QUASARS AT z = 6: THE SURVIVAL OF THE FITTEST

MARTA VOLONTERI AND MARTIN J. REES

Institute of Astronomy, University of Cambridge, Madingley Road, Cambridge CB3 0HA, UK

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ABSTRACT

The Sloan Digital Sky Survey has detected luminous quasars at very high redshift, z > 6. Follow-up observations indicate that at least some of these quasars are powered by supermassive black holes (SMBHs), with masses in excess of $10^9 M_\odot$. SMBHs, therefore, seem to have already existed when the universe was less than 1 Gyr old and the bulk of galaxy formation had yet to take place. Here we investigate the extent to which accretion and dynamical processes influence the early growth of SMBHs. We assess the impact of (1) black hole mergers, (2) the influence of the merger condition, and (3) the negative contribution due to dynamical effects, which can kick black holes out of their host halos (gravitational recoil). We find that if accretion is always limited by the Eddington rate via a thin disk, the maximum allowed radiative efficiency (or spin) to reproduce the luminosity function at $z = 6$ is $\epsilon = 0.12$ (or $\alpha = 0.8$), against the adverse effect of the gravitational recoil. Dynamical effects unquestionably cannot be neglected in studies of high-redshift SMBHs. If black holes can accrete at a supercritical rate during an early phase, reproducing the observed SMBH mass values is not an issue, even in the case that the recoil velocity is in the upper limit range, as the mass ratios of merging binaries are skewed toward low values, where the gravitational recoil effect is very mild. We propose that SMBH growth at early times is very selective, and efficient only for black holes hosted in high density peak halos.

Subject headings: black hole physics — cosmology: theory — galaxies: evolution — quasars: general

1. INTRODUCTION

It seems a challenge for theoretical models to explain the luminosity function of luminous quasars at $z \approx 6$ in the Sloan Digital Sky Survey (SDSS; Fan et al. 2001a). The luminosities of these quasars, well in excess of $10^{47}$ ergs s$^{-1}$, imply that supermassive black holes (SMBHs), with masses $10^9 M_\odot$, were already in place when the universe was only 1 Gyr old. The accretion of mass at the Eddington rate causes the black hole (BH) mass to increase over time as

$$M(t) = M(0) \exp \left( \frac{1 - \epsilon}{\epsilon} \frac{t}{t_{\text{Edd}}} \right), \quad (1)$$

where $t_{\text{Edd}} = 0.45$ Gyr and $\epsilon$ is the radiative efficiency. Among the proposed seed BHs, those more commonly invoked (e.g., Population III star remnants, gravitationally collapsed star clusters) have masses in the range $M(0) = 10^2 - 10^4 M_\odot$, forming at $z = 30$ or less. Given $M(0)$, the higher the efficiency, the longer it takes for the BH to grow in mass by (say) 20 e-foldings (Shapiro 2005). For a Schwarzschild black hole, the standard thin disk radiative efficiency is $\epsilon \approx 0.06$, and there is plenty of time for the BH seed to become supermassive. The timescale to grow from $M(0) = 10^2 - 10^4 M_\odot$ to $\simeq 10^9 M_\odot$ is less than 0.5 Gyr.

If accretion is via a geometrically thin disk, though, the alignment of a SMBH with the angular momentum of the accretion disk tends to efficiently spin holes up (Volonteri et al 2005), and radiative efficiencies can therefore approach 30%–40%. With such a high efficiency, $\epsilon = 0.3$, it takes more than 2 Gyr for the seed to grow to $10^9$ solar masses. It therefore seems difficult to reproduce the observational constraints without invoking exotic processes.

We take here a conservative approach and try to determine the parameter space that allows SMBH growth compatible with observational constraints at $z = 6$, as evinced from the luminosity function of quasars (Fan et al. 2004) and observations of galaxies at $z \approx 6$. In this paper, we critically assess models for the early evolution of SMBHs, exploring the parameter space of the involved processes: accretion rate, radiative efficiency, dynamical processes, and the initial density of massive black hole (MBH) seeds.

At $z < 5$, MBH mergers do not play a fundamental role in building up the mass of SMBHs (Yu & Tremaine 2003), but they may be important at $z > 5$, where we do not have constraints from a Sohn-type argument, which compares the local MBH mass density with the mass density inferred from luminous quasars, as the luminosity function of quasars is not constrained at $z > 6$. Mergers can possibly contribute positively to the buildup of the high-redshift SMBH population (Yoo & Miralda-Escude 2004), as they contribute to make a large black hole from many small seeds.

