

The changing ISM of massive elliptical galaxies and cosmic evolution of radio galaxies and quasars

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ABSTRACT

The recently discovered apparent dramatic expansion in the effective radii of massive elliptical galaxies from $z \simeq 2$ to $z \simeq 0.1$ has been interpreted in terms of either galaxy mergers or the rapid loss of cold gas due to AGN feedback. In examining the latter case we have quantified the extent of the expansion, which is uncertain observationally, in terms of the star formation parameters and time of the expulsion of the cold gas. In either case, the large global decrease in stellar density should translate into a major drop in the ISM density and pressure with cosmic epoch. These cosmological changes are expected to have a major influence on the gas accretion mode, which will shift from ‘cold’ thin disk accretion at high redshifts toward ‘hot’ Bondi fed ADAF accretion at low redshifts. The decline of angular momentum inflow would then lead to a spin down of the black hole, for which we have calculated more precise time scales; a value of about 0.2 Gyr is typical for a $10^9 M_\odot$ central black hole. These results have implications for the different cosmological evolutionary patterns found for the luminosity functions of powerful and weak radio galaxies.

Key words: black hole physics – galaxies: active – galaxies: evolution – galaxies: jets – galaxies: ISM – galaxies: interactions

1 INTRODUCTION

Powerful radio galaxies (RGs) are known to be hosted by the most massive ($10^{11-12} M_\odot$) elliptical galaxies (Matthews, Morgan & Schmidt 1964) at all cosmic epochs (e.g., De Brueck et al. 2002; Rocca-Volmerange et al. 2004; Seymour et al. 2007; Nesvadba et al. 2007). Moreover, the host galaxies are usually found to be either isolated or located in groups as opposed to rich clusters (e.g., Longair & Seldner 1979; Best 2004; Hardcastle, Evans & Croston 2007). Extensive radio and optical studies have confirmed the original inference reached from radio source counts (Longair 1966) that, as compared to the present epoch, the space density of powerful RGs and radio-loud quasars was a factor of $\sim 10^3$ higher during the ‘quasar era’ ($z = 2-3$) (e.g., Dunlop & Peacock 1990; Willott et al. 2001; Grimes et al. 2004). A similarly strong cosmic evolution is exhibited by the optically luminous radio-quiet quasars (RQQs) (e.g., Hartwick & Schade 1990; Wall et al. 2005) which too reside almost exclusively in massive ellipticals (e.g., Falomo et al. 2008).

In contrast, low power radio galaxies, i.e., those with $P_{178\text{MHz}} \leq 1 \times 10^{25} \text{ W Hz}^{-1} \text{ sr}^{-1}$ and typically of Fanaroff-Riley

(1974) morphology class I, exhibit much weaker cosmic evolution, amounting to only a factor of ~ 10 increase in their abundance between $z \sim 0$ and $z \sim 3$ (e.g., Jackson & Wall 1999; Willott et al. 2001). It has been proposed that the remarkable cosmological evolution may be a manifestation of a fundamental change in the dominant gas accretion mode powering the AGN since the quasar era, which has shifted from a predominantly ‘cold’ thin disc accretion at high redshifts toward ‘hot’ Bondi type accretion at low redshifts (e.g., Hardcastle et al. 2007; cf. Cao & Rawlings 2004). The hot accretion (dominant at low z) should funnel rather little angular momentum into the central supermassive black hole (SMBH), leading to the possibility of its spinning down. For such RGs, the ADAF accretion mode (e.g., Narayan & Yi 1995) is usually invoked to explain their observed faint disc emission and low-excitation optical spectra. Note that while such RGs mostly have an FR I radio morphology, this is likely to be dictated mainly by environmental factors, especially the ambient density and its gradient (e.g., Gopal-Krishna & Wiita 1988, 2000; Snellen & Best 2001; Hardcastle et al. 2007; Perucho & Martí 2007; Melliani et al. 2008).

An observational clue linked to the hypothesis of change in the AGN accretion mode with cosmological epoch comes from the drastic decrease in the cool gas content of massive ellipticals since the quasar era. Observations of CO in nearby massive field ellipti-

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cals (within 25 Mpc) have yielded very low detection rates. Curiously, the detection rates are higher for their less luminous counterparts (Sage et al. 2007; Combes et al. 2007; also, Somerville et al. 2008). This counter-intuitive result is, however, fully consistent with the idea that more massive ellipticals preferentially lost their cold gas due to a stronger AGN feedback (e.g., Sazonov et al. 2005; Springel et al. 2005; Croton et al. 2006; Bower et al. 2006; Hopkins et al. 2008). In contrast, large amounts of molecular gas in high- z massive ellipticals have been found from CO detections (e.g., Papadopoulos et al. 2000; De Brueck et al. 2005; Klammer et al. 2005) as well as PAH detections (Lutz et al. 2008).

It is interesting that the fraction of sub-millimetre bright RGs shows a sharp decline from $> 50\%$ at $z > 2.5$ to $< 15\%$ at $z < 2.5$ (Reuland et al. 2004; also, Greve et al. 2006; Seymour et al. 2007). This shows that in RGs the bulk of the expulsion of the cold gas and/or its conversion into stars was largely accomplished during the quasar era.

The available data favour the premise that jet-driven gaseous outflows were common in RGs during the quasar era and that such outflows may have played an important role in the evolution of those massive galaxies (e.g., Nesvadba et al. 2008). The mechanical feedback of $\approx 5\%$ of the QSO's bolometric power during its peak phase can supply the energy needed to induce such intense outflows (e.g., Granato et al. 2004; Hopkins et al. 2005). Powerful outflows, if produced by nearly-Eddington or even super-Eddington accretion onto black holes, have been argued to be fundamentally responsible for: the coordinated growth of the masses of the galactic bulge and the SMBH (e.g., Silk & Rees 1998; Fabian 1999; King 2003); controlling the rate of growth of galaxies (e.g., Bower et al. 2005; Croton et al. 2006; Best et al. 2006); as well as offsetting the cluster cooling flows via heating of their ICM (e.g., Binney & Tabor 1995; Fabian et al. 2003; Birzan et al. 2004).

Whereas the expulsion of huge quantities of cold gas from massive elliptical galaxies over the past ~ 10 gigayears may well have transformed the nature of their AGN activity, it has recently been invoked (Fan et al. 2008) to explain another intriguing result, namely the apparent expansion of the effective radii of massive elliptical galaxies by a factor of $\simeq 3-4$ since the quasar era, as found in several independent studies (e.g., Ferguson et al. 2004; Trujillo et al. 2007; Zirm et al. 2007; van Dokkum et al. 2008; Cimatti et al. 2008; Damjanov et al. 2009). These 'superdense' ellipticals are, however, found to be exceedingly rare in the local universe (e.g., Trujillo et al. 2009). The largest sample of distant galaxies used in these studies involved over 800 galaxies with $0.2 < z < 2$ and masses between $\sim 5 \times 10^{10}$ and $\sim 5 \times 10^{11} M_{\odot}$ (Trujillo et al. 2007). These ellipticals were strongly argued to be much smaller at $z \geq 1.5$, with the more concentrated (spheroidal) galaxies having sizes roughly four times smaller (Trujillo et al. 2007) than those found locally in a large sample drawn from the SDSS catalogue (Shen et al. 2003). The less concentrated (disk dominated) galaxies also showed significant, albeit much more modest, size evolution for similarly high masses.

