Dark matter haloes determine the masses of supermassive black holes

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ABSTRACT
The energy and momentum deposited by the radiation from accretion flows on to the supermassive black holes (BHs) that reside at the centres of virtually all galaxies can halt or even reverse gas inflow, providing a natural mechanism for supermassive BHs to regulate their growth and to couple their properties to those of their host galaxies. However, it remains unclear whether this self-regulation occurs on the scale at which the BH is gravitationally dominant, on that of the stellar bulge, the galaxy or that of the entire dark matter halo. To answer this question, we use self-consistent simulations of the co-evolution of the BH and galaxy populations that reproduce the observed correlations between the masses of the BHs and the properties of their host galaxies. We first confirm unambiguously that the BHs regulate their growth: the amount of energy that the BHs inject into their surroundings remains unchanged when the fraction of the accreted rest mass energy that is injected is varied by four orders of magnitude. The BHs simply adjust their masses so as to inject the same amount of energy. We then use simulations with artificially reduced star formation rates to demonstrate explicitly that BH mass is not set by the stellar mass. Instead, we find that it is determined by the mass of the dark matter halo with a secondary dependence on the halo concentration, of the form that would be expected if the halo binding energy were the fundamental property that controls the mass of the BH. We predict that the BH mass, \( m_{\text{BH}} \), scales with halo mass as \( m_{\text{BH}} \propto m_{\text{halo}}^\alpha \), with \( \alpha \approx 1.55 \pm 0.05 \), and that the scatter around the mean relation in part reflects the scatter in the halo concentration–mass relation.

Key words: hydrodynamics – galaxies: active – galaxies: evolution – galaxies: formation – quasars: general – cosmology: theory.

1 INTRODUCTION
Almost all massive galaxies are thought to contain a central supermassive black hole (BH) and the properties of these BHs are tightly correlated with those of the galaxies in which they reside (e.g. Magorrian et al. 1998; Ferrarese & Merritt 2000; Gebhardt et al. 2000; Tremaine et al. 2002; Häring & Rix 2004; Hopkins et al. 2007b; Ho 2008). It is known that most of the mass of the BHs is assembled via luminous accretion of matter (Soltan 1982). The energy emitted by this process provides a natural mechanism by which BHs can couple their properties to those of their host galaxies. Analytic (e.g. Silk & Rees 1998; Haehnelt, Natarajan & Rees 1998; Fabian 1999; Adams, Graff & Richstone 2001; King 2003; Wyithe & Loeb 2003; Murray, Quataert & Thompson 2005; Merloni & Heinz 2008), semi-analytic (e.g. Kauffmann & Haehnelt 2000; Cattaneo 2001; Granato et al. 2004; Bower et al. 2006) and hydrodynamical (e.g. Springel, Di Matteo & Hernquist 2005; Di Matteo, Springel & Hernquist 2005; Robertson et al. 2006; Sijacki et al. 2007; Hopkins et al. 2007a; Di Matteo et al. 2008; Okamoto, Nemmen & Bower 2008; Booth & Schaye 2009, hereafter BS09) studies have used this coupling between the energy emitted by luminous accretion and the gas local to the BH to investigate the origin of the observed correlation between BH and galaxy properties, and the buildup of the supermassive BH population.

BHs are expected to regulate the rate at which they accrete gas down to the scale on which they are gravitationally dominant. For example, gas flowing in through an accretion disc can become so hot that its thermal emission becomes energetically important. Scattering of the photons emitted by the accreting matter by free electrons gives rise to the so-called Eddington limit. If the accretion rate exceeds this limit, which is inversely proportional to the assumed radiative efficiency of the accretion disc, then the radiative force exceeds the gravitational attraction of the BH and the inflow is quenched, at least within the region that is optically thin to the radiation.

