BLACK HOLE SPIN AND GALACTIC MORPHOLOGY

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ABSTRACT

We investigate the conjecture by Sikora, Stawarz, and Lasota that the observed active galactic nuclei (AGNs) radio loudness bimodality can be explained by the morphology-related bimodality of black hole spin distribution in the centers of galaxies: central black holes (BHs) in giant elliptical galaxies may have (on average) much larger spins than black holes in spiral/disk galaxies. We study how accretion from a warped disk influences the evolution of black hole spins and conclude that within the cosmological framework, where the most massive BHs have grown in mass via merger-driven accretion, one indeed expects most supermassive black holes in elliptical galaxies to have on average higher spin than black holes in spiral galaxies, where random, small accretion episodes (e.g., tidally disrupted stars, accretion of molecular clouds) might have played a more important role.

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

1. INTRODUCTION

It has been known for many years that the radio loudness of AGNs hosted by disk galaxies is on average 3 orders of magnitude lower than the radio loudness of AGNs hosted by giant ellipticals (see Xu et al. 1999, and references therein). However, as shown by Hubble Space Telescope (HST) observations, such a galaxy morphology–radio loudness correspondence is not “one-to-one”: both radio-quiet and radio-loud very luminous quasars are hosted by giant ellipticals (Floyd et al. 2004). On the other hand, the popular version of the so-called “spin paradigm,” asserts that powerful relativistic jets are produced in AGNs with fast rotating black holes (Blandford 1990), implying that BHs rotate slowly in radio-quiet quasars, which represent the majority of quasars. However, such conjecture, at least in its basic interpretation, is in conflict with the high average BH spin in quasars deduced from the high average radiation efficiency of quasars using the “Soțan argument” (Soțan 1982; Wang et al. 2006, and references therein).

Parallel studies of radio emission from X-ray binaries showed that at high accretion rates production of jets is intermittent (Gallo et al. 2003), and that this intermittency can be related to transitions between two different accretion modes (Livio et al. 2003), Nipoti et al. (2005), and Kör ding et al. (2006) to postulate the existence of a similar intermittency of jet production in quasars and formulate an “accretion paradigm” according to which the radio loudness is entirely related to the states of accretion disks. However, Sikora et al. (2007) found that on the radio loudness–Eddington-ratio plane AGNs form two parallel sequences whose occurrence cannot be explained by the “accretion paradigm” (see also Terashima & Wilson 2003; Chiaberge et al. 2005; Panessa et al. 2007). Sikora et al. (2007) therefore proposed a revised version of the spin paradigm, suggesting that giant elliptical galaxies host, on average, black holes with spins larger than those hosted by spiral/disk galaxies.

This morphology-related radio dichotomy breaks down at high accretion rates, where the dominant fraction of luminous quasars hosted by elliptical galaxies is radio quiet. This radio quietness occurs in quasars with high spin values. In such systems with high accretion rates the intermittency is related to the conditions of production of collimated jets, in agreement with what is found in X-ray binaries and with the Soțan argument. It should be emphasized that even if the production of powerful relativistic jets is conditioned by the presence of fast rotating BHs, it also depends on the accretion rate and on the presence of disk MHD winds required to provide the initial collimation of the central Poynting flux dominated outflow.

In this article we will examine under which condition the cosmological evolution of BHs in galaxies may lead to low spins in disk galaxies and high spins in more massive ellipticals. To put our investigation in the relevant context, we will first recall why the value of a black hole’s spin might be of fundamental importance for relativistic jet launching. Assuming that relativistic jets are powered by rotating black holes through the Blandford-Znajek mechanism, Blandford (1990) suggested that the efficiency of jet production is determined by the dimensionless black hole spin, $\hat{a} \equiv J_B/M^2$; where $J_B$ is the angular momentum of the black hole. If true, this could explain the very wide range of radio loudness of AGNs that look very similar in many other aspects by attributing it to a corresponding black hole spin distribution. This so-called “spin paradigm” was explored by Wilson & Colbert (1995), who assumed that the black hole spin evolution is determined mainly by mergers. They claimed that mergers of black holes, following mergers of galaxies, lead to a broad, “bottom-heavy” distribution of the spin, consistent with a distribution of quasar radio loudness. However, this claim was challenged by Hughes & Blandford (2003), who showed that mergers cannot produce the required fraction of black holes with high spins and concluded that accretion of matter is essential in determining black hole spins. In this case, however, as noticed earlier by Moderski & Sikora (1996) and Moderski et al. (1998; hereafter MSL98), one encounters the difficulty of maintaining a sufficient number of black holes at the required low spin, the spin-up by accretion disks being so efficient. MSL98 could match the distribution of radio loudness with the spin distribution only by feeding holes with very small randomly oriented accretion events,

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i.e., by accretion events forming corotating and counterrotating disks with the same probability.