However a very recent alternative analysis, using a tiny but very well observed set of HzRGs, claims that there need be no significant expansion of massive ellipticals and that the multiple earlier claims of this important effect are in works that did not account for the fact that single Sersic profiles are poor matches to the massive ellipticals at low z (Hopkins et al. 2009). These authors assert, firstly, that they do not find the cores of nearby ellipticals to be substantially different in stellar density and size, when compared to their high z counterparts. Secondly, the absence of extended envelopes in the existing optical images of high z ellipticals (similar

to those seen in low z ellipticals) could either be a real effect or merely an artifact of their low surface brightness (Hopkins et al. 2009). The former possibility would be consistent with the recent work showing that minor mergers would expand the outer portion of the elliptical galaxy (e.g., Naab et al. 2009). On the other hand, the latter alternative would imply that no substantive change has occurred in the stellar distribution of the massive ellipticals over the past ~ 10 Gyr that have elapsed since the quasar era. If true, this would further accentuate the difficulty in explaining the enormous cosmic evolution observed both in the cold gas content of massive ellipticals and, even more, in their AGN activity (see above).

If indeed the evolutionary history of RGs is associated predominantly with a transition from a cold disc accretion to a hot accretion mode, as mentioned above, a primary requirement would be getting rid of most of the cold gas content of the massive elliptical hosts of the RGs after the quasar era. The changed accretion mode would then also result in spinning down of the central engine. Therefore, a more detailed modelling of both these processes, namely the cold gas expulsion and BH spin-down, is of considerable importance and the present paper will address these two issues. An improved understanding of the molecular gas expulsion mechanism has now acquired added significance in view of the recent controversy over the claimed cosmological expansion of the massive ellipticals.

To further explore these questions, in Sect. 2.1 we note how the possible expansion of the stellar distributions of massive elliptical galaxies could occur via mergers and in Sect. 2.2 we elaborate upon the cold gas expulsion scenario which can lead to the very large expansion claimed to have occurred throughout the bulk of the galaxies. Different modes of accretion for RGs at different cosmic epochs are considered in Sect. 3. The closely related key issue of the spin evolution of the nuclear SMBH in massive ellipticals is then discussed in Sect. 4. Given the current debate over the reality of the cosmic expansion of massive ellipticals, we shall treat it for the present as an unsettled issue while discussing the cosmic evolution of the AGN population. However, this uncertainty should not affect our main conclusions, which are given in Sect. 5.

2 GALAXY EXPANSION MECHANISMS

2.1 Galaxy expansion from galaxy mergers

Taking the cosmic expansion of massive ellipticals to be a real effect, Trujillo et al. (2007) have invoked major 'dry' mergers (those in which substantial cold gas is absent). These are expected to dominate at moderate redshifts over the 'wet' mergers (those with substantial amounts of cold gas present) that could be more common at $z > 2-3$ (e.g., van Dokkum 2005). Simulations of major dry mergers yield few new stars but they do puff up the effective radius of the galaxy, roughly as $r_e \propto M_{\star}^{0.65-1.3}$, with $M_{\star}$ the stellar mass and with the exponent in this relationship declining as the pericentre distance between the conjoining galaxies increases (Boylan-Kolchin, Ma & Quataert 2006). Thus, two or three nearly equal mass mergers over the past several billion years could account for the claimed observed increase in the galaxy size and might also explain the prevailing age-uniformity found in local massive ellipticals (Trujillo et al. 2007).

But a recent analysis by Bezanson et al. (2009) strongly indicates that major mergers face greater difficulty in producing the reported large expansions and would also probably violate independent constraints on mass growth provided by the evolution of

the galaxy mass function. Brighter cluster galaxies (BCGs) have been reported to evolve in size with redshift even more quickly than most of the early type population; this is most easily understood if BCGs grow from many smaller dry mergers (Bernardi 2009). Yet another study also indicates that dry mergers have not been important for the evolution of most ellipticals out to $z \sim 0.7$ (Scarlata et al. 2007). These different claims may be consistent if results from a detailed study of Virgo cluster ellipticals (Kormendy et al. 2009) are generally applicable. The brighter ellipticals in Virgo all have cuspy cores while the fainter ones do not and these coreless ellipticals have extra light (above a Sersic profile) in their innermost portions. That light could arise from star formation following wet mergers, while the more massive cuspy cored galaxies would have grown through dry mergers (Kormendy et al. 2009). Both minor mergers and dynamical expansion due to mass loss (Sect. 2.2) can yield the large increases in galaxy size and also accommodate the mass function evolution. They both could also lead to significantly more expansion seen in the bulk of the stellar distribution as opposed to just the core (within the central 1 kpc). The “fine tuning” in the amount of mass loss needed for that scenario to be effective may lead to a preference for the minor mergers as the cause for most of the observed expansion (Bezanson et al. 2009).

A detailed simulation of minor mergers onto a massive elliptical was recently conducted by Naab et al. (2009). They find that the dramatic increase in galaxy size, argued to be present by many papers over the past few years, can indeed be explained in this way. Very interestingly, this simulation also shows that the inner portion of the stellar distribution remains dominated by the stars in the original large elliptical, while the more extended outer reaches are mostly filled with the remnants of the smaller galaxies it absorbs (Naab et al. 2009). Another recent study of the evolution of early-type galaxies indicates that variations in the times at which star formation halted for different galaxies can explain about half of the observed expansion and dry mergers can explain the remainder of it (van der Wel et al. 2009). Note however, that a recent very large sample of 150,000 galaxies matched between SDSS and the FIRST radio catalog (Becker et al. 1995), extending to look-back times of $\simeq 2$ Gyr, does not show evidence for mergers or other external environmental factors playing a significant role in triggering nuclear activity in either spiral or elliptical galaxies (Reviglio & Helfand 2009).

The merger events of the central black holes are likely to bring down the net spin of the new hole due to the random addition of the two spins and the spin can be further reduced due to radiation of the angular momentum by gravitational radiation (Hughes & Blandford 2003).

