

Models for jet power in elliptical galaxies: a case for rapidly spinning black holes

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ABSTRACT

The power of jets from black holes is expected to depend on both the spin of the black hole and the structure of the accretion disc in the region of the last stable orbit. We investigate these dependencies using two different physical models for the jet power: the classical Blandford–Znajek (BZ) model and a hybrid model developed by Meier. In the BZ case, the jets are powered by magnetic fields directly threading the spinning black hole while in the hybrid model, the jet energy is extracted from both the accretion disc as well as the black hole via magnetic fields anchored to the accretion flow inside and outside the hole’s ergosphere. The hybrid model takes advantage of the strengths of both the Blandford–Payne and BZ mechanisms, while avoiding the more controversial features of the latter. We develop these models more fully to account for general relativistic effects and to focus on advection-dominated accretion flows (ADAFs) for which the jet power is expected to be a significant fraction of the accreted rest mass energy.

We apply the models to elliptical galaxies, in order to see if these models can explain the observed correlation between the Bondi accretion rates and the total jet powers. For typical values of the disc viscosity parameter $\alpha \sim 0.04$ – 0.3 and mass accretion rates consistent with ADAF model expectations, we find that the observed correlation requires $j \gtrsim 0.9$; that is, it implies that the black holes are rapidly spinning. Our results suggest that the central black holes in the cores of clusters of galaxies must be rapidly rotating in order to drive jets powerful enough to heat the intracluster medium and quench cooling flows.

Key words: accretion, accretion discs – black hole physics – MHD – galaxies: active – galaxies: jets – X-rays: galaxies.

1 INTRODUCTION

Recent high resolution *Chandra* observations of the cores of elliptical galaxies provide dramatic illustration of the impact of jets launched from active galactic nuclei (AGNs) on the interstellar medium of their host galaxies and/or on the intracluster medium if the galaxies reside in clusters. These observations have revealed prominent X-ray surface brightness depressions corresponding to cavities or bubbles inflated by the jets as they interact with the surrounding hot gas (e.g. Allen et al. 2006, hereafter A06; Birzan et al. 2004; Fabian et al. 2006; Rafferty et al. 2006; Taylor et al. 2006). This interaction is presumed to deposit large amounts of energy in their environments (e.g. Churazov et al. 2002; Dalla Vecchia et al. 2004; Sijacki & Springel 2006; Nusser, Silk & Babul 2006), modifying the hierarchy of galaxy formation (e.g. Bower et al. 2006) and altering the evolution of the intracluster medium by counter-

acting radiative losses (e.g. McCarthy et al. 2004). Based on the energetics of creating the observed cavities, the minimum energy associated with the jets range from $\sim 10^{55}$ erg in galaxies, groups, and poor clusters to $\sim 10^{60}$ erg in rich clusters, ranking the outbursts as among the most powerful phenomena in the Universe.

While a detailed understanding of the extragalactic AGN jet phenomena remains elusive, the combination of physically insightful analytic studies (e.g. Blandford & Znajek 1977; Blandford & Payne 1982; Begelman, Blandford & Rees 1984; Punsly & Coroniti 1990; Ferrari 1998; Meier 1999, 2001) and sophisticated general relativistic, magnetohydrodynamic (MHD) numerical simulations (e.g. Koide et al. 2000; Koide 2003; McKinney & Gammie 2004; De Villiers et al. 2005; Komissarov 2005; Hawley & Krolik 2006) are beginning to yield important insights. There is now a general consensus that jets are fundamentally MHD events.

The currently favoured models presuppose an accretion flow threaded by large-scale magnetic fields flowing on to a supermassive black hole. In the neighbourhood of the black hole, the flow settles into a disc-like structure, in which the rate of inward flow of

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matter depends on the efficiency with which its angular momentum can be transferred outward. This transfer of angular momentum is mediated by MHD turbulence. Both MHD turbulence and differential rotation of the plasma in the body of the disc generate and intensify toroidal magnetic fields (Balbus & Hawley 1998). When the pressure associated with the toroidal fields grows strong enough, the field lines escape from the disc forming a rotating helical tower of field lines above and below the disc. Centrifugal forces associated with this rotating magnetic field helix drives any plasma trapped on to the field lines upward and out of the disc, generating outflows. This is the crux of the Blandford–Payne (Blandford & Payne 1982) model for jets, wherein the magnetic fields extract energy from the rotation of accretion disc itself to power the outflows. Depending on the detailed structure of the magnetic fields and the accretion disc, this mechanism is expected to generate outflows ranging from broad, uncollimated winds to highly collimated jets.

In the event that the black hole itself is spinning, the magnetic field can also extract the rotational energy of the central black hole to power the outflows. In the Blandford–Znajek (BZ) model (Blandford & Znajek 1977), the magnetic fields are assumed to be connected directly to the horizon of a spinning black hole while in the Punsly & Coroniti (1990) model, the magnetic fields associated with jet production are anchored to the inflowing plasma inside the black hole’s ergosphere. In both these models, the dragging of inertial frames, relative to an observer at infinity, within the ergosphere of the rotating black hole results in a rotating, tightly wound vertical tower of field lines, and hence powerful outflows. Moreover, Meier (1999) has shown that in the Punsly–Coroniti-like models, the differential dragging of the frames will also act as a dynamo to amplify the magnetic field at the expense of the black hole’s rotational energy and that this will have the effect of further enhancing the jet power. The results of numerical simulation studies (Koide 2003; McKinney & Gammie 2004; De Villiers et al. 2005; Hawley & Krolik 2006) are, in a broad brush sense, consistent with the BZ and Punsly–Coroniti–Meier type of models, indicating that the power of jets depends on both the mass accretion rate as well as the spin of the black hole.

In fact, Meier (2001) argues that understanding the radio-loud and radio-quiet dichotomy in the QSO population requires the jet power to exhibit such dependencies. And while this assertion appears to be indicated by observations of black hole candidate systems in our Galaxy (Cui, Zhang & Chen 1998), very little is known about the prevailing conditions underlying the jet phenomena in extragalactic AGNs. The recent *Chandra* observations of bubbles, however, offer a unique opportunity to remedy this. The X-ray observations not only provide an estimate of the power of the jets emitted by the AGNs but also the rate at which matter is accreting on to the black hole–accretion flow system. Recently A06 analysed the data for nine nearby, X-ray luminous elliptical galaxies and found a remarkably tight correlation between the Bondi accretion rates and the jet powers. Within the context of viable models for the jet–black hole–accretion flow system, such a correlation not only provides insight into the efficiency with which the rest energy of the material accreting on to the black hole is converted into jet energy but also the spin distribution of the black holes powering AGNs. In this paper, we seek to shed light on these issues.

We begin by considering two physical models for the jet power: the BZ model as described above, and a hybrid model proposed by Meier (2001) (see also Punsly & Coroniti 1990; Meier 1999), in which the magnetic field threads the plasma throughout – in the body of the disc where frame-dragging is negligible as well as within the ergosphere. This latter model combines the Blandford–Payne-like

disc acceleration mechanism with the BZ-like scheme that draws upon the rotational energy of the black hole. We improve upon these models by incorporating important general relativistic effects that previously were introduced via highly restricted approximation. Following Meier (2001), we further couple our models for the jet production to the advection-dominated accretion flow model (hereafter ADAF,¹ Narayan, Mahadevan & Quataert 1998, hereafter N98; Narayan 2005; Nemmen et al. 2006). There is a considerable body of work indicating that the launching of jets is most efficient when the accretion flow is advection dominated (e.g. Rees et al. 1982; Meier 2001; Churazov et al. 2005) and that jet production is suppressed in the standard thin accretion disc normally associated with radiatively efficient AGNs (Livio, Ogilvie & Pringle 1999; Meier 2001; Maccarone, Gallo & Fender 2003).

This paper is organized as follows: In Section 2, we describe our models for the accretion flow and the jet power. We summarize the findings of A06 in Section 3. In Section 4, we compare the correlations between the jet power and accretion rate implied by our models with the results of A06 and consider the resulting implications for jet efficiencies and the distribution of black hole spins. We also compare results for accretion models with and without winds and explore the implications of uncertainties in A06 estimates of jet power and mass accretion rates on our findings. Lastly, we highlight our main results in Section 5 and offer some concluding remarks.