MSL98 also addressed the problem of the spin overflipping due to the Bardeen-Petterson effect. When an accretion disk does not lie in the equatorial plane of the BH, that is, when the angular momentum of the accretion disk is misaligned with respect to the direction of \( J_h \), the dragging of the inertial frame causes a precession that twists the disk plane due to the coupling of \( J_h \) with the angular momentum of matter in the disk. The torque tends to align the angular momentum of the matter in the disk with that of the black hole, thus causing the inclination angle between the angular momentum vectors to decrease with decreasing distance from the BH, forcing the inner parts of the accretion disk to rotate in the equatorial plane of the BH (Bardeen & Petterson 1975). Sustained accretion from a twisted disk would align the BH spin (and the innermost equatorial disk) with the angular momentum vector of the disk at large radii (Scheuer & Feiler 1996). If the disk was initially counterrotating with respect to the BH, a complete overflip would eventually occur, and then accretion of corotating material would act to spin up the BH (Bardeen 1970).

MSL98 concluded that the Bardeen-Petterson effect can be neglected, because the alignment time (10 yr; Rees 1978) is longer than the duration of a single accretion event. Later, however, a series of papers put into doubt the validity of Rees’s estimate (Scheuer & Feiler 1996; Natarajan & Pringle 1998). This framework was recently investigated by Volonteri et al. (2005), who argue that the lifetime of quasars is long enough that angular momentum coupling between black holes and accretion disks through the Bardeen-Petterson effect effectively forces the innermost region of accretion disks to align with black hole spins (possibly through spin flips), and hence all AGN black holes should have large spins.

Recently King et al. (2005) pointed out that under some conditions the alignment torque can lead to disk-hole counteralignment reactivating the debate. The counteralignment process was numerically simulated by Lodato & Pringle (2006).

Here we reanalyze the alignment problem in view of all these recent results. We explore the likely outcomes of accretion episodes that grow black holes along the cosmic history, and determine under which conditions black holes in disk galaxies end up having low spins.

2. ASSUMPTIONS

The dynamics involving a spinning black hole accreting matter from a thin disk whose angular momentum is not aligned with the spin axis has been studied in a number of papers (e.g., Papaloizou & Pringle 1983; Pringle 1992; Scheuer & Feiler 1996; Natarajan & Pringle 1998). A misaligned disk is subject to the Lense-Thirring precession, which tends to align the inner parts of the disk with the angular momentum of the black hole. The outer regions of the disk are initially inclined with respect to the hole’s axis, with a transition between alignment and misalignment occurring in between at the so-called “warp radius” (see below). The direction of the angular momentum of the infalling material changes direction as it passes through the warp. In our calculations we assume that the black hole spin evolution is determined by accretion only. Volonteri et al. (2005) have shown that BH mergers play a subdominant role in the global spin evolution.

2.1. Viscosities

Despite many efforts the problem of warped disks, especially in the nonlinear regime, has yet to be solved. Therefore, the characteristic scales of the problem are subject to several uncertainties. The main quantity of interest is the “warp radius” \( R_w \) defined as the radius at which the timescale for radial diffusion of the warp is comparable to the local dragging of inertial frame (“Lense-Thirring” in the weak-field approximation) time \( \left( \alpha c R^2 / R^3 \right)^{-1} \) (Wilkins 1972). The timescale for the warp radial diffusion can be written as

\[
t_w \approx \frac{R^2_{w}}{\nu_2},
\]

where \( \nu_2 \) is a viscosity characterizing the warp propagation, which can be different from the accretion driving viscosity, \( \nu_1 \), which is responsible for the transfer of the component of the angular momentum parallel to the spin of the disk. The relation between \( \nu_1 \) and \( \nu_2 \) is the main uncertainty of the problem, assuming of course that such two-viscosity description is adequate at all. Describing \( \nu_1 \) by the Shakura-Sunyaev parameter \( \alpha \), one can show (Papaloizou & Pringle 1983) that the regime in which \( H/R < \alpha \approx 1 \), (\( H \) being the disk thickness) one has \( \nu_1/\nu_2 \approx \alpha^2 \). In such a case the accretion time \( t_{acc} \approx R_w/\nu_1 \) would be much longer than the warp diffusion time \( t_w \).