2.2 Cold gas expulsion induced galaxy expansion

An alternative explanation for the apparently observed expansion of the effective radii of massive ellipticals by a factor of ~ 3 – 4 since the quasar era invokes rapid expulsion of cold gas due to AGN feedback in the form of winds and jets (Fan et al. 2008; Sect. 1). The presence of large amounts of molecular gas in massive galaxies at such high redshifts is indicated by the CO detections (e.g., Papadopoulos et al. 2000; De Brueck et al. 2005; Klammer et al. 2005; Lutz et al. 2008), corroborated by the result that the median redshift of sub-millimetre selected galaxies also coincides with the quasar era $z = 2$ – 3 (e.g., Chapman et al. 2005). Such massive galaxies are typical hosts for bright QSOs (e.g., Falomo et al. 2008), and observations of sub-mm QSOs indicate that very substantial amounts of cold gas are present in the central regions of their hosts (e.g.,

Omont et al. 2003; Wang et al. 2007; Lutz et al. 2008). Such sub-mm QSOs at redshifts of 2–3 are presumed to be at the transition between very rapid star formation and the beginning of powerful outflows launched by the luminous central engine which, however, may still be shrouded by dust present in the cold gas (e.g., Hopkins et al. 2005). The amount of cold gas found in these galaxies at such stages is comparable to the stellar mass, of at least the inner several kpc (Coppin et al. 2008; Tacconi et al. 2008).

Direct spectroscopic evidence that is consistent with a very powerful outflow in a high z RG (HzRG), with an outflow velocity of $\sim 1000 \text{ km s}^{-1}$ and mass loss rate of $\sim 3,000 M_{\odot} \text{ yr}^{-1}$ has recently been reported (Prochaska & Hennawi 2009). Likewise, good evidence has accumulated for massive outflows of ionized gas from HzRGs ($z = 2$ – 3), possibly due to jet feedback transferring around 10% of its power to the outflowing gas (e.g., Nesvadba et al. 2008; also, Villar-Martin et al. 2007). Also, Tremonti, Moustakas & Diamond-Stanec (2007) found strong outflows exceeding 700 km s^{-1} and large mass loss rates in 10 out of 14 galaxies with fading starbursts and AGN, albeit with $z \sim 0.6$. Under suitable conditions, even several times weaker outflows would be able to expel the cold gas in roughly one Salpeter time ($\simeq 4 \times 10^7 \text{ yr}$) from the visible portion of the galaxy (Fan et al. 2008), though the potential of the dark matter halo may well prevent it from being completely lost from the galaxy. The feedback of only a small fraction of the QSO’s total peak power in winds and jets is sufficient to yield such strong outflows (e.g., Granato et al. 2004; Hopkins et al. 2005). Recall that powerful outflows, if produced by nearly-Eddington or even super-Eddington accretion onto black holes have been argued to be fundamentally responsible for the coordinated growth of galactic bulges and the SMBH mass and for heating of the gas in clusters (e.g., Silk & Rees 1998; King 2003; 2009).

When the mass loss is rapid compared to the dynamical time for the stellar system then it probably causes galaxy expansion over the course of several dozen dynamical times, an effect that has been long studied for stellar clusters (e.g., Hills 1980). Fan et al. (2008) give an approximate calculation of the galaxy expansion for two limiting cases. When the gas ejection timescale, $\tau_{\text{ej}} \ll \tau_{\text{dyn}}$, the dynamical timescale, then the initial and final energies are respectively given by $E = -\beta GM^2/2R$ and $E' = [-\beta GM'^2/R + \beta M'(GM/2R)]$, where β depends on the stellar density profile whose shape is assumed to be the same in the initial and final configurations, because neither the dispersion nor the radius changes during the rapid gas expulsion.

The ratio of the new to original energies is $E'/E = (M'/M)^2[2 - M/M']$ (Biermann & Shapiro 1979). Therefore, if $M/M' > 2$, the system is completely unbound but that is not likely to ever occur for massive galaxies. When $M/M' < 2$ then the system will relax to a new equilibrium system where

$$\alpha(M, M') \equiv R'/R = (2 - M/M')^{-1}, \quad (1)$$

if the new equilibrium system is homologous to the original one. Numerical simulations show that when the system is not disrupted the new configuration does indeed expand by roughly this factor and that it takes about 30–40 initial dynamical times for a new equilibrium to be established (e.g., Goodwin & Bastian 2006). In the other limit, where $\tau_{\text{ej}} \gg \tau_{\text{dyn}}$, the expansion is through the adiabatic orbits of stars, $E = -(\bar{\Phi}\dot{M} + M\Delta\bar{\Phi}/\tau_{\text{dyn}})/2 = 0$, which implies $\Delta\bar{\Phi}/\bar{\Phi} = -\tau_{\text{dyn}}/\tau_{\text{ej}} \approx 0$, where $\bar{\Phi}$ is the mass averaged potential and hence the expansion factor is proportional to the mass loss rate (e.g., Hills 1980) so, $R'/R = M/M'$.

Taking a reasonable value of $\tau_{\text{dyn}} \approx 5 \times 10^7 \text{ yr}$ in the core of a massive elliptical the timescale for relaxation should be ~ 2 Gyr so

that by $z \sim 0.8$ the compact core galaxies seen at $z \sim 2$ would have settled onto the fundamental plane (Fan et al. 2008). A reasonable model would adopt a Sersic profile for the projected stellar density, specialized to $n = 4$, corresponding to a $r^{1/4}$ law (but cf. Hopkins et al. 2009), and a Navarro et al. (1997) profile for the total mass. With these prescriptions and, for the sake of specificity, adopting the evolution model of Granato et al. (2004), the effective radius, r_e , can be found in terms of the virial radius, r_{vir} , the ratio of the stellar dispersion to the dark matter dispersion, f_σ , and the masses of the stars M_* , cold gas, M_{cold} , and the galaxy halo, M_H , (Fan et al. 2008),

$$r_e \approx \frac{0.34}{f_\sigma^2} \frac{M_* + M_{\text{cold}}}{M_H} r_{\text{vir}}. \quad (2)$$

They further argue that reasonable star formation and quasar activity inputs will produce a ratio $M_{\text{cold}}/M_* \sim 2/3$ for massive galaxies nested in halos with $M_H > 10^{12} M_\odot$. If all of this gas is expelled in the fast ejection limit, this would lead to an increase in galaxy size by a factor of ~ 3 from Eqn. (1), since $M = M_* + M_{\text{cold}}$. For less massive galaxies, with $M_H < 10^{12} M_\odot$, the current stellar mass is $< 2 \times 10^{10} M_\odot$; then the nuclear activity is probably too weak to eject the gas until a much later epoch, by which time $M_{\text{cold}} \ll M_*$ and then M_{ej}/M_* where M_{ej} is the mass of the ejected gas, becomes small, too.