2 MODELS OF ACCRETION FLOW AND JET POWER

2.1 Accretion flow structure

There are several lines of evidence that suggest that the accretion flow on to AGNs with strong jets is best described as an ADAF. First, the observed bolometric (radiative) luminosities of radio-loud AGNs are typically many orders of magnitude smaller than the luminosities expected if the mass were flowing on to the central black holes along thin accretion discs (e.g. Di Matteo et al. 2003; Taylor et al. 2006); in the latter case, the bolometric luminosities correspond to roughly 10 per cent of the rest-mass energy of the accreting matter (Shakura & Sunyaev 1973). ADAFs can not only account for the observed radiative quiescence of the jet-emitting AGNs but also reproduce their nuclear X-ray luminosities with accretion disc mass flow rates² comparable to their Bondi rates (Di Matteo et al. 2001; Loewenstein et al. 2001; Di Matteo et al. 2003; Pellegrini 2005; Soria et al. 2006). The exceptional systems (such as the radio-loud quasars) that are both highly luminous and radio-loud are currently an intriguing puzzle: possibly they are high accretion rate systems in which the disc is puffed up by the trapped energy despite the system’s low opacity (e.g. Meier 2001; Maccarone et al. 2003; KÖrding, Jester & Fender 2006; Sikora, Stawarz & Lasota 2007).

Secondly, the magnitude and the structure of the magnetic fields associated with ADAFs are much more conducive to the extraction of spin energy from the hole than those associated with standard thin discs (e.g. Rees et al. 1982; Armitage & Natarajan 1999; Livio et al.

¹ Improved models of ADAFs incorporating winds and convection are also more generally referred to as radiatively inefficient accretion flows (RIAF). We simply use the original acronym ADAF.

² We use the term ‘mass flow rate’ to distinguish the mass flowing through the accretion disc from the mass that is accreted on to the black hole itself. Due to the mass/energy carried away by winds and jets, these two rates may differ significantly.

1999). Thirdly, X-ray binaries with black hole candidates (XRBs) in the ‘low/hard’ state display a strong correlation between the X-ray and radio emissions, which can best be understood in the context of an ADAF–jet system (Gallo, Fender & Pooley 2003; Merloni, Heinz & di Matteo 2003; Falcke, K rding & Markoff 2004). Low-luminosity AGNs with jets show similar correlations and hence, are thought to be scaled-up versions of the galactic XRBs (Merloni et al. 2003; Falcke et al. 2004). One characteristic of the ADAF–jet system is low-mass flow rates. The observations suggest that at high flow rates, the black hole systems switch to the ‘high/soft’ state and the jet activity is suppressed (Maccarone et al. 2003; Greene, Ho & Ulvestad 2006). This observed ‘quenching’ of the jet (as well as the transition from the ‘low/hard’ to the ‘high/soft’) is believed to be due to a change from an ADAF (powerful jets) at low-mass flow rates to a thin accretion disc structure (weak jets) at high-mass flow rates (e.g. Meier 2001).

We describe the structure of the ADAF using the self-similar equations of Narayan & Yi (1995). We are especially interested in the analytical equations (see Appendix A) that describe the vertical half-thickness of the disc H , the angular velocity of the disc Ω' and the magnetic field strength B near the black hole in terms of radius R , black hole mass M_\bullet , accretion rate on to the black hole \dot{M} and advection parameter f (assumed ≈ 1). We require these as inputs to our models for the jet power described in Sections 2.2 and 2.3 below. We note that although the Narayan & Yi (1995) equations are based on an early model for the ADAF that did not allow for mass-loss in the form of winds from the accretion flow, we have verified that these self-similar solutions provide a good approximation to the structure of the inner regions of ADAF models, like that of Blandford & Begelman (1999) (hereafter BB99), that do allow for such losses.

The equations for H , B and Ω' also depend on the properties of the accreting fluid, such as its adiabatic index γ , its viscosity parameter α and the ratio of gas to magnetic pressure β in the fluid. These quantities are not independent. The value of γ depends on β via the relationship $\gamma = (5\beta + 8)/3(2 + \beta)$ (Esin, McClintock & Narayan 1997) and based on the MHD numerical simulations of the evolution of the magnetorotational instability in accretion discs (Hawley, Gammie & Balbus 1995), $\alpha \approx 0.55/(1 + \beta)$. We assume that the magnetic pressure is related to the field strength as $P_{\text{mag}} = B^2/8\pi$.

We constrain the value of α from recent MHD numerical simulations of RIAFs, which take into account self-consistently the role of the Maxwell stresses in establishing its value. The general relativistic simulations (e.g. Hirose et al. 2004; McKinney & Gammie 2004; Hawley & Krolik 2006) find that in the inner portions of the disc $\beta \approx 1$ –10. Although α is not explicitly quoted, we can estimate α from these values of β using the relationship between the two, obtaining $\alpha \approx 0.04$ –0.3. For computational purposes, we will derive results for $\alpha = 0.04$ and 0.3. To put these α -values in context, we note that the simulations around Schwarzschild holes of Proga & Begelman (2003) suggest that near the innermost stable circular orbit, α reaches high values: $\alpha \approx 0.1$ –0.7. Moreover, recent ADAF models of XRBs require values of $\alpha \approx 0.25$ in order to account for the observations (Quataert & Narayan 1999).

To the above model for the ADAF, we introduce three important modifications to take into account general relativistic effects induced by the Kerr metric. First, we take as input for our jet model the values of H , B and Ω' evaluated at R_{ms} , the radius of the marginally stable orbit of the accretion disc. This radius depends sensitively on the dimensionless black hole spin parameter $j \equiv J/J_{\text{max}}$ (‘ a/M ’ in geometrized units, a is the specific angular momentum), where J is the angular momentum of the hole and $J_{\text{max}} = GM_\bullet^2/c$ is the

maximal angular momentum (Bardeen, Press & Teukolsky 1972). Secondly, an observer at infinity will see the disc and the magnetic fields near the black hole rotate, in the Boyer–Lindquist coordinate system, not with an angular velocity Ω' but $\Omega = \Omega' + \omega$, where $\omega \equiv -g_{\phi t}/g_{\phi\phi}$ is the angular velocity, in the same coordinate system, corresponding to the local space–time rotation enforced by the spinning black hole (Bardeen et al. 1972). And thirdly, we take into account the field-enhancing shear caused by frame-dragging when calculating the magnetic field strength in the inner region of the disc, as first suggested by Meier (1999) and tentatively observed in recent MHD numerical simulations (Hawley & Krolik 2006). Following Meier (2001), we relate the amplified, azimuthal component of the magnetic field to the unamplified magnetic field strength derived from the self-similar ADAF solution as $B_\phi = gB$, where $g = \Omega/\Omega'$ is the field-enhancing factor. The amplitude of this factor depends on the black hole spin through the angular velocity of space–time rotation, ω . In the case of a non-rotating black hole, $\omega = 0$ and $g = 1$ (i.e. no field enhancement).

The potential importance of these effects has been recognized previously and Meier (2001) even incorporated them into his jet power model, albeit in the form of simplifications that blunted their impact. For example, he took g to be a free parameter rather than explicitly relate it to space–time rotation induced by the rotating black hole, he adopted a simple approximation for ω , and he evaluated his jet model only for two limiting radii: $7GM_\bullet/c^2$ (corresponding to $j \approx 0$) and $1.5GM_\bullet/c^2$ ($j \approx 1$). In our approach, we are able to capture more fully the dependence of the important model parameters on j and by doing so, are able to explore the sensitivity of the accretion flow and the jet power solutions to variations in the value of the black hole spin parameter.

Finally, we follow Livio et al. (1999) in assuming that the poloidal and azimuthal components of the magnetic field are related to each other as $B_p \approx H/R B_\phi$ where R is the radius in cylindrical coordinates, a relationship based on the assertion that the strength of the poloidal field is limited by the vertical extent of turbulent eddies in the disc. In the case of an ADAF, this yields $B_p \approx B_\phi$ because $H \sim R$.