On the other hand, dynamical processes can disturb the growth of BHs, especially at high redshift (Haiman 2004; Yoo & Miralda-Escude 2004), and make a negative contribution to SMBH growth. In the shallow potential wells of minihalos, the growth of MBHs can be halted by the “gravitational rocket” effect, the recoil due to the nonzero net linear momentum carried away by gravitational waves in the coalescence of two black holes. Also, if MBHs are widespread and binary black holes’ coalescence timescales are long enough for a third MBH to fall in and interact with the central binary, Newtonian three-body interactions can lead to the expulsion, or recoil, of the binary. The accretion history must then be studied jointly with the dynamics involving MBHs.

Yoo & Miralda-Escude (2004) explored the minimum conditions that would allow the growth of seed MBHs up to the limits imposed by the highest-redshift quasar currently known: SDSS J1148+5251. This quasar, at $z = 6.4$, has estimates of the SMBH mass in the range $(2 - 6) \times 10^9 M_\odot$ (Barth et al. 2003; Willott et al. 2003). Yoo & Miralda-Escude showed that the mass of this SMBH can be explained assuming (1) continued Eddington-limited accretion onto MBHs forming in halos with $T_{\text{vir}} > 2000$ K at $z \leq 40$; their model assumes, also, (2) a large
influence of BH mergers in increasing the MBH mass: a contribution by itself on the order of $10^9 M_\odot$. Their investigation takes into account the negative feedback due to the “gravitational rocket” effect (see also Haiman 2004). Recent estimates suggest a more modest recoil velocity, compared with the typical values adopted in the past. We test here how much difference (if any) these new estimates imply for the growth of black holes in pre-galactic halos.

It is also important to understand where $z = 6$ quasars are hosted. By matching the number density in halos more massive than $M_h$ to the space density of quasars at $z = 6$ (Fan et al. 2004), one obtains $M_h = 10^{13} M_\odot$. This assumption corresponds to requiring that the duty cycle of high-redshift quasars is unity, that is, all the BHs inhabiting halos with mass larger than $10^{13} M_\odot$ are active. The only available observation of a quasar host (Walter et al. 2004) shows interesting features. The kinematics of the observed molecular gas implies the lack of a massive bulge around the SMBH but suggests a dark matter halo with mass similar to that of the predicted bulge, on the order of $10^{12} M_\odot$. This is much less than $M_h = 10^{13} M_\odot$ and allows for a much smaller duty cycle. The uncertainties on the dynamical configuration of the gas (e.g., inclination) are still large, however.

In the next section, we review the model for assembly of MBHs in cold dark matter (CDM) cosmogonies. We then review models for MBH growth by accretion in high-redshift halos ($\S$ 3) and discuss the dynamical evolution of MBH binaries in pre-galactic systems ($\S$ 4). The results from the interplay of accretion and dynamical processes are presented in $\S$ 5. Finally, we discuss the implications for the global evolution of the SMBH population ($\S$ 6). All results below refer to a $\Lambda$CDM world model with $\Omega_M = 0.3$, $\Omega_\Lambda = 0.7$, $h = 0.7$, $\Omega_b = 0.045$, $\sigma_8 = 0.93$, and $n = 1$.

## 2. Assembly of Pre-galactic MBHs

The main features of a plausible scenario for the hierarchical assembly, growth, and dynamics of MBHs in a $\Lambda$CDM cosmology have been discussed in Volonteri et al. (2003, 2005) and Madau et al. (2004). Dark matter halos and their associated galaxies undergo many mergers as mass is assembled from high redshift to the present. The halo merger history is tracked backward in time with a Monte Carlo algorithm based on the extended Press-Schechter formalism. “Seed” holes are assumed to form with intermediate masses in the rare high-$\nu$-$\sigma$ peaks collapsing at $z = 20-25$ (Madau & Rees 2001) as end products of the very first generation of stars.

As our fiducial model we take $\nu = 4$ at $z = 24$, corresponding to a mass density parameter in MBHs on the order of $\rho_\star \approx 10^2 M_\odot$ Mpc$^{-3}$, about 3 orders of magnitude smaller than the local SMBH mass density (Aller & Richstone 2002; Merloni et al. 2004; Marconi et al. 2004; Yu & Tremaine 2002). The mass density in seed MBHs cannot be larger than the mass density of SMBHs at $z = 3$, $\rho_{\text{MBH}}(z = 3) \approx 4.5 \times 10^4 M_\odot$ Mpc$^{-3}$ (Merloni 2004).

