Accretion on to black holes formed by direct collapse

Jarrett L. Johnson, Sadegh Khochfar, Thomas H. Greif and Fabrice Durier

Max-Planck-Institut für extraterrestrische Physik, Giessenbachstraße, 85748 Garching, Germany
Max-Planck-Institut für Astrophysik, Karl-Schwarzschild-Straße 1, 85740, Garching, Germany

Accepted 2010 August 5. Received 2010 August 5; in original form 2010 July 22

ABSTRACT

One possible scenario for the formation of massive black holes (BHs) in the early Universe is from the direct collapse of primordial gas in atomic-cooling dark matter haloes in which the gas is unable to cool efficiently via molecular transitions. We study the formation of such BHs, as well as the accretion of gas on to these objects and the high energy radiation emitted in the accretion process, by carrying out cosmological radiation hydrodynamic simulations. In the absence of radiative feedback, we find an upper limit to the accretion rate on to the central object which forms from the initial collapse of hot (\( \lesssim 10^4 \) K) gas of the order of 0.1 \( M_\odot \) yr\(^{-1} \). This is high enough for the formation of a supermassive star, the immediate precursor of a BH, with a mass of the order of \( 10^5 \) \( M_\odot \). Assuming that a fraction of this mass goes into a BH, we track the subsequent accretion of gas on to the BH self-consistently with the high energy radiation emitted from the accretion disc. Using a ray-tracing algorithm to follow the propagation of ionizing radiation, we model in detail the growth and evolution of the H\( \text{\textsc{ii}} \) and He\( \text{\textsc{iii}} \) regions which form around the accreting BH. We find that BHs with masses of the order of \( 10^4 \) \( M_\odot \) initially accrete at close to the Eddington limit, but that the accretion rate drops to \( \sim 10^{-5} \) \( M_\odot \) yr\(^{-1} \) (of the order 1 per cent of the Eddington limit) after \( \sim 10^6 \) yr, due to the expansion of the gas near the BH in response to strong photoheating and radiation pressure. One distinctive signature of the accretion of gas on to BHs formed by direct collapse, as opposed to massive Pop III star formation, is an extremely high ratio of the luminosity emitted in He\( \text{\textsc{ii}} \lambda 1640 \) to that emitted in H\( \alpha \) (or Ly\( \alpha \)), i.e. \( L_{1640}/L_{\text{H}\alpha} \gtrsim 2 \); this nebular emission could be detected by future facilities, such as the James Webb Space Telescope. Finally, we briefly discuss implications for the coevolution of BHs and their host galaxies.

Key words: hydrodynamics – galaxies: formation – galaxies: high-redshift – quasars: general – cosmology: theory – early Universe.

1 INTRODUCTION

How did the seeds for the first supermassive black holes (BHs) form? Observations of high-redshift quasars have shown that BHs can grow rapidly in the early Universe, becoming as massive as \( \gtrsim 10^9 \) \( M_\odot \) within the first billion years of cosmic time (e.g. Willott, McLure & Jarvis 2003; Fan et al. 2004, 2006; see also the recent reviews Haiman 2009; Volonteri 2010). However, it is still an open question how such BHs can grow to be so massive in such a short time. Given the strong connection between galaxies and the BHs that grow in the centres of their host dark matter (DM) haloes (e.g. Magorrian et al. 1998; Richstone et al. 1998; Silk & Rees 1998; Ferrarese & Merritt 2000), the growth of the first massive BHs must be intimately related to the formation of the first galaxies (e.g. Bromm et al. 2009). While the galaxies hosting the first massive BHs at redshifts \( z \gtrsim 10 \) have not yet been probed by observations, the radiation generated by massive BHs accreting in the early Universe will be the target of upcoming telescopes operating at various wavelengths, such as the Energetic X-ray Imaging Survey Telescope\(^1\) (e.g. Grindlay 2005; Grindlay et al. 2010), the James Webb Space Telescope\(^2\) (JWST; e.g. Gardner et al. 2006) and the Atacama Large Millimeter Array\(^3\) (see e.g. Schleicher, Spaans & Klessen 2010a). Therefore, in addition to the question of how the first BHs acquire their mass, it is important and timely to also consider the observational signatures of accretion on to the first massive BHs and whether they will be observable with future facilities.

A growing body of theoretical work has focused on the question of how supermassive BH seeds are formed and grow in the

\(^1\) http://exist.gsfc.nasa.gov/
\(^2\) http://www.jwst.nasa.gov/
\(^3\) http://www.eso.org/sci/facilities/alma/
early Universe, either from the collapse of Population (Pop) III stars (e.g. Haiman & Loeb 2001; Yoo & Miralda-Escudé 2004; Volonteri & Rees 2005; Johnson & Bromm 2007; Alvarez, Wise & Abel 2009; Tanaka & Haiman 2009; see also Devecchi & Volonteri 2009) or from the direct collapse of primordial gas in atomic-cooling haloes (e.g. Bromm & Loeb 2003; Begelman, Volonteri & Rees 2006; Spaans & Silk 2006; Lodato & Natarajan 2006; Regan & Haehnelt 2009a,b; Shang, Bryan & Haiman 2010; Schleicher, Spaans & Glover 2010b). In the former scenario, a BH forms with an initial mass of the order of 100 M⊙, or roughly the mass of its Pop III progenitor; in the latter case, the initial BH mass is likely $\gtrsim 10^3 M_\odot$, thereby providing a substantial jump start towards its growth to $\sim 10^8 M_\odot$ by $z \sim 6$.

The accretion rate and associated radiative feedback from the growth of BHs in the early Universe, in particular, have been studied by numerous authors by carrying out cosmological simulations (see e.g. Li et al. 2007; Pelupessy, Di Matteo & Ciaudi 2007; Di Matteo et al. 2008; Sijacki, Springel & Haehnelt 2009; Booth & Schaye 2009; Levine, Gnedin & Hamilton 2010; see also Sutter & Ricker 2010). Such simulations have commonly used simple prescriptions to estimate the accretion rate on to BHs and to gauge the effect of the radiation emitted during the accretion process (e.g. by injecting thermal energy without a radiative transfer calculation). Recently, however, complementary theoretical investigations have treated the radiative feedback from BH accretion on BH growth in great detail, but have used idealized, non-cosmological initial conditions (e.g. Milosavljević et al. 2009a; Milosavljević, Couch & Bromm 2009b; Park & Ricotti 2010).

In a study complementary to previous cosmological simulations concerning the accretion on to BHs formed from Pop III stars, which has been shown to be inefficient due to radiative feedback both from the preceding Pop III star (e.g. Yoshida 2006; Johnson & Bromm 2007; but see Whalen et al. 2008) and from the accretion of gas on to the BH itself (Alvarez et al. 2009), here we consider the accretion of gas on to BHs formed by direct collapse in an atomic-cooling cosmological halo with a virial temperature of $\sim 10^4$ K at $z \sim 15$, in which H₂ is strongly suppressed by a uniform background photodissociating radiation field. We employ a ray-tracing scheme to track the propagation of H₂ and He II ionizing photons, to account for the photodissociation of H₂ molecules, and to include the effects of radiation pressure on the primordial gas in the vicinity of the BH. Our simulations self-consistently couple the accretion rate of primordial gas on to the BH to the emission of high energy radiation from the accretion disc, the spectrum of which is calculated as a function of the BH mass and accretion rate at every time-step. Our results allow us to accurately calculate the rate of accretion of gas on to the BH and, importantly, to determine the observable signatures of the radiation emitted in the accretion process as it is reprocessed into nebular emission which escapes the host halo.