However, such a description can be questioned on several grounds. First, is the \( \alpha \approx 1 \) appropriate for high-rate accretion onto AGN black holes? There are no reliable estimates of this parameter for AGNs, but outbursts of low-mass X-ray binaries (LMXBs) suggest that in hot accretion disks \( \alpha \approx 0.1 \) (see, e.g., Dubus et al. 2001). In such a case \( \nu_1 \) is comparable to \( \nu_2 \) (Kumar & Pringle 1985). Second, even if the two viscosities are different, is \( t_w \) the relevant time for black hole realignment? This is not clear, since this latter process is very dissipative and could be controlled by accretion and not warp propagation.

2.2. Relevant Radii

During the accretion process, the angular momentum of the disk at the warp location sums up with that of the black hole, so the angle between the angular momentum of the outer disk and the BH spin changes. King et al. (2005) suggest that the condition of alignment or counteralignment can be expressed as a function of the angular momenta of the hole and of the disk: \( J_h \) and \( J_d \). The counteralignment condition depends on the ratio \( 0.5 J_d / J_h \), to be compared with the cosine of the inclination angle, \( \phi \). If \( \cos \phi < -0.5 J_d / J_h \), the counteralignment condition is satisfied. King et al. (2005), however, leave the definition of \( J_d \) vague; indeed, they suggest that \( J_d \) is the angular momentum of the disk inside a certain radius \( R_f \) such that \( J_d(R_f) = J_h \). First, this is not a useful operational definition, because in this case \( \cos \phi < -0.5 J_d / J_h = -0.5 \) is a static condition, which does not depend on the properties of the black hole or of the accretion disk. Second, matter contained within radius \( R_f \) cannot transfer all its angular momentum to \( R_w \) but only a fraction \((R_w/R_f)\frac{1}{2} J_d(R_f)\).

Therefore, a more natural radial scale in the problem is the warp radius \( R_w \), and in the following we will assume that \( J_d = J_d(R_w) \) (note that we share this choice with Lodato & Pringle 2006).

3. METHOD

We explore the dependence of the alignment timescale in a Shakura-Sunyaev disk on viscosity \( \nu_2/\nu_1 \), black hole mass \( M_{BH} \), misalignment angle, Eddington ratio, and accreted mass \( m \). Some articles on this subject use the solution of Collin-Souffrin & Dumont (1990); however, in the view of the basic uncertainties of the problem we decided to use the less refined solution of Shakura & Sunyaev (1973).
Assuming a Shakura-Sunyaev disk ("middle region"), the warp radius (in units of the Schwarzschild radius $R_S$) can be expressed as

$$ R_w = 3.6 \times 10^3 \alpha^{3/8} \left( \frac{M_{\text{BH}}}{10^8 M_\odot} \right)^{1/8} \frac{f_{\text{Edd}}^{1/4}}{f_{\text{Edd}}} \left( \frac{L}{L_{\text{Edd}}} \right)^{-5/8} \alpha^{-1/2}, \quad (2) $$

where $f_{\text{Edd}} \equiv M c^2 / L_{\text{Edd}}$.

We can then define the accretion timescale

$$ t_{\text{acc}} = \frac{R_w}{v_1} = \left( 3 \times 10^6 \text{ yr} \right) \alpha^{-3/2} \left( \frac{\nu_2}{\nu_1} \right)^{-7/8} \times \hat{a}^{7/8} f_{\text{Edd}}^{-3/4} \left( \frac{M_{\text{BH}}}{10^8 M_\odot} \right)^{11/8} \quad (3) $$

(3)

where $v_1 = \alpha H^2 \Omega_K$ was used) and the timescale for warp propagation,

$$ t_w = \frac{v_1}{\nu_2} t_{\text{acc}}. \quad (4) $$

The ratio of angular momenta of the disk at $R_w$, defining $M_d(R_w) = M t_{\text{acc}}(R_w)$, and that of the black hole is

$$ J_d = \frac{M_d}{\hat{a} M_{\text{BH}}} \left( \frac{R_w}{R_S} \right)^{1/2} = 2 \times 10^{-9} f_{\text{Edd}} \left( \frac{t_{\text{acc}}}{1 \text{ yr}} \right) \left( \frac{R_w}{R_S} \right)^{1/2} \hat{a}^{-1}. \quad (5) $$

4. RESULTS

4.1. Single Accretion Episodes

We first discuss the behavior of the disk+BH system during the alignment process. The scheme we adopt is as follows:

1. For a BH with initial mass $M_{\text{BH}0}$, we determine the initial conditions: warp radius, $R_w$, timescale for warp propagation, $t_w = R_w^2 / \nu_2^2$ ($\nu_2$ is chosen either coincident with $\nu_1$ or $\nu_2 = \nu_1 / \alpha^2$), accretion timescale for material at the warp radius, $t_{\text{acc}} = R_w^2 / \nu_1$, angular momentum of the hole and of the disk at $R_w$, $J_h$ and $J_d$.