This choice of seed MBH occupation fraction is similar to that in Volonteri et al. (2003; seed holes in 3.5 $\sigma$ peaks at $z = 20$), ensuring that galaxies hosted in halos with mass greater than $10^{11} M_\odot$ are seeded with a MBH. The assumed threshold allows efficient formation of SMBHs in the range of halo masses effectively probed by dynamical studies of SMBH hosts in the local universe. We will also consider a case in which seed holes are more numerous (e.g., 25 times more than the fiducial case).

We are here interested in the evolution of the most very massive halos present at $z = 6$, that is, very highly biased structures (Diemand et al. 2005). As a consequence, the density of seed holes within the volume occupied by the progenitors of the halo is large enough that the merging of two minihalos both hosting a BH is not a rare event. This is in contrast with the average density of seed holes. In a cosmic volume, the seed holes typically evolve in isolation (cf. Madau et al. 2004).

We generate Monte Carlo realizations (based on the extended Press-Schechter formalism) of the merger hierarchy of an $M_h = 10^{13} M_\odot$ halo at $z = 6$. The halo mass is chosen by requiring that the number density in halos more massive than $M_h$ match the space density of quasars at $z = 6$ (Fan et al. 2004).

## 3. Accretion and Radiative Efficiency

We explore two regimes: in one case, the accretion rate is limited to the Eddington rate; in the second, MBHs are allowed to accrete at supercritical rates (see Volonteri & Rees [2005] for a detailed description of the model). In both cases we assume that accretion episodes are triggered by major mergers (Mihos & Hernquist 1994, 1996; Di Matteo et al. 2005; Kazantzidis et al. 2005), which we define as mergers between halos with a mass ratio of 1 : 10 or higher. Mergers with smaller mass ratios are probably unable to trigger substantial gas inflow (Cox 2004).

In the first case, the hole accretes, at the Eddington rate, a gas rest mass $\Delta m_0$. This leads to a change in the total mass-energy of the hole

$$\Delta m = 7.7 \times 10^5 V_{100}^4 M_\odot,$$

where $V_{100}$ is the circular velocity of the merged system in units of 100 km s$^{-1}$. Adopting equation (2) implies an assumption that the correlations between black hole mass, velocity dispersion, and circular velocity are maintained throughout cosmic time (Tremaine et al. 2002; Ferrarese 2002). The quantities $\Delta m_n$ and $\Delta m_0$ are related through $\Delta m = (1 - \epsilon)\Delta m_0$, where $\epsilon$ is the mass-to-energy conversion efficiency, equal for thin disk accretion to the binding energy per unit mass of a particle in the last stable circular orbit. The MBH’s spin, $S$ (in the second case, we focus only on a subset of halos, those with effective atomic cooling. The cooling curve of metal-free gas has a sharp break at $T < 10^4 K$, so that for halos with $T_{\text{vir}} < 10^4 K$ the only available coolant is molecular hydrogen. $H_2$ is nevertheless a rather inefficient coolant (Madau et al. 2001). Metal-free halos with virial temperatures $T_{\text{vir}} > 10^4 K$ can instead cool even in the absence of $H_2$ by means of neutral hydrogen atomic lines to $\sim 8000 K$. Following Oh & Haiman (2002), we assume that a fraction $f_{\text{disk}}$ of the gas settles into an isothermal, exponential disk, embedded in a dark matter halo described by a Navarro et al. (1997, hereafter NFW) density profile. The mass of the disk can therefore be expressed as $M_{\text{disk}} = f_{\text{disk}}(\Omega_b/\Omega_m) M_h$.  

$^1$ This simple relation would be modified when the thickness of the disk is on the order of its radius and can also be changed by magnetic effects, which allow energy release from within the innermost stable orbit (see e.g., Krolik et al. 2005).
The disk is geometrically thick and has a very high central density (cf. Bromm & Loeb 2003). We refer the reader to Volonteri & Rees (2005) and Lodato & Natarajan (2006) for a detailed description of the model.