In the next section, we describe the setup of our cosmological simulation and the evolution of the primordial gas as it collapses into an atomic-cooling DM halo. In Section 3, we discuss the accretion of gas on to the central collapsed object that forms at the centre of such a halo, prior to its collapse to form a BH. Our calculation of the rate of accretion of gas on to such a newly formed BH and our treatment of the radiation emitted in the accretion process are described in Section 4. We present our results on the accretion rate and associated radiative feedback in Section 5; we consider the observational signatures of such BH accretion in the earliest galaxies in Section 6, and explore the implications of our results for the coevolution of BHs and their host galaxies in Section 7. Finally, in Section 8 we conclude with a brief summary and discussion.

2 COLLAPSE OF PRIMORDIAL GAS IN AN ATOMIC-COOLING HALO

The formation of a BH by direct collapse is predicted to occur when a sufficiently large (e.g. $\gtrsim 10^7 M_\odot$) mass of primordial gas is unable to cool much below $\sim 10^4$ K via molecular transitions and accumulates at the centre of a DM halo with a comparable virial temperature (e.g. Bromm & Loeb 2003; Dijkstra et al. 2008; Shang et al. 2010). Our simulation setup is chosen to realize these conditions. As is common in this scenario for direct BH formation, we invoke a strong Lyman–Werner (LW) background radiation field in order to suppress cooling of the primordial gas by H₂; however, we note that, in principle, other heating sources could also act to keep the gas sufficiently hot (e.g. Sethi, Haiman & Pandey 2010; see also Schleicher et al. 2010c).

2.1 Simulation setup

We carry out our three-dimensional numerical simulations with the parallel version of GADGET (version 1), which includes a tree (hierarchical) gravity solver combined with the smoothed particle hydrodynamics (SPH) method for tracking the evolution of gas (Springel, Yoshida & White 2001; Springel & Hernquist 2002). Along with H₂, H₂⁺, e⁻, He, He⁺ and He++, we have included the five deuterium species D, D⁺, D²⁺, HD and HD⁺, using the same chemical network as in Johnson & Bromm (2006, 2007).

For our simulation of the assembly of a primordial atomic-cooling DM halo at $z \sim 15$, we have employed multi-grid initial conditions which offer higher resolution in the region where the halo forms (e.g. Kawata & Gibson 2003). We initialize the simulation according to the Λ cold dark matter power spectrum at $z = 100$, adopting the cosmological parameters $\Omega_m = 1 - \Omega_k = 0.3, \Omega_b = 0.045, h = 0.7$ and $\sigma_8 = 0.9$, similar to the values measured by the Wilkinson Microwave Anisotropy Probe in its first year (Spergel et al. 2003). We use a periodic box with a comoving size $L = 1 h^{-1}$ Mpc for the parent grid. Our simulations use $N_{DM} = N_{SPH} = 1.05 \times 10^6$ particles for DM and gas, where the SPH particle mass is $m_{SPH} \sim 120 M_\odot$ in the region with the highest resolution. For further details on the technique employed to generate our multi-grid initial conditions, see Greif et al. (2008). The maximum gas density that we resolve is $n_{gas} \sim 10^5$ cm$^{-3}$, higher than in previous work using similar initial conditions (Johnson, Bromm & Bromm 2008; Johnson et al. 2009) owing to the effect of the elevated LW background radiation field that we have included at a level of $J_{1,W} = 10^4 \times 10^{-21}$ erg s$^{-1}$ cm$^{-2}$ Hz$^{-1}$ sr$^{-1}$. Following Shang et al. (2010), we account for the shielding of the gas from this radiation by calculating the column density of H₂ over the local Jeans length and then applying the prescription for H₂ self-shielding provided by Draine & Bertoldi (1996). As also described in Shang et al. (2010), we include the photodissociation of H⁻, assuming a background radiation temperature of $10^4$ K. This is appropriate for a radiation field generated by standard Pop II stars, which may constitute the bulk of the star formation already at $z \lesssim 15$ (see Maio et al. 2010).

2.2 Properties of the gas

Fig. 1 shows the properties of the primordial gas, as a function of the distance from the densest SPH particle in the most massive DM halo in our cosmological box at $z \sim 15$. At this point, $\lesssim 10^6 M_\odot$ of...
The properties of the primordial gas exposed to a molecule-dissociating background radiation field with $J_{\text{W}} = 10^3$, at $z \sim 15$, as a function of distance from the most dense SPH particle, at the location of which a BH is expected to form by direct collapse. Clockwise, from top-left: the number density of hydrogen nuclei, the gas temperature, the H\textsc{ii} fraction and the H$_2$ fraction. Contours denote the distribution of the gas at a fixed radius r, the mass fraction doubling across contour lines. These properties are in rough agreement with those that previous authors have found in similar simulations (see Wise et al. 2008; Regan & Haehnelt 2009b; Shang et al. 2010). Note that due to the generally low H$_2$ fraction, the gas temperature remains of the order of a few times $10^3$ K even at high densities in the core of the halo.

As shown in Figs 1 and 2, the temperature of the gas in the densest regions of the halo spans a range, but is generally well above $\sim 10^3$ K, owing to the low H$_2$ fraction of the gas, $f_{\text{H}_2} \lesssim 10^{-6}$, which leads to cooling times of $\gtrsim 10^6$ yr at densities $\lesssim 10$ cm$^{-3}$ and temperature $\sim 5 \times 10^4$ K. This is in agreement with previous simulations which have tracked the evolution of the gas collapsing into atomic-cooling haloes (Wise, Turk & Abel 2008; Regan & Haehnelt 2009b; Shang, Bryan & Haiman 2010). The density profile of the gas is also in agreement with what these authors have found. Furthermore, we find that the collapsing gas is only mildly turbulent in the centre of the halo, with Mach numbers $\lesssim 2$, as has been found in similar simulations which neglected H$_2$ cooling (Wise & Abel 2007). However, as shown in Fig. 3, we do not find appreciable supersonic infall of gas, i.e. cold flows, owing to the inability of the primordial gas to cool in the central regions of the halo, unlike that found in previous simulations in which H$_2$ cooling was not inhibited by a LW background (Greif et al. 2008).

As is illustrated in Fig. 4, a large portion of the gas within the virial radius, $\sim 600$ pc from the centre of the halo, has been shock heated near the virial temperature of $\sim 1.8 \times 10^4$ K, leading to the partial collisional ionization of hydrogen and to the efficient cooling of the gas through atomic transitions.

Fig. 5 shows the angular momentum distribution of the gas in the central region of the halo, which we also find to be in rough agreement with the previous work (e.g. Regan & Haehnelt 2009b). The angular momentum of the gas has a critical impact on how much mass the central collapsed object, presumably an SMS, can accrete.
before it finally collapses to form a massive BH, as is discussed in the next section.

3 EVOLUTION AND ACCRETION ON TO THE CENTRAL OBJECT

Before the formation of a BH by direct collapse, we expect that the gas in the centre of the halo first collapses to form an SMS, the immediate precursor to the BH. In this section, we study the accretion of gas on to this SMS in the absence of radiative feedback.

