On the spin paradigm and the radio dichotomy of quasars

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Accepted . Received ; in original form .

ABSTRACT
We investigate whether models based on the assumption that jets in quasars are powered by rotating black holes can explain the observed radio dichotomy of quasars. We show that in terms of the “spin paradigm” models, radio-loud quasars could be objects in which the black hole’s rotation rate corresponds to an equilibrium between spin-up by accretion and spin-down by the Blandford-Znajek mechanism. Radio-quiet quasars could be hosting black holes with an average spin much smaller than the equilibrium one. We discuss possible accretion scenarios which can lead to such a bimodal distribution of black hole spins.

Key words: accretion, accretion discs – black hole physics – galaxies: active – galaxies: evolution.

1 INTRODUCTION
Quasars are characterized not only by an intense radiation from the central engine, but also by jets which power large scale radio structures. The ratio of radio luminosity of these structures to the optical luminosity of the central sources shows a bimodal distribution, with only ~ 10 % of quasars belonging to the radio-loud category (see, e.g., Kellerman et al. 1989; Hooper et al. 1995; Falcke, Sherwood, & Perna 1996; Bischof & Becker 1997). Radio-loud quasars have never been found in spiral galaxies, while the hosts of radio-quiet quasars are either spiral or elliptical galaxies (Taylor et al. 1996; Kukula et al. 1998). On the other hand, both radio-quiet and radio-loud quasars have almost identical average IR-optical-UV spectra (Francis et al. 1993; Zheng et al. 1997), which suggests similar accretion conditions in these two samples of objects.

These properties seem to support models which are based on the idea that jets are powered by rotating black holes. However, in order to explain the observed radio loudness bimodality in terms of a bimodal distribution of black hole spins, one has to assume that the population of supermassive black holes is dominated by very low-spin black holes, and one should understand why radio-loud quasars are always hosted by elliptical galaxies. This issue has been addressed by Wilson & Colbert (1995), who proposed that the Blandford-Znajek (B-Z) mechanism, which extracts rotational energy from a black hole (Blandford & Znajek 1977), is not efficient enough to counteract the spinning up of a black hole by the gas accreted from standard α-discs (Shakura & Sunyaev 1973). Low-spin equilibrium states are possible only for very low accretion rates but the time required to establish such an equilibrium is longer than the age of the Universe. This is because the power extracted from the black hole is proportional to the square of intensity of the magnetic field which threads the black hole. The field intensity is limited by the pressure in the accretion flow and, therefore, is very low for very low accretion rates. Hence, once a black hole gets a high spin, it will be rotating fast forever.

Low-spin equilibrium solutions for high accretion rates are possible only for the so called β-discs (Moderski & Sikora 1996). In such discs the viscous stress driving accretion is proportional to the gas pressure only (and not to the total pressure as in the case of α discs). Since the relation between the viscous stress and the pressure in radiation pressure dominated accretion discs is unknown, such discs are a reasonable alternative to α-discs. For high accretion rates the gas pressure is by several orders of magnitude smaller than the radiation pressure (Shakura & Sunyaev 1973). This implies that, for a given accretion rate, β-discs are much denser than α-discs (Sakimoto & Coroniti 1981) and, therefore, can confine much stronger black hole magnetic fields.

Since the power extracted from a black hole via the B-
Z mechanism scales with the square of the magnetic field intensity, equilibrium spins for $\beta$-discs are smaller than for $\alpha$-discs. This, of course, does not mean that the power extracted from a black hole with a $\beta$-disc must be smaller than the power extracted from a black hole with an $\alpha$-disc, because in the expression for the B-Z power a lower equilibrium spin is compensated by a higher intensity of the magnetic field (Moderski & Sikora 1997). Thus, low equilibrium spin black holes can represent radio-quiet quasars only, if the conversion of the extracted black hole energy into jet energy is very inefficient. According to the Wilson-Colbert scenario powerful jets would then exist only in objects where a coalescence of two black holes leads to spins much higher than the equilibrium one. However, powers extracted from black holes with spins much larger than the equilibrium one are so large, that such black holes would be spun-down on time scales two orders of magnitude shorter than the typical lifetime, $\sim 10^8$ years, of radio quasars (see, e.g., Leahy, Muxlow & Stephens 1989).