We estimate the mass accreted by the MBH from the surrounding disk within the Bondi-Hoyle formalism:

\[
\dot{M}_{\text{Bondi}} = \frac{4\pi G^2 M_{\text{BH}}^2 \rho n \eta}{c_s^3} = 4 \times 10^{-5} \left( \frac{M_{\text{BH}}}{10^3 M_\odot} \right)^2 \times \left( \frac{n_0}{10^4 \, \text{cm}^{-3}} \right) \left( \frac{T_{\text{gas}}}{8000 \, \text{K}} \right)^{-1.5} \dot{M}_{\odot} \, \text{yr}^{-1}
\]

(Bondi & Hoyle 1944), where \(n_0\) is the density of the gas, which is on the order of \(10^3 \, \text{cm}^{-3} \leq n_0 \leq 10^5 \, \text{cm}^{-3}\) at the center of the gas disks in high-redshift halos.

The collapsing gas disk likely rotates as a rigid body; rotation is therefore small during the initial collapse phase, and the infall of gas is quasi-radial. The size of the accretion disk, \(r_{\text{in}}\), is on the order of the trapping radius, that is, the radius at which radiation is trapped as the infall speed of the gas becomes larger than the diffusion speed of the radiation:

\[
r_{\text{in}} = R_{\odot} \frac{M}{M_{\text{Edd}}} \propto r_{\text{in}} \dot{M}_{\text{BH}}^{-1},
\]

where \(R_{\odot}\) is the Schwarzschild radius of the black hole. Begelman (1979) and Begelman & Meier (1982) studied supercritical accretion onto a BH in spherical geometry and quiescent thick disks, respectively. Despite the uncertainties, it still seems possible that when the inflow rate is supercritical, the radiative efficiency will drop so that the hole can accept the material without greatly exceeding the Eddington luminosity. The efficiency could be low either because most radiation is trapped and advected inward, or because the flow adjusts so that the material can plunge in from an orbit with small binding energy (Abramowicz & Lasota 1980).

The accretion rate is initially supercritical by a factor of 10 and grows to a factor of about \(10^3\) (Volonteri & Rees 2005), thus making the flow more and more spherical (see eq. [4]). On the other hand, the radius of the accretion disk increases steeply with the mass in the disk. In this case we also follow the evolution of MBHs and their hosts in full detail with a semi-analytical technique (Volonteri et al. 2003, 2005).

4. DYNAMICAL EVOLUTION OF MBHS

To include in this study the dynamics of MBHs, we follow the evolution of MBHs and their hosts in full detail with a semi-analytical technique (Volonteri et al. 2003, 2005).

4.1. MBH Binary Merger Efficiency

The merging—driven by dynamical friction against the dark matter—of two comparable-mass halo-plus-MBH systems (‘‘major mergers’’) drags in the satellite hole toward the center of the more massive progenitor, leading to the formation of a bound MBH binary with separation of \(\sim 1\) pc. In massive galaxies at low redshift, the subsequent evolution of the binary may be largely determined by three-body interactions with background stars (Begelman et al. 1980). Dark matter (DM) particles will be ejected by decaying binaries in the same way as the stars. Another possibility is that gas processes, rather than three-body interactions with stars or DM, may induce MBH binaries to shrink rapidly and coalesce (see, e.g., Gould & Rix 2000; Liu et al. 2003; Escala et al. 2004; Armitage & Natarajan 2005; Dotti et al. 2006; Mayer et al. 2006). If stellar dynamical or gaseous processes drive the binary sufficiently close \((\lesssim 0.01 \, \text{pc})\), gravitational radiation will eventually dominate angular momentum and energy losses and cause the two MBHs to coalesce.

In gas-rich high-redshift halos, the orbital evolution of the central SMBH is likely dominated by dynamical friction against the surrounding gaseous medium. The available simulations (Mayer et al. 2006; Dotti et al. 2006; Escala et al. 2004) show that the binary can shrink to about parsec or slightly subparsec scales by dynamical friction against the gas, depending on the gas thermodynamics. These binary separations are still too large for the binary to coalesce within the Hubble time owing to the emission of gravitational waves. On the other hand, the interaction between a binary and an accretion disk can lead to a very efficient transport of angular momentum, and the secondary MBH can reach the very subparsec separations at which emission of gravitational radiation dominates on short timescales (Liu et al. 2003; Armitage
The viscous timescale depends on the properties of the accretion disk and of the binary:

\[
t_{\text{vis}} = 0.1a_1^{3/2} \left( \frac{h}{r} \right)^{-2} \alpha_{0.1}^{-1/2} \left( \frac{m_1}{10^4 M_\odot} \right)^{-1/2} \text{Gyr}, \tag{5}
\]

where \( a_1 \) is the initial separation of the binary when the secondary MBH starts interacting with the accretion disk in units of parsecs, \( h/r \) is the aspect ratio of the accretion disk, \( h/r = 0.1 \) above, \( \alpha \) is the Shakura & Sunyaev (1973) viscosity parameter, \( \alpha = 0.1 \) above, and \( m_1 \) is the mass of the primary MBH, in solar masses. The emission of gravitational waves dominates the viscous timescales at a separation

\[
a_{\text{GW}} = 10^{-8} \left( \frac{h}{r} \right)^{-16/5} \alpha_{0.1}^{-8/5} q_{0.1}^{3/5} \left( \frac{m_1}{10^4 M_\odot} \right) \text{pc} \tag{6}
\]

(Armitage & Natarajan 2005), where \( q = m_2/m_1 \leq 1 \) is the binary mass ratio. The timescale for coalescence by emission of gravitational waves from \( a_{\text{GW}} \) is much shorter than the Hubble time:

\[
t_{\text{gr}} = \frac{5c^5 a^4(t)}{256G^2 m_1 m_2 (m_1 + m_2)}. \tag{7}
\]

We have assumed here that if an accretion disk surrounds a hard MBH binary, it will coalesce instantaneously owing to interaction with the gas disk. If instead there is no gas readily available, the binary will be losing orbital energy to the stars, if the initial mass function of Population III stars is bimodal and a large population of low-mass stars is formed (Nakamura & Umemura 2001), or to the DM background otherwise. The coalescence timescale typically becomes much longer than the Hubble time (see Madau et al. [2004] for a thorough discussion of this case).

To test the influence of the merger efficiency, we have compared the above merging-efficient model with a conservative merging-inefficient case in which the interaction with gas is neglected and the binaries shrink only by means of scattering from DM and stars, under the assumption that the loss cone stays full.

4.2. Gravitational Recoil

At high redshift, the recoil due to the nonzero net linear momentum carried away by gravitational waves may cause the ejection of MBHs from the shallow potential wells of their hosts (see, e.g., Madau et al. 2004; Madau & Quataert 2004; Merritt et al. 2004). The recoil velocity still has large uncertainties. Early calculations in the Newtonian regime (Fitchett 1983) predicted
at most a recoil velocity of $\approx 100$ km s$^{-1}$. The Newtonian calculations, in the circular case, predict the center of mass to move on circular orbits, spiraling outward while the binary orbit spirals inward, with a velocity

$$v_{\text{CM}} = 1480 \left[ \frac{f(q)}{f_{\text{max}}} \right] \left( \frac{R_s}{R_L} \right)^4 \text{ km s}^{-1}, \quad (8)$$

where $R_s = 2G(m_1 + m_2)/c^2$ is the Schwarzschild radius of the system and the scaling function reads $f(q)/f_{\text{max}}$, with $f(q) = q^2(1 - q)/(1 + q)^3$ and $f_{\text{max}} = 0.01789$. Here $R_L$ represents the radius of the last stable circular orbit, which is $R_L = 6Gm_1/c^2 = 3R_sm_i/(m_1 + m_2)$ in the Schwarzschild metric.

The recoil velocity during the plunge phase probably has the largest contribution. Calculations of the gravitational recoil inside the innermost stable orbit (ISCO) naturally have large uncertainties, but the recoil velocity should be constrained between the upper and lower limits ($V_{\text{upper}}$ and $V_{\text{lower}}$, respectively) suggested by Favata et al. (2004), which span a range between a few and several hundred kilometers per second for a binary with a mass ratio $q \approx 0.1$.

Blanchet et al. (2005) calculated the gravitational recoil to the second post-Newtonian order for nonspinning holes. Their calculations, available at the moment for Schwarzschild holes only, narrow the uncertainties on the final velocity from more than an order of magnitude (Favata et al. 2004) down to about 50%. Damour & Gopakumar (2006), using the effective one-body approach, predict velocities about a factor of 3 less than Blanchet et al. The uncertainties quoted by Blanchet et al. and Damour & Gopakumar make the results of the two calculations incompatible.

