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Photon trapping enables super-Eddington growth of black hole seeds in galaxies at high redshift

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ABSTRACT

We identify a physical mechanism that would have resulted in rapid, obscured growth of seed supermassive black holes in galaxies at $z \gtrsim 6$. Specifically, we find that the density at the centre of typical high-redshift galaxies was at a level where the Bondi accretion rate implies a diffusion speed of photons that was slower than the gravitational infall velocity, resulting in photons being trapped within the accretion flow and advected into the black hole. We show that there is a range of black hole masses ($M_{bh} \sim 10^{3-5} M_\odot$) where the accretion flow traps radiation, corresponding to black holes that were massive enough to generate a photon trapping accretion flow, but small enough that their Bondi radii did not exceed the isothermal scale height of self-gravitating gas. Under these conditions we find that the accretion reaches levels far in excess of the Eddington rate. A prediction of this scenario is that X-ray number counts of active galactic nuclei at $z \gtrsim 6$ would exhibit a cutoff at the low luminosities corresponding to black hole masses below $\sim 10^5 M_\odot$. The super-Eddington growth of $\sim 10^5 M_\odot$ seed black holes at high redshift may have provided a natural acceleration towards the growth of supermassive black holes at $z \sim 6-7$, less than a billion years after the big bang.

Key words: galaxies: formation – galaxies: high-redshift – galaxies: nuclei – quasars: general – cosmology: theory.

1 INTRODUCTION

Exceptionally bright quasars with redshifts up to $z \sim 7$ have been discovered (e.g. Fan et al. 2001, 2003; Mortlock et al. 2011). These quasars are thought to be powered by the thin disc accretion of gas onto supermassive black holes at the centres of galaxies. Their maximum (Eddington) luminosity depends on the mass of the black hole, and the brighter quasars are inferred to have black holes with masses of more than a few billion solar masses. Since their discovery (Fan et al. 2001, 2003), the existence of such massive black holes at $z \gtrsim 6$ has posed a challenge to models for their formation. This is because an $\sim 10^9 M_\odot$ black hole accreting at the Eddington rate with a radiative efficiency of 10 per cent requires almost the full age of the Universe at $z \sim 6$ to grow from a stellar mass seed. Many authors have therefore discussed solutions to this apparent mystery by including a significant build-up of mass through mergers (Haiman & Loeb 2001), collapse of low-spin systems (Eisenstein & Loeb 1995) and suppression of molecular line cooling via a large Lyman–Werner flux (Dijkstra et al. 2008). Other authors have also attempted to explain the fast growth of black holes at high redshifts based on supermassive stars (Loeb & Rasio 1994; Bromm & Loeb 2003) and more recently, quasi-stars (Begelman 2010; Volonteri & Begelman 2010).

Volonteri & Rees (2005) discussed the possibility of super-Eddington accretion of black hole seeds in high-redshift galaxies. They pointed out that seed black holes are located at the centres of isothermal discs where the conditions for quasi-spherical Bondi accretion should be prevalent. At high redshift these disc centres are sufficiently dense that the Bondi accretion rate greatly exceeds the Eddington rate. Volonteri & Rees (2005) point out that this super-Eddington accretion provides a route by which a large fraction of the mass e-foldings needed to grow a supermassive black hole by redshift 6 could be accommodated within a small fraction of the age of the Universe. However, the calculations in Volonteri & Rees (2005) ignored feedback effects like gas heating, which may raise the sound speed, and hence lower the density and therefore the Bondi accretion rate. For example, Milosavljević et al. (2009) show that photoheating and radiation pressure from photoionization significantly reduce the steady-state accretion rate and potentially render a quasi-radial accretion flow unsteady and inefficient. They find that the time-averaged accretion rate is always a small fraction of the Bondi accretion rate. Thus, the very high accretion rates implied by the Bondi accretion in the centre of a high-redshift isothermal disc might never be reached. On the other hand, if the accretion rate is sufficiently high that the emergent photons are trapped within the
accretion flow, then these feedback effects cannot operate (Begelman 1979), and so the accretion rate can reach arbitrarily high levels.

In this paper we find that at sufficiently high redshift, the central densities of galaxies imply Bondi accretion rates that exceed the rate required to trap radiation and advect it into a black hole (Begelman 1979). Thus, we find that there were periods in the growth of black holes at high redshift where the growth was super-Eddington and feedback mechanisms could not halt the accretion flow. Our goal is not to make a self-consistent model for both the transport of material from large galactic radii and central black hole accretion. Rather, we note that there is a significant literature looking at the problem of super-Eddington accretion assuming large mass-delivery rates to the region of the black hole, and identify the cosmological conditions that would provide sufficient accretion rates to allow this by trapping radiation. We begin in Section 2 with a description of our simple model, before presenting our results in Section 3. We finish with a discussion in Section 4 and conclusions in Section 5. In our numerical examples, we adopt the standard set of cosmological parameters (Komatsu et al. 2011), with values of $\Omega_m = 0.04$, $\Omega_m = 0.24$ and $\Omega_{\Lambda} = 0.76$ for the matter, baryon and dark energy fractional density, respectively, $h = 0.73$, for the dimensionless Hubble constant and $\sigma_8 = 0.82$.

2 MODEL

The basis of this paper is a comparison between the Bondi accretion rate and the accretion rate required to trap photons within the accretion flow. We discuss these in turn.

