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Star-forming accretion flows and the low-luminosity nuclei of giant elliptical galaxies

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ABSTRACT

The luminosities of the centres of nearby elliptical galaxies are very low compared to models of thin disc accretion on to their black holes at the Bondi rate, typically a few hundredths to a few tenths of a solar mass per year. This has motivated models of inefficiently radiated accretion that invoke weak electron–ion thermal coupling, and/or inhibited accretion rates due to convection or outflows. Here we point out that, even if such processes are operating, a significant fraction of the accreting gas is prevented from reaching the central black hole because it condenses into stars in a gravitationally unstable disc. Star formation occurs inside the Bondi radius (typically $\sim$100 pc in giant ellipticals), but still relatively far from the black hole in terms of Schwarzschild radii. Star formation depletes and heats the gas disc, eventually leading to a marginally stable, but much reduced, accretion flow to the black hole. We predict the presence of cold ($\sim$100 K), dusty gas discs, containing clustered Hα emission and occasional Type II supernovae, both resulting from the presence of massive stars. Star formation accounts for several features of the M87 system: a thin disc, traced by Hα emission, is observed on scales of about 100 pc, with features reminiscent of spiral arms and dust lanes; the star formation rate inferred from the intensity of Hα emission is consistent with the Bondi accretion rate of the system. Star formation may therefore help to suppress accretion on to the central engines of massive ellipticals. We also discuss some implications for the fuelling of the Galactic Centre and quasars.

Key words: accretion, accretion discs – black hole physics – stars: formation – galaxies: active – galaxies: elliptical and lenticular, cD – galaxies: individual: M87.

1 INTRODUCTION

Optical spectroscopy and photometry provide good evidence that black holes with masses of 10^8–10^{10} $M_{\odot}$ reside at the centres of giant elliptical galaxies (e.g. Kormendy & Richstone 1995). X-ray observations of these galaxies allow an inference of the interstellar gas temperature and thus an estimate of the Bondi accretion radius, where the dynamics of the gas starts to be dominated by the potential of the black hole. The density of the gas is also derived from these observations, so the Bondi accretion rate can be estimated. The high resolution of the *Chandra X-ray Observatory* has proved to be particularly useful for such studies.

The bolometric luminosities of the black hole engines can be estimated from multiwavelength observations, particularly in the infrared. These luminosities are typically several orders of magnitude fainter than predicted by models of thin disc accretion at the Bondi rate in which about 10 per cent of the rest-mass energy is radiated (e.g. Fabian & Rees 1995). A particularly well-studied case is M87 (3C 274, NGC 4486; Di Matteo et al. 2000, 2003; see also Reynolds et al. 1996a), but similar results have also been found for NGC 1399, 4472 and 4636 (M49) (Di Matteo et al. 2000; Loewenstein et al. 2001), for NGC 4697 (Di Matteo et al. 2000; Sarazin, Irwin & Bregman 2001), for NGC 4697 and for NGC 4649 (M60) (Di Matteo et al. 2000; Randall, Sarazin & Irwin 2004).

In order to explain such quiescent accretors, two-temperature collisionless accretion models have been invoked (e.g. Ichimaru 1977; Narayan & Yi 1995; Quataert & Gruzinov 1999; Narayan 2002). These models typically assume that the energy exchange between electrons and ions proceeds by Coulomb collisions. The associated accretion flows do not radiate efficiently if (1) the dissipation primarily heats the protons and (2) the electron–ion equilibration time is longer than the radial infall time. The gravitational energy is then carried by ions to the event horizon of the black

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hole, in what are known as advection-dominated accretion flows (ADAFs).

For the ellipticals, minimalistic standard ADAF models predict too high a luminosity (e.g. Di Matteo et al. 2000). Variations of ADAFs are, however, still in the running: ADAFs can be combined or replaced with models involving a radially dependent accretion rate and winds (Blandford & Begelman 1999; Quataert & Narayan 2000). Alternatively, ADAFs unstable to convection – called convection-dominated accretion flows (CDAFs; Narayan, Igumenshchev & Abramowicz 2000; Quataert & Gruzinov 2000) – have the outward transport of angular momentum inhibited by the convection and thus their accretion rate on to the central engine reduced.