The latest estimate of the recoil comes from fully relativistic numerical simulations (Baker et al. 2006) following the dynamical evolution of a black hole binary within the ISCO. These simulations, carried out for a mass ratio $q = 0.2$, predict a recoil midway between the Blanchet et al. (2005) and Damour & Gopakumar (2006) predictions. The recoil predicted by Baker et al. is still large enough to eject the merging binary from small pregalactic structures. As shown by Yoo & Miralda-Escudé (2004), if the recoil effect is mild (e.g., the Damour & Gopakumar [2006] or Favata et al. [2004] lower limits), it is easy to fulfill the constraints on quasars at $z = 6$. As we are considering here the most pessimistic (although realistic) conditions under which it is still possible to develop the population of $z = 6$ quasars, we adopt as a zero point the Baker et al. (2006) results. We rescale the recoil to different mass ratios adopting Fitchett’s scaling function. Although Fitchett derived the scaling in the perturbative regime, outside the strong-gravity region, both Blanchet et al. (2005) and Damour & Gopakumar (2006) find a dependence on the mass ratio in very good agreement with Fitchett’s scaling.

One remaining issue is the effect of black hole spins, and the spin-orbit coupling, which however seems to be mild. Favata et al. (2004) suggest a modification of Fitchett’s scaling function to take into account spin-orbit coupling: $f(q) = f(q)|1 + (7/29)\tilde{a}_1/(1 - q)|/[1 + (7/29)\tilde{a}_1/(1 - 0.127)]$. Here $\tilde{a}_1$ is the physical spin parameter of the large hole ($S_1 = \tilde{a}_1GMm_i/c$, $0 \leq \tilde{a}_1 \leq 1$) and $\tilde{a}_i$ is the “effective” spin parameter of the binary system. Damour suggests $\tilde{a}_i = \hat{a}_i(1 + 3\hat{a}_i/4)/(1 + q)^2$ in the post-Newtonian limit. Figure 2 compares different theoretical estimates for the recoil velocity.

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2. We performed one calculation adopting the Favata et al. (2004) lower limits and checked that the fraction of displaced binaries amounts to about 5%.
Figure 3 compares the MBH merger rates in two different models accounting for the orbital evolution of MBH binaries in the phase preceding emission of gravitational waves. If gas processes (efficient merging) drive the MBH orbital decay, MBHs start merging efficiently at very early times, when host DM halos are still small. Although the absolute number of mergers is larger, more MBHs are ejected from DM halos as a result of the rocket effect, because halos are smaller at earlier times. The two effects, increased mergers and increased ejections, compensate (see also Fig. 5 below), thus leading to similar SMBH masses in the main halo at $z = 6$. Note that the MBH merger rate in such a highly biased volume is not the average cosmic rate, as we are following only the hierarchy leading to extremely high mass halos at $z = 6$ (see Diemand et al. [2005] for similar considerations at lower redshift). MBHs in a high-overdensity volume experience a much higher number of mergers with respect to the typical MBH evolving in a less overdense region.

Figure 4 summarizes our results in terms of MBH binary ejection. We have explored a large parameter space, and we selected here a few representative examples. Accretion and dynamical histories are strongly intermingled, as the kick velocity depends on the binary mass ratio. The rocket recoil is substantial only for mass ratios of order $q \approx 0.1$, which is believed to be the most probable mass ratio for MBH binaries (because of the inefficiency of dynamical friction in bringing the satellite MBH within the influence of the main one in the case of minor mergers; see also Sesana et al. 2005). Interestingly, the Volonteri & Rees (2005) model for supercritical accretion skews the mass ratios of MBH binaries toward lower values, as the conditions for supercritical accretion to happen are typically fulfilled by only one of two merging systems. This latter result supports a biased and selective growth of high-$z$ BHs, as the conditions for supercritical accretion appear to be fulfilled only in halos with $T_{\text{vir}} > 10^6$ K, which represent 3 $\sigma$ or 4 $\sigma$ peaks in the field of density fluctuations. If only a tiny fraction of the MBHs, those hosted in the most massive halos at this early time, undergo rapid growth, the conditions for a mild effect of the recoil are naturally met, regardless of which detailed kick-velocity calculation is considered. The distribution of merging-MBH mass ratios is in fact skewed toward low mass ratios ($q = m_2/m_1 \ll 1$), where the expected kick velocity is low (Fig. 2).
5.2. MBH Accretion

The mass growth history of a MBH ending up as a SMBH in an $M_h = 10^{13} M_\odot$ halo at $z = 6$ is shown in Figure 5 for a series of different models. MBH mass increase by mergers is taken into account, as well as mergers’ “negative feedback” onto the MBH growth due to the possibility of binary ejection following coalescence.