Now we improve upon the estimate of Fan et al. (2008) for $\alpha(M, M')$. To find the factor M/M' we note that initially the mass would include the mass in the cold gas M_{cold} ; after condensation and star formation, part of that mass gets added to M_* . Then the ratio of the initial to the final mass, after expulsion of the gas at time t_x , would be given by

$$\frac{M}{M'} = \frac{M_{\text{cold}}(t_x) + M_*(t_x)}{M_*(t_x)}. \quad (3)$$

The time t_x begins at the epoch when virialization of the halo is essentially complete. To make further progress, we invoke the semi-analytic model of Mao et al. (2007) using the framework of Granato et al. (2004). We have calculated this mass ratio by integrating the Eqns A1–A6 of Mao et al. (2007), while making the assumption that the expulsion event (which could be either AGN or supernovae driven, both of which are accommodated in the formulation) results in the entire remaining cold gas exiting from the galaxy. With Δt_{burst} the length of star formation, the most useful independent variable is $\tau \equiv t_x/\Delta t_{\text{burst}}$. Then the mass ratio is found to be

$$\frac{M}{M'} = 1 + \frac{1 - \exp[-(s\gamma - 1)\tau\delta]}{(s - 1/\gamma) \exp(\tau\delta) - s - \exp[(s\gamma - 1)\tau\delta]/\gamma}, \quad (4)$$

where

$$\delta(M_{12}) \equiv \frac{\Delta t_{\text{burst}}}{t_{\text{cond}}} = \frac{5}{8} \mathcal{F}(M_{12}) M_{12}^{-0.2}, \quad (5)$$

with M_{12} the total galactic mass in units of $10^{12} M_\odot$, t_{cond} the gas condensation and star formation time, and $\mathcal{F}(x) = 1$ for $x \geq 1$ and $\mathcal{F}(x) = x^{-1}$ for $x \leq 1$. The other quantities in the expression for the mass ratio are $\gamma = 0.7 + 0.6 M_{12}^{-2/3} [4/(1+z)]$ and $s \equiv t_{\text{cond}}/t_* \approx 5$ for isothermal spheres, with t_* defined by the rate of star formation, $\dot{M}_* \approx \frac{M_{\text{cold}}}{t_*}$. So M/M' is primarily a function of redshift z , the total mass of the galaxy M_{12} , and τ , which is the time of the gas expulsion expressed in units of the starburst duration, which is taken to begin at the time the halo became virialised (and gas condensation and star formation begin). Plausible values have been adopted for the other star formation and condensation related parameters (see Mao et al. 2007).

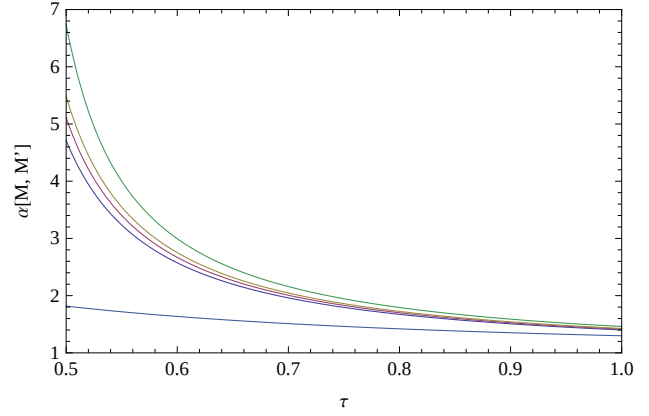


Figure 1. Values of α , the expansion factor, as functions of the time of the cold gas expulsion (measured from the start of star formation in units of the starburst duration) for the case of rapid expulsion ($\tau_{\text{ej}} \ll \tau_{\text{dyn}}$). The four upper curves are for $z = 5, 3, 2.5, 2$, starting from the top. The lowest curve is for the case of adiabatic slow mass loss ($\tau_{\text{ej}} \gg \tau_{\text{dyn}}$, for the case of $z = 3$) where τ_{ej} is the gas exit time scale. For all curves a halo mass of $M_{12} = 1$ is assumed.

We would like to point out that Eqn (4), without adding 1, is an improvement over Eqn. (5) of Fan et al. (2008) which is an approximation. The deviation is significant for low τ , where additional factors in the numerator and denominator of the expression become important. As a result, the expansion factor seems to have been underestimated in the model of Mao et al. (2007). In Fig. 1, the expansion factors for the case of rapid mass ejection, $\alpha(\tau)$ are shown for a range of the halo virialization redshift, $z = 5, 3, 2.5, 2$; one case of slow (adiabatic) mass loss is also shown in Fig. 1. It is seen that for a typical value of $\tau = 0.5 - 0.6$, and an initial gas-to-stellar mass ratio of $2/3$, an expansion factor of about 3–4 is expected for the rapid expulsion case and less than half as much for the adiabatic case.

We see that while $\alpha > 4$ is possible it requires rather high M_{cold} to M_* ratios or very small values of τ . Values of $\tau < 0.4$ imply ejection of much of the gas before most of the stars form; however, those circumstances are unlikely to yield a massive elliptical galaxy. So while this process certainly can puff up the galaxies substantially, it requires a large, yet not too large (otherwise the galaxy is disrupted), quantity of cold gas to be expelled very rapidly, underscoring the “fine tuning” problem also noted by Bezanson et al. (2009).

Should the expansion factor be indeed modest with $\alpha \leq 1.5$ (Hopkins et al. 2009) there are two alternative possibilities for the cold gas depletion. First, as seen from Fig. 1, if the explosive event is delayed by about 0.1 in τ , which is in units of $\Delta t_{\text{burst}} = 6 \times 10^8 \left[\frac{4}{1+z} \right]^{1.5} \mathcal{F}(M_{12})$ years, while the starburst has been in progress, this would lead to a lower α . There could be a variety of reasons for this: either the QSO switches on late or massive stars are formed slightly later depending on the details of the clumpiness of the gas distribution (we have taken it to be isothermal here). A second possibility envisaged is a steady adiabatic mass loss which results in the lower curve of Fig. 1. Either of these two situations would be consistent with the negative results of Hopkins et al. (2009).

3 COSMIC EVOLUTION OF THE ACCRETION MODE

While the mechanical feedback from the AGN might be responsible for the expulsion of the bulk of cold/warm gaseous phase of massive ellipticals during the quasar era (Sect. 2), the existing hot gas is likely to be largely retained and even replenished by continuing stellar mass loss. Thus, independent of the role of cold gas expulsion in causing any expansion of the galaxy, a Bondi type accretion of this hot ISM phase onto the central SMBH could become the favored mechanism for powering the low-redshift/luminosity RGs (mostly of edge-darkened or FR I RGs), particularly because any emission lines seen in their optical spectra are usually of the low-excitation type (e.g., Hardcastle et al. 2007; Balmaverde, Baldi & Capetti 2008; also, Baum et al. 1995).