However, observations indicate that the time-averaged accretion rate is far below the Eddington rate (Kollmeier et al. 2006), suggesting the presence of processes acting on larger scales. Indeed, the existence of tight correlations between the mass of the BH

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and the properties of the stellar bulge indicates that self-regulation may happen on the scale of the bulge (∼1 kpc; Adams et al. 2001; Hopkins et al. 2007a), far exceeding the radius within which the BH is gravitationally dominant. However, since galaxy-wide processes such as galaxy mergers can trigger gas flows into the bulge (Sanders et al. 1988; Mihos & Hernquist 1994), it is conceivable that BHs could regulate their growth on the scale of the entire galaxy (∼10 kpc; Haehnelt et al. 1998; Fabian 1999; Wyithe & Loeb 2003) or even on that of the dark matter (DM) haloes hosting the galaxies (∼10^7 kpc; Silk & Rees 1998; Ferrarese 2002). Finally, it is possible, perhaps even likely, that self-regulation takes place simultaneously on multiple scales. For example, frequent, short, Eddington-limited outbursts may be able to regulate the inflow of gas on the scale of the bulge averaged over much longer time-scales.

In this Letter we investigate, using self-consistent simulations of the co-evolution of the BH and galaxy populations, on what scale the self-regulation of BHs takes place. In Section 2, we describe the numerical techniques and simulation set employed in this study. In Section 3, we demonstrate that BH self-regulation takes place on the scale of the DM halo and that the BH mass is determined by the binding energy of the DM halo rather than by the stellar mass of the host galaxy. Throughout we assume a flat Λ cold dark matter (ΛCDM) cosmology with the cosmological parameters: \( \{ \Omega_m, \Omega_b, \Omega_\Lambda, \sigma_8, n_s, h \} = \{ 0.238, 0.0418, 0.762, 0.74, 0.951, 0.73 \} \), as determined from the Wilkinson Microwave Anisotropy Probe 3-yr data (Spergel et al. 2007).

## 2 Numerical Methods

We have carried out a set of cosmological simulations using smoothed particle hydrodynamics (SPH). We employ a significantly extended version of the parallel PMTree-SPH code gadget III (last described in Springel 2005), a Lagrangian code used to calculate gravitational and hydrodynamic forces on a particle-by-particle basis. The initial particle positions and velocities are set at \( z = 127 \) using the Zeldovich approximation to linearly evolve positions from an initially glass-like state.

In addition to hydrodynamic forces, we treat star formation, supernova feedback, radiative cooling, chemodynamics and BH accretion and feedback, as described in Schaye & Dalla Vecchia (2008), Dalla Vecchia & Schaye (2008), Wiersma, Schaye & Smith (2009a), Wiersma et al. (2009b) and BS09, respectively. For clarity we summarize here the essential features of the BH model, which is itself a substantially modified version of that introduced by Springel et al. (2005).

### 2.1 The black hole model

Seed BHs of mass \( \dot{m}_{\text{seed}} = 10^{-3} m_g \) – where \( m_g \) is the simulation gas particle mass – are placed into every DM halo that contains more than 100 DM particles and does not already contain a BH particle. Haloes are identified by regularly running a friends-of-friends group finder on-the-fly during the simulation. After forming, BHs grow by two processes: accretion of ambient gas and mergers. Gas accretion occurs at the minimum of the Eddington rate, \( \dot{m}_{\text{edd, min}} = 4 \pi G M_{\text{BH}} m_p / \sigma_T c^2 \), and \( \dot{m}_{\text{seed}} = 4 \pi G^2 m_g^2 \rho (c_s^2 + v^2)^{1/2} \), where \( m_p \) is the proton mass, \( \sigma_T \) is the Thomson cross-section, \( c \) is the speed of light, \( c_s \) and \( p \) are the sound speed and density of the local medium, respectively, \( v \) is the velocity of the BH relative to the ambient medium and \( \alpha \) is a dimensionless efficiency parameter. The parameter \( \alpha \), which was set to 100 by Springel et al. (2005), accounts for the fact that our simulations possess neither the necessary resolution nor the physics to accurately model accretion on to a BH on small scales. Note that for \( \alpha = 1 \), this accretion rate reduces to the so-called Bondi–Hoyle (Bondi & Hoyle 1944) rate.