2.1 The Bondi accretion rate

We begin with the expression for the Bondi accretion rate (Bondi 1952) on to a central black hole of mass $M_{\text{bh}}$,

$$M_{\text{Bondi}} = 4\pi r^2 v_{\text{ff}} \rho_{\text{Bondi}} = 4\pi r^2 v_{\text{ff}} \rho_{\text{Bondi}} \left( \frac{GM_{\text{bh}}}{r_{\text{Bondi}}} \right)^{1/2}$$

(1)

where $v_{\text{ff}}$ is the free-fall time at the Bondi radius

$$r_{\text{Bondi}} = \left( \frac{GM_{\text{bh}}}{c_s^2} \right)^{1/2}$$

(2)

and $c_s$ is the sound speed, which for an isothermal gas, we assume, corresponds to a temperature of $10^4$ K. To evaluate $M_{\text{Bondi}}$ we specify the central number density of a self-gravitating disc (Schaye 2004),

$$n_0 = \frac{GM_{\text{disc}}^2}{12\pi c_{\text{m}}^2 R_0^4 \mu_m m_p}$$

(3)

Here $R_0$ is the characteristic radius of an exponential disc of surface density profile,

$$\Sigma(r) = \Sigma_0 \exp(-r/R_0),$$

(4)

with $\Sigma_0 = M_{\text{disc}}/2\pi R_0^2$, and the disc scalelength is given by

$$R_0 = \sqrt{\frac{\lambda}{2\nu}} r_{\text{vir}},$$

(5)

where $\lambda$ is the dimensionless spin parameter of the halo. In equation (3), $M_{\text{disc}} = m_p M_{\text{halo}}$, is the disc mass, $m_p$ is the mass of hydrogen and $\mu = 1.22$ is the molecular weight of primordial neutral gas. At the high redshifts of interest, most of the virialized galactic gas is expected to cool rapidly and assemble into the disc. We therefore assume $m_p = 0.17$. The corresponding mass density is $\rho_0 = m_p n_0$. The virial radius of a halo with mass $M_{\text{halo}}$ is given by the expression

$$r_{\text{vir}} = 0.784 h^{-1} \text{kpc} \left( \frac{M_{\text{halo}}}{10^8 M_\odot h} \right)^{1/3} \left[ \frac{\nu(z)}{10} \right]^{-1},$$

where $\nu(z)$ is close to unity and defined as $\nu \equiv \left[ \frac{\Omega_m}{\Omega_m^*} \left( \Delta_c/18\pi^2 \right) \right]$, $\Omega_m \equiv \left[ 1 + (\Omega_{\Lambda}/\Omega_m) (1 + z)^{3} \right]^{-1}$, $\Delta_c = 18\pi^2 + 82d - 39d^2$ and $d = \Omega_m - 1$ (see equations 22–25 in Barkana & Loeb 2001, for more details). From equations (3) and (6) we see that the central density $n_0$ and hence the Bondi accretion rate $M_{\text{Bondi}}$ scales as $(1 + z)^{1/2}$, and as a result accretion rates are expected to be much larger at high redshift.

Dormant central black holes are ubiquitous in local galaxies (Magorrian et al. 1998). The masses of these supermassive black holes scale with the physical properties of their hosts (e.g. Magorrian et al. 1998; Merritt & Ferrarese 2001; Tremaine et al. 2002). However, at high redshift the relations observed in the local Universe may not be in place. We therefore do not impose a model for the relation between black hole and halo mass in this paper, and instead explore a range of values. Indeed, our results indicate that feedback, which is thought to drive the black hole–halo mass relation, would not be effective at early times. The grey curves in Fig. 1 show the Bondi accretion rate as a function of redshift for different values of halo and black hole mass. Here we assume $\lambda = 0.05$ corresponding to the mean spin parameter for dark-matter haloes (Mo, Mao & White 1998).

2.2 Photon trapping by the accretion flow

If the diffusion velocity of photons at a radius $r$ is smaller than the free-fall velocity of the material at radius $r$, then photons become trapped in the accretion flow (Begelman 1979). In such cases, the black hole would be obscured. In this section, we estimate the accretion rate required to achieve this photon trapping at radius $r$.

The free-fall time from radius $r$ to a smaller radius $r_{\text{min}}$ is

$$t_{\text{ff}} \sim \frac{(r - r_{\text{min}})}{v_{\text{ff}}},$$

(6)

where $v_{\text{ff}} \sim \sqrt{GM_{\text{bh}}/r}$ within the Bondi radius. This should be compared with the diffusion time of

$$t_{\text{diff}} \sim \frac{(r - r_{\text{min}})}{\mu},$$

(7)

Here the optical depth is given by

$$\tau \sim \frac{\sigma_T m_p}{m_p} (r - r_{\text{min}}),$$

(8)

where $\sigma_T$ is the Thomson cross-section, $F_{\text{sig}}$ is the ratio of the scattering cross-section of the accreting gas to the emergent radiation in units of $\sigma_T$ and $\mu$ is the line-of-sight averaged density between radii $r_{\text{min}}$ and $r$. The density within the Bondi radius scales as $\rho(r) = \rho_0 (r/r_{\text{Bondi}})^{-1.5}$, yielding

$$\mu \equiv \frac{2 \rho_0}{r_{\text{Bondi}}} \frac{r_{\text{ff}}}{r_{\text{min}}} \left[ \left( \frac{r_{\text{min}}}{r_{\text{Bondi}}} \right)^{-0.5} - \left( \frac{r}{r_{\text{Bondi}}} \right)^{-0.5} \right]$$

$$= 2 \rho_0 \frac{r_{\text{Bondi}}}{r_{\text{min}}} \left[ \left( \frac{r_{\text{min}}}{r_{\text{Bondi}}} \right)^{0.5} - 1 \right] \text{ for } r < r_{\text{Bondi}}.$$
in which we have used the relation $M_{\text{Bondi}} = 4\pi r^2 v_g$. We find that the accretion rate required for photon trapping has a minimum value (i.e. $\frac{dM_{\text{min}}}{dr} = 0$) at a radius of $r = 4r_{\text{min}}$, yielding a minimum required accretion rate for photon trapping of

$$\dot{M}_{\text{min}} = \left(\frac{8\pi \sigma T_c}{F_{\text{sig}} \sigma T_c}\right) r_{\text{min}}.$$  

Since the Eddington accretion rate at an efficiency $\epsilon$ is $\dot{M}_{\text{Edd}} = \frac{4\pi G M_{\text{bh}} M_p}{(c \sigma T_c \epsilon)}$, we find

$$\dot{M}_{\text{min}} = 2 \left(\frac{\epsilon}{F_{\text{sig}} \sigma T_c}\right) \frac{r_{\text{min}}}{r_g} \dot{M}_{\text{Edd}},$$

where $r_g = GM_{\text{bh}}/c^2$.