Powering jets in radio-loud quasars by a black hole can last as long as $10^8$ years, only if losses of angular momentum due to the B-Z mechanism are compensated by gains of angular momentum from an accretion disc. This would give a correlation between radio and optical luminosities, in accordance with observations (Serjeant et al. 1997). Then, however, one would have to explain why the majority of super-massive black holes have spin values much lower than the equilibrium one – the condition for the existence of radio-quiet quasars. As suggested by Moderski, Sikora & Lasota (1997), black holes in most objects could be forced to rotate slowly by multi-accretion events with random orientations of the angular momentum vector. In such a scenario, quasars which become radio-loud are only those which undergo major accretion events, induced, i.e., by a merger of two big galaxies. Following this, a black hole can easily double its mass and reach an equilibrium state. In this paper we explore this possibility and derive conditions the model must satisfy in order to explain the observed radio properties of quasars.

This paper is organized as follows. In Section 2 we present equilibrium spin solutions, obtained for a variety of accretion disc models. In Section 3, we discuss multi-event accretion scenarios. In Section 4 we use our results to derive conditions which must be satisfied by quasar evolution models in order to obtain bimodal distribution of black hole spins.

## 2 EQUILIBRIUM STATES

The evolution of a black hole, as determined by a gain of energy and angular momentum from an accretion disc and a loss of energy and angular momentum via the B-Z mechanism, is described by equations (Moderski & Sikora 1996a)

$$\frac{dA}{dt} = (j_{in} - 2\alpha c\dot{e}_{in}) \frac{\dot{M}}{M} - \left(\frac{2\dot{\Omega}_h}{kA} - 2A\right) \frac{P}{Mc^2},$$

$$\frac{d\ln M}{dt} = \dot{e}_{in} \frac{\dot{M}}{M} - \frac{P}{Mc^2},$$

where $M$ and $J$ are the energy-mass and the angular momentum of a black hole, respectively; $A \equiv cJ/GM^2 = J/J_{max}$ is the dimensionless angular momentum of the black hole; $\dot{\Omega}_h = \dot{E}/\dot{M} - \Omega$ is the angular velocity of a black hole - see Appendix A; and

$$\dot{M} \equiv \dot{m} \dot{M}_{Edd}$$

is the accretion rate ($\dot{M}_{Edd}$ is the Eddington accretion rate defined by $c^2 \dot{M}_{Edd} \equiv L_{Edd} \simeq 10^{38}M/M_\odot [\text{erg sec}^{-1}]$); $\dot{e}_{in} = \dot{e}_{in}/c^2$ and $j_{in} = \dot{e}_{in}/GM$ are dimensionless specific energy and angular momentum of matter at inner edge of the accretion disc; $\dot{e}_{in} = \dot{e}_{in}/c^2$ and $j_{in} = \phi_{in}/GM$ are dimensionless specific energy and angular momentum of matter at inner edge of the accretion disc; $\dot{\Omega}_h = \dot{\Omega}_h/c^2$ is the angular velocity of the magnetic field lines threading the horizon and $\Omega_h$ is the angular velocity of a black hole - see Appendix A; and

$$P = \frac{k(1-k)}{32} G^2 c^5 \dot{e}_{in}^2 A^2 B^2 \dot{M}^2$$

is the power extracted by the B-Z mechanism, where $B_\perp$ is the poloidal component of the magnetic field threading the horizon.

