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BLACK HOLE FORAGING: FEEDBACK DRIVES FEEDING

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ABSTRACT

We suggest a new picture of supermassive black hole (SMBH) growth in galaxy centers. Momentum-driven feedback from an accreting hole gives significant orbital energy, but little angular momentum to the surrounding gas. Once central accretion drops, the feedback weakens and swept-up gas falls back toward the SMBH on near-parabolic orbits. These intersect near the black hole with partially opposed specific angular momenta, causing further infall and ultimately the formation of a small-scale accretion disk. The feeding rates into the disk typically exceed Eddington by factors of a few, growing the hole on the Salpeter timescale and stimulating further feedback. Natural consequences of this picture include (1) the formation and maintenance of a roughly toroidal distribution of obscuring matter near the hole; (2) random orientations of successive accretion disk episodes; (3) the possibility of rapid SMBH growth; (4) tidal disruption of stars and close binaries formed from infalling gas, resulting in visible flares and ejection of hypervelocity stars; (5) super-solar abundances of the matter accreting on to the SMBH; and (6) a lower central dark-matter density, and hence annihilation signal, than adiabatic SMBH growth implies. We also suggest a simple subgrid recipe for implementing this process in numerical simulations.

Key words: accretion, accretion disks – black hole physics – galaxies: evolution – quasars: general

Online-only material: color figure

1. INTRODUCTION

The relation between supermassive black holes (SMBHs) and their host galaxies is a major theme of current astrophysics. The scaling relations (Ferrarese & Merritt 2000; Gebhardt et al. 2000; Häring & Rix 2004) between the SMBH mass M and the velocity dispersion σ and mass $M_{\rm bulge}$ of the host spheroid strongly suggest that the hole's enormous binding energy affects the host in important ways. A credible picture of this process is gradually emerging (e.g., Silk & Rees 1998; Fabian 1999; King 2003, 2005; Zubovas & King 2012). But we are still far from a deterministic theory of SMBH–galaxy coevolution, because we have no cogent picture of how the host affects the hole, i.e., of what causes SMBH mass growth. We know that this must largely occur through accretion of gas: the Soltan (1982) relation implies that mass growth produces electromagnetic radiation with accretion efficiency $\eta \simeq 0.1\times$ rest-mass energy, at least at low redshifts. This rules out dark-matter accretion as a major contributor, and direct accretion of stars through tidal disruption is inefficient (Frank & Rees 1976).

Because all gas has angular momentum, accretion on to the hole at the smallest scales must be through an accretion disk. But these scales must indeed be small: the viscous timescale

$$t_{\rm visc} = \frac{1}{\alpha}\left(\frac{R}{H}\right)^2\left(\frac{R^3}{GM}\right)^{1/2} \qquad (1)$$

approaches a Hubble time at scales of only a few times 0.1 pc if the accreting gas can cool, so that the disk aspect ratio $H/R \ll 1$ (e.g., King & Pringle 2006, 2007; $\alpha \lesssim 1$ is the standard Shakura & Sunyaev (1973) viscosity parameter). However, if $M_{\rm disk}/M \gtrsim (H/R) \sim 0.003$, the disk is self-gravitating and forms stars instead of accreting. Therefore, for efficient black hole growth, $\gtrsim 10^{2-3}$ individual accretion events are required, each of which contributes only a small fraction of M and lasts $\lesssim 10^6$ yr (King et al. 2008), implying that the accretion disks have radii $R_{\rm disk} \lesssim 0.003$ pc. Yet the gas that the hole must eventually accrete, which can be of order $10^{8-9}\,M_\odot$, must occupy a far larger region $R_{\rm gas} \sim$ 10–100 pc.

So the missing element in current treatments is a connection between these scales, telling us how gas falls from a region of size $R_{\rm gas}$ to make a succession of disks at scales $\sim R_{\rm disk}$. In numerical simulations of galaxy evolution, Bondi (1952) accretion is a popular choice, but has several critical drawbacks.