2. Using the King et al. (2005) condition for misalignment, we determine if the BH and the inner disk are aligned or counteraligned (counteraligned if $\cos \phi < -0.5 J_d / J_h$).

3. Over time steps $\Delta t = t_{\text{acc}}(R_w)$ we compute the necessary quantities at the end of each step: increase in black hole mass due to accretion, new BH spin (following Bardeen 1970; where the counteralignment or alignment conditions are taken into consideration, i.e., BHs can be spun down or up), new $J_h$, new $R_w$, new $J_d$, and new angle between $J_h$ and $J_d$ (vectorial sum). In every time step the disk within $R_w$ is consumed.

Figures 1 and 2 give examples of single accretion episodes for different initial angles between the angular momentum vector of the outer (not warped) portions of the accretion disk and the black hole spin. They show the evolution of the spin magnitude and inclination as computed for $\nu_2 = \nu_1$ (Fig. 1) and $\nu_2 / \nu_1 = 1 / \alpha^2$ (Fig. 2).

As we can see, the alignment timescale is basically independent of the misalignment angle. To significantly modify the BH spin, one has to bring to $R_w$ an amount of angular momentum comparable to $J_h$. Therefore, if $J_h > J_d(R_w)$, then

$$ t_{\text{align}} \propto J_h \left( \frac{f_{\text{acc}}(R_w)}{J_d(R_w)} \right). \quad (6) $$

Since $J_h \propto \hat{a} M_{\text{BH}}(R_w)^{1/2}$, $J_d(R_w) \propto M_d(R_w)(R_w)^{1/2}$, and $M_d(R_w) = M t_{\text{acc}}(R_w)$, equation (6) gives (Rees 1978)

$$ t_{\text{align}} \propto \hat{a} \frac{M_{\text{BH}}}{M} \left( \frac{R_s}{R_w} \right)^{1/2}. \quad (7) $$

Defining the mass accreted during $t_{\text{align}}$ as $m_{\text{align}} = t_{\text{align}} \dot{M}$, one gets

$$ m_{\text{align}} \propto M_{\text{BH}} \hat{a} \left( \frac{R_s}{R_w} \right)^{1/2}. \quad (8) $$

Therefore, a series of many randomly oriented accretion events with accreted mass $m < m_{\text{align}}$ should result in a black hole’s spin oscillating around zero. For the opposite case of $m \gg m_{\text{align}}$ the black hole will be spun up to large positive spins; for $m \sim M_{\text{BH}}$ the hole will be spun up to $\sim 1$. Finally, note that since it is reasonable to assume that $m_{\text{align}} \ll M_{\text{BH}}$ the existence of AGNs hosting black holes with $\sim 1$ is rather unlikely.

Our calculations show that it is difficult to avoid high spin for the most massive black holes. For large BH masses, the accretion timescale is very long, and consequently the warp radius and the angular momentum within the warp radius, $J_d$, are large. If $J_d > 2 J_h$, then the value $0.5 J_d / J_h > 1$ and the counteralignment condition cannot be satisfied for any angle. This condition corresponds to

$$ M_{\text{BH,max}} > \left( 6.2 \times 10^8 M_\odot \right) \alpha^{28/23} \left( \frac{\nu_2}{\nu_1} \right)^{19/23} f_{\text{Edd}}^{-2/23} \hat{a}^{-3/23}. \quad (9) $$

If $\nu_2 = \nu_1$, $M_{\text{BH,max}}$ is of order $10^{-7} - 10^{-8} M_\odot$, for most sensible choices of $\alpha$ and $f_{\text{Edd}}$. In this case the most massive black holes force accretion to occur from aligned disks, therefore causing a systematic spin-up. If the warp propagation is instead better described by $\nu_2 = \nu_1 / \alpha^2$, $M_{\text{BH,max}}$ becomes exceedingly high and large accretion events can still act to spin down the black hole, provided $m < m_{\text{align}}$.

The condition expressed in equation (9) is true only if there is enough mass to fill the warp radius, that is, if the total mass of the disk is larger than

$$ M_{\text{d,min}} > 6.5 \times 10^5 M_\odot \left( \frac{M_{\text{BH}}}{10^8 M_\odot} \right)^{19/8} \alpha^{-3/2} f_{\text{Edd}}^{1/4} \hat{a}^{7/8} \left( \frac{\nu_2}{\nu_1} \right)^{-7/8}. \quad (10) $$

If the mass to be accreted by the BH in an episode is smaller than $M_{\text{d,min}}$, then $J_d < J_h$, and both alignment or counteralignment can happen.