2.2 The BZ jet model

According to the BZ model, the total power of the resulting jet is given by (e.g. Macdonald & Thorne 1982; Thorne, Price & Macdonald 1986):

$$P_{\text{jet}}^{\text{BZ}} = \frac{1}{32} \omega_F^2 B_\perp^2 R_H^2 j^2 c, \quad (1)$$

where $R_H = [1 + (1 - j^2)^{1/2}] GM_\bullet/c^2$ is the horizon radius, B_\perp is the strength of the magnetic field normal to the horizon, and the factor $\omega_F \equiv \Omega_F(\Omega_H - \Omega_F)/\Omega_H^2$ depends on the angular velocity of the field lines Ω_F relative to that of the hole, Ω_H . Following the usual practice, we assume that $\omega_F = 1/2$, which maximizes the power output (e.g. Macdonald & Thorne 1982; Thorne et al. 1986). As for B_\perp Livio et al. (1999) has argued that the field threading the horizon ought to be comparable in strength to the field threading the inner regions of the accretion flow. Hence, we take the field strength at the horizon to be the same as that at the radius of the marginally stable orbit of the accretion disc; that is, $B_\perp \approx B_p(R_{\text{ms}}) \approx g(R_{\text{ms}})B(R_{\text{ms}})$. These field configurations are consistent with those seen in the in the jet launch regions in numerical simulations of accretion on to Kerr black holes (e.g. Hirose et al. 2004; McKinney & Gammie 2004). Finally, since the accretion rate enters the jet power only through the field strength, the jet power only depends on the accretion rate measured at the

marginally stable orbit of the accretion flow $\dot{M}_{\text{ms}} \equiv \dot{M}(R_{\text{ms}})$. If the mass is conserved in the disc as in the early version of the ADAF model (Narayan & Yi 1995, N98) then \dot{M} is constant with radius, but contemporary ADAF models often allow for mass-loss from the accretion disc in the form of winds and outflows (see Section 2.4).

The primary advantage of the BZ model is its simplicity. The picture describes a black hole whose event horizon is equivalent to a rotating conducting surface with surface charges and currents. In effect, the horizon is a unipolar inductor on to which the magnetic fields are attached. The magnetic fields are torqued by the rotation of the black hole and in turn, drive an outflow. There are, however, concerns about the model. Punsly & Coroniti (1990) have argued that the very notion of flux being emitted from the event horizon renders the model unphysical because the event horizon is casually disconnected from events upstream. Disturbances generated at the event horizon can only propagate inward. Moreover, the BZ model is also formulated on the assumption of zero accretion that underlies the BZ model; the BZ model is derived in the limit of vanishing plasma density and pressure (Hawley & Krolik 2006). On the other hand, McKinney & Gammie (2004) have reported good agreement between several aspects of the BZ model and their numerical simulation results, and for this reason, we consider this model in the present work. We do note, however, that McKinney & Gammie (2004) also reported a growing difference between the model and their results with increasing black hole spins, a result that can partly be attributed to the fact that the BZ model is derived in the limit of slowly rotating black holes.

2.3 Hybrid model

Contrary to the premise of very low plasma densities that underlies the BZ model, accretion of matter is presumed to be a key element of real AGN systems. In fact, numerical simulations show that the coupling between the accretion flow and the magnetic fields is an essential element of jet production. For this reason, we consider a second jet model, the hybrid model of Meier (2001). As noted previously, this model is constructed so that large-scale magnetic fields thread the accretion disc outside the ergosphere as well as rotating plasma within the ergosphere that is flowing on to the black hole. Hence, the model is able to draw upon both the rotational energy of the accretion disc as well as the spinning black hole in order to drive outflows, though the extraction of the black hole rotational energy occurs indirectly through field lines anchored to the plasma subject to space–time rotation imposed by the black hole. This model also takes into account the effects of field amplification by both the differential rotation of the plasma in the body of the disc and the differential frame-dragging. As discussed by Meier (2001), the jet power in this model is a strong function of the thickness of the accretion disc and the black hole spin; strong magnetic fields and rapid rotation, the necessary ingredients for the launching of powerful jets, only arise when the disc is thick and the hole is spinning rapidly.

Following Meier (2001), the total jet power for the hybrid model is given by

$$P_{\text{jet}}^{\text{disc}} = (B_{\phi} H R \Omega)^2 / 32c, \quad (2)$$

where $B_{\phi} = gB$ and $\Omega = \Omega' + \omega$. All quantities are evaluated at $R = R_{\text{ms}}$, which is also assumed to be the approximate characteristic size of the jet-formation region. We note that we denote the jet power for this model with the superscript ‘disc’ in order to highlight that in addition to drawing upon the rotational energy of the spinning black hole, this model also draws energy from the accretion disc.

As a matter of interest, we note that there is considerable support for the Punsly–Coroniti–Meier component of the hybrid model from recent numerical simulations studies (Koide 2003; McKinney & Gammie 2004; De Villiers et al. 2005). On the other hand, the relevance of the Blandford–Payne mechanism, which describes the extraction of jet power from the rotation of the accretion disc and which is incorporated in our hybrid model, is very much a matter of debate. For example, Hawley & Krolik (2006) have recently argued that this mechanism is not at all important for understanding the collimated jet outflows seen in their numerical simulations while Blandford (2005) has suggested that such conclusions are premature and that the simulation results are dependent on the initial conditions adopted.

2.4 Jet models properties

There are several interesting properties of our jet models that ought to be highlighted. First, the jet power for neither the BZ nor the hybrid models (equations 1 and 2, respectively) depends on the black hole mass. Moreover, the dependence on the accretion rate enters the jet power only through the field strength (see Appendix A) and for the combined accretion flow–jet power model under consideration here, the jet powers for both models can be expressed as $P_{\text{jet}}(\alpha, j, \dot{M}_{\text{ms}}) \propto \dot{M}_{\text{ms}}$; that is, the jet power depends linearly on the mass accretion rate. However, both have a complicated dependence on j and α (see Appendix A). Fig. 1 illustrates the spin dependence of both the BZ and the hybrid P_{jet} models, for viscosity parameters $\alpha = 0.04$ and 0.3 , and a fiducial value of $\dot{M}_{\text{ms}} = 10^{-3} M_{\odot} \text{ yr}^{-1}$. This corresponds to one tenth of the mean Bondi accretion rate of the elliptical galaxies studied by A06.

As Fig. 1 shows, both the BZ and the hybrid P_{jet} models behave similarly and give comparable results for intermediate and high values of j . The slight difference in the jet power between the two models is not significant given the uncertainty in our understanding of the detailed mechanisms underlying jet formation. At low-spin

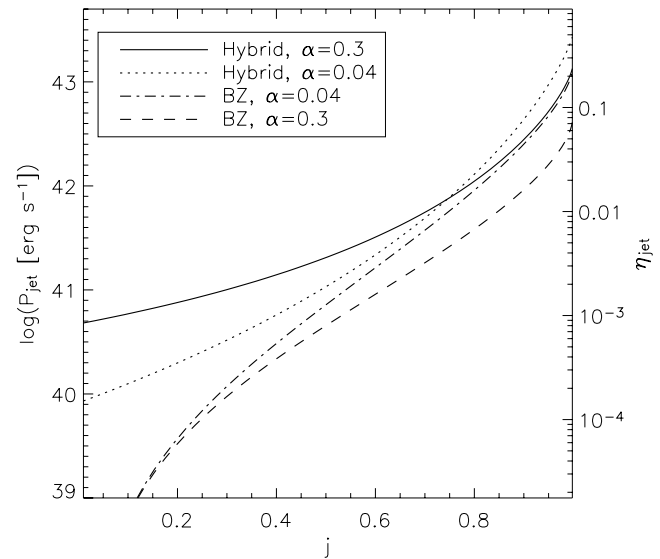


Figure 1. The jet powers for our modified BZ (equation 1) and the hybrid model (equation 2) as a function of black hole spin parameter j . The curves are computed assuming a fiducial value of $\dot{M}_{\text{ms}} = 10^{-3} M_{\odot} \text{ yr}^{-1}$, and for two different values of the viscosity parameter: $\alpha = 0.04$ and 0.3 . The right-hand axis shows the jet efficiency $\eta_{\text{jet}} \equiv P_{\text{jet}} / \dot{M}_{\text{ms}} c^2$. Since the jet powers scale linearly with \dot{M}_{ms} , η_{jet} depends only on j and α .

values, however, the jet power in the hybrid model is considerably higher than that for the BZ case. The drop-off in the BZ case reflects the decline in the black hole rotational energy available to power the jet while in the hybrid case, the decline is limited by contributions from the accretion disc. The most important point made by Fig. 1 is that the jet power for both the BZ and the hybrid models is a strong function of the black hole spin, spanning a range of three orders of magnitude as the black hole spin varies from 0 to 1. This strong dependence is largely due to our improvements over previous models, which did not incorporate carefully the physics of the Kerr metric (Ghosh & Abramowicz 1997; Armitage & Natarajan 1999; Meier 2001). This type of j -dependence for the jet power has previously only been seen in complex numerical MHD simulations of jet formation (Hawley & Krolik 2006; McKinney 2005).