There are, however, some unresolved issues with the physics of these models. Regarding CDAFs, the parameter space allowed for solutions that reduce the outward transport of angular momentum becomes restricted when magnetic fields are present and seems to require that the field saturates at a very low value. Otherwise the magnetorotational instability dominates the outward angular momentum transport (Balbus & Hawley 2002; Narayan et al. 2002). It should also be noted that virtually all flavours of optically thin ADAFs and CDAFs require maintenance of two-temperature collisionless flows. The unanswered plasma physics questions about the viability of this situation have also been an important area of ongoing research (Begelman & Chiueh 1988; Bisnovatyi-Kogan & Lovelace 1997; Quataert 1998; Gruzinov 1998; Blackman 1999; Quataert & Gruzinov 2000; Pariev & Blackman 2005).

For these reasons, it is still of interest to consider other contributors to the quiescence of elliptical galaxy centres. Here we propose a very simple alternative (or companion process) to these accretion models. We consider the possibility that most of the accreting gas forms stars, shutting off the supply of fuel to the engine. In Section 2 we make a careful estimate of the Bondi accretion rate in elliptical galaxies. In Section 3 we then describe the expected properties of the disc that forms inside the Bondi radius, concluding that it is gravitationally unstable and should form stars. In Section 4 we consider general properties of star-forming 'accretion' discs around black holes, describing the signatures and feedback from star formation, its efficiency and the amount of gas present in the disc in steady state. We discuss the application of this model to M87, other ellipticals, the Galactic Centre, quasars and active galactic nuclei in Section 5, and conclude in Section 6.

2 THE ACCRETION RATE

X-ray observations show that elliptical galaxies contain hot gas approximately at the virial temperature (e.g. Mathews & Brighenti 2003). Where high-resolution data are available, there is often evidence for a temperature gradient such that the central temperature is cooler than the outer regions, e.g. for M87 (Di Matteo et al. 2003) and for NGC 4472, 4636 and 5044 (Mathews & Brighenti 2003, and references therein). Typical temperatures on scales of several hundred parsecs from the galaxy centres are $kT \sim 1$ keV and so the gas is fully ionized with a mean particle mass of $\mu m_{\rm H}$ with $\mu \simeq 0.6$. We define $\mu_{0.6} \equiv \mu/0.6$. This gas has an effective isothermal sound speed of $c_{\rm s} \equiv (P/\rho)^{1/2} = 400\mu_{0.6}^{-1/2}(kT_\beta/{\rm keV})^{1/2}\ {\rm km\,s^{-1}}$, where the total pressure is $P \equiv nkT_\beta$ and ρ is the density. It is possible that thermal pressure accounts for only a fraction, β, of the total pressure (the remainder may include magnetic, turbulent and cosmic ray pressures) and we account for this by defining $T_\beta \equiv T/\beta$ to be an effective temperature that is a factor β^{-1} higher than the actual gas temperature, T, i.e. as inferred from X-ray observations. However, to simplify the equations, henceforth we shall assume $\beta = 1$ and use $T_{\rm keV}$ in place of $kT/(\beta\,{\rm keV})$. To allow for non-thermal pressure, one can simply substitute $T_{\rm keV} \rightarrow T_{\rm keV}/\beta$.

An object of mass M begins to have a significant gravitational influence on the gas at a distance $r_{\rm B}$, the Bondi radius. Note that, in the case of supermassive black holes in giant elliptical galaxies, M may also include a contribution from a concentrated stellar cluster around the black hole. Following the notation of Shu (1992), the Bondi radius is

$$r_{\rm B} \equiv \frac{GM}{c_{\rm s,\infty}^2} = 26.9\mu_{0.6,\infty}\frac{M_9}{T_{\rm keV,\infty}}\ \ {\rm pc}, \tag{1}$$

where $M_9 = M/10^9\,{\rm M_\odot}$ and subscript ∞ refers to a location far from the black hole, i.e. at least several $r_{\rm B}$.