4.2. Multi-Accretion Events

We then run a series of simulations in which we explore different parameters. We start with a small BH, $M_{\text{BH}0} = 10^5 M_\odot$, and have it grow by a series of accretion episodes. The accreted mass $m$ is randomly extracted from only one of two different distributions: (1) a distribution flat in $m$, with $m < 0.1 M_{\text{BH}0}$; (2) a distribution flat in $m$, with $m < 0.01 M_{\text{BH}0}$. The angle $\phi$ is extracted...
from a flat distribution $0 < \phi < \pi$ at the beginning of every accretion episode. Every simulation is composed by a large number of accretion episodes until one of the following conditions are met: $M_{\text{BH}} > 10^9 M_\odot$ or $t_{\text{tot}} > 10^{10}$ yr, that is, the total simulation time (total time a BH accretes to grow from the initial $M_{\text{BH}0} = 10^5 M_\odot$ mass to its final mass) is shorter than the age of the universe. During an episode where the BH accretes counteraligned material, the BH is spun down. If the black hole is spun down until its spin is zero, any subsequently accreted matter acts to spin the BH up again, although the direction of the spin axis is now reversed and aligned with the angular momentum of the disk. We ran 100 simulations for every parameter sets choice, and we traced the spins at the end of every accretion episode for all the accretion episodes in the simulations.

We have explored a wide range of parameters, and here we summarize our findings. We have varied the accretion rate from $f_{\text{Edd}} = 0.05$ to 1. If the accretion rate is low, the main caveat is that black holes do not reach high masses within the Hubble time; however, the efficiency of alignment is not strongly dependent on $f_{\text{Edd}}$ (see eq. [8]).

We have considered different black hole spins at birth from $\dot{a} = 10^{-3}$ to 0.9. After the BHs have changed their initial mass by about 1 order of magnitude, the distributions are indistinguishable from each other. During the first few $e$-foldings, however, the spin distribution is peaked around the black hole spin at birth.

We have also varied the viscosity parameter $\alpha$ (see § 2.1) and the relation between the viscosity characterizing the warp propagation ($\nu_2$) with respect to the viscosity responsible for the transfer of the component of the angular momentum parallel to the spin of the disk ($\nu_1$). When $\alpha$ is varied, but $\nu_2/\nu_1$ is kept fixed, the differences between the spin distributions are not large. A smaller $\alpha$ skews the distribution toward higher spins (cf. eqs. [2] and [8]).

One of the main parameters influencing the spin distribution is the relation between $\nu_2$ and $\nu_1$. If $\nu_2/\nu_1 = 1$, after the BHs have reached $m_{\text{BH}} \sim 10^6 M_\odot$, the spin distribution is dominated by rapidly spinning black holes. Equation (9) also shows that the most massive black holes force accretion to occur from aligned disks, therefore causing a systematic spin-up, unless very small parcels of material are accreted at every single accretion episode. If the warp propagation is instead better described by a high $\nu_2 = \nu_1/\alpha^{-1}$, $M_{\text{BH,max}}$ becomes exceedingly high, and all sorts of accretion events can still act to spin down the black hole, provided that $m < m_{\text{align}}$.  

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**Fig. 1.** Evolution of the misalignment angle between the angular momentum vector of the outer accretion disk and the BH spin (top panels) and of the magnitude of the BH spin (bottom panels) due to accretion of aligned material (spin-up) or counteraligned material (spin-down). The initial BH mass is $M_{\text{BH}0} = 5 \times 10^6 M_\odot$, the initial spin is $\dot{a} = 0.5$, $\nu_2 = \nu_1$, $\alpha = 0.1$, and the accretion rate is $f_{\text{Edd}} = 1$. The four curves show different initial misalignment angles (top to bottom: $\phi = 3, 2, 1, 0.5$ rad).
In fact, we confirm the results by MSL98, that is that the main parameter governing the distribution of BH spins is the amount of material accreted in a single accretion episode. This result is clear from Figures 3 and 4, which refer to different choices for the distributions of \( m \). Only if the mass accreted in one episode is smaller than \( m_{\text{align}} \) can the distribution of black hole spins remain flat. In Section 4.3 we discuss the likelihood of different \( m \) distributions in the light of evolutionary models for the BH population in a hierarchical cosmology.