The right-hand axis of Fig. 1 shows the jet efficiency factor, defined as $\eta_{\text{jet}} \equiv P_{\text{jet}}/\dot{M}_{\text{ms}}c^2$. Since our model jet powers scale linearly with \dot{M}_{ms} , the corresponding jet efficiency factor depends only on the black hole spin parameter j and the viscosity parameter α . For high values of j , the efficiencies predicted by our two jet models are comparable to the Novikov–Thorne thin disc radiative efficiency, which is related to the binding energy of the innermost stable circular orbit (Novikov & Thorne 1973), and the jet carries away a significant fraction of the rest mass energy flowing through the disc. We note that for the maximal black hole spin, $j = 0.998$, as implied by the analysis of Thorne (1974), the jet efficiencies are $\eta = 0.22$ (BZ model) and 0.48 (hybrid model) for $\alpha = 0.04$. On the other hand, if $\alpha = 0.3$, the efficiencies drop to $\eta = 0.07$ (BZ model) and 0.24 (hybrid model).

3 THE EMPIRICAL \dot{M} – P_{jet} CORRELATION

Given a distribution of gas about a central black hole, the most simple configuration describing the accretion of the gas on to the black hole is the Bondi flow model (Bondi 1952), which assumes a non-luminous central source and a spherically symmetric flow with negligible angular momentum. The resulting Bondi accretion rate can be written as $\dot{M}_{\text{Bondi}} = \pi\lambda c_s \rho r_A^2$, where $r_A = 2GM_{\bullet}/c_s^2$ is the accretion radius, G is the gravitational constant, M_{\bullet} is the black hole mass, c_s is the sound speed of the gas at r_A , ρ is the density of gas at r_A and λ is a numerical coefficient that depends on the adiabatic index of the gas. This estimate is frequently used in studies of the central X-ray emitting gas in elliptical galaxies (e.g. Di Matteo et al. 2003; Pellegrini 2005).

Recently, A06 derived the Bondi accretion rates and jet powers from *Chandra* X-ray observations of nine nearby, X-ray luminous giant elliptical galaxies that show evidence of jet-inflated cavities in their central regions, and found a tight correlation between the two. The jet powers were estimated from the energies and time-scales required to inflate cavities observed in the surrounding X-ray emitting gas such that $P_{\text{jet}} = E/t_{\text{age}}$, where E is the energy required to create the observed bubbles and t_{age} is the age of the bubble. A06 estimates are based on the assumption that the X-ray bubbles are inflated slowly. The values of \dot{M}_{Bondi} for the systems were calculated from the observed gas temperature and density profiles. In most cases, the Bondi radius is not observed directly, and the appropriate density and temperature are determined by extrapolating the observed data, typically by a factor of 3 or greater, in radius (cf. Rafferty et al. 2006). The black hole masses were deduced from the optical velocity dispersion measurements using the correlation between central black hole mass and velocity dispersion of Tremaine et al. (2002). Since A06 presents a state of the art analysis of current data, we accept the estimates of the Bondi accretion rate and jet

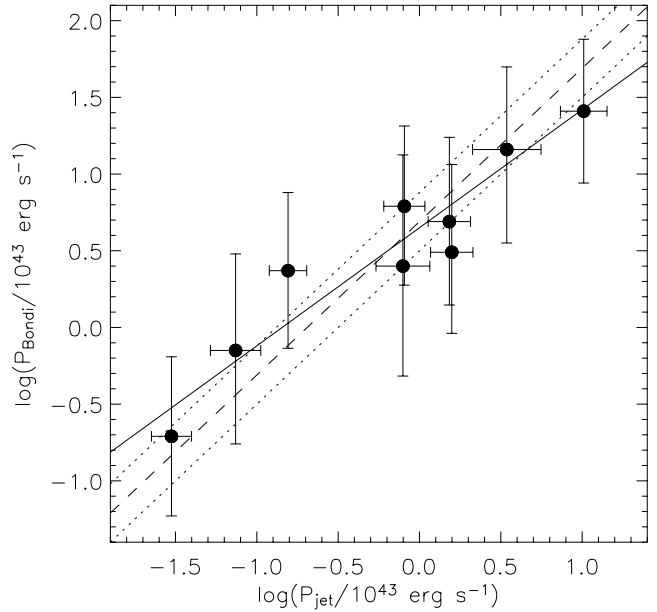


Figure 2. The empirical relationship between the Bondi accretion power ($P_{\text{Bondi}} = 0.1\dot{M}_{\text{Bondi}}c^2$) and jet power (P_{jet}) for nine nearby, X-ray luminous giant elliptical galaxies derived by A06. The error bars plotted include a systematic uncertainty of 0.46 dex in $\log P_{\text{Bondi}}$. The fitted power-law model predicted by our jet models ($A = 0.69 \pm 0.19$ and $B = 1$) is represented by the dashed (best fit) and dotted (error bars) lines; the best-fitting power-law model determined by A06 is shown as the solid line.

power at face value in what follows. Clearly it will be important to see how these observational constraints improve with future X-ray missions. In Section 4.2, we consider the impact of uncertainties in the calculation of the jet power from the physical properties of cavities observed in the *Chandra* data.

The A06 results for the nine systems are shown in Fig. 2, where P_{Bondi} is the total accretion power released for an efficiency of 10 per cent, $P_{\text{Bondi}} = 0.1\dot{M}_{\text{Bondi}}c^2$. The error bars plotted include a systematic uncertainty of 0.46 dex in $\log P_{\text{Bondi}}$, implied by the intrinsic dispersion of 0.23 dex in the $\log M_{\bullet}$ – $\log \sigma$ relation (see A06).

The correlation between \dot{M}_{Bondi} (or equivalently, P_{Bondi}) and jet power P_{jet} was expressed by A06 as a power-law model of the form

$$\log \frac{P_{\text{Bondi}}}{10^{43} \text{ erg s}^{-1}} = A + B \log \frac{P_{\text{jet}}}{10^{43} \text{ erg s}^{-1}}, \quad (3)$$

with $A = 0.65 \pm 0.16$ and $B = 0.77 \pm 0.20$. This power-law model is shown in Fig. 2 as a solid line.

4 \dot{M} – P_{jet} RELATION: IMPLICATIONS FOR BLACK HOLE SPIN AND ACCRETION RATE

The A06 results are interesting in two regards. First, they show that the central AGNs in these systems are extremely sub-Eddington; that is, they have extremely low accretion rates compared to the Eddington rate, $\dot{M}_{\text{Bondi}}/\dot{M}_{\text{Edd}} \lesssim 10^{-3}$, where $\dot{M}_{\text{Edd}} \equiv 22M_{\bullet}/(10^9 M_{\odot}) M_{\odot} \text{ yr}^{-1}$. Our ADAF model for the accretion flow is only valid for such highly sub-Eddington flows. Secondly, as noted by A06, the correlation implies that a non-negligible fraction [$P_{\text{jet}}/(\dot{M}_{\text{Bondi}}c^2) = 2.2^{+1.0}_{-0.7}$ per cent for $P_{\text{jet}} = 10^{43} \text{ erg s}^{-1}$] of the energy associated with the rest mass of the gas entering r_A is channelled into jet power.