Beginning from the point of the simulation at $z \sim 15$, the state of which is shown in Figs 1–5, we calculate the rate at which gas accretes on to the SMS. To this end, we employ the same sink particle routine as used in Johnson et al. (2008), and we adopt a simple prescription for the accretion of SPH particles on to the sink particle, which represents the unresolved region surrounding the SMS in our simulation and is created at the location of the most dense SPH particle with an initial mass of $\sim 10^4 M_\odot$, roughly equal to the mass in the kernel. Defined as the density at which the local Jeans mass becomes comparable to the SPH kernel mass, the maximum
Figure 5. The specific angular momentum of the gas about the most dense SPH particle at $z \sim 15$, as a function of distance (right-hand panel) and of the enclosed mass (left-hand panel). As in Fig. 1, contours denote the distribution of the gas at a fixed radius $r$ and at fixed mass coordinate $M_{\text{enc}}$, in the right- and left-hand panels, respectively. The specific angular momentum needed for circularization of the gas, defined as $L_{\text{enc}} = (GM_{\text{enc}}r) \times 1/2$, is shown by the orange lines. At the radii resolved in our simulation, the gas is in general not rotationally supported and of the order of $10^3 M_\odot$ of gas may be accreted on to the sink particle representing the collapsed massive object in the centre of the halo. However, the flattening of the $l$ profile at $r \lesssim 10$ pc indicates that a rotationally supported disc is likely to form at radii that are not resolved in our simulation (see also Regan & Haehnelt 2009b).

density that is resolved is $n_{\text{res}} \sim 10^4$ cm$^{-3}$; in turn, the minimum resolved length-scale is roughly $(m_{\text{SPH}}/(m_{\text{H}}m_{\text{H}_{2}}))^{1/2} \sim 2$ pc, where $m_{\text{H}}$ is the mass of the hydrogen atom. Therefore, conservatively, we allow SPH particles to be accreted on to the sink particle if they come within 2 pc of the sink and if the semimajor axis of their orbit about the sink is $< 2$ pc. This allows us to resolve the evolution of the gas surrounding the sink particle and to obtain an upper limit for the mass of the SMS as it grows in time. We track the accretion of SPH particles for 2.2 Myr, roughly equal to the $\sim 2$ Myr maximum lifetime of an SMS (see e.g. Begelman 2010).

Fig. 6 shows the rate at which mass accretes on to the sink particle and the total mass accreted, as functions of the time since the formation of the sink. The accretion proceeds relatively smoothly, with variations in the accretion rate reflecting variations in the density of the accreting gas, shown in Figs 1 and 4. The mass of the sink particle increases continuously at a rate $0.1 M_\odot$ yr$^{-1} \lesssim M_{\text{sink}} \lesssim 0.3 M_\odot$ yr$^{-1}$. The total mass of the sink particle after 2 Myr is $\sim 3 \times 10^3 M_\odot$, consistent with the fraction, $\sim 2/3$, of the inner $10^3 M_\odot$ of gas which was moving on infalling trajectories at speeds approaching the sound speed, just before the creation of the sink particle (see Fig. 3).

It is important to note that the accretion rates that we find here are upper limits, for two reasons. First, the sink particle that is created exerts no pressure on the surrounding gas, which may lead to somewhat artificially enhanced infall velocities for the neighbouring SPH particles. Secondly, at this stage we neglect any radiative feedback on the gas, which could act to inhibit accretion. How important radiative feedback is in slowing the accretion of gas on to an SMS is not entirely clear. However, at the high accretion rates that we find the infalling matter may be optically thick, trapping the ionizing radiation emitted by the SMS and so limiting the radiative feedback at the scales we resolve.

Some fraction of the gas accreted on to the SMS is expected to end up in the BH that forms at its centre; this fraction may be up to $\sim 90$ per cent (see Shibata & Shapiro 2002), but may be perhaps an order of magnitude lower due to the copious high-energy radiation emitted from the BH as it grows in the centre of the star, as discussed in Begelman (2010). Therefore, we shall consider that a range of initial BH masses, from $\sim 10^4$ to $\sim 10^5 M_\odot$, is possible. In the following sections, we study the accretion of gas on to the BH that forms for different choices of its initial mass.

4 MODELLING BLACK HOLE ACCRETION AND RADIATIVE FEEDBACK

In this section, we describe our methodology for calculating the accretion rate on to BHs, as well as our treatment of the radiative output from the accretion disc, which is coupled in detail to the surrounding gas.

4.1 The accretion rate

The Bondi radius, which must be resolved in a simulation in order for the Bondi accretion rate on to a BH to be properly estimated, is given by

$$R_{\text{Bondi}} \sim 0.3 \text{ pc} \left( \frac{M_{\text{BH}}}{10^3 M_\odot} \right) \left( \frac{T}{10^9 \text{K}} \right)^{-1},$$

where $M_{\text{BH}}$ is the mass of the BH and $T$ is the temperature of the gas surrounding the BH. We compare this to the smallest scale that can be resolved in our simulation; specifically, we compute the minimum length-scale $R_{\text{res}}$ of the SPH kernel for which the Jeans mass is larger than the mass of the kernel ($\sim 50 M_{\text{SPH}}$), given by

$$R_{\text{res}} \sim 0.2 \text{ pc} \left( \frac{M_{\text{SPH}}}{120 M_\odot} \right) \left( \frac{T}{10^9 \text{K}} \right)^{-1}.$$
where $M_{\text{SPH}}$ is the mass of an SPH particle; in our case, $M_{\text{SPH}} = 120 \, M_\odot$. Thus, for the cases that we consider, with $M_{\text{BH}} \gtrsim 10^4 \, M_\odot$, the Bondi radius will always be resolved and we can have confidence in the accretion rates that we find, at least at the scales we resolve at which the gas is not rotationally supported, as shown in Fig. 3 (see Hopkins & Quataert 2010; also Booth & Schaye 2009). However, it should be noted that at smaller scales some fraction of the gas may instead go into an accretion disc wind or otherwise be accreted on longer time-scales than we assume here, such as over the viscous time-scale of the accretion disc (see e.g. Power, Nayakshin & King 2010). In this sense, the accretion rates that we find, calculated as described below, are likely to be upper limits.

The accretion rate on to the BH is calculated directly as the Bondi (1952) accretion rate:

$$\dot{M}_{\text{BH}} \sim \frac{4\pi(r_{\text{BH}})^2 \rho_{\text{gas}}}{(c_s^2 + v_{\text{BH}}^2)^{\frac{1}{2}}},$$

(3)

where the gas sound speed $c_s$ and gas density $\rho_{\text{gas}}$ are taken as the average of these values over all SPH particles within the smoothing length of the sink particle which marks the position of the BH; finally, $v_{\text{BH}}$ is the square of the velocity of the sink particle with respect to the SPH particles within its smoothing length.

We track the total accreted mass on to the BH and every time a mass equal to the mass of an SPH particle ($120 \, M_\odot$) is accreted, a random SPH particle lying within a smoothing length of the BH is accreted, using the same sink particle routine as discussed in the previous section.