As long as we resolve the scales and physics relevant to the Bondi–Hoyle accretion, we could set \( \alpha = 1 \). If a simulation resolves the Jeans scales in the accreting gas, then it will also resolve the scales relevant for the Bondi–Hoyle accretion on to any BH larger than the simulation mass resolution (BS09). We therefore generally set \( \alpha \) equal to unity. However, this argument breaks down in the presence of a multiphase interstellar medium (ISM), because our simulations do not resolve the properties of the cold, molecular phase, and as such the accretion rate may be orders of magnitude higher than the Bondi–Hoyle rate predicted by our simulations for star-forming gas. We therefore use a power-law scaling of the accretion efficiency such that \( \alpha = (n_{\text{HI}}/n_{\text{HI}}^0)^{\beta} \) in star-forming gas, where \( n_{\text{HI}}^0 = 0.1 \text{ cm}^{-3} \) is the critical density for the formation of a cold, star-forming gas phase. The parameter \( \beta \) is a free parameter in our simulations. We set \( \beta = 2 \), but note that the results shown here are insensitive to changes in this parameter when \( \beta \geq 2 \) (see BS09), because in that case the growth of the BHs is limited by feedback.

Energy feedback is implemented by allowing BHs to inject a fixed fraction of the rest mass energy of the gas they accrete into the surrounding medium. The energy deposition rate is given by

\[
\dot{E} = \epsilon_i \dot{m}_{\text{acc}} c^2 = \frac{\epsilon_i \dot{m}_{\text{BH}}}{\epsilon_i - \epsilon_f} m_{\text{BH}} c^2,
\]

where \( \epsilon_i \) is the radiative efficiency of the BH, \( \dot{m}_{\text{acc}} \) is the rate at which the BH is accreting gas and \( m_{\text{BH}} \) is the rate of BH mass growth.

We set \( \epsilon_i \) to be 0.1, the mean value for radiatively efficient accretion on to a Schwarzschild BH (Shakura & Sunyaev 1973). We vary \( \epsilon_i \) but use \( \epsilon_f = 0.15 \) as our fiducial value. It was shown in BS09 that, for \( \epsilon_i = 0.15 \), simulations identical to these reproduce the observed redshift zero \( m_{\text{BH}} - m_* \) and \( m_{\text{BH}} - \sigma \) relations, where \( \sigma \) is the one-dimensional velocity dispersion of the stars and \( m_* \) is the galaxy stellar mass. Energy is returned to the surroundings of the BH ‘thermally’, that is, by increasing the temperature of the BH’s neighbouring SPH particles by at least \( \Delta T_{\text{min}} \). A BH performs no heating until it has built up enough of an energy reservoir to heat by this amount. The use of an energy reservoir is necessary in these simulations as otherwise gas will be able to radiate away the energy every time-step. Imposing a minimum temperature increase ensures that the radiative cooling time is sufficiently long for the feedback to be effective. In our fiducial model, we set \( N_{\text{heat}} = 1 \) and \( \Delta T_{\text{min}} = 10^8 \text{ K} \) but the results are insensitive to the exact values of these parameters (see BS09).

### 2.2 The simulation set

The simulations employed in the current work use cubic boxes of size 12.5 and 50 Mpc \( h^{-1} \) and assume periodic boundary conditions. Each simulation contains either 128^3 or 256^3 particles of both gas and collisionless CDM. Comoving gravitational softenings are set to 1/25 of the mean interparticle separation down to \( z = 2.91 \), below which we switch to a fixed proper scale. The 12.5 Mpc \( h^{-1} \) (50 Mpc \( h^{-1} \)) boxes are evolved as far as redshift 2 (zero). The numerical parameters of the simulations used in this study are summarized in Table 1. All results presented in this Letter are derived from the 50.0 Mpc \( h^{-1} \), 256^3 particle simulations, with the other box sizes and particle numbers employed to demonstrate numerical convergence.