4.1 Reconsidering the mass accretion rate

Presumably less matter than the amount predicted by the Bondi rate gets down to the black hole for several reasons. As the gas in an ADAF has angular momentum, accretion is driven not just by gravity as in the Bondi flow but also by the rate of angular momentum transport characterized by the parameter α in our ADAF model; given a certain ambient density in the external medium, the accretion rate predicted by the ADAF model is lower than the Bondi accretion rate: $\dot{M}_{\text{ADAF}} \sim \alpha \dot{M}_{\text{Bondi}}$ (e.g. N98; Proga & Begelman 2003), with $\alpha < 1$. Furthermore, some part of the gas may be prevented from being accreted due to winds and/or convection (e.g. BB99; Quataert & Gruzinov 2000; Proga & Begelman 2003; Igumenshchev, Narayan & Abramowicz 2003) occurring in the ADAF, reducing even more \dot{M}_{ms} compared to \dot{M}_{Bondi} . We allow \dot{M}_{ms} to be smaller than \dot{M}_{Bondi} by introducing the parameter ϵ_{Bondi} such that $\dot{M}_{\text{ms}} = \epsilon_{\text{Bondi}} \dot{M}_{\text{Bondi}}$. The parameter ϵ_{Bondi} represents the fraction of material supplied by the external medium at the Bondi radius that ultimately reaches the innermost stable circular orbit of the accretion flow and gets accreted afterwards. Therefore this parameter encompasses our ignorance about the possible physical processes that may modify the density profile of the accretion flow with respect to the ADAF solution. We note that here we are only considering the simplest situation where the density profile is affected. Winds too will reduce the rate of energy and angular momentum accretion on to the central black hole. We defer the discussion of such complications to Section 4.3.

Including this modification into our models gives jet power relations of the form

$$P_{\text{jet}}(P_{\text{Bondi}}, \alpha, \epsilon_{\text{Bondi}}, j) \propto P_{\text{Bondi}}. \quad (4)$$

This relationship has the same functional form as the correlation found by A06 (equation 3) and we can now use the above equation to constrain the main model parameters (α , ϵ_{Bondi} and j) using A06 results.

Comparing equations (3) and (4), it follows that our models predicts the slope $B = 1$, which is somewhat higher than the value obtained by A06, $B = 0.77 \pm 0.20$, although the difference is not statistically significant (see Fig. 2). As A and B in equation (3) presumably are strongly correlated, we fit the data of A06 to a power-law model with the fixed value $B = 1$ to find the corresponding value of A . Using the χ^2 fit statistics, which accounts only for errors in the values of P_{Bondi} , we find $A = 0.69 \pm 0.19$ with $\chi^2 = 2.4$ for eight degrees of freedom, which indicates that our $B = 1$ fit also provides a good description of the data, and hints that the systematic errors in the data points are not fully independent. The power-law model with $A = 0.69$ and $B = 1$ is plotted in Fig. 2 as the dashed line; the dotted lines delimitate the corresponding statistical uncertainty in this power law.

In the $B = 1$ model, the value of the parameter A is a measure of the jet efficiency η_{jet} . Setting $B = 1$ in equation (3) we have

$$A = \log \left(\frac{P_{\text{Bondi}}}{P_{\text{jet}}} \right) = \log \left(\frac{0.1}{\eta_{\text{jet}} \epsilon_{\text{Bondi}}} \right), \quad (5)$$

which can be recast as

$$\log(\eta_{\text{jet}}) = -(1 + A + \log \epsilon_{\text{Bondi}}). \quad (6)$$

Fig. 3 shows the jet efficiency factor corresponding to the observed value of $A = 0.69 \pm 0.19$ ($B = 1$), as a function of the mass accretion parameter ϵ_{Bondi} . The value of $\eta_{\text{jet}} \approx 2$ per cent derived by A06 is, in fact, the minimum value and corresponds to $\epsilon_{\text{Bondi}} = 1$ or $\dot{M}_{\text{ms}} = \dot{M}_{\text{Bondi}}$. Since ϵ_{Bondi} is expected to be significantly less than unity, the required jet efficiency will be higher.

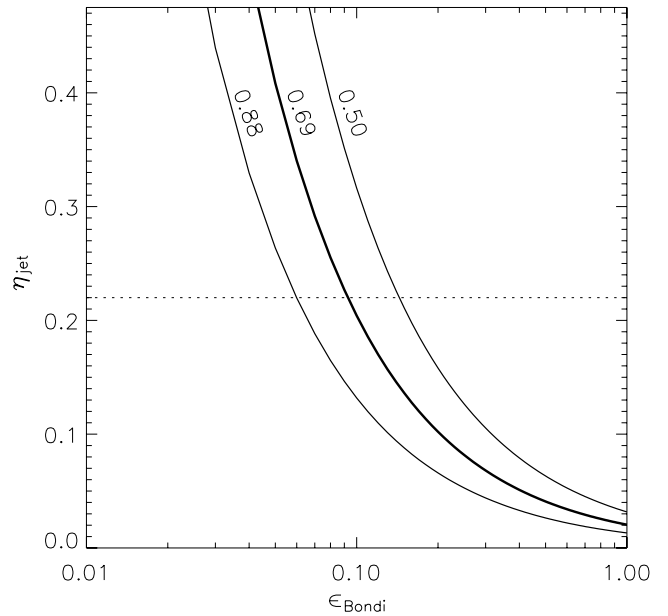


Figure 3. Relation between the parameters η_{jet} and ϵ_{Bondi} required to reproduce the observed correlation between \dot{M}_{Bondi} and P_{jet} . Each line corresponds to a different value of A (shown above the lines), the thicker line represents the best-fitting value of A . We limit the value of efficiency to the theoretical maximum for the hybrid model with $\alpha = 0.04$ (48 percent). The dotted line corresponds to the maximum efficiency for the BZ model with $\alpha = 0.04$. These correspond to maximally spinning black holes with $j = 0.998$. The inferred efficiencies decline with increasing α since the plausible values of ϵ_{Bondi} increase. For $\alpha = 0.3$, the maximal efficiencies are $\eta = 0.07$ (BZ model) and 0.24 (hybrid model).

Since in our models for the jet power, the jet efficiency factor is a function of the viscosity parameter α and the black hole spin j , the above equations imply that the value of the parameter A is a function of the parameters α , ϵ_{Bondi} and j . And therefore, for a given value of α , we can use the observed value of $A = 0.69 \pm 0.19$ to explore the range of possible values of ϵ_{Bondi} and j . As described previously, various different analysis and arguments suggest that α ranges between 0.04 and 0.3. In the discussion below, we will derive results for the two values that bound the range: $\alpha = 0.04$ and 0.3.

Fig. 4 shows the parameter space (j , ϵ_{Bondi}) corresponding to the observed values of A for the BZ (dashed line, Section 2.2) and the hybrid (solid line, Section 2.3) jet models. The two panels show the results for the two different values of α near the innermost stable circular orbit, $\alpha = 0.3$ ($\beta \sim 1$, left-hand panel) and $\alpha = 0.04$ ($\beta \sim 10$, right-hand panel). We remind the reader that β is the ratio of the gas pressure to the magnetic pressure near the radius of the marginally stable orbit. The label beside each contour is the value of A for that line. We show the effect of considering an uncertainty of ± 0.19 in A . Note that the range we have adopted assumes that the uncertainty in each data point is independent. A would be smaller if, for example, the jet power was systematically underestimated in all systems (see Section 4.2).