### 4.2 Radiation emitted from the accretion disc

At every time-step, we compute the spectrum of the radiation emitted from the accretion disc of the BH, as a function of the mass of the BH and the accretion rate, as follows. The temperature $T$ of the accretion disc, as a function of distance $r$ from the BH, is given by

$$T(r) = \left(\frac{3}{8\pi} \frac{G M_{\text{BH}} M_{\text{disc}}}{\sigma_{\text{SB}} r^2}\right)^{\frac{1}{4}},$$

(4)

where $M_{\text{BH}}$ is the mass of the BH, $\dot{M}_{\text{BH}}$ the accretion rate and $\sigma_{\text{SB}}$ the Stefan–Boltzmann constant (Pringle 1981). For simplicity, we have taken the disc to be a thin disc, such that each annulus radiates as a blackbody of temperature given by the above equation. Integrating the flux over the surface of the disc, from $r_{\text{inner}}$ to $r_{\text{outer}} = 10^4 r_{\text{inner}}$, at which distance the contributions to both the photodissociating and ionizing fluxes are negligible, we find the total emitted flux as a function of the BH mass and its accretion rate. We take for the inner-most radius of the disc

$$r_{\text{inner}} \sim 2 \text{km} \left(\frac{M_{\text{BH}}}{M_\odot}\right),$$

(5)

corresponding to a high value for the BH spin parameter $a \gtrsim 0.9$ (e.g. Makishima et al. 2000; Vierdayanti, Watari & Mineshige 2008), which we expect considering that at its formation the BH accretes stellar material which can have a large angular momentum (see e.g. Volonteri & Begelman 2010).

Fig. 7 provides several examples of the spectra that we obtain using this treatment, for different choices of the BH mass and accretion rate. For a given BH mass, a higher accretion rate results in a higher bolometric luminosity and also in emission at higher energies. However, for accretion of gas at a given fraction of the Eddington rate, the peak of the spectrum is at higher energy for a less massive BH, owing to the smaller inner radius of the accretion disc at which the temperature is higher.

Recent observations of the X-ray emission from an accreting BH with a mass $10^5 \lesssim M_{\text{BH}} \lesssim 10^8 \, M_\odot$ (Farrell et al. 2009), bracketing the range of BH masses that we consider in the present study, suggest that the bulk of the emission at energies $\gtrsim 1 \text{ keV}$ can be accounted for by the presence of an accretion disc radiating as a blackbody, as we assume here, while above this energy a power-law component is shown to provide an appropriate fit to the observed spectrum. For simplicity, we do not include a power-law component in the spectrum emitted from the accreting BH. Including one, however, would not likely make a significant difference in the strength of the radiative feedback on the gas within the host halo, as the optical depth $\tau(\nu)$ for photons with energy $h\nu$ is

$$\tau(\nu) \sim \sigma(\nu) \Omega_\odot M_{\text{halo}} / (4\pi r_{\text{vir}}^2 m_H) \sim 10^{-2} (h\nu/1 \text{ keV})^{-3},$$

(6)

where $\sigma(\nu) \sim 4 \times 10^{-20} \text{ cm}^2 (h\nu/0.1 \text{ keV})^{-3}$ is the photoionization cross-section for neutral primordial gas (e.g. Kuhlen & Madau 2005), $\Omega_\odot$ is the cosmological baryon fraction, $M_{\text{halo}} \sim 3.5 \times 10^7 \, M_\odot$ is the total mass of the halo, $r_{\text{vir}} \sim 600 \, \text{pc}$ is the virial radius of the halo and $m_H$ is the mass of the hydrogen atom. At $h\nu \gtrsim 1 \text{ keV}$, $\tau \lesssim 10^{-2}$; hence, most of these photons would not interact with the gas in the halo and would contribute little to the photoionization of the gas. We point out, though, that this does have implications for our estimate of the escape fraction of ionizing photons, as discussed in Section 5.4.

### 4.3 Radiative feedback on the gas

To model the propagation of the high energy radiation emitted from the accretion disc, we use a modified version of the ray-tracing routine described in Johnson et al. (2009) and Greif et al. (2009). We refer the interested reader to those works for a detailed description of the scheme that we use to treat the propagation of ionizing radiation by solving, at every time-step, for the $H\alpha$ and $H\beta$...
regions surrounding the accreting BH; here we will briefly describe the improvements and modifications that have been made to that scheme.

The main modification in the calculation of the photoheating and photoionization rates is that here we adopt a multi-colour disc spectrum appropriate for a BH accretion disc, as described in Section 4.2, instead of a stellar spectra as in previous work. We integrate over the accretion disc radius \( r \) and over the frequency of the thermal radiation \( \nu \) emitted from the disc to calculate the photoionization and associated photoheating rates for both the photoionization of both \( \text{H} \) and \( \text{He} \), as well as to obtain the total numbers of ionizing photons, \( Q_{\text{H}} \) and \( Q_{\text{He}} \), emitted per second from the disc for each of these reactions, respectively.

As well, we estimate the photodissociation rate of \( \text{H}_2 \) by finding the total flux emitted from the accretion disc at 12.87 eV (e.g. Abel et al. 1997). We account for the self-shielding of \( \text{H}_2 \) in our ray-tracing routine by calculating the column density of molecules along each ray and using the prescription provided by Draine & Bertoldi (1996); however, we find that the column densities realized in our simulations are too low for self-shielding to have an effect in lowering the photodissociation rate.

Finally, as follows, we account for the transfer of momentum to the gas due to the speed of light in the outer regions of the disc. As this could effectively shield a substantial portion of the surrounding gas from the high energy radiation, the overall strength of the radiative feedback on the gas in our simulations is likely to be an overestimate.

\[ P_{\text{rad},i} \sim \frac{L_{\text{acc}} \sigma T n_{\text{elec},i} V_i^4}{4 \pi r_i^2 c}, \quad (7) \]

where \( L_{\text{acc}} \) is the total luminosity emitted from the accretion disc, \( n_{\text{elec},i} \) is the electron number density of the gas represented by the SPH particle, \( r_i \) is the distance from the BH to the particle, \( c \) is the speed of light and \( \sigma T \) is the Thomson cross-section. The approximate volume \( V_i \) of gas represented by the particle is estimated as \( V_i \sim M_i / \rho_i \), where \( M_i \) is the mass of the particle and \( \rho_i \) is its density.

The radiation pressure due to photoionization is estimated as

\[ P_{\text{ion},i} \sim \frac{(E_i) (k_{\text{H},\text{II}} n_{\text{H},\text{II}} + k_{\text{He},\text{II}} n_{\text{He},\text{II}}) V_i^2}{c}, \quad (8) \]

where \( (E_i) \) is the average energy of ionizing photons emitted from the accretion disc, \( n_{\text{H},\text{II}} \) and \( n_{\text{He},\text{II}} \) are the number densities of \( \text{H}_{\text{II}} \) and \( \text{He}_{\text{II}} \), and \( k_{\text{H},\text{II}} \) and \( k_{\text{He},\text{II}} \) are the photoionization rates of \( \text{H}_{\text{II}} \) and \( \text{He}_{\text{II}} \). These two pressures added together yield the total radiation pressure on particle \( i \), \( P_{\text{rad},i} = P_{\text{elec},i} + P_{\text{ion},i} \). In general, we find that the photoionization pressure is much greater than the electron scattering pressure, due to the non-negligible fractions of \( \text{H}_{\text{II}} \) and \( \text{He}_{\text{II}} \) in photoionized regions.