The value of $B_\perp$ depends on the accretion disc model and on the efficiency of the diffusion of the external magnetic field into a disc. If the black hole magnetic field does not penetrate the disc but instead interacts with it by inducing surface currents, then its energy density can be as high as maximum total pressure in a disc, i.e. $B^2 \perp \sim 8\pi\rho_{\text{rot}, \perp}$.

The dependence of $\rho_{\text{rot}, \perp}$ on accretion rate, viscosity and spin for $\alpha$- and $\beta$-discs is illustrated in Figure 1 and respective formulas are specified in Appendix B. Using Figure 1 for estimations of $P$ in Eq. 1 one should remember that the disc pressure depends also on the black hole mass: for radiation pressure dominated $\alpha$-discs it is $\propto 1/M$; for radiation pressure dominated $\beta$-discs $\propto 1/M^{4/3}$; and for gas pressure dominated discs $\propto 1/M^{9/10}$. Much smaller $B_\perp$ are confined by a disc, if the black hole magnetic field diffuses into a disc. Then, as was argued by Ghosh & Abramowicz (1997), the energy density of magnetic field which threads
the black hole cannot exceed the energy density of the magnetic field in a disc. Assuming that the angular momentum transport in the disc is dominated by the Maxwell stress, \( t_{r\phi} = B_t B_\phi / 4\pi \) (Balbus & Hawley 1991), one can find, using definitions of the \( \alpha \)-disc (\( t_{r\phi} = \alpha \rho_{\text{tot}} \)) and of the \( \beta \)-disc (\( t_{r\phi} = \beta \rho_{\text{gas}} \)), that for \( \alpha \)-disks \( B^2_\perp / 8\pi \sim \alpha \rho_{\text{tot}, \text{max}} \), and for \( \beta \)-disks \( B^2_\perp / 8\pi \sim \beta \rho_{\text{gas}, \text{max}} \).

For a given accretion rate, the rate of angular momentum transport is model independent. This transport rate is given by the integral of \( t_{r\phi} \) over the disc height. In a radiation pressure dominated accretion disc in which electron scattering is the only source of opacity the disc semi-thickness is equal to \( H = (3/4)M r_S \), where \( r_S = 2GM/c^2 \) is the Schwarzschild radius. Therefore, \( B^2_\perp \sim t_{r\phi} \) is the same for both models and is numerically equal to \( 8\pi \rho(a_{\text{tot}, \text{max}}) \) for discs with \( \alpha = \beta = 1 \). The maximum pressure for such discs is \( \sim 10^5 / M_8 \) dyne cm\(^{-2} \) which gives \( P_{\text{max}} \sim 3 \times 10^{14} M_8 \) ergs s\(^{-1} \), as calculated from Eq. (3) for \( k = 1/2 \) and \( A = 1 \). Assuming that efficiency of radio production is 10% \( (L_r \sim 0.1P) \) and noting that most luminous radio sources have \( L_r \sim 10^{45} \) ergs s\(^{-1} \) one needs \( M \sim 3 \times 10^5 M_8 \) to explain them. The existence of such massive black holes in some galactic nuclei is confirmed by the evaluation of the black mass in M87 (Harms et al. 1994). In a model with non-diamagnetic discs it is difficult, however, to obtain the radio-loudness of radio-loud quasars, defined as the ratio of the radio-loud quasars corresponds roughly to \( P/L \sim 0.1 \) (Rawlings & Saunders 1991), where \( L_d = (1-\tilde{e}_n)MC^2 \) is the luminosity of an accretion disc. Such values are reachable by models with non-diamagnetic discs only if \( L_d \leq 0.01L_{\text{edd}} \) and it is not clear whether such objects produce UV bumps and emission lines which are attributes of both radio-quiet and radio-loud quasars. Accretion flows in such systems would be rather dominated by advection (see Lasota 1998; Narayan, Mahdevan & Quataert 1998 for recent reviews). The presence of such a flow in M87 was proposed by Reynolds et al. (1996).