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Table 1. Numerical parameters of the simulations. From left to right: simulation identifier, comoving box size (Mpc $h^{-1}$), number of both gas and DM particles, final redshift, gas particle mass ($10^7$ $M_\odot$ $h^{-1}$), DM particle mass ($10^7$ $M_\odot$ $h^{-1}$) and maximum physical gravitational softening (kpc $h^{-1}$). Each simulation was run multiple times using different values of $\epsilon_f$.

| Name        | L     | $n_{part}$ | $z_f$ | $m_g$ | $m_{DM}$ | $\epsilon_{max,phys}$ |
|-------------|-------|------------|------|-------|----------|------------------------|
| L050N256    | 50.0  | 256        | 0    | 8.7   | 41.0     | 2.0                    |
| L050N128    | 50.0  | 128        | 0    | 69.6  | 328.0    | 4.0                    |
| L012N256    | 12.5  | 256        | 0    | 0.1   | 0.6      | 0.5                    |

3 RESULTS AND DISCUSSION

It is instructive to first consider under what conditions BHs can regulate their growth. To regulate its growth on a mass scale $M_{sr}$, a BH of mass $m_{BH}$ must be able to inject energy (or momentum) at a rate $\epsilon_f$ that is sufficient to counteract the force of gravity on the scale $M_{sr}$, averaged over the dynamical time associated with this scale. The mass $M_{sr}$ could, for example, correspond to that of the BH, the stellar bulge or the DM halo. If the BH cannot inject energy sufficiently rapidly, then gravity will win and its mass will increase. Provided that the maximum rate at which it can inject energy increases with $m_{BH}$ (as is e.g. the case for Bondi–Hoyle and Eddington-limited accretion with a constant radiative efficiency) and provided this rate increases sufficiently rapidly to counteract the growth of $M_{sr}$, the BH will ultimately reach the critical mass $m_{BH, crit}(M_{sr})$ required to halt the inflow on the scale $M_{sr}$. If, on the other hand, $m_{BH} \gg m_{BH, crit}(M_{sr})$, then the BH will quickly quench the accretion flow and its mass will consequently remain nearly unchanged. The BH will in that case return to the equilibrium value $m_{BH, crit}(M_{sr})$ on the time-scale which characterizes the growth of $M_{sr}$.

If the BH regulates its growth on the mass scale $M_g$ and if $m_{BH} \ll M_g$, then the critical rate of energy injection required for self-regulation is independent of the mass of the BH. It then follows from equation (1) that $m_{BH} \propto \epsilon_f^{-1}$, which implies

$$m_{BH} - m_{seed} \propto \epsilon_f^{-1},$$

where $m_{seed}$ is the initial mass of the BH. Hence, if the self-gravity of BHs is negligible on the maximum scale on which they regulate their growth and if $m_{BH} \gg m_{seed}$, then we expect $m_{BH} \propto \epsilon_f^{-1}$.

The black diamonds plotted in Fig. 1 show the predicted global mass density in BHs at redshift zero as a function of $\epsilon_f$, the efficiency with which BHs couple energy into the ISM, normalized to the density obtained for $\epsilon_f = 0.15$. Similarly, the black plus signs indicate the normalization of the $m_{BH, crit} - \sigma$ relation divided by that for the $\epsilon_f = 0.15$ run. The feedback efficiency, $\epsilon_f$, is varied, in factors of 4, from $\epsilon_f = 9.2 \times 10^{-6}$ to 9.6, which implies that the fraction of the accreted rest mass energy that is injected (i.e. $\epsilon_f \approx 1$) varies from $9.2 \times 10^{-7}$ to 0.96. BH mass is clearly inversely proportional to the assumed feedback efficiency for $10^{-4} < \epsilon_f < 1$. For $\epsilon_f > 1$ the trend breaks down because the BH masses remain similar to the assumed seed mass, in accord with equation (2). If we had used a lower seed mass, then the trend would have extended to greater values of $\epsilon_f$. The deviation from inverse proportionality that sets in below $\epsilon_f = 10^{-4}$ is more interesting. Such low values yield BH masses that are more than $0.15/10^{-4} \approx 10^3$ times greater than observed, in which case they are no longer negligible compared to the masses of their host galaxies. In that case, the critical rate of energy deposition will no longer be independent of $m_{BH}$ and we do not expect equation (2) to hold.