4.3. Merger-Driven Accretion

We first present an evolutionary track for BH spins, where a BH grows by a sequence of randomly oriented accretion episodes in a merger-driven scenario. The BH mass evolutionary tracks are extracted from semianalytical simulations of BH growth that have been shown to reproduce the evolution of the BH population as traced by the luminosity function of quasars (Marulli et al. 2006, 2007; Volonteri et al. 2006). We focus here on two specific tracks, for a putative BH in an “elliptical” (E) galaxy, and one in a putative “disk” (D) galaxy (Fig. 5). Here the morphological classification is purely based on the frequency of major mergers, i.e., mergers between comparable mass galaxy systems, which are believed to contribute mainly to the spherical component of galaxies. A BH hosted in an elliptical galaxy should have experienced a major accretion event in connection with the last high-redshift major merger, which formed the elliptical galaxy as we see it now. Afterward, the galaxy (BH) has not grown in mass due to merger-driven star formation (accretion). In the case of the BH hosted in a disk galaxy, a small number of minor mergers might have happened after the last major mergers. These minor mergers are believed to be responsible for rebuilding the galaxy disk. In conjunction with these minor mergers, a small infall of gas can produce a relatively minor accretion episode onto the BH as well.

Along the evolutionary tracks for our BHs, we trace the joint evolution of accretion onto the BH, the dynamics of the accretion disk, and the consequences for the spin. The scheme we adopt is similar to the one described in Section 4.1. During an episode where the BH accretes counteraligned material, the BH is spun down until the spin is zero, and subsequently any accreted matter acts to spin the BH up again, although the direction of the spin axis is now reversed and aligned with the angular momentum of the disk.

In case “E,” the last accretion episode caused a large increase in the BH mass, following the major merger that created the elliptical itself (Hopkins & Hernquist 2006). During this episode the spin increased significantly as well, up to very high values. Let us remember here that in the extreme event of a maximally rotating hole spun down by retrograde accretion, the BH is braked...
after its mass has increased by the factor \( \sqrt{3}/2 \). Any mass accreted afterward spins up the black hole, and if the final mass increase is by a factor 3, the BH will end up maximally rotating again.

Let us now consider a “D”-type evolution. A BH would experience a series of small accretion episodes (triggered possibly by minor mergers), extending for a longer period of times. If these episodes are uncorrelated, that is if the inflow during a given episode is not aligned with the orientation of the spin of the BH, the randomization of the angle \( \phi \) over the (few) accretion episodes tends to spin down the BH.

We run a statistical sample of “E” and “D” track, for BHs hosted in large (i.e., Andromeda-size systems) galaxies. We find that, if only the merger-driven evolution is taken into account, BHs in elliptical galaxies are left with large spins. BHs in disk galaxies have, on average, slightly lower spins, however, the distribution is still peaked at large values (Fig. 6).

4.4. Short-Lived Accretion Events

Even minor mergers tend to trigger inflows of matter that are too large to lead to the series of short lived accretion events necessary to leave BHs with small spins (cf. the discussion in MSL98). Moreover, several observations suggest that single accretion events last \( \approx 10^5 \) yr in Seyfert galaxies, while the total activity lifetime (based on the fraction of disk galaxies that are Seyfert) is \( 10^8-10^9 \) yr (e.g., Kharb et al. 2006; Ho et al. 1997).

This suggests that accretion events are very small and very “compact.”

One type of random event that leads to short-lived accretion episodes is the tidal disruption of stars. One expects disks formed by stellar debris to form with a random orientation. Stellar disruptions would therefore contribute to the spin-down of BHs. Let us consider the maximal influence that feeding via tidal disruption of stars can have on spinning down a BH. The number of tidal disruptions of solar type stars in an isothermal cusp per billion years can be written as

\[
N_\ast = 4 \times 10^5 \left( \frac{\sigma}{60 \text{ km s}^{-1}} \right) \left( \frac{M_{\text{BH}}}{10^6 M_\odot} \right)^{-1}.
\]

Assuming that BH masses scale with the velocity dispersion, \( \sigma \), of the galaxy (we adopt here the Tremaine et al. 2002 scaling), we can derive the relative mass increase for a BH in 1 billion years:

\[
\frac{M_{\ast}}{M_{\text{BH}}} = 0.37 \left( \frac{M_{\text{BH}}}{10^6 M_\odot} \right)^{-9/8}.
\]

The maximal level of spin-down would occur assuming that all the tidal disruption events form counterrotating disks, leading to retrograde accretion (note that the mass of the debris disk is
much smaller than $M_{d, \text{min}}$, cf. eq. [10], so that counteralignment is allowed for any BH mass). Equation (12) shows that a small (say $10^6 M_\odot$) BH starting at $\dot{a} = 0.998$ would be spun down completely; on the other hand, the spin of a larger (say $10^7 M_\odot$) BH would not be changed drastically.