From equations (1) and (3) one gets that at equilibrium, \( dA/dt = 0 \), the radio loudness for a thin, Keplerian disc is expressed by

\[
P/L_d = \frac{\tilde{J}\tilde{n} - 2A\tilde{e}_n}{(\tilde{r}A - 2A)(1 - \tilde{e}_n)}
\]

and is independent on the disc model chosen. The equilibrium radio loudness as a function of the equilibrium spin is presented in Figure 3.

The evolution of the black hole spin, as derived from equations (1) and (3) can be characterized by two phases. During the first, \( A \) makes its way towards the equilibrium value \( A_{eq} \) given by \( dA/dt = 0 \), the second is described by \( dA/dt \approx 0 \) and Eq. (4). Examples of such evolution for \( \alpha \)- and \( \beta \)-disks were presented by Moderski & Sikora (1996a; 1997). \( A_{eq} \) solutions as a function of \( \tilde{m} \) are shown in Figure 2. For high accretion rate \( \alpha \)-disks these solutions are independent of \( M_{\text{BH}} \). For \( \beta \)-disks and gas pressured dominated disc \( A_{eq} \) depends on \( M_{\text{BH}} \), but very weakly, \( \propto M_{\text{BH}}^{-1/10} \), respectively. One can see in Figure 2 that for \( \alpha \)-disc models, low spin equilibrium solutions (\( A < 0.1 \)) exist only for very low accretion rates and only for very small \( \alpha \) parameters. Such states cannot be achieved during a Hubble time (see the dashed lines which mark the time scales during which black holes reach \( A_{eq} \)) and, therefore, black holes once spun-up to a high spin value will rotate fast forever. Black holes accreting from \( \alpha \)-disks can have low spins only if they are formed with low spins and then accrete very little in comparison with the mass they collected during the formation process (Moderski, Sikora & Lasota 1997), or if some processes lead to random changes of a disc angular momentum. These possibilities are discussed in next Sections.

Black holes with low equilibrium spins for the high accretion rates exist only if the accretion disc is a \( \beta \)-disc. For the viscosity parameter \( \beta \sim 0.1 \) and \( \tilde{m} \sim 10 \) such models give \( A_{eq} \sim 0.05 \), and if black holes start from different initial spins they reach the equilibrium spin in \( \sim 10^7 \) years. However, as can be found from Figures 2 and 3, \( P/L_d \)’s for such parameters correspond to radio-loudness of the radio-loud quasars rather than to that of radio-quiet quasars (see Section 4). This problem could eventually be over-passed by postulating very low efficiency of conversion of energy extracted from a black hole to the jet. Then, black holes with a low equilibrium spin would represent radio-quiet quasars, while black holes reaching \( A \sim 1 \gg A_{eq} \), e.g. following coalescence of two black holes (Wilson & Colbert 1993), would represent radio-loud quasars. However, as Figure 4 shows, such black holes spin-down in less than \( 10^6 \) years, i.e. in two orders shorter time scale than typical lifetime of radio sources in quasars.

All these considerations show that, in the framework of the spin paradigm, it is impossible to account for the observed radio dichotomy of quasars if the angular momentum accreted during their evolution is always of the same sign.
3 SWITCHING BETWEEN PRO- AND RETROGRADE ACCRETION DISCS