Two of these are crucial. The first is that in reality all gas has significant angular momentum, and so cannot fall in radially, in the way envisaged for Bondi accretion. Angular momentum is the main barrier to accretion. However, since $R_{\rm gas} \lesssim$ scale height of the interstellar medium (ISM), the cold gas in this region is probably not in large-scale rotation, i.e., has a distribution of (partly) opposing angular momenta with a small net angular momentum. Therefore, a way of canceling these opposing angular momenta would greatly enhance accretion.

A second serious problem in using the Bondi formula is its implication that gas falls toward the black hole because of the destabilizing influence of its gravity. But the hole's mass is so small compared to that of even a small region of the galaxy that this is implausible. As we remarked above, the property of the hole which is highly significant for the galaxy is not its mass M, but its binding energy $\eta c^2 M$, where $\eta \simeq 0.1$. In mass terms, the hole is typically only one part in about 10^{-3} of the galaxy bulge stellar mass $M_{\rm bulge}$ (Häring & Rix 2004). But for binding energies the situation is reversed: a hole of mass $10^8\,M_\odot$ has $\eta c^2 M \sim 10^{61}$ erg, while the bulge binding energy is $\sim\sigma^2 M_{\rm bulge} \sim 10^{58}$ erg for a typical velocity dispersion $\sigma \simeq 200\,{\rm km\,s^{-1}}$ (this disparity is even bigger for smaller SMBH if these follow the scaling relations).

This suggests that the cause of black hole accretion ultimately involves its effects on the galaxy, i.e., feedback. We already know quite a lot about black hole feedback in galaxies, and how it produces the SMBH–galaxy scaling relations. What is important for our purposes here is that the feedback is carried by quasi-spherical winds driven by radiation pressure; these are detected via blueshifted X-ray iron absorption lines (e.g., Pounds et al. 2003a, 2003b; Tombesi et al. 2010, 2011). The winds have momentum scalars $\dot{M}_{\rm out} v \simeq L_{\rm Edd}/c$, where $L_{\rm Edd}$ is the Eddington luminosity of the hole, $\dot{M}_{\rm out}$ is the wind outflow rate, and $v \sim \eta c$ its velocity (King & Pounds 2003). The winds interact with the host galaxy by shocking against its interstellar gas, giving initial post-shock temperatures $\sim 10^{10}$ K. While the SMBH is growing, these shocks lie close to the hole. Here the much cooler ($\sim 10^7$ K) radiation field produced by accretion removes most of the shock energy through the inverse Compton effect (King 2003).

So only the wind ram pressure, i.e., the momentum rate $L_{\rm Edd}/c$ mentioned above, is communicated to the host ISM (these are called "momentum-driven" flows). This thrust can push the host ISM only modestly outward, and is apparently unable to prevent the hole from growing. But once the hole mass reaches the M–σ scaling relation, i.e.,

$$M = M_\sigma = \frac{f_{\rm g}\kappa}{\pi G^2}\sigma^4, \tag{2}$$

with $f_{\rm g}$ the local gas fraction, the wind shocks are able to move far away from the hole (King 2003, 2005), beyond the critical radius $R_{\rm cool} \sim 0.5$ kpc where the radiation field of the accreting black hole becomes too dilute to cool the shocked wind. This now expands adiabatically ("energy-driven" flow), sweeping the host ISM before it at high speed (~ 1000 km s^{-1}) and largely clearing the galaxy bulge of gas (Zubovas & King 2012). This terminates black hole growth, leaving the hole near the mass M_σ (Equation (2)).

2. FEEDBACK CAUSES FEEDING

This sequence shows that the growth of the SMBH toward the M–σ relation is characterized by quasi-spherical momentum-driven outflow episodes which push the interstellar gas out, but do not unbind it. This changes the dynamical state of the ISM in two important ways. First, the SMBH driven wind does not transfer angular momentum to the gas, but increases its gravitational energy. This results in a *decrease* of the typical pericentric radius of the gas. Second, gas with differing angular momenta is pushed together, leading to (partial) cancellation.