In contrast, powerful RGs (FR IIs) and RQQs found at high redshifts are believed to be powered by a sustained thin disc accretion of cold gas, which is both capable of spinning up the SMBH (e.g., Rees & Volonteri 2007; Berti & Volonteri 2008) and also producing intense disc emission, in addition to ejecting powerful relativistic jets whenever conditions near the central engine become propitious for jet collimation (e.g., Sikora, Stawarz & Lasota 2007; Livio, Pringle & King 2003; Meier 2001). In contrast, the hot accretion (dominant at low z) should funnel only little angular momentum into the central BH; for such RGs the ADAF accretion mode (e.g., Narayan & Yi 1995) can explain their observed faint disc emission with low-excitation optical spectra. Evolutionary scenarios along these lines are in vogue (e.g., Somerville et al. 2008 and references therein), particularly since such central engines can explain the radiatively inefficient “low-excitation RGs” which characterize the bulk of the FR I population and some less powerful FR II sources as well (e.g., Hardcastle et al. 2007; Best et al. 2006; also Allen 2006). Only rare instances involving ‘wet’ mergers at low z would push massive ellipticals into the high-excitation AGN mode attributable to cold gas accretion. Evidence now exists for the hot gas accretion to be the dominant process for FR I RGs across three decades in radio power (e.g., Balmaverde, Baldi & Capetti 2008); that paper reports a nearly linear scaling between the accretion power and the jet power so that $\simeq 1\%$ of the rest mass energy goes into jet power.

Thus, it appears that the massive elliptical galaxies that had hosted powerful AGN (both radio-loud and radio-quiet) during the quasar era will, in most cases, be able to produce only their low-luminosity versions at small redshifts. We argue that observational support for massive ellipticals undergoing transformation from hot to cold accretion mode emerges from the fact that the co-moving density of powerful RGs, mostly the edge-brightened FR II’s, peaks at $\phi \simeq 10^{-6.2} \text{ Mpc}^{-3}$ for $P_{151\text{MHz}} \simeq 10^{26.5} \text{ W Hz}^{-1} \text{ sr}^{-1}$ between $2 < z < 3$, while the weaker radio sources found at $z \simeq 0$ (essentially all FR I’s) have a comparable density of $\phi \simeq 10^{-5.8} \text{ Mpc}^{-3}$ for $P_{151\text{MHz}} \simeq 10^{24} \text{ W Hz}^{-1} \text{ sr}^{-1}$ (Willott et al. 2001; converted to Λ CDM cosmology by Grimes et al. 2004). Note that this radio luminosity is typical of FR I RGs, being nearly an order of magnitude below the value for the most powerful FR I sources (Best et al. 2005). So, the above scheme for accretion mode change would posit that a substantial fraction of the ellipticals hosting the FR I RGs at low z are probably descendants of the galaxies that had hosted FR II RGs at $z > 2$, as also hinted by Best et al. (2005). This scenario also allows us to understand why the decline in the abundance with cosmic time is so much less steep for low power radio RGs than for powerful RGs.

It is further expected that hot accretion, and thus most FR I RGs, would preferentially be associated with more massive galax-

ies at low z . This is because the Bondi accretion rate scales approximately as $M_{\text{BH}}^{1.5}$, (e.g., Best et al. 2007). Clearly, hot accretion would be favoured by deeper potential wells and this fits with the well known result that FR I RGs are usually associated with the larger cluster galaxies (e.g., Longair & Seldner 1979; Best 2004; Hardcastle et al. 2007) which must also be also richer in hot ISM (e.g., Mathews & Brighenti 2003). In the next section we discuss some important ramifications of the jet production that arise from the switchover to the hot accretion mode at later cosmic epoch.

4 JET PRODUCTION

We adopt the flexible paradigm that the jets are either powered totally by black hole spin, essentially by the Blandford-Znajek (1977; BZ) process or also partially from a magnetized accretion disk, as in the hybrid model of Meier (1999; 2001). Note that even though jets may also form in RQQs, they are liable to be quenched at a nascent stage (e.g., Gopal-Krishna, Mangalam & Wiita 2008), and hence we would consider only the accretion disk output in the case of RQQs. The jet powers in these two cases can be written as (McDonald & Thorne 1982; Meier 1999)

$$\mathcal{L}_{\text{jet}} = \begin{cases} \mathcal{L}_{\text{BZ}} = \frac{1}{32} \omega_F^2 B_{\perp}^2 R_H^2 j^2 c & \text{BZ process} \\ \mathcal{L}_{\text{H}} = \frac{1}{32c} (B_{\phi} R_{ms} \Omega)^2 & \text{Hybrid process} \end{cases} \quad (6)$$

where in the first expression, $\omega_F \equiv \Omega_F (\Omega_H - \Omega_F) / \Omega_H^2$ depends on the angular velocity of the field, Ω_F , relative to that of the hole, Ω_H , $B_{\perp}$ is the component of the magnetic field perpendicular to the hole, $R_H = r(j)m \equiv (1 + (1 - j^2)^{1/2}) GM_{\bullet}/c^2$ is the radius of the event horizon of the BH, where $m \equiv GM_{\bullet}/c^2$, and $j = a/m$ is the dimensionless angular momentum of the BH; for the second expression, B_{ϕ} is the azimuthal component of the magnetic field in the inner portion of the disk, R_{ms} is the radius of the marginally stable orbit, the height of the disk is taken to be $H \simeq R_{ms}$, and Ω is the angular velocity of the disk at the marginally stable orbit.

The BZ model is technically derived in the limit of zero accretion but the hybrid models, on the other hand, do require at least a modest level of accretion, most likely in the ADAF mode (e.g., Narayan & Yi 1995). As mentioned above, an ADAF type of accretion disk close to the SMBH, fueled by essentially Bondi accretion at larger scales, is now the common premise for the weaker “radio mode” for AGN, where jets, but little optical emission, are produced (e.g., Somerville et al. 2008).

Although some general relativistic magnetohydrodynamical (GRMHD) simulations indicate that only an essentially BZ type of process is likely to yield a relativistic jet while hybrid type models do not work well (McKinney 2005), other GRMHD simulations do seem to show the flows develop essentially along the line of the hybrid models (e.g., Koide et al. 2002; Nishikawa et al. 2005); it is fair to say that no consensus has yet emerged on this point. Even the BZ process for a massive, rapidly spinning BH does not guarantee that a powerful jet will emerge, since it has been argued that successful formation of relativistic jets requires additional collimation by MHD outflows from accretion disks (Sikora et al. 2007). When the accretion rate is close enough to the Eddington rate for a “standard” (optically thick but geometrically thin) accretion disc to form (e.g., Shakura & Sunyaev 1973), the optical and UV luminosity will be high, but powerful jets will only rarely emerge.