A spinning black hole can be decelerated if it accretes angular momentum of the opposite sign. Since a simple version of retrograde accretion was discussed in details in one of our previous papers (Moderski & Sikora 1996b), here we present only results which are directly connected to the possibility of spinning down of black holes. As was shown by Moderski and Sikora (1996b), it suffices to accrete \( \sim 0.2 \) of the initial black hole mass to decelerate a black hole from its maximum spin to zero. However, accretion of an larger amount larger than 0.2 of the initial mass will spin-up a black hole again, so that to maintain a low value of the time-averaged spin, the evolution of a black hole must be governed by many “small” accretion events with random orientation of the accreted angular momentum. In the calculations we assume that portions (supplied e.g. by giant molecular clouds or by the capture of dwarf galaxies) all have equal masses. Because of the Bardeen-Petterson effect (Bardeen & Petterson 1975) the disc, near the black hole, rotates in its equatorial plane we randomly choose only between direct or retrograde accretion and we did not consider more complicated alignments. Some examples of the evolution of the black hole spin for several values of the mass accreted at each accretion event are presented in Figure 5. We can see from Figure 5 that spinning down of a black hole with alternating direct–retrograde discs is very efficient and leads to black holes with spin values fluctuating around zero. To check whether this effect depends on the mass of the accreted portion of matter we performed a number of numerical evolutions for different ratios of the accreted mass to the initial mass of the black hole, \( \Delta m \). We traced the evolution of 50 thousands systems evolving from \( A = 1 \) until accreted mass exceeded four times the initial mass of the black hole and studied how the black hole spin changed with the number of accretion events. An example

Figure 3. \( P/L_d \) vs. \( A_{eq} \) for geometrically thin accretion disc (inner edge of the disc on the marginally stable orbit) and for maximal efficiency of energy extraction \((k = 1/2)\).

Figure 4. Evolution of the black hole spin (left panels) and of the radio-loudness (right panels) for the \( \beta \) accretion disc model. Initial values of spin are \( A_0 = 0.0, 0.4, 1.0 \). Solid lines are for discs with \( \beta = 1 \); dotted lines are for \( \beta = 0.1 \); and dashed lines are for \( \beta = 0.01 \). The upper panels are for \( \dot{m} = 1 \) and the lower panels are for \( \dot{m} = 10 \). Evolutions calculated for initial black hole mass \( M_{BH} = 10^8 M_\odot \).

Figure 5. Examples of spin evolution for different masses accreted each time. All evolutions are from initial spin \( A_0 = 1 \). From the upper left to bottom right we have \( \Delta m = 0.005, 0.01, 0.05, 0.1 \), where \( \Delta m \) is the ratio of the mass of the accreted portion to the initial mass of the black hole \( M_0 \).
of spin distributions in the simulated population for three different numbers of accretion episodes and for $\Delta m = 0.01$ is presented in Figure 3. We found that value of the peak does not depend on the mass $\Delta m$ but only on total accreted mass, and that the spread in distribution is wider for higher $\Delta m$. In Figure 3 we show the mean spin value and the spread in the population as a function of accreted mass. It is worth noting that results do not change significantly if we consider geometrically thick discs, in which the inner disc radius can be close to the marginally bound orbit (Abramowicz & Lasota 1980). We checked that masses smaller by only 10% are needed to establish the same population of black holes as in the case of thin discs.

One should mention that results presented in this Section were obtained assuming that $P = 0$. With this assumption, the black hole evolution is described by a simple equation

$$\frac{dA}{d\ln M} = \frac{2}{\tilde{e}_{\text{in}}} - 2A,$$

(5)

which for a disc edge located at the marginally stable ($r_{\text{ms}}$) or marginally bound ($r_{\text{mb}}$) orbit has an analytical solution (see Bardeen 1970 and Moderski & Sikora 1996b). Our results can be used to describe evolution of the black hole also for non-zero $P$, provided we follow only these parts of evolutionary tracks for which $A < A_{\text{eq}}$, because at such $A$'s the evolution is dominated by accretion processes, rather than by the B-Z mechanism and, therefore, the results are insensitive to the value of $P$.