When a black hole accretion episode ends, the outward thrust supporting the gas against gravity drops, and it must fall back from the radius $R_{\rm shell}$ of the swept-up region. Clearly, this infall is unlikely to be spherically symmetric. Instead, individual clumps or high-density regions fall on ballistic orbits. Because of the cancellation of angular momentum and the increase of gravitational energy during the outflow phase, these orbits are highly eccentric with pericenters much closer to the hole than the radii from which the gas was originally swept up during the wind feedback phase. On such eccentric orbits, any gas cloud is likely to be tidally stretched, forming a stream, in particular near pericenter.

We now estimate the resulting density of clouds/streams on such orbits. Consider a population of clouds/streams orbiting with the same peri- and apocentric radii $R_\pm$, and hence with the same orbital energy and specific angular momentum:

$$E = \frac{R_+^2\Phi_+ - R_-^2\Phi_-}{R_+^2 - R_-^2}, \qquad L^2 = \frac{2R_+^2R_-^2(\Phi_+ - \Phi_-)}{R_+^2 - R_-^2}. \tag{3}$$

Here, $\Phi_\pm \equiv \Phi(R_\pm)$, where $\Phi(R) = -GMR^{-1} + \Phi_{\rm bulge}(R)$ is the total gravitational potential. Neglecting collisions and internal shocks, the phase-space density of clouds/streams is conserved and simply the product of delta functions in E and L^2. Integrating it over all velocities yields the spatial density:

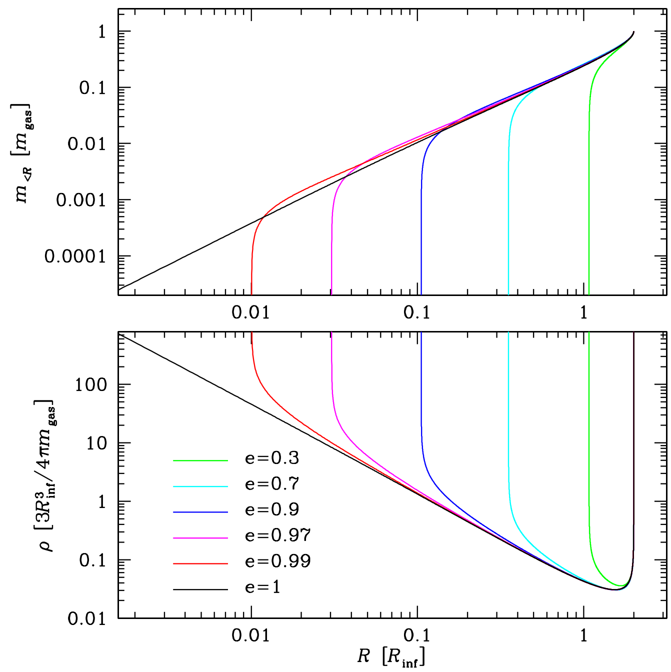

Figure 1. Enclosed mass (top) and mean density (bottom) of a population of clouds/streams orbiting the hole with the same apocentric radius $R_+ = 2R_{\rm inf}$ (corresponding to $M \approx M_\sigma/2$, for other choices the picture is very similar) but different eccentricities $e = (R_+ - R_-)/(R_+ + R_-)$ as indicated. The bulge was modeled as an isothermal sphere.

(A color version of this figure is available in the online journal.)

$$\rho(R) = \frac{m\,C}{R\sqrt{2R^2(E - \Phi(R)) - L^2}} \tag{4}$$

with

$$C^{-1} \equiv 4\pi \int_{R_-}^{R_+} \frac{R\,dR}{\sqrt{2R^2(E - \Phi(R)) - L^2}}, \tag{5}$$

where m is the total gas mass. We identify the apocenter with the radius of the initially swept-up shell, $R_+ = R_{\rm shell}$, and numerically evaluate C and the mass $m_{<R} = 4\pi\int_{R_-}^{R}\rho\,R^2dR$ enclosed at any time within radius R. The resulting density and enclosed mass are plotted in Figure 1 for various pericenters but with the apocenter fixed at $R_+ = 2R_{\rm inf}$ with

$$R_{\rm inf} \equiv GM/\sigma^2, \tag{6}$$

the radius of the hole's sphere of influence. Because eccentric orbits have a long residence time near apocenter, the density is maximal there and most of the gas is now further from the hole than before, in a kind of thick shell near $R_{\rm shell}$. However, the infalling gas creates a second density maximum near pericenter, where the clouds/streams tend to collide with probability $\propto \rho^2$ and with significant relative velocity. Near apocenter, on the other hand, collisions are not only less likely (because the orbiting clouds simply return near to their initial position, avoiding each other) but also have modest relative velocities and thus do not lead to cancellation of angular momentum.