A potentially very interesting outcome of the cold gas expulsion for both the BZ and hybrid models would arise if flux conservation were to be assumed for all the gas that feeds the central

engine so that $B \propto R^{-2}$. Then the magnetic field contribution to the jet power, which is $\propto B^2$ in either case, is enhanced by a factor of about $\alpha^4 \sim 100$. However, it is more likely that the gas feeding the disk, at least initially, is practically within the sole gravitational purview of the SMBH and therefore would not partake of the overall gas expulsion. This idea is supported by the observations indicating much less stellar density evolution in the inner ~ 1 kpc of massive ellipticals than in the outer portion of their stellar populations (Bezanson et al. 2009). In such an event any effect due to magnetic field changes is expected to be much more modest until the disk matter is fully accreted by the BH and its spin-down sets in.

4.1 Spin down time scale

The formulation of the BZ model in Macdonald & Thorne (1982) leads to the following forms for the jet power, $\mathcal{L}$, and torque, $\mathcal{G}$ (details are given in Appendix A),

$$\mathcal{L} = \mathcal{L}_0 j^2 r^2(j) \text{ where } \mathcal{L}_0 = \frac{m^2 c}{32} B_\perp^2 g, \quad (7)$$

$$\mathcal{G} = \mathcal{G}_0 j r^3(j) \text{ where } \mathcal{G}_0 = \frac{m^3}{8} B_\perp^2 f, \quad (8)$$

where we have dropped the subscript BZ. The geometric factors, g and f , are the results of angle averaging over the horizon of the magnetic flux and the spin of the field and are model dependent; expressions for them are given in Eqns (A6) and (A7). When the maximum power is transferred from the BH to the jet, as is typically assumed, these factors are of order unity.

The angular momentum budget and the rotation energy budget are given by, respectively

$$\mathcal{J} = \mathcal{J}_0 j \text{ where } \mathcal{J}_0 = c M_\bullet m \text{ and} \quad (9)$$

$$\mathcal{E} = \mathcal{E}_0 \left(1 - \frac{r^{1/2}(j)}{\sqrt{2}} \right) \text{ where } \mathcal{E}_0 = M_\bullet c^2. \quad (10)$$

For reference we give the numerical values of the various quantities used in cgs units:

$$\mathcal{J}_0 = c M_\bullet m = 9 \times 10^{64} M_8^2 \text{ (g cm}^2/\text{s);} \quad (11)$$

$$\mathcal{L}_0 = \frac{m^2 c}{32} B_\perp^2 g = 2 \times 10^{43} g B_4^2 M_8^2 \text{ (erg/s);} \quad (12)$$

$$\mathcal{G}_0 = \frac{m^3}{8} B_\perp^2 f = 4 \times 10^{46} f B_4^2 M_8^3 \text{ (erg).} \quad (13)$$

This lets us compute a spin-down time from angular momentum conservation

$$\tau_{j,BZ} = \frac{\mathcal{J}_0}{\mathcal{G}_0} \int_{j_f}^{j_i} \frac{dj}{r^3(j)} = 7.0 \times 10^8 \text{ yrs } \frac{[(\kappa(j_i, j_f)/0.1)]}{B_4^2 M_9 f} \quad (14)$$

where, j_i and j_f are the initial and final spins, respectively, $\kappa(j_i, j_f)$ is the value of the integral, $B_4 = B/10^4$ Gauss and $M_9 = M_\bullet/10^9 M_\odot$. The spin down time can be found by calculating the time for the rotational energy to reduce by a factor ϵ ; see Fig. 2. A detailed evaluation of this factor is given in Appendix B. The time corresponding to $\epsilon = 1/e$ for $j_i = 0.5$ is calculated to be 0.5 Gyr for $M_9 = 1, B_4 = 1$.

Another estimate of the spin down time that is more relevant observationally is the e-folding time scale of the jet power. The j_f at a time when a reduction by a factor q is reached can be calculated from

$$q j_i^2 r^2(j_i) = j_f^2 r^2(j_f). \quad (15)$$

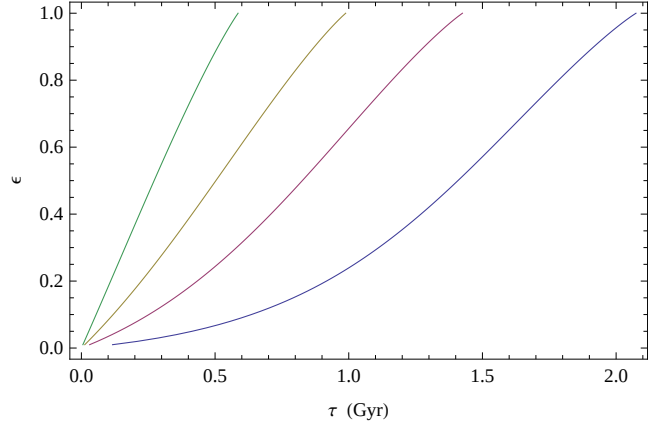


Figure 2. The fraction of BH rotational energy, ϵ , lost by the BZ process as a function of time is shown for various initial values of the BH spin parameter $j = 0.2, 0.4, 0.6, 0.8$ from right to left. The values assumed for other parameters are $B_4 = 1, M_9 = 1$.

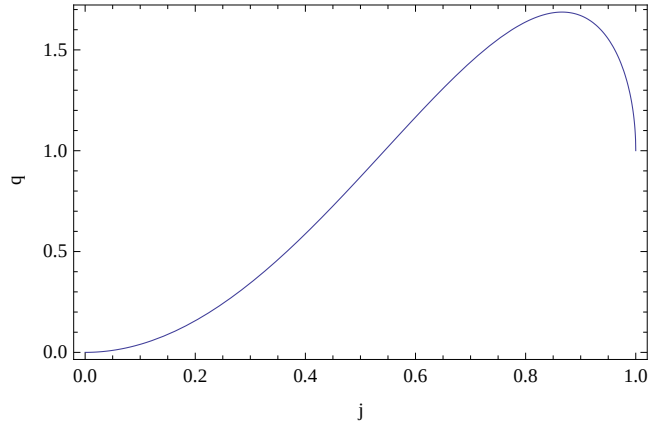


Figure 3. The power of the BZ jet as a function of the BH spin parameter j , normalized to the value at $j = 1.0$. The power goes through a maximum at $j = \sqrt{3}/2$ before dropping; this arises from the competition between the horizon radius and the spin of the BH.

The above equation leads to a quartic in j_f ,

$$j_f^4 - 2q j_f + q^2 = 0, \quad (16)$$

whose roots are expressible analytically. The positive real root between 0 and j_i is found to be $j_f(q, j_i)$, which is fed into Eqn (B3) to calculate the time in terms of τ_j from Eqn (14). The $\mathcal{L}(j)$ goes through a maximum at $j = \sqrt{3}/2$ (Fig. 3). The evolution of the power as a function of time is given in Fig. 4. The time corresponding to $q = 1/e, j_i = 0.5$ is calculated to be 0.5 Gyr for $M_9 = 1, B_4 = 1$. The spin down time scale is likely to be reduced further when the mass of the hole increases by Bondi accretion (with negligible net angular momentum) as it spins down. The effective timescale is estimated to be $\tau_{spin} \approx \frac{(\tau_a/2)\tau_j}{(\tau_j + \tau_a/2)}$, where

the accretion time scale is $\tau_a \equiv \frac{M_\bullet c^2}{L_E} = 0.45$ Gyr, where L_E is the Eddington luminosity. As a result, $\tau_{spin} \approx 0.2$ Gyr for the typical case $j_i = 0.5, M_9 = 1, B_4 = 1$. Note that the timescales derived above are inversely proportional to the BH mass (Eqn 14).