The limiting accretion rate depends on the emergent radiation spectrum. As described below in Section 3.2, we find that the condition for photon trapping is more readily achieved for X-ray photons than for optical/ultraviolet (UV) photons. We therefore frame our discussion around trapping of optical/UV photons. In the case of optical/UV photons the cross-section would be dominated by dust at radii larger than a sublimation radius $r_{\text{min}} = r_{\text{sub}}$, beyond which the opacity of the gas to optical/UV photons could be larger than the Thomson opacity by as much as two or three orders of magnitude due to the presence of dust (Laor & Draine 1993). The existence of dust can be questioned for high-redshift galaxies (e.g. Bouwens et al. 2010); however, high metallicities are inferred from the broad emission lines of all quasars out to $z = 7.1$ and so metal enrichment (due to star formation) is known to precede the growth of the black hole in the galactic nuclei of interest (Hamann & Simon 2010). On the other hand, the diffusion time may be lessened if the gas is clumpy, corresponding to a lower effective opacity. We therefore take a typical value for the cross-section that is $F_{\text{sig}} = 100$ times larger than the Thomson cross-section (additional cases are presented below in Section 2.3). Defining $r_{\text{min}} = F_{\text{min}} r_g$, we assume a typical value of $F_{\text{min}} = 10^3$ to describe the sublimation radius (Netzer & Laor 1993).

Limiting accretion rates corresponding to these default values are plotted in the left-hand panels of Fig. 1 (dark lines) for the case where the vertical structure of the disc is set by the sound speed of the gas. The upper and lower panels correspond to cases of halo masses of $M_{\text{halo}} = 10^9$ and $10^{10} M_\odot$. The larger value corresponds to the lower end of the inferred halo masses of the Lyman-break (Ly-break) population at $z \gtrsim 6$ (Trenti et al. 2010).

### 2.3 Disc structure

Before proceeding we first note that we have so far assumed the thickness of the gaseous disc at the centre of the galaxy to be set by the sound speed of gas at $10^8$ K. However, from equation (13) we see a strong dependence on the value for the effective sound speed.

In the dense regions at the centres of high-redshift galaxies, we expect that there may be rapid cooling and fragmentation of the gas, with associated star formation. Isothermal gas at $10^4$ K may
hole in a $10^{10} \, M_{\odot}$ halo will result in trapped radiation at $z > 4$. In this section, we discuss the range of halo and black hole masses that result in photon trapping.

### 3.1 Scaling relations

The conditions for the halo mass, black hole mass and redshift that conspire to provide accretion rates that trap the optical/UV radiation can be obtained by combining equations (1) and (10). Evaluating the Bondi radius, we get

$$c_T = \left( \frac{G \Sigma_{\text{bh}}}{\pi} \right)^{0.25}.$$  

This value of $c_T$ is the maximum value possible for a disc at large radius (as a higher $c_T$ would imply an unphysical disc with $h > r$). Therefore, for a $Q = 1$ disc the turbulent velocity should decrease towards small $r$. At the small radii corresponding to the Bondi radius we find $c_T < c_s$, implying that the centres of discs are likely to be denser and thinner than in the isothermal $10^4 \, \text{K}$ case. Thus, although the turbulent velocity can be significant at large radii in the disc, we find that when evaluated at the Bondi radius, the turbulent velocity is smaller than the sound speed unless the halo is very massive and at high redshift ($z > 4$).

On the other hand, there may also be sources of energy injection that heat the gas to temperatures that correspond to velocities much larger than $c_T$ (or $c_s$), although in order to remain bound the velocity of the gas must be smaller than the virial velocity of the halo.

We therefore show two cases in addition to $c_s \sim 10 \, \text{km s}^{-1}$, namely $c_T$ and $c_s \sim f v_{\text{inf}}$, where $f \sim 0.5$ in order to bracket the range of interest, with corresponding limiting accretion rates plotted in the central and right-hand panels of Fig. 1 (dark lines). In the case of turbulent velocity with $Q = 1$ we note that increased clumpiness of the gas would greatly lower the cross-section of the infalling gas. We choose $F_{\text{sig}} = 1$ as our fiducial example. Similarly, heated gas in the case of a disc height set by $c_s$, we expect lower clumping and so take $F_{\text{sig}} = 1000$. The results are sensitive to which velocity sets the disc height. However, the expected difference in the clumpiness modifies the expected cross-section $F_{\text{sig}}$ in such a way that partially counteracts the effect of the disc density. In the remainder of this paper, we show examples for both the cases of an isothermal disc and a turbulent $Q = 1$ disc, which we expect to bracket the range of likely physical conditions.

We also note that the gas discs in which the black holes are embedded are not static, but rotating. As a result, Bondi accretion does not strictly apply. Previous studies have shown that accretion flows with low angular momentum have an accretion rate that is lower than the Bondi rate (e.g. Proga & Begelman 2003; Moscibrodzka & Proga 2008). However, Narayan & Fabian (2011) showed that gas with sub-Keplarian rotation, as expected for gas at the centre of a large galactic disc of $10^7 \, \text{K}$ gas, accretes at a factor of only a few less than the Bondi rate. Thus, we neglect the disc rotation for the purposes of our simple calculations in this paper.