We then calculate the acceleration, in the direction opposite to that of the BH, of SPH particles lying within the photoionized region as

\[ a_{\text{rad},i} \sim \frac{P_{\text{rad},i} V_i^2}{M_i}, \quad (9) \]

which is expressed more precisely as

\[ a_{\text{rad},i} = (m_i c)^{-1} \left[ \frac{L_{\text{acc}} \sigma T f_{\text{elec}}}{4 \pi r_i^2} + (E_i) (k_{\text{H},\text{II}} n_{\text{H},\text{II}} + k_{\text{He},\text{II}} n_{\text{He},\text{II}}) \right], \quad (10) \]

where \( f_{\text{elec}} \) is the fraction of the free electron fraction, \( \text{H} \) fraction and \( \text{He} \) fraction, respectively. Finally, the effect of the radiation pressure is realized by imparting a velocity kick of magnitude \( \delta v_{\text{rad},i} \) to the SPH particle in the direction opposite to the BH:

\[ \delta v_{\text{rad},i} = a_{\text{rad},i} \delta t, \quad (11) \]

where \( \delta t \) is the length of the time-step taken in the simulation, over which the radiative pressure is applied.

We emphasize that in this treatment we have assumed that the radiation is emitted isotropically from the BH accretion disc, and in this we have not accounted for the absorption of ionizing radiation in the outer regions of the disc. As this could effectively shield a substantial portion of the surrounding gas from the high energy radiation, the overall strength of the radiative feedback on the gas in our simulations is likely to be an overestimate.

5 RESULTS ON BLACK HOLE ACCRETION AND RADIATIVE OUTPUT

In this section, we present the results of our simulations of BH accretion and the associated radiative feedback. The evolution of the gas within the host halo, the accretion rate of gas on to the BH, the radiative output from the accretion disc and the escape of ionizing radiation from the host halo are all closely interconnected. We discuss each of these in turn.

Using as initial conditions the output from the cosmological simulations described in Section 2 after the creation of the central sink particle, which we take to represent the accreting BH, we ran three separate simulations, each with a different initial BH mass. These initial BH masses are \( M_{\text{BH,init}} = 10^4 \), \( 2.5 \times 10^4 \) and \( 5 \times 10^4 M_\odot \).

5.1 Evolution of the gas

With the start of BH accretion, the radiative output from the accretion disc dramatically impacts the state of the surrounding gas. Figs 8, 9 and 10 show the properties of the gas, for the simulations with initial BH masses of \( 5 \times 10^4 \), \( 2.5 \times 10^4 \) and \( 10^5 M_\odot \), respectively, after 3 Myr since the formation of the BH. Comparing each of these figures, it is clear that accretion on to the more massive BHs leads to stronger radiative feedback on the gas. To illustrate the main effects, we shall thus focus on the case of the most massive BH, shown in Fig. 8.

As shown in the top panels in Fig. 8, there is a two-phase medium that develops within \( \sim 100 \) pc of the BH, with the gas that is photoionized having, in general, lower densities (\( \lesssim 10 \) cm\(^{-3} \)) and higher temperatures (\( \gtrsim 10^5 \) K), while the more dense gas that is shielded from the ionizing radiation remains at \( \lesssim 10^4 \) K. The photoheated gas experiences a boost in both gas pressure and radiation pressure, which result in the gas being pushed outwards at speeds up to \( v_r \sim 80 \) km s\(^{-1} \), as shown in the bottom-left panel.

As shown in the bottom-middle panel of Fig. 8, the gas pressure is generally much higher than the radiation pressure, except in photoionized regions where the fractions of \( \text{H} \) and \( \text{He} \) are relatively high, as it is here that the photoionization rates per unit volume are highest (see equation 8). In particular, this occurs at \( \sim 100–200 \) pc from the BH, as shown in the panels on the right, where the radiation pressure therefore plays an important role in driving the gas out from the centre of the halo. Interestingly, at distances \( \gtrsim 100 \) pc from the BH we find that the temperature of the ionized gas jumps to \( \lesssim 10^5 \) K; this occurs where the \( \text{He} \) fraction is highest, leading to an enhanced photoheating rate through the reaction \( \text{He} \; \gamma \rightarrow \text{He} \; + e^- \). This increased heating rate leads to temperatures high enough that hydrogen is almost completely collisionally ionized, leading in turn to a decreased cooling rate by atomic hydrogen;
Figure 8. Properties of the gas at the end of the simulation, 3 Myr after the formation of the BH, as a function of physical distance from the BH, for the case of a $5 \times 10^4 \, M_\odot$ initial BH mass, with the contours denoting the relative mass fraction of all of the gas shown. Clockwise from top-left: number density of hydrogen nuclei, temperature, neutral hydrogen (HI) fraction, singly ionized helium (He II) fraction, gas and radiation pressure, and radial velocity. Strong photoheating and photoionization pressure drive the gas out from the centre of the halo at up to $\sim 80 \, \text{km s}^{-1}$, although some relatively cold neutral gas remains in the vicinity of the BH, as it is shielded from the photoionizing radiation. While the radiation pressure is higher than the gas pressure where the neutral fraction is highest (see equation 8), in general the gas pressure is dominant. In the photoionized region at $r \geq 100 \, \text{pc}$ from the BH, the primary coolant is He II instead of HI, yielding higher temperatures in this region than in the higher density central region of the halo.

Figure 9. Just as Fig. 8, but for the case of a $2.5 \times 10^4 \, M_\odot$ initial BH mass. The gas properties are similar to those found in the $5 \times 10^4 \, M_\odot$ case, although the temperature of the photoionized gas at large radii is somewhat lower in this case, as is the speed at which the gas is pushed outwards from the centre of the halo, largely owing to the initially much higher accretion rate and ionizing photon output of the more massive BH (see Figs 12 and 13).

in effect, the gas approaches thermal equilibrium with He II as the primary coolant at $r \gtrsim 100 \, \text{pc}$, whereas HI is the primary coolant closer to the BH where the gas density is higher. A similar effect has been reported by Meiksin, Tittley & Brown (2010).

5.2 Accretion rate and Eddington fraction

The accretion rate, as described in Section 4.1, is directly coupled to the properties of the gas in the vicinity of the BH. Fig. 11 shows the evolution of the temperature $T_{\text{Bondi}} = 3 m_H c_s^2 / k_B$ and density $n_{\text{Bondi}} = \rho_{\text{gas}} / m_H$ in the vicinity of the BH that are used to compute the accretion rate $\dot{M}_{\text{BH}}$ (see equation 3), as a function of the time since the formation of the BH. Comparing the top and bottom panels, it is clear that the temperature of the accreted gas varies, at most, by a factor of $\lesssim 2$, whereas the density of the gas decreases by at least an order of magnitude in all cases, owing to the radiative feedback on the gas. Therefore, we conclude that it is largely the evolution of the density which dictates the evolution of the accretion rate.