As a reference, we have considered an upper limit to the radiative efficiency of $\epsilon = 0.16$, or $a = 0.9$ adopting the standard conversion, inspired by the Gammie et al. (2004) simulations. Figure 5 shows, however, that MBH masses on the order of $10^9 M_\odot$ can be reached by $z = 6$ only if dynamical effects are not too destructive. In fact, if we adopt the Baker et al. (2006) scaling (eq. [9]) for the gravitational recoil velocity, MBH masses are always less than $10^9$ solar masses. The constraints imposed by $z = 6$ quasars can be met if either (1) the radiative efficiency is lower than $\epsilon_{\text{max}} = 0.12$ (corresponding to a spin parameter $a = 0.8$) or (2) the number of seed MBHs is much more than 25 times larger (i.e., seeds inhabit density peaks $\leq 3.5 \sigma$ at $z = 24$; note that in this case, massive black holes would be expected also in dwarf galaxies with total masses well below $M_h = 10^{11} M_\odot$ today), or, alternatively, (3) if the MBHs go through a phase of supercritical accretion.

If the dynamical influence of the “gravitational rocket” is better described by, for example, Favata et al. (2004) or Damour & Gopakumar (2006), though, the observational constraints can be met more easily, allowing for higher radiative efficiencies or rarer seeds. It is worth noting that a radiative efficiency $\epsilon > 0.16$ can be accommodated only if the influence of mergers is large (see also Shapiro 2005; Yoo & Miralda-Escudé 2004), for example, seeds in lower $\sigma$-peaks with efficient merging of binary MBHs.

We have compared theoretical luminosity functions (LFs) with the most recent determination of the quasar blue LF from the SDSS (Fan et al. 2004). The SDSS samples only the very bright end of the LF; Fan et al. (2001b) fitted a single power law to the data, which span luminosities larger than $10^{13} L_\odot$ in the blue band. Upper limits on the faint end of the luminosity function, though, can be derived from the nondetection of $z = 6$ quasars.

Fig. 5.—Averaged mass growth history of a MBH at the center of a $10^{13} M_\odot$ halo at $z = 6$. Top left, Eddington accretion rate only, maximum radiative efficiency $\epsilon = 0.16$; top right, Eddington accretion rate only, maximum radiative efficiency $\epsilon = 0.12$; bottom left, supercritical accretion, in a fat disk with $f_d = 0.1$; bottom right, Eddington accretion rate only, maximum radiative efficiency $\epsilon = 0.16$ and a 25 times larger abundance of seed holes. The crosses mark the locus of SDSS J1148+5251.
quasars in the Chandra Deep Field–North (Barger et al. 2003) and in the Canada-France High-z Quasar Survey (Willott et al. 2005).

We have adopted here the parameterization of the LF given by Willott et al. and compared our results with two representative cases for the faint-end \((\alpha, \beta) = (-1, -2.5)\) and bright-end \((\alpha, \beta) = (-2, -3.5)\). Letting the luminosity break \(L_B\) vary between \(7 \times 10^{11}\) and \(2 \times 10^{12}\) \(L_\odot\) leads to indistinguishable curves. Figure 6 shows that quasars bright enough to reproduce the observed LF cannot be created under the assumption that accretion is limited at the Eddington rate, occurring via a disk that can spin black holes up to \(\hat{a} = 0.9\) (corresponding using the standard conversion to \(\epsilon = 0.16\)). The constraints from the faint end of LF are still very weak (cf. Fig. 5 of Willott et al. 2005), so we cannot rule out either a model with a lower \(\epsilon_{\text{max}}\) or a model with supercritical accretion based on these results only.

6. DISCUSSION

The detection of luminous quasars at \(z \simeq 6\), suggesting an early growth of SMBHs, deserves a special investigation within the hierarchical scenario for galaxy formation. Following a series of papers tracing the seeds of the SMBH population observed today from the first stars, we have here focused on the constraints set by the presence of SMBHs already in place when the universe was less than 1 Gyr old.