We can express the actual accretion rate as being some fraction $f_{\rm B}$ of the Bondi rate

$$\begin{aligned}\dot{M} &= f_{\rm B}\dot{M}_{\rm B}\\ &= f_{\rm B}\lambda_\gamma 4\pi\rho_\infty (GM)^2\big/c_{\rm s,\infty}^3\\ &= 0.0150 f_{\rm B}\frac{\lambda_\gamma}{0.272}n_\infty\mu_{0.6,\infty}^{5/2}M_9^2T_{\rm keV,\infty}^{-3/2}\ \ {\rm M_\odot\ yr^{-1}},\end{aligned} \tag{2}$$

where λ_γ is a dimensionless factor that depends on the effective equation of state of the gas near the Bondi radius ($\lambda_\gamma = 0.116, 0.272, 0.377, 1.12$ for $\gamma = 5/3, 3/2, 7/5, 1$; Shu 1992, p. 437). The factor $f_{\rm B}$ takes account of effects such as the inhibition of accretion from certain directions because of an outflow from the galactic central engine.

The expected luminosity of the nuclear region inside $r_{\rm B}$ can be expressed as

$$\begin{aligned}L &= \epsilon\dot{M}c^2\\ &= 8.50\times 10^{43}\,\epsilon_{0.1}f_{\rm B}\frac{\lambda_\gamma}{0.272}\mu_{0.6,\infty}^{5/2}\frac{n_\infty M_9^2}{T_{\rm keV,\infty}^{3/2}}\ \ {\rm erg\ s^{-1}},\end{aligned} \tag{3}$$

where $\epsilon_{0.1} \equiv \epsilon/0.1$, with this canonical value being appropriate for thin disc accretion. The fact that luminosities of the nuclear regions of many giant ellipticals are less than the above value has motivated models of either inefficient radiation (small ϵ) or inefficient accretion (small $f_{\rm B}$). We can relate this luminosity to the Eddington value:

$$\frac{L}{L_{\rm E}} = 6.81\times 10^{-4}\,\epsilon_{0.1}f_{\rm B}\frac{\lambda_\gamma}{0.272}\mu_{0.6,\infty}^{5/2}\frac{n_\infty M_9}{T_{\rm keV,\infty}^{3/2}}. \tag{4}$$

Now consider the heating and cooling rates at the Bondi radius. Their relative values determine which value of γ is most appropriate in the equation of state. We assume that the heating is only due to compression. The heating rate per unit volume of gas in a spherical shell at $r_{\rm B}$ is

$$\begin{aligned}\Gamma &= 2P_{\rm B}u_{r,{\rm B}}/r_{\rm B}\\ &= 2f_\rho^\gamma n_\infty kT_\infty\beta_{\rm B}^{-1}f_{\rm u}c_{\rm s,\infty}/r_{\rm B}\\ &= 1.54\times 10^{-21}\ f_\rho^\gamma f_{\rm u}\mu_{0.6,\infty}^{-3/2}T_{\rm keV,\infty}^{5/2}n_\infty M_9^{-1}\ \ {\rm erg\ cm^{-3}\ s^{-1}},\end{aligned} \tag{5}$$

where subscript B refers to the Bondi radius, $f_\rho \equiv \rho_{\rm B}/\rho_\infty = 1.44$, 1.55, 1.64, 2.45 for $\gamma = 5/3$, 3/2, 7/5, 1, and $f_{\rm u} \equiv u_{r,{\rm B}}/c_{\rm s,\infty} =$ 0.78, 0.72, 0.69, 0.46 for $\gamma = 5/3$, 3/2, 7/5, 1, and where in the last expression we have assumed $\beta_{\rm B} = \beta_\infty = 1$. The cooling rate for gas under conditions typical of the centres of ellipticals has been calculated by Sutherland & Dopita (1993). For 0.0274 keV $< kT$ $<$ 2.74 keV and for metallicities that are about a factor of 3 greater than solar, the cooling is dominated by resonance-line emission from