Indeed, as shown in Fig. 12, at the outset of the simulations the accretion rates are higher for the more massive BHs, and the Eddington fractions are comparable in all three cases. However, as the hydrodynamic response of the gas to the radiative feedback develops and the density drops, the accretion rates drop dramatically. The accretion rate of the $5 \times 10^4 \, M_\odot$ BH plummets by two orders of magnitude over the 3 Myr time-scale of the simulation. The hydrodynamical response is not as swift in the other two cases, and the $10^4$ and the $2.5 \times 10^4 \, M_\odot$ BHs experience brief periods of accretion at near the Eddington rate before the radiative feedback on the gas drives the gas density down, ultimately resulting in the accretion rate dropping by an order of magnitude. For the majority of the time, all three BHs accrete at well below 10 per cent of the Eddington rate.
Figure 10. Just as Fig. 8, but for the case of a $10^4 M_\odot$ initial BH mass. Only a small fraction of the gas near the BH are ionized in this case, owing to the lower number of ionizing photons emitted from the accretion disc (see Fig. 13). As a result, the density structure of the gas closest to the BH is less disrupted than in the cases of the more massive BHs (see also Fig. 14).

Figure 11. The density (top panel) and the temperature (bottom panel) in the vicinity of the BH that are used to compute the Bondi accretion rate, as functions of time, for the three different cases of BH mass in our study. The temperature varies much less than the density over the 3 Myr of our simulations; hence, it is largely variations in the density of the accreted gas that affect changes in the accretion rate (see Fig. 12) and in the strength of the associated radiative feedback (see Fig. 13).

5.3 Ionizing and H$_2$-dissociating photon output

We integrate over the spectrum emitted from the accretion disc to compute the total number of ionizing photons emitted per unit time. The numbers of H i- and He ii-ionizing photons emitted per unit time, $Q_{\text{H i}}$ and $Q_{\text{He ii}}$, are defined here to be the total number of photons emitted with energies above 13.6 and 54.4 eV, the ionization thresholds of H i and He ii, respectively. The numbers of ionizing photons emitted per second are shown in the top two panels of Fig. 13. Comparing Figs 12 and 13, it is clear that the ionizing photon production rate is strongly correlated with the accretion rate.

Indeed, we see the same trends in the production rate of ionizing radiation as for the accretion rate, with the more massive BHs initially emitting photons at a higher rate and the overall emission rates dropping by at least an order of magnitude for all three BH masses.

Due to the hard spectra emitted from the accretion discs (see Fig. 7), for all three initial BH masses, the production rate of He ii-ionizing photons is almost as high as that of H i-ionizing photons, as shown in the top two panels of Fig. 13. Indeed, the fraction $Q_{\text{He ii}}/Q_{\text{H i}}$ is considerably higher than is expected for even very...
massive Pop III stars (see e.g. Bromm, Kudritzki & Loeb 2001; Schaerer 2002). As we discuss further in Section 6, this translates into a distinctive observational signature of the accretion of primordial gas on to BHs in the early Universe.

In the bottom panel of Fig. 13 is shown the flux of molecule-dissociating (LW) radiation \( J_{1W} \), in the standard units of \( 10^{-21} \text{ erg s}^{-1} \text{ cm}^{-2} \text{ Hz}^{-1} \text{ sr}^{-1} \), normalized to a physical distance of 1 kpc from the BH. Interestingly, the photodissociating flux is almost always lower than the background flux assumed in our simulation (\( J_{1W} = 10^3 \)) by at least an order of magnitude, at the fiducial distance of 1 kpc. Therefore, the radiative output from the BH significantly impacts only the gas at much smaller distances from the BH. However, the flux near the centre of the halo, where the most dense gas resides, is easily high enough to suppress the cooling of gas by \( H_2 \), and so to effectively suppress Pop III star formation.

Fig. 14 highlights the effects that the ionizing radiation have on the state of the gas (see also Figs 8–10). The more massive BHs, initially emitting more ionizing photons, have a stronger dynamical impact on the gas, as shown in the left-hand panels, causing more disruption of the dense gas surrounding the BH. In turn, the lower gas densities lead to lower recombination rates in the central dense regions, allowing the ionizing radiation to propagate further and heat the gas at larger distances from the BH. As well, the limited amount of ionizing radiation generated by the \( 10^4 \) M\(_\odot\) BH is evident in the bottom panels, as the \( H\alpha \) region is bottled up at small radii and thus the radiation affects only a small region surrounding the BH.

Also evident in Fig. 14 is the anisotropic expansion of the \( H\alpha \) regions surrounding the BHs, which arise from variations in the density of the gas in the centre of the halo. Even a relatively small difference in the density of the gas encountered by the ionization front as it moves out from the central BH can translate into a large difference in the recombination rate, which scales as density squared, and so in the rate at which the ionizing photons are absorbed by the gas. Therefore, as the ionization front encounters a large range of densities of the gas at a give radius (see Figs 8–10) as it expands outward, in directions in which the density is lower it extends out further than in directions in which the density is higher. This is a generic feature of \( H\alpha \) regions formed in a cosmological density field, and has also been found for the case of \( H\alpha \) regions generated by stars (see e.g. Alvarez, Bromm & Shapiro 2006; Abel, Wise & Bryan 2007).

### 5.4 Escape fraction of ionizing photons

Interwined with the emission of ionizing radiation from the accretion disc and the hydrodynamic response of the gas to the radiation is the fraction of ionizing photons which are able to escape from the host halo. To calculate this fraction, at every time-step we add up the total number of photons which propagate beyond 1 physical kpc from the BH; as this is well beyond the virial radius of the host halo (~600 pc) where the gas densities approach those of the intergalactic medium (IGM) (see Fig. 14), this provides a good estimate of the fraction of ionizing photons that can escape and contribute to the reionization the general IGM.

Fig. 15 shows the escape fraction \( f_{\text{esc}} \) of ionizing photons for the two most massive BHs; no ionizing radiation escapes the host halo in the case of the \( 10^4 \) M\(_\odot\) BH, the \( H\alpha \) region surrounding this BH being confined to the central regions of the halo, as shown in Figs 10 and 14. We find that only a small fraction, never exceeding \( \sim 0.1 \), of the ionizing photons are able to escape the host halo; indeed, the time-averaged escape fraction is \( \lesssim 10^{-2} \), for all the cases we studied.

We emphasize, however, that the escape fractions that we calculate here do not account for photons with energies higher than those emitted from the thermal multi-colour disc from which we derive our source spectrum. In particular, X-ray photons with energies e.g. \( \gtrsim 1 \) keV, such as may be emitted in a non-thermal power-law spectrum, could easily escape from the host halo and contribute to the reionization and photoheating of the IGM (e.g. Ricotti & Ostriker 2004; Kuhlen & Madau 2005; Volonteri & Gnedin 2009). However, even accounting for the long mean-free path of such higher energy emission, we expect that the majority of ionizing photons are still emitted from the thermal disc that we consider (see the observed spectra in e.g. Farrell et al. 2010), and hence that the error in our calculation of the escape fraction is not large.