We have thus confirmed that feedback enables BHs to regulate their growth. Moreover, we demonstrated that this self-regulation takes places on scales over which the gravitational influence of the BHs is negligible, provided that the fraction of the accreted rest mass energy that is coupled back into the ISM is $\gtrsim 10^{-3}$.

To test whether it is the stellar or the DM distribution that determines the mass of BHs, we compare the BH masses in two simulations that are identical except for the assumed efficiency of star formation. One uses our fiducial star formation law, but in the other simulation we reduced its amplitude by a factor of 100, making the gas consumption time-scale much longer than the age of the Universe. Because changing the amount of stars would imply changing the rate of injection of supernova energy, which could affect the efficiency of BH feedback, we neglected feedback from star formation in both runs. In the simulation with ‘normal’ star formation the central regions of the galaxies are dominated gravitationally by the baryonic component of the galaxy, whereas in the simulation with reduced star formation the DM dominates everywhere. Fig. 2 shows the $m_{BH} - m_{halo}$ and $m_{BH} - m_{DM}$ relations at redshift zero. While the two runs produce nearly identical BH masses for a fixed halo mass, the $m_{BH} - m_{halo}$ relation is shifted to lower stellar masses by more than an order of magnitude in the model with reduced star formation. The insensitivity of the relation between $m_{BH}$ and $m_{halo}$ to the assumed star formation efficiency demonstrates that the BH mass is not set by the gravitational potential on the scale of the galaxy. We have verified that the same result holds at redshift 2 for the simulations with 64 times better mass resolution. Clearly, stellar mass does not significantly influence the relation between the mass of the BH and that of its host halo. This implies that BH self-regulation occurs on the scale of DM haloes.

If the rate by which the BHs inject energy is independent of the assumed feedback efficiency, then we expect the same to be true for
the factor by which BH feedback suppresses star formation. This is confirmed by comparison of the global star formation rates in runs with different values of $\epsilon$ (see fig. 6 of BS09).

Fig. 3 compares the predicted $\log_{10}(m_{\text{BH}}/M_\odot) - \log_{10}(m_{\text{halo}})$ relation with observation (Bandara, Crampton & Simard 2009). The agreement is striking. The slope and normalization of the observed $\log_{10}(m_{\text{BH}}/M_\odot) - \log_{10}(m_{\text{halo}})$ relation are $1.55 \pm 0.31$ and $8.18 \pm 0.11$, respectively, whereas the simulation predicts $1.55 \pm 0.05$ and $8.01 \pm 0.04$. Note that the simulation was only tuned to match the normalization of the relations between $m_{\text{BH}}$ and the galaxy stellar properties.

If the energy injected by a BH is proportional to the halo gravitational binding energy, then, for isothermal models (Silk & Rees 1998), $m_{\text{BH}} \propto m_{\text{halo}}^{5/3}$. Here we extend these models to the more realistic universal halo density profile (Navarro, Frenk & White 1997), whose shape is specified by a concentration parameter, $c$ (we assumed $c \propto v_{\text{max}}^2/v_{\text{halo}}^2$, where $v_{\text{max}}$ and $v_{\text{halo}}$ are the maximum halo circular velocity and the circular velocity at the virial radius, respectively). It is known that concentration decreases with increasing halo mass, $c \propto m_{\text{halo}}^{-0.1}$ (Bullock et al. 2001; Duffy et al. 2008), which then affects BH mass through the dependence of halo binding energy on concentration. If the total energy injected by a BH of a given mass is proportional to the energy required to unbind gas from a DM halo (Lokas & Mamon 2001) out to some fraction of the virial radius, $r_{ej}/r_{\text{halo}}$, then