The halos that we choose for our investigation, \(M_h > 10^{13} M_\odot\) at \(z = 6\), are 5.5 \(\sigma\) density fluctuations. They are therefore not representative of the typical halo mass at \(z = 6\). On the other hand, the masses are chosen by requiring that their number density match the number density of quasars derived by Fan et al. (2004). That the highest massive halos experienced strong evolution at early times is in line with the “antihierarchical” evolution of galactic structures suggested by high-redshift galactic surveys (e.g., Kodama et al. 2004). A simple exercise can help us quantify the global importance of the black holes hosted in the selected halos.

If the correlation between black hole mass and halo circular velocity holds up to \(z = 6\), we can use the Press-Schechter formalism to determine the MBH density as a function of halo mass at \(z = 6\). The MBH mass density in halos with mass larger than \(M_h = 10^{13} M_\odot\) is \(\sim 20 M_\odot\) \(\text{Mpc}^{-3}\), while for \(M_h = 10^{12} M_\odot\) we have \(\lesssim 10^3 M_\odot\) \(\text{Mpc}^{-3}\), and for \(M_h = 10^{11} M_\odot\) we have \(\lesssim 10^4 M_\odot\) \(\text{Mpc}^{-3}\). So, if all halos with, for example, \(M_h = 10^{11} M_\odot\) at \(z = 6\) host a SMBH whose mass scales with the local \(M_{\text{BH}}-V_c\) relation, the SMBH density at \(z = 6\) is similar to that at \(z = 3\), leaving very little room for accretion between \(z = 6\) and \(z = 3\). This implies
that either the $M_{\rm BH}-V_c$ relation is redshift dependent, and in the opposite sense from that suggested by the Walter et al. (2004) observation, or the occupation fraction of black holes is not unity at all redshifts. We define the MBH occupation fraction as the average number of halos hosting a MBH over the total number of halos in a comoving volume. Volonteri et al. (2003; see also Marulli et al. 2006) shows that models in which SMBHs evolve from seeds forming in biased regions (e.g., remnants of Population III stars) predict a SMBH occupation fraction that decreases with halo mass.

If the occupation fraction is a declining function of the halo mass, we expect MBH-MBH interactions to be more common for massive halos. MBHs hosted in low-mass halos grow in isolation, while MBHs lurking in massive halos have a much higher probability of experiencing a MBH-MBH merger. So, on one hand MBHs hosted in low-mass halos have a larger probability of being ejected by a kick, and on the other hand their probability of experiencing a merger in the first place is smaller. We might therefore expect that a SMBH can grow almost undisturbed by $z = 6$ up to a billion solar masses, if hosted in, say, a $10^{11} M_\odot$ halo. A requirement in this case is however that the mass of the black holes, in these high-redshift systems, be able to grow as $M_{\rm BH} \propto z$, as described above. The above scenario is thus independent of the accretion of gas into the SMBH.

To summarize, if the SMBHs powering $z = 6$ quasars are hosted in halos with $M_\Delta \lesssim 10^{12} M_\odot$, they can grow almost in isolation, and accretion can be limited to the Eddington rate, provided the radiative efficiency is not too high, but the mass of the SMBHs must be allowed to exceed the nominal $M_{\rm BH}-V_c$ correlation. This is the opposite of what is seen in numerical simulations (Robertson et al. 2006) but in agreement with the Walter et al. (2004) results.

The emerging picture, therefore, points to a Darwinian natural selection scenario. At early times, only a small fraction of MBHs, those hosted in the highest-density peaks ($T_{\rm vir} > 10^4 K$), evolve rapidly and efficiently (see also Johnson & Bromm 2006). The effects are twofold: not only do these rare BHs quickly become supermassive owing to efficient accretion, but the effects of dangerous dynamical interactions are also softened, as the merging of low mass ratio MBH binaries is favored. Alternative models for MBH seed formation (e.g., Koushiappas et al. 2004) tend to select as sites of black hole formation halos with $T_{\rm vir} > 10^4 K$ as well.

Clarification of the early evolution of SMBHs can come from detection of the gravitational wave signal from their inspiraling, as different models for the MBH early evolution predict different mass ratio distributions for the merging MBHs. Reliable calculations narrowing the uncertainties in the kick velocity due to gravitational emission for spinning MBH binaries will help to clarify the threat that kicks represent for the early evolution of the MBH population.

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