The energy loss time-scale is quite similar for the hybrid model, although the dependence on BH mass is much more grad-

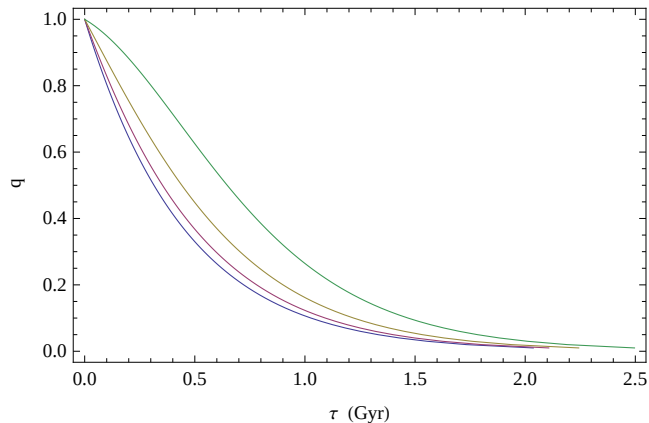


Figure 4. The power reduction factor, q , of the BZ jet as a function of time is shown for various initial values of the BH spin parameter $j = 0.2, 0.4, 0.6, 0.8$ from bottom to top. The values assumed for other parameters are $B_4 = 1$, $M_9 = 1$.

ual. Using the rotational energy from Eqn (10) above in the limit of low spin, and Eqn (12) for the power in the MHD jet given by Meier (1999), we have the time scale

$$\tau_{L, \text{hybrid}} = 5 \times 10^8 \text{ yr } \alpha_{-2}^{-1/10} M_9^{-1/10} \dot{m}_{-3}^{-4/5} \zeta_{-1}^{-1}, \quad (17)$$

where α is the disk viscosity parameter, $\dot{m}$ is the accretion rate in units of the rate that would produce an Eddington luminosity, and ζ is a duty cycle parameter expected to be around 10^{-1} (Meier 1999). All quantities are normalized to the power of ten that is indicated by the subscript and are chosen to be typical for an ADAF type of flow. Thus we find that in either case, for a BH of mass about $\sim 10^9 M_\odot$ the time-scale for spin-down is a few times 10^8 yr . Note that in the hybrid model the spin-down time scale is extremely insensitive to $M_\bullet$ (Eqn 17), while it scales as $M_\bullet^{-1}$ for the BZ model (Eqn 14) where the angular momentum is directly extracted from the hole as opposed to the hybrid case, where it is partly extracted from the accretion disk. An important caveat to note is that the spin-down time scales derived above (eg. Eqn (14)) depend on the assumption that the magnetic field strength is constant during the spin-down.

5 DISCUSSION AND CONCLUSIONS

In this paper we have focused attention on the transition period connecting the quasar era ($z = 2 - 3$) to the present era, between which the co-moving space density of powerful radio galaxies and QSOs has declined by almost three orders of magnitude. As described in the foregoing sections, this “transitional” era is probably marked by a changeover in the dominant mode of AGN activity from one fuelled by accretion of cold gas onto the SMBH to another which is dominated by Bondi accretion of hot gas from the interstellar medium of the galaxy (Sect. 3, Somerville et al. 2008 and references therein). Only the cold accretion mode can support a strong disc emission and a high spin of the central supermassive black hole (SMBH) which can therefore eject powerful jets under suitable conditions. An important consequence of this paradigm of accretion mode change is that the massive elliptical galaxies that were able to host powerful AGNs at high redshifts would usually be no longer able to do so at recent cosmic epochs (Sect. 3).

The postulated basic operational flip of the central engines in massive ellipticals since the quasar era underscores the need to

take a closer look at the physical processes governing the transitional era. Clearly, the two dominant processes underlying this transformation are the elimination of vast reservoirs of molecular gas present in the massive ellipticals at high- z and the consequent spinning down of their SMBH. In this paper we have addressed both these points quantitatively by presenting an improved theoretical modelling of the physical mechanisms involved.

Based on the available evidence we have argued in favour of the AGN wind or jet driven expulsion as being the principal mechanism by which massive ellipticals have lost their vast reservoirs of molecular gas since the quasar era. Recently, this issue of cold gas ejection has attracted a great deal of attention, following the claims of an observed $\sim 3 - 4$ fold increase in the sizes of massive ellipticals since the quasar era, along with a very recent counter-claim (Sect. 1).

We have obtained a detailed, improved expression for the expansion factor, Eqn (4), which allows us to distinguish between the expulsion histories of galaxies in terms of the star formation parameters. The distinction is expressed in our model as the expansion factor would be lower if more stars form in the core before the gas expulsion process begins. Alternatively, a steady adiabatic mass loss leads to a modest expansion factor. In the former case, the core would be more luminous as more star formation occurs and less gas is expelled. More detailed modelling along these lines can explore the space of the star formation parameters, the timing of gas expulsion, and the cases between adiabatic and rapid mass loss, in order to make better connections with observations.

Irrespective of the potential role in causing the reported dramatic size evolution of massive elliptical galaxies, the cold gas expulsion will only have a delayed effect on the output of the central engine (accretion disk plus the SMBH). This is because the cold gas associated with the nuclear gas cloud that is gravitationally bound to the SMBH is expected to be only weakly subject to the AGN induced cold gas expulsion. This resulting residual AGN activity from the cold accretion era would largely define the transitional era and be itself determined by the duration of the spin-down phase which the SMBH would trace once the nuclear molecular cloud is fully accreted and, consequently, a dramatic fading of the central engine has set in.

We have focused largely on the Blandford-Znajek mechanism of extracting rotational energy from the BH in the absence of accretion torque. We have derived two time scales using energy loss and power reduction as criteria for determining the spin down, which turns out to be about $0.5/(B_4^2 M_9)$ Gyr. If Bondi accretion proceeds, it will add to the mass of the BH but little to its angular momentum and thus reduce the spin-down time scale to about $0.2/(B_4^2 M_9)$ Gyr. The evolution of the jet power indicates an increase before a gradual decline if the initial spin, $j > \sqrt{3}/2$, as a result of the hole’s increasing size. This naturally has implications for the evolution of the jet. We plan to expand our work to hybrid models in greater detail, and to thus explore disc accretion models that explicitly involve angular momentum transport from the hole to the disc.