4 DISCUSSION

Radio dichotomy of quasars was discovered many years ago (Strittmatter et al. 1980; Kellerman et al. 1989), but is still waiting for a theoretical explanation. As for now, the consensus concerns only one aspect of the problem: it is clear that jets in quasars must be formed near a supermassive black hole. This follows from the energetics of quasar jets, since no other known sources could power jets at a rate reaching $10^{46}$ ergs s$^{-1}$ for millions of years (Rawlings & Saunders 1991; Leahy et al. 1989). Independent argument for the formation of extra galactic jets in the vicinity of supermassive black holes is provided by direct VLBI observations of nearby radio galaxies. In particular, in 3C 274 (M87) the jet is seen down to $10^{15}$ cm from the center (Junor & Biretta 1993), which corresponds to 100 gravitational radii for the $3 \times 10^6 M_\odot$ (Harms et al. 1994) central black hole. There, very deep in the gravitational potential well, jets could be powered either by the innermost parts of an accretion disc (Blandford & Payne 1982; Park & Vishniac 1994; Contopoulos 1995; Begelman 1995) or by a rotating black hole (Blandford & Znajek 1977; Rees et al. 1982). However, no jet production model can be successful, if it fails to explain why only a small fraction of quasars is radio loud and why radio-loudness has a bimodal distribution.

In terms of the $R = F_\nu / F_\nu^o$ ratio, where $F_\nu$ and $F_\nu^o$ are the monochromatic fluxes measured at frequencies $\sim 10^{10}$ Hz and $\sim 10^{15}$ Hz, respectively, radio-quiet quasars cluster around $R \sim 0.3$ and radio-loud quasars cluster around $R \sim 300$ (Kellermann et al. 1989; Falcke, Sherwood & Paumk 1996). Thus, the average radio-loudness of the two quasar populations differs by a factor $10^3$ and this number, together with typical radio luminosities of radio loud quasars, $L_\nu \sim 10^{45}$ ergs s$^{-1}$, provides the basic quantitative conditions which should be satisfied by any unified model of quasars. These conditions, together with our results discussed in the two previous sections are used below to test spin based models of a jet activity in quasars. The predic-
tions of such models should also satisfy such observationally established trends, as
- radio-loud quasars avoid disc-galaxies and have UV-luminosities $\gtrsim 10^{46}$ ergs s$^{-1}$;
- radio-quiet quasars are present both in spiral and elliptical galaxies (Taylor et al. 1996) and their radio properties do not depend on the galaxy morphology (Kukula et al. 1998);
- radio properties of radio-quiet quasars suggest that they are, like in radio loud quasars, related to the jet production by a central engine.

Assuming that the efficiency of conversion of jet energy into radio emission is $\sim 10\%$, the typical jet in radio-loud quasars should have $P \sim 10^{46}$ ergs s$^{-1}$. Similar jet powers are deduced by calculating the total energy content of extended radio sources and dividing it by the age of the source (Leahy et al. 1989), or from energetics of $\gamma$-ray production in sub-parsec jets (see, e.g., Sikora 1997). Largest powers which can be extracted from rotating black holes are given by equation (3). For $A = 1$ and $B_L = 8 r_{\text{pot}}$ we obtain $P_{\text{max}} \sim 3 \times 10^{41} M_8 r_{\text{pot}, s}^2$ ergs s$^{-1}$, where $r_{\text{pot}, s} = p_{\text{pot}, s}/10^8$ dyne cm$^{-2}$ and $M_8 = M/10^8 M_\odot$. One can see from Figure 1 that for high accretion rates ($\dot{m} > 0.1$, say) and black hole masses $\sim M = 10^8 M_\odot$, a pressure $\geq 10^8$ dyne cm$^{-2}$ is provided by $\alpha$-discs with $\alpha \leq 0.1$, and by all $\beta$-discs. Thus, the B-Z mechanism is efficient enough to power jets in radio-loud quasars, provided the black hole magnetic field is supported by the total disc pressure. If the latter is not true and, as Ghosh and Abramowicz (1997) argued, the energy density of the black hole magnetic field cannot exceed the energy density of the maximum magnetic field in a disc, then the maximum pressure of the black hole's magnetic field is numerically equal to the total pressure in an $\alpha = 1$ disc. In this case black hole masses $\sim 3 \times 10^9 M_\odot$ are required in order to get $P \sim 10^{46}$ ergs s$^{-1}$. The case of M87 seems to prove that such black holes are not necessarily exceptional (Harms et al. 1994). However, one should note here, that the question of the diffusion of an external magnetic field into an accretion disc is still open (see, e.g., Wang 1995; Bardou & Heuvel 1996).