### 3 PHOTON TRAPPING

In cases where the Bondi accretion rate is larger than the critical accretion rate, rest-frame optical/UV photons are trapped, and the active galactic nucleus (AGN) can be obscured. Thus the cross-over of the limiting and Bondi accretion rate curves in Fig. 1 represents the redshift beyond which accretion traps radiation for an AGN powered by the particular black hole masses shown. If the vertical disc structure is set by the sound speed, we find that a $10^7 \, M_{\odot}$ black hole in a $10^{10} \, M_{\odot}$ halo will result in trapped radiation at $z > 4$. In this section, we discuss the range of halo and black hole masses that result in photon trapping.

#### 3.2 Trapping of X-rays

Up until now we have found that the critical rate at which the accretion flow traps optical/UV photons can be reached in the centres...

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of high-redshift galaxies owing to gas with a large effective opacity (described by $F_{\text{sig}} > 1$) beyond a radius $F_{\text{min}} r_g$. The results of Laor & Draine (1993) indicate that while $F_{\text{sig}} > 1$ is expected for UV photons travelling through smoothly distributed gas, the X-ray component of the spectrum will see an opacity to the inflowing gas that is set by the Thomson cross-section (i.e. $F_{\text{sig}, X} = 1$). Moreover, if the X-ray photon component of the spectrum were not trapped, it could halt the accretion flow via radiation pressure if it exceeded the Eddington rate by itself, thus preventing trapping of the optical/UV radiation. On the other hand, X-rays see this Thomson opacity at radii much smaller than the sublimation radius, with $F_{\text{min}} r_g < 1$, corresponding to the innermost stable circular orbit, we expect that an accretion flow which traps the X-rays encounter opacity ($r_{\text{min}, X} = F_{\text{min}, X} r_g$) is described by a value of $F_{\text{min}, X} > 50$. To see this we note that the ratio of the Bondi accretion rate to the rate needed to trap X-rays is

$$\frac{M_{\text{Bondi}}}{M_{\text{lim}, X}} = \left( \frac{F_{\text{min}}}{F_{\text{lim}, X}} \right) \left( \frac{1}{F_{\text{sig}}} \right) \frac{M_{\text{Bondi}}}{M_{\text{lim}}}.$$  \hspace{1cm} (17)

Since we expect $F_{\text{lim}, X} \sim 6$, corresponding to the innermost stable circular orbit, we expect that an accretion flow which traps the optical/UV radiation will also trap the X-rays. As a result we do not consider the effect of X-rays on the accretion flow for the remainder of this paper.

### 3.3 Critical spin parameter

Equation (13) shows that the trapping of photons is very sensitive to the value of the spin parameter $\lambda$ which governs the density of the galactic disc. In particular, photons may be more easily trapped within discs of low-spin parameter. The assembly of dark-matter haloes leads to a distribution of spin parameters. To illustrate the parameter space we recast equation (13). Photon trapping at all...
wavelengths occurs for spin parameters $\lambda < \lambda_c$, where
\[
\lambda_c \approx 0.12 \left( \frac{M_{\text{bh}}}{10^8 M_\odot} \right)^{1/4} \left( \frac{M_{\text{halo}}}{10^{10} M_\odot} \right)^{1/6} \left( \frac{1+z}{7} \right) \left( \frac{m_d}{0.17} \right)^{1/2} \\
\times \left( \frac{F_{\text{min}}}{10^3} \right)^{-1/4} \left( \frac{c_s}{10 \text{ km s}^{-1}} \right)^{-5/4} \left( \frac{F_{\text{sig}}}{10^5} \right)^{1/4}.
\] (18)

In the left-hand panels of Fig. 3 we plot contours of $\lambda_c$ as a function of black hole and halo mass at $z = 6$. Two scenarios are shown corresponding to an isothermal $10^4$ K gas having $c_s \approx 10$ km s$^{-1}$ with $(F_{\text{sig}}/10^2)(F_{\text{min}}/10^5)^{-1} = 1$ (upper panel) and a turbulent $Q = 1$ disc having $c_s = c_T$ with $(F_{\text{sig}}/10^2)(F_{\text{min}}/10^5)^{-1} = 0.01$ (lower panel). The examples assume $m_d = 0.17$. The mean spin parameter $\lambda = 0.05$ is shown in black, and the grey shaded regions illustrate the parts of parameter space where trapping is not possible for a mean disc. In the isothermal case, Fig. 3 shows that emergent photons from black holes with $M_{\text{bh}} \gtrsim 10^{11} M_\odot$ will be trapped at all wavelengths in the centres of high-redshift galaxies (\$10^8-10^9 M_\odot\$) that are at the mean of the spin-parameter distribution. Conversely, in the turbulent $Q = 1$ disc case we find that black holes smaller than $\sim 10^7 M_\odot$ would have their radiation trapped.

We note the differing behaviour between the isothermal and $Q = 1$ cases, both in terms of the reversed trend of $\lambda_c$ with $M_{\text{bh}}$ at fixed $M_{\text{halo}}$ and of $\lambda_c$ with $M_{\text{halo}}$ at fixed $M_{\text{bh}}$. This difference results from the fact that in the $Q = 1$ case there is a dependence of the effective sound speed (which we have evaluated at the Bondi radius) on the black hole mass, whereas in the isothermal case there is no dependence. As can be seen from equations (12) and (18), larger values of $M_{\text{bh}}$ result in larger values of sound speed, and hence smaller values of $\lambda_c$. This dependence dominates the effect of $M_{\text{bh}}$ on the value of $\lambda_c$ over the explicit $M_{\text{bh}}^{1/4}$ dependence in equation (18), leading to the reversed behaviour relative to the isothermal case.