Based on the arrangement of the contours in this figure, we find that the observed tight correlation between accretion rates and jet powers implies a narrow range of black hole spins for the elliptical galaxies of the sample of A06, irrespective of the value of α and the specific model adopted, with the main result being that the central black holes powering the jets in these systems must be rapidly

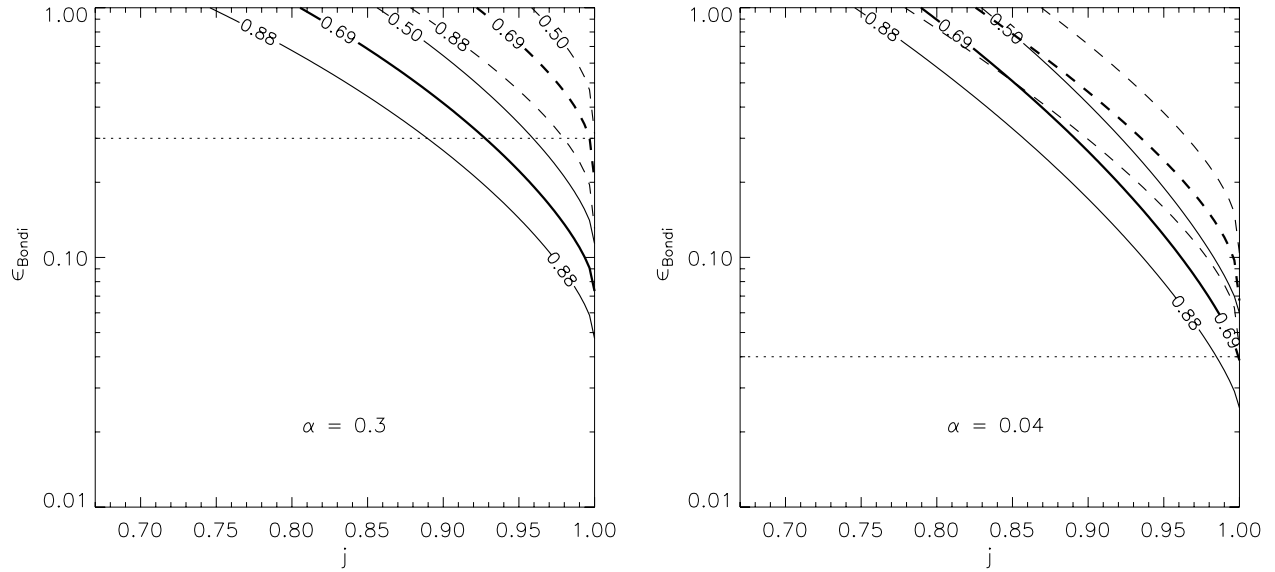


Figure 4. Range of values for the parameters ϵ_{Bondi} and j from the jet models, which reproduce the measured values of A of the empirical correlation between \dot{M}_{Bondi} and P_{jet} (equation 5). The dashed lines show the predictions of the BZ model and the solid lines show those of the hybrid model. The left-hand panel corresponds to $\alpha = 0.3$ and the right-hand panel to $\alpha = 0.04$, the two values that bound the range of plausible values for the viscosity parameter. The label beside each line is the value of A for that line. The thickest lines correspond to the best-fitting value of A . The horizontal dotted lines indicate the value $\epsilon_{\text{Bondi}} = \alpha$ (i.e. $\dot{M}_{\text{ms}} = \dot{M}_{\text{ADAF}}$).

spinning. Specifically, we find that if $\alpha \approx 0.3$ (Fig. 4a), as required to ensure an agreement between ADAF model predictions for XRBs and the observations (Quataert & Narayan 1999), then the hybrid model requires $j \gtrsim 0.75$, while the BZ model implies $j \gtrsim 0.88$. If the value of the viscosity parameter is on the low side, $\alpha \approx 0.04$ (Fig. 4b), both jet models require $j \gtrsim 0.75$. In all cases, this corresponds to $\eta_{\text{jet}} \gtrsim 2$ per cent.

These lower limits for both j and η_{jet} correspond to the situation where rate of mass accretion on to the black hole equals the Bondi rate ($\epsilon_{\text{Bondi}} \sim 1$). As discussed above, only a fraction of the Bondi rate is likely to make its way to the black hole and realistically, $\dot{M}_{\text{ms}} \lesssim \dot{M}_{\text{ADAF}}$. This restricts the allowed region of parameter space in Fig. 4 to that below the horizontal dotted lines ($\epsilon_{\text{Bondi}} = \alpha$) corresponding to $\dot{M}_{\text{ms}} = \dot{M}_{\text{ADAF}}$.

Subject to these constraints on the mass accretion rate, the observed correlations between mass accretion rate and jet power imply that in the case of the hybrid model, the central black holes must have spins $j \gtrsim 0.89$ if $\alpha \approx 0.3$. The corresponding range for the mass accretion rate on to the black hole is $\dot{M}_{\text{ms}} \approx (0.05\text{--}0.3)\dot{M}_{\text{Bondi}} \approx (0.16\text{--}1.0)\dot{M}_{\text{ADAF}}$. If α is on the low side, then the black holes in the hybrid model must be spinning rapidly ($j \gtrsim 0.98$) and the corresponding range for the mass accretion rate is $\dot{M}_{\text{ms}} \approx (0.02\text{--}0.04)\dot{M}_{\text{Bondi}} \approx (0.5\text{--}1.0)\dot{M}_{\text{ADAF}}$. For a given value of j , the corresponding jet efficiency factor can be read off from Fig. 3. For example, if $\alpha = 0.3$ and $j = 0.89$, $\eta_{\text{jet}} = 4.5$ per cent, giving an overall efficiency of $\eta_{\text{jet}\epsilon_{\text{Bondi}}} = 1.3$ per cent.

In the case of the BZ model, the spin distribution corresponds to $j \gtrsim 0.98$ ($\alpha = 0.3$) and $j > 1$ ($\alpha = 0.04$). The latter is unphysical and indicates that values of viscosity parameters as low as $\alpha = 0.04$ are disfavoured. For this line of reasoning, we find that the BZ model requires $\alpha \geq 0.2$ in order to guarantee $j \leq 1$ and $\dot{M}_{\text{ms}} \leq \dot{M}_{\text{ADAF}}$. It is also possible these ‘problems’ reflect the limitations of the BZ model. As indicated previously, the BZ model is, strictly speaking, applicable only in the limit of slow rotation where the unipolar approximation is reasonable. For this reason, and for the clarity of

the discussion, we shall hereafter only discuss the results in the context of the hybrid model though we shall continue to show the results for both models.

Since our results indicate high-spin values for jet-emitting black holes, it is useful to bear in mind that recent numerical relativistic MHD simulations of thick-discs of Gammie, Shapiro & McKinney (2004) suggest that the accreting plasma will bring the black hole in spin equilibrium not at the maximal value of $j = 0.998$ (Thorne 1974) but at a lower value of $j \approx 0.93$. This is not a trivial difference. At high-spin rates, a 7 per cent reduction in j translates into a factor of ~ 3 reduction in the jet power. In the specific case of the A06 elliptical galaxies, we find that black holes in these galaxies must be fed at rates $\epsilon_{\text{Bondi}} \gtrsim 0.05$, depending on the value of α and the mechanism of jet powering at work. In the particular case of a hybrid model with $\alpha \approx 0.3$, the black holes cannot be fed at rates much lower than \dot{M}_{ADAF} (i.e. $0.2 \lesssim \epsilon_{\text{Bondi}} \lesssim 0.3$, where the lower limit is due to the reduced efficiency associated with $j = 0.93$ and corresponds to $\dot{M}_{\text{ms}} = \dot{M}_{\text{ADAF}}$), limiting the potential role of mass-loss through winds.

4.2 Reconsidering the jet power estimates

A06 uses the X-ray observations of the intracluster medium cavities to estimate the jet power. To derive the energy E required to create the observed bubbles in the X-ray emitting gas, A06 assume that the cavities are inflated slowly and obtain the relationship $E = 4PV$ for $\gamma = 4/3$ (relativistic plasma), where P is the thermal pressure of the surrounding X-ray emitting gas, V is the volume of the cavity and γ is the mean adiabatic index of the fluid within the cavity. This is the minimum energy required to inflate the cavities and does not take into account any additional energy that may have gone into heating the intracluster medium. A more realistic scenario is likely to involve overpressurized bubbles, which upon injection expand rapidly to reach pressure equilibrium with their surroundings, and in

the process generate weak shock waves that heat the X-ray emitting gas. Nusser et al. (2006) calculate the total energy E' deposited by the jets in such cases and find that

$$E' \approx 3PV \left(\frac{P_i}{P} \right)^{1/4}, \quad (7)$$

where P_i is the thermal pressure of the surrounding gas at the location where the bubble is injected. Equation (7) suggests that if a bubble is injected with overpressure $P_i/P \gtrsim 10$ then $E' \gtrsim 6PV$. We consider the possibility that as an extreme case the bubble energies are twice the value assumed by A06 ($E' = 8PV$), and calculate the impact of this on the observed jet powers and on the results derived from our models.