Assuming, as before, that the fraction of the jet energy converted into radiation is $10\%$, and that the bolometric corrections for jet radiation at $\sim 10^{43}$ Hz and for accretion disc radiation at $\sim 10^{15}$ Hz are of the same order, we obtain $P/L_d \sim 10 L_r/L_o \sim 10 F_{\nu_F}/F_{\nu_o} \sim 10^{-4} R$. Thus, for radio-loud quasars models should predict $P/L_d \sim 0.1$, while radio-quiet quasars should cluster around $P/L_d \sim 10^{-4}$. As is seen from Figures 2 and 3, $P/L_d \sim 0.1$ nicely corresponds to black hole equilibrium spin solutions for all but $\alpha > 0.1$ disc models. One can also check, that there are no equilibrium spin solutions which would correspond to radio-loudness of radio-quiet objects. For them $A < 0.3$ is required, provided that radio luminosity scales linearly with $P$. A population of such low spin black holes can exist only if black holes are born with very low spin and then accrete very little (Moderski, Sikora & Lasota 1997), or if black hole evolution is determined by multi-accretion events with random angular momenta.

As one can deduce from results presented in Figure 7, hundreds of accretion events per object are required in order to have more than $90\%$ of black holes with $A < 0.1$ at any given moment. This is too much to be obtained by accretion events induced by capture of dwarf galaxies, but can be achieved by accretion of molecular clouds. Molecular cloud accretion events were recently proposed by Sanders et al. (1998) to explain some properties of Sgr A$^*$ and other AGNs. This scenario is supported by the random orientation of central engines vs. the orientation of galactic discs, as deduced from observations of “UV” cones (Wilson & Tsvetanov 1994; McLeod & Rieke 1995) and radio axis (Clarke, Kinney & Pringle 1998) in Seyfert galaxies. Here we should note, that in our simplified treatment of the multi-accretion scenario (Section 3), we didn’t take into account the coupling between the spin of the black hole and the orbital angular momentum of the approaching molecular clouds. Such coupling supposedly leads to random wondering of the black hole spin vector.

One can now speculate that changes of orientation of the black hole spin could be interrupted and the black hole could be spun-up to very high spins following a merger process. This process could induce a massive and long lasting accretion event. If during such an event the accretion proceeds from a fixed plane and at least doubles the black hole mass, the black hole spin reaches the equilibrium spin and the object becomes a typical radio-loud quasar (Moderski, Sikora & Lasota 1997).

Since mergers happen mostly in groups and clusters of galaxies, where the population of galaxies is dominated by ellipticals, this could explain why radio loud quasars avoid spiral galaxies. Observational arguments for such a scenario are exactly the same as those used by Wilson and Colbert (1993). The only difference is, that they postulated formation of high spin black holes via coalescence of two supermassive black holes, while in our scenario high spins result from an accretion process. Note, however, that a coalescence of two black holes, if it happens, does not have to affect much our scenario. If the coalescence involves two black with very different masses, the final spin will be determined by the accretion process, otherwise both processes lead to similar spins.

What are the perspectives for an observational test of the assumption that radio-quiet objects have low spins? A possibility to measure the spin of supermassive black holes is provided by the detailed studies of profiles of the X-ray fluorescent iron line produced in the surface layer of the innermost parts of accretion discs. Such lines are detected in many Seyfert galaxies, which represent the low luminosity branch of radio-quiet quasars. For at least one of such objects the line profile was claimed to be consistent with the kinematics given by the rotation of a disc around a black hole in fast rotation (Kawasawa et al. 1996). However, as demonstrated by Reynolds and Begelman (1997), similar line profiles can be produced around non-rotating black holes, provided that a large part of the line emission comes from below the marginally stable orbit. Therefore, much more detailed theoretical models and sensitive observations are required to get conclusive diagnostics from this type of investigations.