### 3.4 Breakout of the Bondi radius

In the previous sections, we have illustrated that the large densities expected at the centres of high-redshift galaxies lead to conditions where accretion rates may be sufficient to trap photons within the accretion flow. These calculations are based on the density at the centre of a pressure supported self-gravitating disc. In this section, we point out that the calculation applies only to black hole masses for which the Bondi radius is smaller than the scale height at the disc centre. Moreover, we note that the trapping of radiation cannot be realized once the black hole grows sufficiently that its Bondi radius exceeds the scale height.
We again utilize the self-gravitating disc model. The scale height at the disc centre is

$$z_0 = \frac{c_s^2}{\pi G \Sigma_0}.$$  

(19)

This expression ignores the gravitational contribution from the black hole which would serve to reduce the scale height, and is valid when $r_{\text{Bondi}} < z_0$. Photon trapping and obscuration are only possible in this regime. We therefore calculate the ratio of the Bondi radius to the central disc scale height as

$$r_{\text{Bondi}} \approx \frac{G \Sigma_0 m_\bullet}{c_s^2 \lambda_0^2 R_{\text{mt}}^3}.$$  

(20)

Putting in characteristic values we obtain

$$r_{\text{Bondi}} \approx 2.6 \left( \frac{M_{\bullet}}{10^5 M_\odot} \right) \left( \frac{M_{\text{halo}}}{10^{10} M_\odot} \right)^{1/3} \left( \frac{1 + z}{7} \right)^2 \times \left( \frac{c_s}{10 \text{ km s}^{-1}} \right)^{-4} \left( \frac{\lambda}{0.05} \right)^{-2}.$$  

(21)

Thus, within $M \sim 10^5 M_\odot$ haloes at $z \sim 6$, black holes in excess of $M_{\bullet} \sim 10^5 M_\odot$ have Bondi radii which are larger than the disc scale height and so may not be observed. We note that the ratio is very sensitive to the value of the sound speed, with a value larger than the fiducial $10 \text{ km s}^{-1}$ significantly reducing the ratio, allowing larger black holes to accrete in the photon trapping mode. A small gas fraction also reduces the ratio. However, the breakout of the Bondi radius implies that high-mass black holes such as those observed within the SDSS quasars (e.g. Fan et al. 2001, 2003) could not have their emergent radiation trapped.

To better understand the constraints imposed on black hole masses where photon trapping can occur, we evaluate the critical value of the spin parameter $\lambda_{c, B}$ at which $r_{\text{Bondi}} = z_0$,

$$\lambda_{c, B} \sim 0.08 \left( \frac{M_{\bullet}}{10^5 M_\odot} \right)^{1/2} \left( \frac{M_{\text{halo}}}{10^{10} M_\odot} \right)^{1/2} \left( \frac{c_s}{10 \text{ km s}^{-1}} \right)^{1/2} \left( \frac{m_\bullet}{0.17} \right)^{1/2}.$$  

(22)

In the central panels of Fig. 3 we plot contours of $\lambda_{c, B}$ as a function of the black hole and halo mass at $z = 6$. As before, two scenarios are shown corresponding to an isothermal $10^5$ K gas having $c_s \sim 10 \text{ km s}^{-1}$ with $(F_{\text{sig}}/10^5)(F_{\text{min}}/10^5)^{-1} = 1$ (upper panel) and a turbulent $Q = 1$ disc having $c_s = c_T$ with $(F_{\text{sig}}/10^5)(F_{\text{min}}/10^5)^{-1} = 0.01$ (lower panel). We assume $m_\bullet = 0.17$. The mean spin parameter $\lambda$ is 0.05 is shown in black, and the grey shading illustrates the region of parameter space where photon trapping is not possible for the mean disc because the Bondi radius is larger than the scale height. Fig. 3 shows that we expect black holes with $M_{\bullet} \gtrsim 10^5 M_\odot$ within high-redshift galaxies ($\sim 10^{10} M_\odot$) having the mean of the spin parameter cannot form a photon trapping accretion flow.

### 3.5 When and where could photon trapping occur?

The right-hand panels of Fig. 3 show the probabilities that the spin parameter $\lambda$ lies between the two critical values needed for the conditions of (i) photon trapping and (ii) a Bondi radius that is contained within the disc scale height (i.e. $\lambda_{c, B} < \lambda < \lambda_c$). To calculate this probability

$$P = \frac{1}{\sqrt{2\pi \sigma_\lambda}} \int_{\lambda_{c, B}}^{\lambda_c} \exp \left[ -\frac{(\ln \lambda - \ln \lambda_c)^2}{2\sigma_\lambda^2} \right] d\lambda,$$  

(23)

we assume that the distribution of spin parameters is Gaussian in the natural logarithm $\ln \lambda$, with variance $\sigma_\lambda = 0.5$ and a mean at $\lambda = 0.05$ (Mo et al. 1998). The probability is $P = 0$ if $\lambda_{c, B} < \lambda_{c, B}$, indicating a black hole–halo mass combination that cannot produce a photon trapping accretion flow. Contours are shown that represent black hole–halo mass combinations for which $P = 1$, 10 and 50 per cent of discs would have densities that result in photon trapping. As before, two examples are shown corresponding to an isothermal $10^5$ K gas having $c_s \sim 10 \text{ km s}^{-1}$ with $(F_{\text{sig}}/10^5)(F_{\text{min}}/10^5)^{-1} = 1$ (upper panels) and a turbulent $Q = 1$ disc having $c_s = c_T$ with $(F_{\text{sig}}/10^5)(F_{\text{min}}/10^5)^{-1} = 0.01$ (lower panels). We assume $m_\bullet = 0.17$. The mean disc at $z \sim 6$ has a central density that leads to photon trapping for black hole masses up to $M_{\bullet} \sim 10^5 M_\odot$ within haloes up to $M_{\text{halo}} \sim 10^5 M_\odot$.