These high-impact-velocity collisions near pericenter (which are neglected in Figure 1) lead to accretion-disk formation because the gas loses energy much faster than angular momentum, a process familiar from accretion in close binary systems. The colliding gas must shock and lose much of its orbital energy to cooling. In addition, the collisions may cancel some, potentially most, of the angular momentum, creating a cascade of ever smaller but less eccentric orbits. Ultimately, gas on the innermost orbits circularizes and forms a disk. If more gas penetrates to this radius, the disk is destroyed but quickly replaced by an even smaller one. Moreover, any misalignment of the disk angular momentum with the black hole spin results in disk tearing, when angular-momentum cancellation leads to a further reduction of the inner disk radius by a factor 10–100 (Nixon et al. 2012).

This whole process is rather complex and chaotic, but certainly has the potential to transfer some of the gas from $R_{\rm gas} \sim$ 10–100 pc into an accretion disk at $R_{\rm disk} \sim 0.001$–0.01 pc, where standard viscosity-driven accretion physics takes over the mass transport, and feeds the SMBH on a timescale of $\sim 10^6$ yr.

3. THE FEEDING RATE

The fundamental feature of our picture is that once central accretion (and hence feedback) slows, gas is no longer supported against gravity. This suggests that during the chaotic infall phase, gas feeds a small-scale accretion disk around the SMBH at some fraction of the dynamical infall rate:

$$\dot{M}_{\rm feed} \lesssim \dot{M}_{\rm dyn} \simeq \frac{f_{\rm g}\,\sigma^3}{G}. \tag{7}$$

For an SMBH with mass M close to M_σ, this exceeds the Eddington accretion rate $\dot{M}_{\rm Edd}$ by factors $\sim$10–100 at most (King 2007).

This feeding rate should characterize the rapid growth phases for the SMBH. For gas fractions $\gtrsim 0.1$ it implies disk feeding at rates a few times $\dot{M}_{\rm Edd}$. This is likely to result in the following scenario (cf. King & Pringle 2006; 2007). The outer parts of the disk may become self-gravitating and form stars, while the remaining gas flows inward under the disk viscosity at slightly super-Eddington rates. This leads to SMBH accretion at about $\dot{M}_{\rm Edd}$, and similar mass outflow rates, with momentum scalars $\dot{M}_{\rm out} v \simeq L_{\rm Edd}/c$ (King & Pounds 2003). This fits self-consistently with the feedback needed to give the observed M–σ scaling relation (King 2003, 2005).

Once central accretion stops, the SMBH should be quiescent for the sum of the infall timescale R_+/σ and the viscous timescale (Equation (1)). In general infall is more rapid, so the controlling timescale is probably viscous and depends critically on the radius $R_{\rm disk}$ at which the chaotic infall process places the disk.

We note that in our picture, both the precise value of the mass feeding rate and its duty cycle are determined by essentially stochastic processes. This makes it difficult to go beyond the simple estimates given here either analytically or numerically. We return to this problem in the last section.

4. BLACK HOLE OBSCURATION

We expect this same mechanism to produce the putative accretion "torus" at radii larger than $R_{\rm disk}$. This structure is postulated (Antonucci & Miller 1985; Antonucci 1993) to cover a large solid angle, obscuring the hole along many lines of sight, and so accounting for the populations of unobscured (Type I) and obscured (Type II) active galactic nuclei (AGNs). The main problem in understanding the torus in physical terms is that it must consist of cool material, which by its nature cannot form a vertically extended disk or torus. However, a large solid angle is natural if much of this obscuring gas is not yet settled into a disk, but still falling in on a range of orbits of very different inclinations. The column density $\Sigma = \int \rho\, dR$ of a population of gas clouds/streams with total mass m and common apo- and pericentric radii is

$$\Sigma \sim \frac{m}{2\pi (R_- + R_+)\sqrt{R_- R_+}}. \tag{8}$$

(using Equation (4) with $\Phi_{\rm bulge} = 0$). This diverges for small pericentric radii R_-, so the black hole must be obscured either completely or, more probably, for many lines of sight and/or extended periods of time. In fact, the obscuring matter may not be in form of a torus at all but merely a collection of clouds/streams orbiting the hole on eccentric orbits.