While the ionizing photon production rate (see Fig. 13) is, at least initially, comparable to what could be produced by a cluster...
Accretion on to black holes

Figure 14. The projected number density (left-hand column), density-weighted temperature (middle column) and density-weighted H II fraction (right-hand column) of the gas in the vicinity of the accreting BH after 3 Myr since the formation of the BH, for the three cases of BH mass that we consider: $10^4 \, M_\odot$ (bottom row), $2.5 \times 10^4 \, M_\odot$ (middle row) and $5 \times 10^4 \, M_\odot$ (top row). The radiation emitted from the accretion discs of the more massive BHs has a stronger impact on the gas in the halo, leading to generally higher temperatures and to expansion of the gas causing expansion and disruption of the dense gas in the centre of the halo from which the BH feeds (see also Figs 8, 11–13). Note also that more ionizing radiation escapes the halo for the cases of the more massive BHs (see Fig. 15).

of very massive Pop III stars in the centre of a similar halo, the escape fraction that we find is much lower than in the case of such a stellar cluster, which can approach unity (Johnson et al. 2009; see also Wise & Cen 2009). This is a product of the self-regulation of the radiative output from accreting BHs, as the high energy radiation they emit heats and disperses the gas, in turn limiting the accretion rate and the emission of radiation. In the case of stellar clusters, once they are formed their radiative output is not coupled to the properties of the surrounding gas, allowing them to emit copious amounts of ionizing radiation even after dispersing the gas surrounding them; this leads to generally higher escape fractions, one of the most extreme cases being that of massive Pop III stars formed in minihaloes (see e.g. Whalen, Abel & Norman 2004; Alvarez, Bromm & Shapiro 2006).
6 AN OBSERVATIONAL SIGNATURE OF BLACK HOLE ACCRETION

In this section, we investigate a key signature of the observable radiation emitted from haloes hosting BHs accreting primordial gas in the early Universe: the strength of the emission in several prominent recombination lines.

6.1 Recombination line luminosities

The luminosity emitted in a given recombination line is given by (e.g. Schaerer 2002; Osterbrock & Ferland 2006)

\[ L = \alpha(1 - f_{\text{esc}})Q, \]  

(12)

where \( \alpha \) is the energy emitted in the line for each recombination, \( Q \) is the number of photons emitted per unit time that ionize the chemical species which produces the emission, and \( f_{\text{esc}} \) is the fraction of these photons that escape the halo. For our case, we use either \( Q_{\text{HI}} \) or \( Q_{\text{HeII}} \), shown in Fig. 13. Also, we do not differentiate between the escape fraction of \( \text{H}\alpha \) and He II-ionizing photons, as we find that the \( \text{H}\alpha \) and the He II regions surrounding the BH are almost completely overlapping, due to the high ratio of \( Q_{\text{HeII}} / Q_{\text{H}\alpha} \); in any case, the escape fraction is so low as to not be critical for our estimate of the recombination line luminosities.

The luminosities that we find, using the values for \( \alpha \) given in Schaerer (2002), for Ly\( \alpha \), H\( \alpha \), and He II \( \lambda 1640 \) are shown in Fig. 16. Due to the low values of \( f_{\text{esc}} \), the luminosities we find here mirror almost exactly the ionizing photon production rates shown in Fig. 13; in turn, variations in the luminosities also closely track those in the BH accretion rate.

The luminosities in He II \( \lambda 1640 \) relative to those in H\( \alpha \) or Ly\( \alpha \), for all three cases, are extremely high. The ratios of the observable flux in He II \( \lambda 1640 \) to that in H\( \alpha \), to which the IGM is optically thin even before reionization is complete, are shown in Fig. 17. In all cases, this flux ratio exceeds \(~2\), substantially higher than that expected even for very massive Pop III stars (i.e. \( \gtrsim 100 \ M_{\odot} \)); see e.g. Johnson et al. 2009). This suggests that a distinct observable signature of BHs accreting primordial gas in the early Universe could be such extremely high flux ratios, ultimately due to the very hard spectra emitted from the BH accretion disc.

\[ F_{1640} = \frac{L_{1640}}{4\pi D_{L}^{2}} \sim 10^{-20} \text{ erg s}^{-1} \text{ cm}^{-2} \left( \frac{L_{1640}}{10^{40} \text{ erg s}^{-1}} \right) \left( \frac{1+z}{10} \right)^{-2} \]  

(13)

where \( L_{1640} \) is the luminosity emitted in the line and \( D_{L}(z) \) is the luminosity distance to redshift \( z \).

The flux limit for \( 3\sigma \) detection of the He II \( \lambda 1640 \) line emitted from a source at \( z \sim 10 \), with 100 h of observation by the Near-Infrared Spectrograph (NIRSpec) that will be aboard the JWST, is of the order of \( 10^{-19} \text{ erg s}^{-1} \text{ cm}^{-2} \). Therefore, at the peak luminosities in this line, shown in Fig. 16, emission due to both the \( 2.5 \times 10^{4} \) and the \( 5 \times 10^{4} \ M_{\odot} \) BHs would be marginally detectable at \( z \lesssim 10 \). Given a small optical depth to Ly\( \alpha \) through the IGM, or to sufficiently high velocity of the emitted gas relative to the IGM.

\( \text{http://www.stsci.edu/jwst/science/sensitivity/} \)
For $M_\alpha$ during the time that the BHs are accreting; using an average BH, $M_\alpha \sim 2 \times 10^9 M_\odot$. Gas accretes on to a central SMS of the order of $0.1 M_\odot$. The ratio $M_\alpha$ used both in the $J$-band $F_{\nu,J}$, as shown in Fig. 7.

When we track the growth of these BHs for only $3 \text{ Myr}$ in the present work, we can extrapolate our findings to roughly estimate how the BH to host halo mass ratio may evolve at later times. We find in our simulations that the host DM halo grows at a rate of $0.5 M_\odot \text{ yr}^{-1}$ during the time that the BHs are accreting; using an average BH accretion rate from Fig. 12, $\sim 10^{-3} M_\odot \text{ yr}^{-1}$, along with this rate, we find that the ratio $M_\text{BH}/M_\text{halo}$ may drop to $\sim 10^{-4}$ over $\sim 10^3 \text{ yr}$.

Another important question with regard to detecting BHs formed by direct collapse is the cosmological environment in which they are most likely to be contained. The competition between accretion and the high LW background flux needed for the formation of such BHs, Dijkstra et al. (2008) argue that the host haloes of these BHs may lie within $\sim 10$ physical kpc of another halo with mass $\sim 10^9 M_\odot$. Thus, it may be the case that these haloes are found very close to other, more luminous star-forming galaxies. However, this distance and the proximity of the BH to the host halo is sensitive to the mass of the BH, such that the emission due to more massive BHs than considered here will likely be more easily detected.

We note, however, that the accretion rates needed for these BHs could also be detectable at the same redshift. However, due to lower instrument sensitivity at the relevant wavelengths and to its lower luminosity, the flux in the He II line would be too low for detection.

Another important question is whether the accretion rates are zero for $3 \text{ Myr}$ after the formation of the BHs; at later times, with the growth of the host halo by continuous accretion and mergers, the BHs could accrete gas at high rates again. Also, as we have shown here, the luminosity of the radiation emitted due to the accretion process is sensitive to the mass of the BH, such that the emission due to massive BHs from mergers is more easily detected.