Thus, we would expect photon trapping to be common in galaxies hosting $M_{\bullet} \lesssim 10^5 M_\odot$ black holes at $z \sim 6$. Moreover, since from equation (13) we see that this growth is super-Eddington, we find that photon trapping provides a mechanism by which rapid black hole growth could proceed at high redshift, helping to explain how supermassive black holes grow less than a billion years after the big bang.

In Fig. 4 we explore how the conclusions regarding the black hole and halo mass ranges that produce photon trapping accretion flows are affected by redshift and the parameter combination $(F_{\text{sig}}/10^5)(F_{\text{min}}/10^5)^{-1}$. We choose two values of redshift, $z = 6$ and 10, and assume the case of a disc structure set by isothermal $10^5$ K gas with $m_\bullet = 0.17$. These cases are shown in the upper and lower rows, respectively. In each case we show examples with $(F_{\text{sig}}/10^5)(F_{\text{min}}/10^5)^{-1} = 0.3, 1$ and 3 (left-hand, central and right-hand panels, respectively). In all panels, we show contours for the probabilities that the spin parameter $\lambda$ lies between the two critical values needed for the conditions of (i) photon trapping and (ii) a Bondi radius that is contained within the disc scale height (i.e. $\lambda_{c, B} < \lambda < \lambda_c$). As before, contours are shown that represent black hole–halo mass combinations for which $P = 1$, 10 and 50 per cent of discs would have central densities that lead to photon trapping.

In Fig. 5 we mirror the results of Fig. 4, this time for the case of a turbulent $Q = 1$ disc. Here, lower values of $(F_{\text{sig}}/10^5)(F_{\text{min}}/10^5)^{-1} = 0.003, 0.01$ and 0.03 are used (left-hand, central and right-hand panels, respectively), reflecting our assumption that the cross-section of the gas will be much lower than in the isothermal case.

In both scenarios we find that larger values of $(F_{\text{sig}}/10^5)(F_{\text{min}}/10^5)^{-1}$ lead to more massive black holes in smaller haloes having their emergent radiation trapped. At redshift $z \sim 6$–10 we find that black holes with masses up to $M_{\bullet} \sim 10^9 M_\odot$ have their radiation trapped in 10–50 per cent of cases. Thus, photon trapping in high-redshift ($z \gtrsim 6$) galaxies is likely to be dominated by low-mass ($M_{\bullet} \lesssim 10^5 M_\odot$) black holes within $M_{\text{halo}} \sim 10^5 M_\odot$ haloes.

### 3.6 Comparison with $M_{\bullet} - M_{\text{halo}}$ models

As noted in Section 1, the relations observed between black hole mass and galaxy properties in the local Universe may not be in place at high redshift. For this reason, we have not imposed a model for the relation between black hole and halo mass in this paper, and instead have explored a range of values. However, it is interesting to compare the range of black hole mass founds to be accreting in the photon trapping mode with expectations of the simple models.
relating black hole and halo masses that have been successful in describing some of the properties of high-redshift quasars.

Motivated by local observations (Ferrarese 2002), we consider a model in which the central black hole mass is correlated with the halo circular velocity. This scenario is supported by the results of Shields et al. (2003) who studied quasars out to \( z \sim 3 \) and demonstrated that the relation between black hole masses and the stellar velocity dispersion does not evolve with redshift. This is expected if the mass of the black hole is determined by the depth of the gravitational potential well in which it resides, as would be the case if growth is regulated by feedback from quasar outflows (e.g. Silk & Rees 1998; Wyithe & Loeb 2003). Expressing the halo virial velocity, \( v_c \), in terms of the halo mass, \( M_{\text{halo}} \), and redshift, \( z \), the redshift-dependent relation between the supermassive black hole and halo masses may be written as

\[
M_{\text{bh}} = \epsilon_{\text{bh}} M_{\text{halo}} \left( \frac{M_{\text{halo}}}{10^{12} M_\odot} \right)^{2/3} \left[ \xi(z) \right]^{5/6} \left( \frac{1 + z}{7} \right)^{5/2}.
\]

(24)

The normalizing constant in this relation has an observed value of \( \epsilon_{\text{bh}} \approx 10^{-4.3} \) based on calibration of equation (24) at \( z = 0 \) (Ferrarese 2002), where we take the underlying assumption that the halo mass profile resembles a singular isothermal sphere. In Wyithe & Loeb (2005), this model was shown to be consistent with the clustering and luminosity function data of the 2dF quasar redshift survey (Croom et al. 2002).

Interestingly, the predicted value for supermassive black hole masses in galaxy mass haloes (as shown by the thick grey curves in Figs 4 and 5) is comparable to the range in which we would expect photon trapping in galaxies with \( z \gtrsim 6 \). Thus, prior to obtaining masses where the self-regulating mechanisms thought to be responsible for the relations between black hole mass and halo properties take effect we would expect black holes in high-redshift galaxies to reach the level where super-Eddington accretion would become a natural part of their evolution.

### 4 DISCUSSION

The findings in this paper point to some potentially important implications for the growth of high-redshift supermassive black holes. Recent simulations by Li (2011) show that self-gravity overcomes radiative feedback and that accretion on to intermediate mass black holes reaches the Eddington rate. However, at high enough accretion rates, a spherical inflow is not subject to the Eddington limit (though the emergent radiation is), and the Bondi accretion rate can take arbitrarily large values in high-density galactic centres (Begelman 1979). The conditions for spherical accretion may only rarely be realized at the centres of high-redshift galaxies. On the other hand, there is a class of slim disc models (e.g. Abramowicz et al. 1988; Chen et al. 1995; Ohsuga et al. 2005; Watarai 2006;...
Figure 5. Contours of the probability in black hole–halo mass parameter space that accretion will result in photon trapping, corresponding to the spin parameter $\lambda$ lying between the two critical values, i.e. $\lambda_c < \lambda < \lambda_c^\prime$. In this figure we assume the gas to have $Q = 1$ with a disc structure set by turbulent velocity. Contours are shown for $P = 1, 10$ and $50$ per cent. Examples are shown for two values of redshift, $z = 6$ and $10$ (top and bottom). We choose $m_g = 0.17$. In each case we show examples with $(\text{F}_{\text{sig}}/10^2)(\text{F}_{\text{min}}/10^3)^{-1} = 0.003, 0.01$ and $0.03$ (from left to right). The thick grey curve is the predicted value for the black hole mass–halo mass relation (equation 24).