Following A06, we calculate the modified jet powers as $P'_{\text{jet}} = E'/t_{\text{age}} = 2P_{\text{jet}}$ and refit the power-law model with $B = 1$ to the modified data, taking into account the proper error propagation on the values of P'_{jet} . We obtain $A = 0.39 \pm 0.19$ with $\chi^2 = 2.4$ for eight degrees of freedom. As the assumed value of E' implies higher jet powers, our models need higher values of j to reproduce the increased values of P'_{jet} . In particular, if $\alpha \approx 0.3$ these values of A imply $j_{\text{min}} \approx 0.84$ for the hybrid model and $j_{\text{min}} \approx 0.95$ for the BZ model. Therefore, a narrower range of still large spins is required to explain P'_{jet} and the implied numerical values of ϵ_{Bondi} and ϵ_{ADAF} also increase slightly.

One important uncertainty in the above calculation (and also in that of A06) is the estimate of the age of the bubbles. A06 calculate the ages using the formula $t_{\text{cs}} = D/c_s$, where D is the distance of the bubble centre from the black hole and c_s is the adiabatic sound speed, but as discussed by Birzan et al. (2004) (see also Rafferty et al. 2006) there are two other ways of estimating the age of the cavities (equations 3 and 4 of Birzan et al. 2004) which result in longer time-scales when compared to t_{cs} . An underestimate of the time-scales involved impacts the estimate of jet power, as determined above, by artificially enhancing the jet power above the ‘true’ value. To take this into account, we consider the possibility that the A06 ages are too low by a factor of ~ 2 (Birzan et al. 2004; Rafferty et al. 2006), such that the modified age is $t_{\text{age}} \approx 2t_{\text{cs}}$, implying a smaller jet power $P'_{\text{jet}} = 1/2 P_{\text{jet}}$. With reference to the original jet power estimates of A06, we find that revising the time-scales as suggested here results, not surprisingly, in a systematic decrease in lower limits of j -values. For example, in the case of the $\alpha = 0.3$ hybrid model, we find that $j_{\text{min}} \approx 0.6$.

In spite of the above-mentioned concerns with the A06 estimates of the total jet energy and the associated time-scales, it seems likely that their end-product, the estimates of the jet power in the nine elliptical galaxies, is reasonable. The impact of underestimating the total jet energy very nearly cancels out the impact of underestimating the time-scales involved and the analysis that we have presented thus far will only be minimally affected.

4.3 Reconsidering ADAF models: the effect of accretion disc winds

Finally, we have alluded previously to the possibility that the accretion flow may be modified by mass-loss in the form of winds. Winds from the accretion flow are likely to be the norm rather than the exception (e.g. Blandford 2005) and these winds will remove mass, angular momentum and energy from the flow. As a consequence, ADAF-like accretion flows with winds (advection-dominated inflow–outflow solution, ADIOS model, BB99) will have a different dynamical structure than ADAFs without winds. For in-

stance given the same value of \dot{M} near the black hole, the angular velocity, scale height, total pressure and magnetic field strength predicted by the ADAFs with and without winds will be different. And since our models for the jet power are linked to the accretion flow model through the quantities H , Ω' and B , it is not inconceivable that jet powers linked to no-wind ADAFs could be very different from jet powers linked to ADIOS models.

The ADIOS solution has three parameters in addition to those characterizing the no mass-loss ADAF model: p_w , λ_w and ϵ_w . These parameters describe how much mass, energy and angular momentum, respectively, the wind removes from the accretion flow and specific choice of values for these three parameters leads to different types of winds (for more information see BB99). To assess the implications for jet power if the true underlying model is more correctly described by the ADIOS solution, we begin by fixing the parameters common to the two models (the ADIOS and the no-wind model) as follows: $\alpha = 0.1$, $\gamma = 1.5$ and $M_{\bullet} = 10^9 M_{\odot}$ (we have confirmed that varying these particular variables does not affect our conclusions). We then consider various ADIOS solutions arising from varying the wind parameters across the following range: $p_w = 0-1$, $\lambda_w = 0.1-0.75$ and $\epsilon_w = 0.1-0.5$. This range of values encompasses all the interesting types of winds.

We found that while some features of the two models are very different, the solutions for H , Ω' and B at the radius of marginally stable orbit in the ADIOS model are, to first order, comparable to the corresponding solutions derived from the no-wind model. In detail, the differences are such that the ADIOS jet powers are always smaller than the no-wind values by a factor of order unity, regardless of the combination of values of the wind parameters. Lower jet powers in the ADIOS model, in turn, imply both a narrower range of spins for jet-emitting black holes and that these black holes must be spinning faster than those embedded in the corresponding no mass-loss ADAF in order to reproduce the correlations derived by A06. In other words, if winds are indeed the norm, then the central black holes must be spinning even more rapidly than suggested by our analysis above.

4.4 Comparison with previous estimates of jet efficiency

A06 reported an efficiency of conversion of \dot{M}_{Bondi} into jet power of $P_{\text{jet}}/\dot{M}_{\text{Bondi}}c^2 \approx 2$ per cent for $P_{\text{jet}} = 10^{43}$ erg s $^{-1}$. In our modelling, this corresponds to the case where $\epsilon_{\text{Bondi}} = 1$ (i.e. $\dot{M}_{\text{ms}} = \dot{M}_{\text{Bondi}}$). Since it is more likely that $\dot{M}_{\text{ms}} \lesssim \dot{M}_{\text{Bondi}}$, our models suggest that the efficiency of conversion of the accreting matter into jet power is considerably higher. For instance, η_{jet} may reach values as high as ≈ 50 per cent for high-spin rates. The extraction of spin energy from the holes is responsible for this noticeable increase in the jet efficiency. We note that the increase in the jet power with black hole spin has been verified in numerical simulations of jets (e.g. De Villiers et al. 2005; McKinney 2005; Hawley & Krolik 2006).

The upper limit we have obtained for η_{jet} is more than an order-of-magnitude larger than that reported by Armitage & Natarajan (1999). This result is mainly due to the fact that we have included a Kerr metric shear-driven dynamo (Meier 1999), which enhances the field strength and was not included by these authors (see Section 2.1 and Appendix A). Based on the calculations of Armitage & Natarajan (1999), Cao & Rawlings (2004) have recently asserted that accretion flow–jet models of the kind we have considered cannot account for the jet powers of a sample of low accretion rate 3CR FR I radio galaxies observed using the *Hubble Space Telescope* (jet powers in the range $\sim 10^{41}-10^{45}$ erg s $^{-1}$). As a result of much higher

efficiencies, adopting the critical accretion rate $\dot{M}_{\text{crit}} \sim \alpha^2 \dot{M}_{\text{Edd}}$ (Esin et al. 1997) above which the ADAF solution ceases to be valid, and taking $j = 0.998$ and $\alpha = 0.3$, our jet models yield $P_{\text{jet}} \approx 10^{45-46} \text{ erg s}^{-1}$, which is more than enough to account for the observed jet power in the sample used by Cao & Rawlings (2004).

The adopted model for the jet-formation mechanism predicts a linear relation between accretion rate and jet power, while the measured best-fitting slope of the relation measured by A06 is marginally better fit by a non-linear relationship. More than likely, this marginal discrepancy is due to the small number of systems in the A06 sample. However, we note that if ϵ_{Bondi} has a slight dependence on j , as suggested by recent numerical simulations of Hawley & Krolik (2006), then ϵ_{Bondi} will also exhibit a weak dependence on P_{jet} and the agreement of the model with the observed correlation will be improved. We verified that if $\epsilon_{\text{Bondi}} \propto P_{\text{jet}}^{0.2}$ the slope predicted by the model agrees with the observed best-fitting slope. In the numerical simulations, the dependence of \dot{M}_{ms} on j arises due amplified magnetic fields that act to transfer angular momentum from the hole to the accretion disc. In addition, it is not inconceivable that the dependence $\epsilon_{\text{Bondi}}(P_{\text{jet}})$ might also emerge as a consequence of feedback effects of the jet on the material fuelling the accretion flow.