The BH by direct collapse is a high ratio of these fluxes, although this is less valid for the ratios of the less massive BHs, owing to the harder spectrum of radiation emitted from their accretion discs, for a given accretion rate. For instance, Ferrarese (2002) suggests a relation between the host halo mass and BH mass of $\frac{M_\text{BH}}{M_\text{halo}} \lesssim 10^{-5}$ for $M_\text{halo} \lesssim 10^{12} M_\odot$, whereas the ratio we have in the present work is two orders higher than this at the time of BH formation. Furthermore, from their simulations of BH growth, Booth & Schaye (2010) find that $M_\text{BH} \sim 10^9 M_\odot (\frac{M_\text{halo}}{10^{13} M_\odot})^{0.55}$, hence, the $\gtrsim 10^4 M_\odot$ BHs that we consider forming in $\lesssim 10^9 M_\odot$ haloes would lie roughly four orders of magnitude above the relation that these authors find.

Due to the strong radiative feedback from the BHs, we find that star formation is not likely to take place in the central dense regions of the host halo, unless the BH accretion rate drops even further than we find in our simulations or the host halo is able to grow massive enough that some gas can become shielded from the photodissociating radiation emitted from the accretion disc (see also Oh & Haiman 2002). If star formation does occur, the accretion rate may be boosted by the transfer of angular momentum outward due to supernova explosions, thereby driving more gas to the centre of the halo (see Chen et al. 2009; Kumar & Johnson 2010). If star formation does not occur, then BHs formed by direct collapse may remain unaccompanied by stars in their host halo, at least until the host halo merges with another halo in which stars may have already formed. Indeed, it may be that the host halo ultimately merges with the halo or haloes in which reside the sources of radiation that are responsible for the high background LW radiation field that led to the formation of the BH to begin with (see e.g. Dijkstra et al. 2008).

8 SUMMARY AND DISCUSSION

We have presented cosmological simulations of the formation and growth of BHs formed by direct collapse in an atomic-cooling halo at $z \sim 15$. In these simulations, the accretion of gas and the detailed interaction of the gas with the high energy radiation emitted in the accretion process were modelled self-consistently, allowing us to arrive at a number of novel conclusions, which are summarized below.

During the growth of the host halo prior to the formation of the BH by direct collapse, we find little evidence for supersonic ‘cold flows’ into the centre of the halo, due to the inability of the gas to cool under the influence of the elevated photodissociating background radiation field, which is a likely prerequisite for this scenario of BH formation. None the less, in the absence of radiative feedback, we find an upper limit to the rate at which the warm ($\gtrsim 10^4$ K) gas accretes on to a central SMS of the order of $0.1 M_\odot \text{ yr}^{-1}$, high enough for it to grow to a mass of $\gtrsim 10^5 M_\odot$ before it collapses to form a BH.

With the formation of the BH and a hot accretion disc, photo-heating of the gas in the host halo, as well as the radiative pressure to which it is subjected, leads to its expansion, which in turn leads to accretion rates which generally fall with time. While in all three

7 IMPLICATIONS FOR GALAXY AND BLACK HOLE COEVOlUTION

Numerous observations of high-redshift galaxies suggest that the central BHs they host are more massive than would be expected from the observed BH–host galaxy relations at lower redshift (e.g. Walter et al. 2004; McLure et al. 2006; Peng et al. 2006; Bennert et al. 2010; Merloni et al. 2010; Greene, Peng & Ludwig 2010; Decarli et al. 2010). It is interesting to note that the expected mass range of BHs (i.e. $10^5$–$10^6 M_\odot$) formed by direct collapse in atomic-cooling haloes is at least qualitatively in line with this trend.
cases of initial BH mass that we simulate the accretion rate initially approaches the Eddington limit, within the 3 Myr of our simulations the accretion rate drops by at least an order of magnitude. The time-averaged accretion rates that we find are \( \lesssim 10 \) per cent of the Eddington rate.

In part due to these low accretion rates, the time-averaged ionizing photon escape fractions are very low (\( \lesssim 0.01 \)); this suggests that BHs formed by direct collapse contribute little to the reionization of the IGM in the vicinity of their host haloes. These low escape fractions, in combination with the hard spectra emitted from the accretion disc, lead to an extremely high ratio of the nebular recombination radiation emitted in He \( \alpha \lambda 1640 \) to that emitted in H\( \alpha \) or Ly\( \alpha \). This is a distinctive observational signature of BH accretion in the early Universe, which may be detected by the JWST. Another potential observational signature is that BHs formed by direct collapse may not be accompanied by stars formed in their host halo, due to the strong radiative feedback from the accretion process. Finally, the large ratio of the masses of these BHs to the masses of their host haloes may provide some insight into the origin of observed BHs at high redshift which have higher masses, for a given host halo mass, than is observed in the local Universe.

We emphasize that the accretion rates that we report here are likely to be upper limits, for two reasons. The first is that we do not consider the radiative feedback from the SMS that may form as the immediate predecessor to the massive BH (see e.g. Bromm & Loeb 2003; Volonteri & Begelman 2010), which may act to decrease the accretion rate of the BH once it forms by driving gas away from the centre of the halo (see e.g. Johnson & Bromm 2007; Alvarez et al. 2009). Secondly, as mentioned in Section 4.1, not all of the gas that is accreted at the scale of the Bondi radius is likely to be accreted on to the BH, as we have assumed.

However, we also note that at scales below those that we resolve in our simulations the radiative feedback on the gas may be limited by shielding of the gas in the shadow of the accretion disc of the BH, as mentioned in Section 4.3. Such a reduction in the radiative feedback on the accreting gas would likely allow for somewhat higher accretion rates. Thus, while both the strength of the radiative feedback and the accretion rates that we find here may be overestimated separately, how each of these would change in response to more detailed modelling of the other remains to be fully understood.

To better estimate the accretion rates of BHs in the early Universe, higher resolution cosmological simulations which follow the formation and evolution of the accretion disc around the BH would be of great interest.

The sub-Eddington accretion rates that we find may pose a significant challenge to models that call for the formation of \( \gtrsim 10^8 \, M_\odot \) BHs by \( z \sim 6 \) by accretion of gas on to BHs formed by direct collapse at \( 10 \lesssim z \lesssim 15 \), and also for models which envision similar massive BHs as forming from Pop III stars fuelled by dark matter annihilation (see e.g. Freese et al. 2010). It is important to note, though, that the single host halo that we have considered in this study represents only a \( \sigma \) fluctuation of the cosmological density field, while the most massive BHs at \( z \sim 6 \) likely reside in much more rare, faster growing haloes (e.g. Barkana & Loeb 2001). Therefore, while radiative feedback is certain to hinder the accretion of gas even in such rare overdense regions to some extent, the faster growth of the host halo would undoubtedly lead to generally higher average accretion rates (see Li et al. 2007). However, such a rare, fast growing halo would likely have a progenitor which hosted the formation of a Pop III star at very early times (i.e. at \( z \gtrsim 30 \)), before a substantial LW background radiation field could have been established (see e.g. Johnson et al. 2008). While this would likely preclude the formation of a BH by direct collapse in the same halo at a later time, such Pop III relic BHs could well be the seeds for the most massive BHs observed to power high-redshift quasars.

ACKNOWLEDGMENTS

We are grateful to the support staff of the SFC supercomputer cluster at the Rechenzentrum Garching of the Max Planck Society, on which the simulations presented here were carried out. JLJ would also like to thank Mitch Begelman, Dominik Schleicher and Zoltán Haiman for helpful discussions, as well as the anonymous referee for comments which improved the clarity of this work.

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