Ohsuga & Mineshige (2011), for which super-Eddington accretion flows are possible. These models include a large viscosity, and so rapid transport of material through the disc within a small multiple of the free-fall time (e.g. Watarai 2006), and are optically thick and advection dominated (Chen et al. 1995). Photon trapping plays an important role within these discs, even those with complex flows, leading to very inefficient energy conversion and a luminosity that is independent of accretion rate (e.g. Ohsuga et al. 2005). Simulations indicate that at sufficient densities, the mass accretion rate can reach hundreds of Eddington, although the photon luminosity does not (e.g. Ohsuga & Mineshige 2011). These super-Eddington accretion discs will generate outflows (which cannot be generated in spherical accretion; Begelman 1979) that are likely to be channelled along the poles, and so will not suppress the accretion along the equatorial plane. The calculations presented in this paper provide the boundary condition for super-Eddington flows. These rapid accretion events may provide the seeds for supermassive black hole growth.

A plausible scenario therefore includes two episodes of black hole growth. Initially the accretion may have been spherical or a slim disc with high viscosity, so that the Eddington rate was greatly exceeded by the accretion rate. We find that radiation would have been trapped and advected into the central black hole by the very large Bondi accretion rates at high redshift. This phase of growth would be obscured at both optical/UV and X-ray wavelengths. Once the Bondi radius became larger than the scale height of the galaxy disc, the accretion rate dropped, and higher angular momentum gas would have settled into a thin disc accretion mode, in which the accretion time was much longer than the free-fall time, allowing the radiation to escape. For a pressure supported disc, we have found that the Bondi radius is expected to exceed the scale height of the disc for black holes in excess of $\sim 10^5 M_{\odot}$. The luminous quasars discovered at $z \gtrsim 6$ with black hole masses in excess of $10^8 M_{\odot}$ are therefore thought to be shining in this later mode, unobscured in the optical and with accretion rates close to, but smaller than Eddington. However, a portion of their prior growth, when the black hole mass was $\sim 10^4$–$5 M_{\odot}$, would have been in the photon trapping mode.

Super-Eddington accretion rates arising from the large densities in the cores of high-redshift galaxies were previously considered by Volonteri & Rees (2005). As noted in Section 1, the calculations in Volonteri & Rees (2005) neglected feedback effects like gas heating, which may lower the Bondi accretion rate (e.g. Milosavljević et al. 2009). In this paper, we note that if the accretion rate is sufficiently high that the emergent photons are trapped within the accretion flow, then feedback effects cannot operate (Begelman 1979). Our model also makes two different assumptions which modify our conclusions relative to Volonteri & Rees (2005). First, Volonteri & Rees (2005) assume that once the gas is enriched, metal line cooling allows the gas to cool to temperatures much lower than $10^4 K$, so that it fragments to form stars and the super-Eddington accretion episode is ended. This assumption was necessary in order that super-Eddington accretion might not lead to black hole densities in excess of those observed. However as noted earlier, simulations by Hopkins et al. (2012) and Joung et al. (2009) show that feedback from star
formation causes a clumpy, self-regulated two-phase medium which has a Toomre’s Q of approximately unity. Observationally, Genzel et al. (2010) inferred a Toomre’s Q value of $Q = 1$ in ULIRGS. Thus, in this case the growing black hole would be embedded in the centre of a turbulent pressure supported disc in which Bondi accretion can continue. Rather than making an arbitrary assumption that growth stops when the size of the accretion disc grows by a factor of 5, we instead assume that the super-Eddington photon trapping accretion mode would be regulated by the time when the Bondi radius exceeds the scale height of the pressure supported disc. We find that this condition prevents super-Eddington accretion on to high-mass black holes.

4.1 Obscured accretion

Our results have some relevance to the recent discussion surrounding obscured accretion in high-redshift galaxies. While luminous optical quasars in the most massive haloes ($\gtrsim 10^{12} \, M_{\odot}$) dominate the observations of high-redshift supermassive black holes, Treister et al. (2011) recently presented evidence that most of the black hole accretion at $z \gtrsim 6$ is actually optically obscured, and in galaxies below halo masses of $\sim 10^{10} \, M_{\odot}$. Since high-redshift Ly-break galaxies are thought not to be dusty (e.g. Bouwens et al. 2010), our results might have provided a mechanism by which the AGN can be obscured even in the absence of a large dusty component. However, our results do not support a photon trapping explanation for this result. First, X-ray photons would be more strongly trapped by the accretion flow than the UV photons, indicating that we would not expect to observe X-rays without optical detection. In addition, we also find that photon trapping is only expected in a fraction of galaxies rather than in all galaxies as implied by Treister et al. (2011).