5 CONCLUSIONS

We have employed two physical models for the black hole–accretion flow–jets with features that are broadly consistent with the results of numerical simulation to understand the empirical correlation between accretion rates and jet powers of X-ray luminous elliptical galaxies derived by A06: the classical BZ model and a hybrid model proposed by Meier (2001). In the BZ model the energy is extracted electromagnetically from the spinning black hole; in the hybrid model the magnetic field threads both the inner and the outer regions of the accretion flow and the energy that powers the outflow is extracted from both the rotation of the disc as well as from the rotation of the black hole, via the frame-dragged accretion flow inside the hole’s ergosphere. We assume that the accretion flow is advection dominated (ADAF) and take into account general relativistic effects not fully appreciated before in these models.

In the absence of disc winds and feedback, the model suggests that the jet power should be linearly dependent on the accretion mass flow rate. The normalization of this relation is dependent on the disc viscosity and the black hole spin, with the dispersion around the best-fitting correlation of A06 being caused by different values of the black hole spin, j , and the ratio of the disc mass flow rate to the accretion rate on to the black hole, ϵ_{Bondi} . We find that the jet efficiency ($\eta_{\text{jet}} = P_{\text{jet}}/\dot{M}_{\text{ms}}c^2$) can exceed 10 per cent so that the jet may carry away an appreciable fraction of the rest mass energy of the accreted material.

We compared our jet power models to the jet power estimates made by A06. Adopting typical values of the viscosity parameter $\alpha \sim 0.04\text{--}0.3$, the \dot{M}_{Bondi} versus P_{jet} correlation implies a narrow range of spins $j \approx 0.75\text{--}1$ and accretion rates $\dot{M}_{\text{ms}} \approx (0.04\text{--}1) \dot{M}_{\text{Bondi}}$. If we further demand that the mass accretion rate be restricted to $\dot{M}_{\text{ms}} \lesssim \dot{M}_{\text{ADAF}} \sim \alpha \dot{M}_{\text{Bondi}}$, as is both likely to be the case and consistent with ADAF model expectations, we find that the observed correlations require $j \gtrsim 0.9$; that is, the correlations imply rapidly spinning black holes in all of the target galaxies. If additionally, the ADAF accretion flows also experience mass, energy and angular momentum loss via winds as in the ADIOS model proposed by BB99, the correlations indicate nearly maximally spinning black holes (see Section 4.3). It is reassuring that semi-analytic cosmolog-

ical simulations of the spin evolution of black holes through mergers and gas accretion (Volonteri et al. 2005) and estimates of the radiative efficiencies of global populations of quasars based on Soltan-type arguments (e.g. Soltan 1982; Yu & Tremaine 2002; Wang et al. 2006) suggest that most nearby massive holes are rapidly rotating.

The relatively small scatter in the correlation between P_{Bondi} and P_{jet} of A06, combined with the strong dependence of jet power on black hole spin implies large black hole spins is probably a general result valid for all elliptical galaxies. Of course the sample considered by A06 is relatively small, but if this result holds up in larger samples, and can be shown to also apply to the larger central radio galaxies of galaxy clusters (Rafferty et al. 2006), our results suggest that they have the most powerful jets because they have sufficient black hole mass to host relatively high-mass flow rates while the disc remains in the ADAF state (and hence is able to generate strong poloidal magnetic fields near the black hole). In this picture, all central ellipticals have the capacity to produce powerful jets, and the resulting jet power is determined by the structure of the accretion disc. Our results reveal a potentially fundamental connection between black holes and the formation of the most massive galaxies: the central holes in the cores of galaxy clusters must be rapidly rotating, in order to make the radio jets powerful enough to provide an effective feedback mechanism and quench the cooling flows, therefore preventing star formation and explaining the observed galaxy luminosity function (e.g. Bower et al. 2006; Croton et al. 2006) as well as accounting for the observed X-ray luminosity–temperature and X-ray luminosity–Sunyaev–Zeldovich effect correlations for the intracluster medium (e.g. Babul et al. 2002; McCarthy et al. 2004).

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APPENDIX A: DERIVATION OF THE JET POWER

We list the equations we used to compute the dependence of the jet power on α , j and \dot{M} using the BZ model (Section 2.2) and the hybrid model (Section 2.3). The jet power is given by equation (1) (BZ model) and equation (2) (hybrid model). The code that implements the equations described in this work and returns the jet power is available at the URL <http://www.if.ufrgs.br/~rms/jetpower.htm>.

The following equations describe the self-similar ADAF structure (Narayan & Yi 1995), where we use the black hole mass in solar units ($m = M_{\bullet}/M_{\odot}$), accretion rates in Eddington units ($\dot{m} = \dot{M}/\dot{M}_{\text{Edd}}$, \dot{M}_{Edd} is the Eddington accretion rate defined in Section 3) and radii in Schwarzschild units [$r = R/(2GM_{\bullet}/c^2)$]:

$$\Omega' = 7.19 \times 10^4 c_2 m^{-1} r^{-3/2} \text{ s}^{-1}, \quad (\text{A1})$$

$$B = 6.55 \times 10^8 \alpha^{-1/2} (1 - \beta)^{1/2} c_1^{-1/2} c_3^{1/4} m^{-1/2} \dot{m}^{1/2} r^{-5/4} \text{ G}, \quad (\text{A2})$$

$$H/R \approx (2.5c_3)^{1/2}. \quad (\text{A3})$$

The constants c_1 , c_2 and c_3 are defined as

$$\begin{aligned} c_1 &= \frac{5 + 2\epsilon'}{3\alpha^2} g'(\alpha, \epsilon'), \\ c_2 &= \left[\frac{2\epsilon'(5 + 2\epsilon')}{9\alpha^2} g'(\alpha, \epsilon') \right]^{1/2}, \\ c_3 &= c_2^2 / \epsilon', \\ \epsilon' &\equiv \frac{1}{f} \left(\frac{5/3 - \gamma}{\gamma - 1} \right), \\ g'(\alpha, \epsilon') &\equiv \left[1 + \frac{18\alpha^2}{(5 + 2\epsilon')^2} \right]^{1/2}. \end{aligned}$$

The relations among α , γ and β are given in Section 2.1. The angular velocity of the field seen by an outside observer at infinity in the Boyer–Lindquist frame is $\Omega = \Omega' + \omega$, where the angular velocity of the local metric is given by (Bardeen et al. 1972)

$$\omega \equiv -\frac{g_{\phi t}}{g_{\phi\phi}} = \frac{2aM_{\bullet}}{a^2(R + 2M_{\bullet}) + R^3}, \quad (\text{A4})$$

using geometrized units ($G = c = 1$).

We estimate the field-enhancing shear caused by the Kerr metric following Meier (2001) as $g = \Omega/\Omega'$, such that the azimuthal component of the field is given by $B_{\phi} = gB$ (see Section 2.1). The poloidal component is related to the azimuthal component following Livio et al. (1999) as $B_p \approx H/R B_{\phi} \approx B_{\phi}$.

Lastly, we insert all the quantities defined above into equations (1) and (2), and then evaluate the resulting equations at the marginally stable orbit of the accretion disc R_{ms} , given by (Bardeen et al. 1972)

$$\begin{aligned} R_{\text{ms}} &= M_{\bullet} \{3 + Z_2 - [(3 - Z_1)(3 + Z_1 + 2Z_2)]^{1/2}\}, \\ Z_1 &\equiv 1 + (1 - j^2)^{1/3}[(1 + j)^{1/3} + (1 - j)^{1/3}], \\ Z_2 &\equiv (3j^2 + Z_1^2)^{1/2}. \end{aligned} \quad (\text{A5})$$

Taking $R = R_{\text{ms}}$ in the BZ model corresponds to assume that the strength of the magnetic field at the horizon of the hole is very similar to the corresponding strength at the marginally stable orbit. This is a reasonable assumption according to recent numerical simulations of jet formation (e.g. Hirose et al. 2004; McKinney & Gammie 2004).

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