Early-type disks typically host faint bulges characterized by steep density cusps, both inside (Bahcall & Wolf 1976; Merritt & Szell 2006) and outside (Faber et al. 1997) the sphere of influence of the BH. In this environment, the rate of stars that are tidally disrupted by BHs (Hills 1975; Rees 1978) less massive than $10^8 M_\odot$ is non-negligible (Milosavljevic et al. 2006). Since in elliptical galaxies the central relaxation timescale is typically longer than the Hubble time, and the central density profile often displays a shallow core, tidal disruption of stars is unlikely to play a dominant role.

An additional feeding mechanism might be at work in gas-rich galaxies with active star formation. Compact self-gravitating cores of molecular clouds (MC) can occasionally reach subparsec regions, and may do it with random directions provided that the galactic disk is much thicker than the spatial scale of the BH gravity domination region (I. Shlosman 2007, private communication). Although the rate of such events is uncertain, we can adopt the estimates of Kharb et al. (2006) and assume that about $10^4$ of such events happen. We can further assume a lognormal distribution for the mass function of MC close to galaxy centers (based on the Milky Way case, e.g., Perets et al. 2007). We do not distinguish here giant MC and clumps, and for illustrative purpose we assume a single lognormal distribution peaked at $\log (M_{\text{MC}}/M_\odot) = 4$, with a dispersion of 0.75.

Figure 7 shows the possible effect that accretion of molecular clouds can have on spinning BHs. The result is, on the whole, similar to that produced by minor mergers of black holes (Hughes & Blandford 2003), that is a spin-down in a random walk fashion. The larger the BH mass, the more effective the spin-down.

In a gas-poor elliptical galaxy, however, substantial populations of molecular clouds are lacking (e.g., Sage et al. 2007), thus hampering the latter mechanism for the short-lived accretion events proposed.

5. DISCUSSION AND CONCLUSIONS

We have investigated the evolution of BH spins driven by accretion from disks with angular momentum vectors that can be misaligned with respect to the spin axis. We have assumed that accretion disks can be described by Shakura-Sunyaev $\alpha$-disks, and that when the angular momentum of the accretion disk is not aligned with the spin of the BHs, the disk itself is warped. The inner portions of the disks experience a Lense-Thirring torque, which tends to align the inner parts of the disk. The timescale of the Lense-Thirring precession increases faster with distance from
a BH than the timescale of warp propagation, and they equate at the so-called "warp radius," where a transition occurs from alignment to misalignment. King et al. (2005) pointed out that for highly misaligned disks, counteralignment rather than alignment can occur. The co- or counteralignment of the accretion disks has important consequences on the spin of BHs. A black hole accreting from prograde orbits (i.e., alignment case) is spun up by the coupling between the angular momentum of the infalling material and its spin (Bardeen 1970). If, instead, an initially spinning hole accretes from retrograde orbits (i.e., counteralignment case), it is spun down. An initially nonrotating BH gets spun up to a maximally rotating state ($\hat{a} = 1$) after reaching the mass $M_{\text{BH}} = \sqrt{6} M_{\text{BH}_0}$. A maximally rotating hole ($\hat{a} = 1$) gets spun down by retrograde accretion to $\hat{a} = 0$ after reaching the mass $M_{\text{BH}} = \sqrt{3/2} M_{\text{BH}_0}$. A 180° flip of the spin of an extreme Kerr hole will occur after $M_{\text{BH}} = 3 M_{\text{BH}_0}$.

It is therefore necessary that accretion episodes increase the mass of a BH by less than $M_{\text{BH}} = \sqrt{6} M_{\text{BH}_0}$, in order to keep the spin at low values if accretion preferentially occurs from prograde orbits. Natarajan & Pringle (1998) suggested that accretion indeed occurs from aligned disks (i.e., prograde orbits), as the timescale for disk alignment is much shorter than the timescale of the BH mass growth by $M_{\text{BH}} = \sqrt{6} M_{\text{BH}_0}$. King et al. (2005) suggested, however, that when the initial misalignment angle is large and $m$ is sufficiently small, counteralignment, rather than alignment, occurs and BHs can be spun down in a large fraction of the accretion episodes.

We have quantified here the likelihood of counteralignment and spin-down as claimed by King et al. (2005). We identify two main parameters influencing the distribution of BH spins: the distribution of the accreted mass, $m$, with respect to the mass of the BH, and the relation between $\nu_2$ and $\nu_1$, where $\nu_2$ is the viscosity characterizing the warp propagation, and $\nu_1$ is responsible for the transfer of the component of the angular momentum parallel to the spin of the disk. The viscosity $\nu_2$ can in principle differ from $\nu_1$.