$$m_{\text{BH}} \propto \frac{c}{[\ln(1+c) - c/(1+c)]^2} \times \left( \frac{1}{1 + c e_{\text{halo}}/r_{\text{halo}}} - \frac{2 \ln \left( 1 + c e_{\text{halo}}/r_{\text{halo}} \right)}{1 + c e_{\text{halo}}/r_{\text{halo}}} \right) m_{\text{halo}}^{5/3} \ .$$

(3)

Inserting $c \propto m_{\text{halo}}^{-0.1}$ and computing the logarithmic derivative with respect to $m_{\text{halo}}$ in the mass range $10^{10} M_\odot < m_{\text{halo}} < 10^{14} M_\odot$, we find that the slope is a weak function of $r_{ej}/r_{\text{halo}}$ that varies from 1.50 at $r_{ej} = 10^{-3} r_{\text{halo}}$ to 1.61 at $r_{ej} = r_{\text{halo}}$. The close match between theory, simulation and observation suggests that the halo binding energy, rather than halo mass, determines the mass of the BH.

The residuals from the $m_{\text{BH}} - m_{\text{halo}}$ relation ($\delta m_{\text{BH}}$) are correlated with halo concentration (Spearman rank correlation coefficient $\rho = 0.29$; probability of significance $P = 0.9998$) as would be expected if $m_{\text{BH}}$ is sensitive to the halo binding energy. The residuals are also correlated with galaxy stellar mass, though much less strongly ($\rho = 0.09$; $P = 0.96$). Taken together, these correlations tell us that, at a given halo mass, galaxies with BHs more massive than the average will also contain a larger than average number of stars and are hosted by more concentrated haloes. This suggests that the galaxy stellar mass is also determined by the halo binding energy. Thus, outliers in the $m_{\text{BH}} - m_{\text{halo}}$ relation may still lie close to the mean of $m_{\text{BH}} - m_{\text{halo}}$ relation. Furthermore, higher concentrations imply earlier formation times and spheroidal components do indeed typically host old stellar populations.
In addition to the ‘quasar mode’ of feedback discussed in this work, it has recently become clear that a second ‘radio mode’ may be required to quench cooling flows in galaxy groups and clusters (see e.g. Cattaneo et al. 2009, for a review). Although we do not explicitly include a ‘radio mode’ in this work, the active galactic nucleus feedback prescription explored here is capable of suppressing cooling flows, at least on group scales, providing excellent matches to observed group density and temperature profiles as well as galaxy stellar masses and age distributions (McCarthy et al. 2009). It is known that BHs obtain most of their mass in the ‘quasar mode’ (Soltan 1982), so any discussion of what determines the masses of BHs must focus primarily on this mode of accretion. Finally, the ability of a BH to quench cooling flows in the ‘radio mode’ is expected to be closely related to the virial properties of the hot halo (Cattaneo et al. 2009) and would therefore provide an additional link between BHs and DM haloes over and above what we discuss here and so serve to make any fundamental connection between the BH mass and the properties of the DM halo even stronger.

We conclude that our simulation results suggest that in order to effectively halt BH (and galaxy) growth, gas must not return to the galaxy on a short time-scale. This requires that the BH injects enough energy to eject gas out to scales where the DM halo potential is dominant. The mass of the BH is therefore determined primarily by the mass of the DM halo with a secondary dependence on halo concentration, of the form that would be expected if the BH mass were controlled by the halo binding energy. The tight correlation between $m_{\text{BH}}$ and $m_*$ is then a consequence of the more fundamental relations between the halo binding energy and both $m_{\text{BH}}$ and $m_*$. 

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