The transition from cold to hot accretion dominated phase in the cosmic evolutionary history of the AGN population is marked by a period when the SMBH would continue to be fuelled (and spun up) by the accretion of the cold gas located within the sphere of influence of the SMBH. This gas is likely to survive the otherwise efficient cold gas expulsion due to the intense AGN activity during the quasar era. But, even if such nuclear gas cloud has a radius as large as 1 pc, it would sustain cold disc accretion phase for no more than 10^9 years (e.g., King & Pringle 2007). Moreover,

the spin-down of the central engine would also occur in $\sim 10^9$ years (also, Meier 2001) for reasonable values of mass accretion rate and the duty cycle parameter (Eqn 17). Thereafter, i.e., during the past 6-7 Gyr, the occurrences of cold accretion dominated AGN would become rare, mostly sustained by occasional ‘wet’ mergers of captured galaxies. Thus, the AGN activity over the past several gigayears would be increasingly marked by Bondi accretion of hot gas powering the central engines and thereby producing mostly FR I RGs (Best et al. 2006; Hardcastle et al. 2007).

To summarize, we have envisaged a scheme characterized by the following sequence of events related to the evolution of massive elliptical galaxies. At $z \simeq 3$ they contain an abundant supply of cold gas that yields both prolific star formation and luminous thin disk accretion. The latter usually leads to a fast spinning BH. Once a good fraction of the cold gas has been accreted, the resulting fast spinning BHs would become capable of ejecting powerful jets and forming the luminous radio sources that were so much more abundant during the quasar era. But these powerful jets and disc winds can easily expel most of the cold gas reservoir, possibly resulting in a substantial expansion of the host galaxy. A likely outcome of the cold gas expulsion is the eventual dramatic weakening of the central engine itself through a spinning down of the SMBH, usually on a timescale less than ~ 1 Gyr.

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APPENDIX A: FORMULAE FOR POWER AND TORQUE

We derive the Eqns (7–13) given in Sect. 4.1. We start from Eqn (7.19) in Macdonald & Thorne (1982) for the luminosity through a ring at latitude θ

$$\Delta\mathcal{L} = \frac{\Omega_F(\Omega_H - \Omega_F)}{4\pi c} \varpi^2 B_n \Delta\psi, \quad (\text{A1})$$

where $\varpi(\theta) = R_H(j) \sin \theta$ is the ring radius, $\Omega_H(j)$ is the angular velocity of the hole, Ω_F is the angular velocity of the field, $B_n = \mathbf{B} \cdot \mathbf{n}$ is the field perpendicular to the hole surface and $\Delta\psi(\theta)$ is the flux through a ring at θ of arc length $R_H(j)\Delta\theta$, which is given by

$$\Delta\psi(\theta) = 2\pi\varpi(\theta)R_H(j)B_n\Delta\theta. \quad (\text{A2})$$

The torque over the same ring is given by

$$\Delta\mathcal{G} = \frac{\Delta\mathcal{L}}{\Omega_H - \Omega_F}. \quad (\text{A3})$$

Now when the above equations for the power and torque are integrated over θ we obtain

$$\mathcal{L} = \frac{B_\perp^2 R_H(j)^4 \Omega_H(j)^2}{8c} g \quad \text{and} \quad (\text{A4})$$

$$\mathcal{G} = \frac{B_\perp^2 R_H(j)^4 \Omega_H(j)}{4c} f, \quad (\text{A5})$$

where the geometric factors f and g are found to be

$$f = \int_0^\pi \left(\frac{(\mathbf{B} \cdot \mathbf{n})^2}{B_\perp^2} \right) \left(\frac{2\Omega_F}{\Omega_H} \right) \sin^3 \theta \, d\theta \quad (\text{A6})$$

$$g = \int_0^\pi \left(\frac{(\mathbf{B} \cdot \mathbf{n})^2}{B_\perp^2} \right) \left(\frac{4\Omega_F}{\Omega_H} \right) \left(1 - \frac{\Omega_F}{\Omega_H} \right) \sin^3 \theta \, d\theta, \quad (\text{A7})$$

where $\mathbf{B}$ and Ω_F could be functions of θ and $B_\perp^2 \equiv \int_0^\pi (\mathbf{B} \cdot \mathbf{n})^2 \sin^3 \theta \, d\theta$ is the angle averaged root mean squared value of the normal component of the field on the surface of the hole. These factors are of order unity; if the maximum power is transferred, then $\Omega_F = \Omega_H/2$, and both of these integrals are unity. Then Eqns (7, 8) follow after substituting for $R_H(j)$ and $\Omega_H(j)$, which are respectively the black hole’s radius and angular velocity and are given by

$$R_H(j) = r(j)m = (1 + (1 - j^2)^{1/2}) GM_\bullet / c^2, \quad (\text{A8})$$

$$\Omega_H(j) = \left(\frac{j}{2} \right) \frac{c}{R_H(j)}. \quad (\text{A9})$$

Eqns (9) and (10) then follow as the fundamental equations for the black hole angular momentum and rotational energy in terms of the irreducible mass of the hole. The coefficients for these physical quantities and the luminosity are easily found to be Eqns (11–13).

APPENDIX B: CALCULATION OF THE SPIN DOWN FACTOR

We evaluate the energy radiated when the spin reduces from $j_i \rightarrow j_f$ to be

$$\Delta\mathcal{E} = \int \mathcal{L} dt = \mathcal{L}_0 \int r^2(j) j^2 \left[\frac{dj}{dt} \right]^{-1} dj \quad (\text{B1})$$

As a fraction of the rotational energy budget, this is found to be

$$\epsilon(j_i, j_f) \equiv \frac{\Delta\mathcal{E}}{\mathcal{E}} = \left[\int_{j_i}^0 \frac{j}{r(j)} dj \right]^{-1} \int_{j_i}^{j_f} \frac{j}{r(j)} dj \quad (\text{B2})$$

The integral κ defined in Eqn (14) is calculated to be

$$\kappa(j_i, j_f) = \left[\left(\frac{1}{16} \right) \ln \left(\frac{2-w}{w} \right) + \left(\frac{3w^2 + 3w - 4}{24w^3} \right) \right]_{w_f}^{w_i} \quad (\text{B3})$$

where $w_i = r(j_i)$ and $w_f = r(j_f)$. Using similar techniques the integral

$$\chi(j_i, j_f) \equiv \int_{j_i}^{j_f} \frac{j}{r(j)} dj = [\ln(w) - w]_{w_i}^{w_f}. \quad (\text{B4})$$

Now from Eqn (B2), the value of w_f for $\epsilon = 1/e$ and $j_i = 0.5$ is solved for and fed into Eqn (B3). Then using Eqn (14) the time to spin down is found to be $\simeq 0.5$ Gyr.

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