Whatever the geometry of the obscuring matter, our model renders the standard geometrical explanation for AGN unification (Antonucci 1993) time-dependent, since the orientation of that matter changes randomly over time and because we expect cyclically recurring inflow phases. This is in line with observational evidence of occasional changes between Seyfert types (e.g., Alloin et al. 1985; Shappee et al. 2013).

5. THE CENTRAL BUBBLE

Our discussion so far has not specified the physical scale $R_{\rm shell}$ where the momentum-driven outflows are typically halted. Our feeding mechanism works independently of this scale, but it may set the duty cycle and orientation of the individual accretion disk episodes. We note that King & Pounds (2013) have recently suggested that radiation pressure from the central active nucleus tends to create a shell of gas at a characteristic radius $R_{\rm tr} \sim 50\,(\sigma/200\,{\rm km\,s^{-1}})^2$ pc, at which the gas becomes transparent to the radiation from the accretion disk.

This is larger than the radius (Equation (6)) of the sphere of influence by a factor M_σ/M and the shell's mass is comparable with the final mass M_σ of the hole. In this picture, momentum-driven outflows must be halted here, as their inertia is of course far smaller. This means that $R_{\rm shell} \simeq R_{\rm tr}$. This idea agrees with observations of warm absorbers, which can be interpreted as arrested momentum-driven outflows.

6. DISCUSSION

We have suggested that black hole feeding is ultimately caused by feedback. By elongating the gas orbits and promoting collisions, this causes cancellation of opposed specific gas angular momenta, allowing accretion disks to form at small distances from the black hole, where they can feed the hole on timescales close to Salpeter (1964). This is different from a situation where the gas is initially pressure supported, when cooling and collisions of the resulting condensations can lead to turbulent infall (Gaspari et al. 2013). Our picture explains a number of other aspects.

As we have shown above, a near-toroidal topology for obscuring gas is a natural result. It is also clear that the orientation of the accretion structure (disk + "torus") cannot be constant over time, but must be essentially random. This is just the situation envisaged in the picture of chaotic accretion suggested by King & Pringle (2006, 2007), which results in relatively low black hole spins. This implies rapid mass growth and low gravitational-wave recoil velocities for merging black holes. The impact of the black hole wind on the gas which ultimately falls in may cause some of it to form stars, and this can also happen in the collisions during gas infall. Of course, any gas converted to and/or heated by stars is prevented from participating in the black hole feeding. However, at each feeding cycle only a small fraction of the gas within R_{shell} is required to reach R_{disk}, and only gas locked in stellar remnants and dwarfs is ultimately prevented from accreting.

Because angular momentum has been largely canceled, such newly formed stars fall in on near-parabolic orbits. This has several consequences. First, stars coming too close to the hole create visible tidal disruption events (Rees 1988); second, tidal dissociation of close binaries produces hypervelocity stars (Hills 1988); finally, massive stars which escape these fates inject metal-enriched gas into their surroundings. In any plausible picture most of this gas remains near to the hole, and could undergo repeated star formation. This may be the origin of the high chemical enrichment observed in AGN spectra (Shields 1976; Baldwin & Netzer 1978; Hamann & Ferland 1992; Ferland et al. 1996; Dietrich et al. 1999, 2003a, 2003b; Arav et al. 2007).