If the accreted mass, $m$, is much smaller than the mass of the BH (e.g., $m < 0.01 M_{\text{BH}}$), the distribution of black hole spins is flat, as the timescale for spin overflipping due to the Bardeen-Petterson effect is longer than the timescale to accrete the whole $m$. If instead $m \approx M_{\text{BH}}$, BHs can align with the angular momentum of the accretion disk and accrete enough mass to be spun up. In this case the distribution of BH spins is dominated by rapidly rotating systems.

Understanding if the description of the warp propagation is correctly described by a different viscosity with respect to the one responsible for the radial propagation of the angular momentum...
is beyond the scope of this work. We have therefore explored a wide range of possible viscosities, and here we simply report our results. If $\nu_{2}/\nu_{1} = 1$, the timescale for alignment is short, and the spin of a BH increases rapidly. If the warp propagation is instead better described by a high $\nu_{2} = \nu_{1}/\alpha^{2}$, a substantial fraction of black holes of all masses can have small spins, provided $m \ll M_{\text{BH}}$.

However, both semianalytical models of the cosmic BH evolution (Volonteri et al. 2005) and simulations of merger-driven accretion (di Matteo et al. 2005) show that most BHs increase their mass by an amount $\gg m_{\text{align}}$, if the evolution of the luminosity function (LF) of quasars is kept as a constraint. These high $m$ values are likely characteristic of the most luminous quasars and most massive black holes, especially at high redshift.

We therefore expect that bright quasars at $z > 3$ have large spins, in contrast to the suggestion of King & Pringle (2006). High spins in bright quasars are also indicated by the high radiative efficiency of quasars, as deduced from observations applying the SoItan argument (Soltan 1982; Wang et al. 2006, and references therein).

If the mass of a BH need to reach $10^{6} M_{\odot}$ by $z = 3$, or even more strikingly by $z = 6$, so that they can represent the engines of quasars with luminosity $L > 10^{46}$ erg s$^{-1}$, BHs need to grow from typical seed masses (e.g., Madau & Rees 2001; Koushiappas et al. 2004; Begelman et al. 2006; Lodato & Natarajan 2006) by at least 3–4 orders of magnitude in $10^{6}$–$10^{9}$ yr. The necessity of long and continuous accretion episodes therefore imply that for these BHs $m \gg m_{\text{align}}$.

Smaller BHs, powering low-luminosity active galactic nuclei, can instead grow by accreting smaller packets of material, such as tidally disrupted stars (for BHs with mass $< 2 \times 10^{6} M_{\odot}$; Milosavljevic et al. 2006), or possibly molecular clouds (Hopkins & Hernquist 2006). For these black holes the spin distribution is more probably flat, or skewed toward low values. This latter result is in agreement with Sikora et al. (2007), who find that disk galaxies tend to be weaker radio sources than elliptical hosts. In the hierarchical framework we might expect that the BH hosted by an elliptical galaxy had, as last major accretion episode, a large increase in its mass following the major merger that created the elliptical itself (Hopkins & Hernquist 2006). During this episode the spin increased significantly as well, possibly up to very high values. Subsequently the black hole might have grown by swallowing the occasional molecular cloud, or by tidally disrupting stars. If the total contribution of these random episodes represents a small fraction of the BH mass, the spin is, however, kept at high values.

Black holes in spiral galaxies, on the other hand, probably had their last major merger (i.e., last major accretion episode), if any, at high redshift, so that enough time elapsed for the galaxy disk to re-form. Most of the latest growth of the BH should have happened through minor events, which have likely contributed to the BH spin-down.

Our results are supported also by the recent finding by Capetti & Balmaverde (2006; 2007) that radio bimodality correlates with bimodality of stellar brightness profiles in galactic nuclei. The inner regions of radio-loud galaxies display shallow cores (star deficient). Cores, in turn, are preferentially reside in giant ellipticals (see Lauer et al. 2007, and references therein). Radio-quiet galaxies, including nearby low-luminosity Seyferts, have instead power-law (cuspy) brightness profiles and preferentially reside in S0 and spiral galaxies.

Hence, noting that core nuclei result from merging BHs following galaxy mergers (Ebisuzaki et al. 1991; Milosavljevic & Merritt 2001; Milosavljevic et al. 2002; Ravindranath et al. 2002; Volonteri et al. 2003), Balmaverde & Capetti’s discovery is consistent with our conjecture that spin bimodality is determined by diverse evolutionary tracks of BH spins in disk galaxies (random small mass accretion events) and giant elliptical galaxies (massive accretion events which follow galaxy mergers). Both tidal disruption of stars and accretion of gaseous clouds is unlikely in shallow, stellar dominated galaxy cores. Therefore, it is conceivable that the observed morphology-related bimodality of AGN radio loudness results from bimodality of central black holes’ spin distribution.

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