While other mechanisms may lead to buried accretion, the findings of Treister et al. (2011) have also been disputed by a number of authors (Cowie, Barger & Hasinger 2012; Fiore, Puccetti & Mathur 2011; Willott 2011), who do not observe the same level of X-ray emission. These authors find limits on the average X-ray luminosity in the rest-frame 0.5–2 keV band of $L_{X,0.5–2} < 4 \times 10^{41} \, erg \, s^{-1}$ for $z \sim 6.5$ dropout galaxies. This luminosity can be related to black hole mass as

$$L_{0.5–2} \sim 3 \times 10^{41} \left( \frac{M_{bh}}{10^7 M_{\odot}} \right) \eta \, erg \, s^{-1},$$

where $\eta = 1$ is the fraction of the Eddington accretion rate, and we have assumed the spectral energy distribution of Elvis et al. (1994). Observed luminosities must have $\eta < 1$ even if the accretion rate is super-Eddington (Begelman 1979). The observed luminosity limit therefore corresponds to observed black hole masses of $M_{bh} \lesssim 1.3 \times 10^5 \, M_{\odot}$. We do not find that black holes with masses $M_{bh} \gtrsim 10^5 \, M_{\odot}$ produce photon trapping accretion flows, and so photon trapping does not explain the lack of observed X-ray sources among the $z \sim 6.5$ dropouts (Cowie et al. 2012; Fiore et al. 2011; Willott 2011). Since the stacked observations in these studies are based on only $\sim 10^2$ galaxies, this lack of detection could follow from the low duty cycle of AGN (which is likely a few per cent; Wyithe & Loeb 2003). However, our results do suggest that 90–100 per cent of discs with black holes below this mass would be in the photon trapping mode. Thus, we would expect that deeper and wider field X-ray observations using future X-ray observatories, to display a cutoff in the X-ray luminosity function at about $L_{0.5–2} \sim 3 \times 10^{41} \, erg \, s^{-1}$. Conversely, the discovery of X-ray AGN with luminosities an order of magnitude lower than current limits would therefore rule out photon trapping accretion as a mechanism for rapid growth of early black holes. Finally, we note that since super-Eddington accretion at high redshift is obscured at both optical and X-ray wavelengths, rapid growth of seed black holes could not provide a significant source of X-ray for reionization of the IGM (Volonteri & Rees 2005).

4.2 Seed black hole growth

The photon trapping mode is likely to be important for the rapid growth of seed supermassive black holes with masses of $\sim 10^{6–7} \, M_{\odot}$. Because the Bondi accretion rates in these high-redshift galaxies could be orders of magnitude larger than the Eddington rate, the photon trapping mechanism helps alleviate the difficulty of growing supermassive black holes of more than a billion solar masses (corresponding to the most distant quasars; Mortlock et al. 2011) within the first billion years of the Universe’s age. This point was made in detail in Volonteri & Rees (2005). To illustrate we note that accretion at the Eddington rate (with $\epsilon = 0.1$) leads to an e-folding time of $t = 4 \times 10^7 \, yr$. Assuming that a black hole accretes with a duty cycle of unity, the number of e-folding times available by $z \sim 7$ is therefore $\sim 20$. This should be compared with the 20 e-folds needed to grow a $1 M_{\odot}$ black hole seed up to a mass of $10^7 \, M_{\odot}$. Thus, there is only just enough time during the age of the Universe at $z \gtrsim 6$ for a stellar black hole seed to grow to a supermassive black hole. We suggest that the period of obscured growth in some galaxies would provide a path towards growing these supermassive black holes.

We can make an estimate of the time-scale over which photon trapping and hence super-Eddington accretion might occur. From Figs 4 and 5 we see that super-Eddington accretion may grow the black hole by $N \sim 4–10$ e-folds in mass. Combining this result with equation (16) we find that this growth will occur over a time-scale of

$$t_{\text{SE}} \sim 4 \times 10^7 \, yr \left( \frac{N}{4} \right) \left( \frac{M_{\text{BH}}}{10^7 M_{\odot}} \right)^{-1} \left( \frac{M_{\text{halo}}}{10^7 M_{\odot}} \right)^{-2/3} \left( \frac{1 + z}{7} \right)^{-4} \left( \frac{m_{\text{g}}}{0.17} \right)^{-2} \left( \frac{c_s}{10 \, \text{km \ s}^{-1}} \right)^{5/4} \left( \frac{\lambda}{0.05} \right)^{4}.$$  

This value is comparable to the free-fall time from the characteristic radius of a galactic disc at $z \sim 6–10$ (e.g. Wyithe & Loeb 2003), and an order of magnitude less than the lifetime of massive stars. Compared with the age of the Universe at $z \sim 6$, equation (26) implies that the duty cycle of super-Eddington accretion is only $\sim 0.5$ per cent.

5 CONCLUSION

In this paper we have determined the cosmological regime in which photons produced through accretion on to a central black hole are trapped by infalling material, so that (i) radiation feedback on the infall of gas outside the dust sublimation radius is suppressed, allowing accretion rates far in excess of the Eddington limit, (ii) AGN appear obscured and (iii) the growth time of black hole is short. Specifically we find that a large fraction of galaxies at $z \gtrsim 6$ with masses up to those of the observed Ly-break population (halo masses of $\sim 10^{10–11} \, M_{\odot}$) exhibit Bondi accretion rates on to $M_{\text{BH}} \sim 10^{1–3} \, M_{\odot}$ black holes that are sufficiently high to trap the resulting photons.

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rest-frame optical/UV/X-ray radiation. The obscuration due to photon trapping is found only to occur for black holes with masses up to $\sim 10^5 \, M_\odot$ because larger black holes have Bondi radii that exceed the scale height of the disc from which they accrete gas, so that a photon trapping mechanism cannot operate. As a result, we find a natural distinction between obscured, photon trapping accretion on to $\sim 10^5 \, M_\odot$ black holes in galaxies of halo mass $\lesssim 10^{10} \, M_\odot$ and the luminous accretion seen in the brightest quasars with black hole masses of $\sim 10^8-10^9 \, M_\odot$ within haloes of mass $\sim 10^{11}-10^{12} \, M_\odot$.

Our results indicate that super-Eddington accretion of mass to form seed black holes of $\sim 10^5 \, M_\odot$ provided a mechanism by which supermassive black holes were able to form prior to $z \sim 6$.

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