We note that the idea of feedback-stimulated feeding opens the possibility of runaway growth: the black hole forages for its own food, and grows still faster. Given an abundant food supply (i.e., $f_{\mathrm{g}} \gtrsim 0.1$) this growth is stopped only as the hole reaches the limiting M–σ mass and drives all the food away. A runaway SMBH like this would of course have a tendency to grow at the Eddington rate for most of its (short) feeding frenzy. This may explain very massive SMBH observed at high redshifts (e.g., Barth et al. 2003; Willott et al. 2003; Fan et al. 2003; Mortlock et al. 2011). Here the close proximity of all galaxies means that many are likely to be gas-rich (i.e., $f_{\mathrm{g}} \gtrsim 0.1$) because of mergers, so runaways are favored.

One interesting aspect of the proposed mechanism is the mutual dependence of feeding and feedback on each other. Clearly, this whole process must be started by some initial accretion which was *not* triggered by feedback, but by sufficient gas coming within $\lesssim$0.001 pc of the infant hole. Such an event could be triggered by a galactic merger, but must be relatively rare. This implies that early SMBH formation may be somewhat random, but more likely in frequently perturbed/merging galaxies.

Conversely, if the SMBH's neighborhood at $R \lesssim R_{\mathrm{gas}}$ acquires some net rotation, for example, during a merger, then the distribution of angular momenta is unlikely to allow for angular-momentum cancellation. In such a situation, the SMBH suffers from starvation. Despite sitting tantalizingly close to its food, it cannot reach it nor bring it down easily. However, if the rotating gas can cool, it will form a disk (and possibly stars), clearing most of the space and opening the possibility for re-starting the feeding cycle.

Also, our proposed feeding mechanism will not work efficiently if the feedback is dominated by a collimated jet rather than wide-angle outflows. This is obvious if the impact shocks are efficiently cooled (momentum-driven flow) as the jet simply carves a narrow hole in the gas it impacts. If the shocks do not cool (energy-driven), their effect is wider but still unlikely to cause feeding in the way described here.

Finally, we note that the episodic inflows and outflows of a fraction $f_{\mathrm{g}} \sim 0.1$ of matter at velocities well above σ entail abrupt changes in the gravitational potential in the inner $R_{\mathrm{shell}} \sim 10$–100 pc. Therefore, the growth of the SMBH is not an adiabatic process for the dynamics of collisionless matter on these scales. Instead, the abrupt variations in the potential redistribute the orbital actions. This renders the central dark-matter density smaller than current estimates (by, e.g., Young 1980; Quinlan et al. 1995) based on the adiabatic assumption, though possibly still larger than in absence of a SMBH. This implies a significant reduction in the expected dark-matter annihilation signal from SMBH hosting galaxy centers.

7. A SUBGRID RECIPE

We have suggested that feeding of SMBHs may in many cases be stimulated by feedback. A practical question is how one might implement this process in simulations of galaxy formation which cannot resolve the hole's sphere of influence, let alone the dynamics and cooling of infall and outflow, and instead must use a subgrid recipe. Clearly, any Bondi-like subgrid recipe adapted to account for the angular momentum of the gas at $\gtrsim R_{\mathrm{gas}}$ cannot adequately describe these dynamics. Instead a completely different approach is required.

We have seen that feedback-induced feeding generally occurs at a fraction of the dynamical infall rate (Equation (7)) when it operates. This is generally slightly super-Eddington (for $f_{\mathrm{g}} \gtrsim 0.1$). This in turn makes the SMBH grow at about the Eddington rate, and rejects the remainder of the mass in a wind, which is what causes the feedback. If $M < M_\sigma$ we know in reality this will result in momentum-driven feedback, which keeps the accretion going, and does not blow the gas away. Once $M \geqslant M_\sigma$, the feedback changes character to energy-driven and terminates SMBH mass growth.

Given the discussion above, a suitable subgrid recipe is as follows. Grow M from surrounding gas at the rate $\dot{M} = \min\{\epsilon \dot{M}_{\mathrm{dyn}}, \dot{M}_{\mathrm{Edd}}\}$ (see Equation (7)) with $\epsilon \sim 0.1$. If $M < M_\sigma$, *neglect* feedback. If $M \geqslant M_\sigma$, deposit *energy* into the surrounding gas at the rate $(\eta/2)c^2\dot{M} \simeq 0.05c^2\dot{M}$ (Zubovas & King 2012).

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