\text{acc}}/M_{\text{b}}^{\text{max}}$, and hence higher ρ_{acc} . The temperature of the accreting gas is now evaluated at ρ_{acc} , and can be higher than the value $\rho_{\text{vir}}/3$ of model A. A value of $\beta = 2/3$ improves the predictions of model B over that of the simpler model A. With this value of β , models A and B accrete gas at the

same density for a halo with mass M_c , which has a baryon fraction $\langle f_b \rangle / 2$.

APPENDIX B: FILTERING MASS

Gnedin & Hui (1998) found that growth of the density fluctuations in the gas is suppressed for comoving wavenumber $k > k_F$, where the critical wavenumber k_F is related to the Jeans wavenumber k_J by

$$\frac{1}{k_F^2(t)} = \frac{1}{D(t)} \int_0^t dt' \frac{\ddot{D}(t') + 2H(t')\dot{D}(t')}{k_J^2(t')} \int_{t'}^t \frac{dt''}{a^2(t'')}. \quad (\text{B1})$$

The physical Jeans wavenumber k_J is defined as

$$k_J \equiv \frac{a}{c_s} (4\pi G \langle \rho_{\text{tot}} \rangle)^{1/2}. \quad (\text{B2})$$

Here, $D(t)$ and $H(t)$ are the linear growth factor, and Hubble constant as functions of cosmic time t , respectively, and the sound speed c_s is defined as

$$c_s = \left(\frac{5 k_B T_0}{3 \mu m_p} \right)^{1/2}. \quad (\text{B3})$$

Note that Gnedin (2000) used the volume-averaged IGM temperature, $\langle T \rangle_V$, instead of T_0 . He compared the corresponding filtering mass

$$M_F = \frac{4\pi}{3} \langle \rho_{\text{tot}} \rangle \left(\frac{2\pi a}{k_F} \right)^3 \quad (\text{B4})$$

with the characteristic mass measured in cosmological simulations of reionization. He found a good agreement between the characteristic mass and the filtering mass down to $z = 4$.

In the upper panel of Fig. B1, we show evolution of T_0 and the volume-averaged temperatures in the ‘reference’ and ‘low-resolution’ simulations. Since larger objects form in the ‘low-resolution’ simulation owing to its larger box size, the IGM temperature in the ‘low-resolution’ simulation at low redshift becomes slightly higher than that in the ‘reference’ simulation because of the shock heating. In the lower panel, we compare the filtering mass calculated from both T_0 and $\langle T \rangle_V$ with the characteristic mass. Clearly, the filtering mass significantly overestimates the characteristic mass at low redshift, while there is a reasonably good agreement at high redshift ($z > 6$) between them.

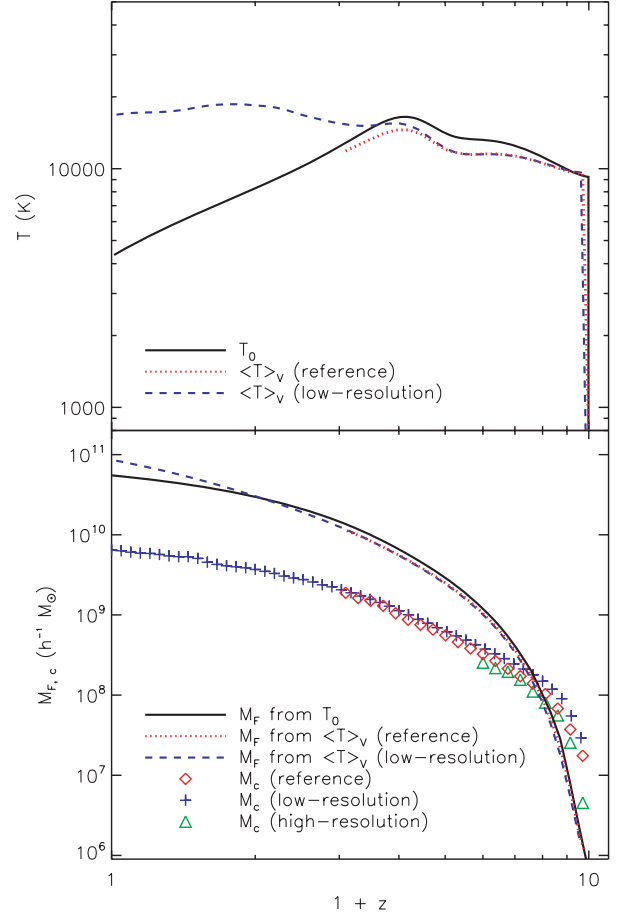


Figure B1. Upper panel: evolution of the IGM temperature at mean density (solid line) and the volume-averaged temperatures in the ‘reference’ (dotted line) and ‘low-resolution’ (dashed line) simulations. These temperatures are used to compute the filtering mass shown in the lower panel. Lower panel: comparison of filtering mass, M_F , obtained from equation (B4), and characteristic mass, M_c , obtained from the simulations. The filtering mass $M_F \approx 10M_c$ by $z = 0$, illustrating the fact that the filtering mass strongly overestimates the effect of photoheating on the formation of small galaxies.

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