The remarkable discovery of relativistic jets in several Galactic X-ray sources (cf. Mirabel & Rodriguez 1994; Hjellming & Rupen 1995; Newell, Spencer & Garrett 1997) suggests that the radio-dichotomy exists for Galactic compact objects as well. As was argued recently by Zhang, Cui & Chen (1997), jet activity in these sources can also be conditioned by the value of the black hole spin.
ACKNOWLEDGMENTS

RM and MS acknowledge support from KBN grant 2P03D01209. During his stay in Meudon RM was supported by the Réseau Formation Recherche of the French Ministère de l’Enseignement Supérieure et de la Recherche, he also acknowledges the support from Foundation for Polish Science Fellowship. JPL thanks the Isaac Newton Institute for hospitality in April 1998.

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APPENDIX A: DIMENSIONLESS QUANTITIES IN THE KERR METRIC

The black hole radius in the Kerr metric is
\( r_h = c^2 r_h^3 / G M = 1 + (1 - A^2)^{1/2} \), (A1)
The black hole angular velocity is
\[ \Omega_h = \frac{GM}{c^3 r_h} \]
\( \Omega_h = A / 2 r_h^3 \) (A2)

The marginally stable orbit is (upper signs are for direct accretion and lower sign for retrograde accretion)
\[ r_{ms} = c^2 r_{ms} / G M = 3 + Z_2 - \frac{1}{2}[(3 - Z_1)(3 + 2Z_1)]^{1/2} \]
where
\[ Z_1 = 1 + (1 - A^2)^{1/3} [(1 + A)^{1/3} + (1 - A)^{1/3}] \]
and
\[ Z_2 = (3A^2 + Z_1^2)^{1/2} \]

and inversely:
\[ A = \pm \frac{1}{3} r_{ms}^{1/2} (4 - (3r_{ms} - 2)^{1/2}) \]
The specific energy and specific angular momentum of a particle on the marginally stable orbit are, respectively,
\[ e_{ms} = e_{ms} / c^2 = \left( 1 - \frac{2}{3r_{ms}^3} \right)^{1/2} \]
and
\[ j_{ms} = e_{ms} c / G M = + \frac{2}{3 \sqrt{3}} \left[ 1 + 2(3r_{ms} - 2)^{1/2} \right] \]

APPENDIX B: PRESSURE IN \( \alpha \) - AND \( \beta \) -DISCS

For high accretion rates pressure in geometrically thin discs around supermassive black holes is dominated by radiation and can be approximated by formulas (Novikov & Thorne 1973; Sakimoto & Coroniti 1981)
\[ p_{\nu \nu}^{(\alpha)} = 3 \times 10^8 \alpha^{-1} M_\bullet^{-1} \tilde{r}^{-3/2} \frac{B^2 \nu}{A^2} \] (B1)
for $\alpha$-discs and

$$p_{\text{rad}}^{(\beta)} = 1.6 \times 10^{14} \dot{m}^{8/5} \beta^{-4/5} M_8^{-4/5} \bar{r}^{-18/5} \frac{Z^{8/5}}{B^{8/5} D^{4/5}} \quad (B2)$$

for $\beta$-discs, where $\dot{m} = \dot{M} c^2 / L_{\text{Edd}}$, and $A$, $B$, $D$, $E$, $Z$ are relativistic corrections and can be found in Novikov & Thorne (1973).

For low accretion rates the pressure is dominated by a gas and is

$$p_{\text{gas}} = 2.5 \times 10^{13} \beta^{-9/10} \dot{m}^{4/5} M_9^{9/10} \bar{r}^{-51/20} \frac{B^{1/5} E^{1/2} Z^{4/5}}{A D^{2/5}}. \quad (B3)$$