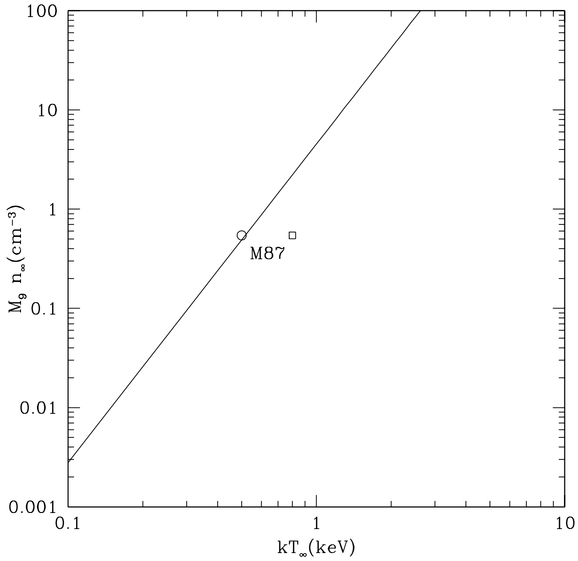

Figure 1. Parameter space for the balance of heating and cooling at the Bondi radius. The solid line shows the condition when the heating rate due to adiabatic compression and the cooling of three times solar metallicity gas are in balance, so that an accretion solution with $\gamma = 1$ would be a reasonable approximation. Systems below this line are unable to cool to maintain the isothermal condition and therefore $\gamma > 1$. The reported data for M87 are shown by the square (Di Matteo et al. 2003). A power-law extrapolation of the observed temperature gradient with radius to the location of the Bondi radius leads to the position marked with a circle. Given existing observational data, consideration of a range of equations of state from $\gamma \simeq 1$–1.5 seems appropriate.

O, Ne and Fe and can be approximated as

$$\Lambda = 6.38 \times 10^{-23}\,\frac{Z}{3.2\,\mathrm{Z}_\odot} T_{\mathrm{keV}}^{-0.71} n^2\ \ \mathrm{erg\,cm^{-3}\,s^{-1}}. \tag{6}$$

We can use equations (5) and (6) to evaluate when heating balances cooling at the Bondi radius, so that the isothermal $\gamma = 1$ solution is most applicable. This condition is

$$n_\infty M_9 = 24.1 f_\rho^{\gamma-2+0.71(\gamma-1)} f_{\mathrm{u}} T_{\mathrm{keV},\infty}^{3.21}\ \ \mathrm{cm}^{-3}, \tag{7}$$

which is also shown in Fig. 1. If at a given temperature the combination $n_\infty M_9$ is less than this isothermal condition, then equations of state with $1 < \gamma < 5/3$ are more appropriate, as they account for the inability of the gas to cool fast enough to maintain the isothermal state. Note that, while we have included the β dependence in the heating rate, we have not allowed for any 'cooling' of the non-thermal component.

Now, as a specific example, consider M87. It is thought to harbour a black hole of mass $(3.2 \pm 0.9) \times 10^9\,\mathrm{M}_\odot$ [this is the dynamical mass inside ~5 pc (Macchetto et al. 1997), but from the mass-to-light ratio stars are unlikely to contribute significantly (Harms et al. 1994)], which is embedded in gas with a density near the Bondi radius of about $0.17 \pm 0.01\,\mathrm{cm}^{-3}$ and temperature 0.80 ± 0.01 keV (Di Matteo et al. 2003). This temperature is an average of the inner kiloparsec, which is a much larger scale than the Bondi radius. Using these values, the heating rate is a factor of 14.6 and 4.1 times larger than the cooling rate for $\gamma = 3/2$ and 1, respectively (see also Fig. 1). However, the data shown by Di Matteo et al. (2003) show a temperature gradient so that the inner regions are cooler. If this trend continues all the way to the Bondi radius, then the cooling rate would balance heating, even keeping the same value of the density. Thus the existing data are not able to distinguish clearly whether conditions are closer to isothermal or adiabatic, and we must consider a range of possible values of γ. For the M87 system, the expected accretion rate is $0.036 f_{\mathrm{B}}(\lambda_\gamma/0.272)\,\mathrm{M}_\odot\,\mathrm{yr}^{-1}$ and the accretion luminosity is $(2.07 \pm 1.2) \times 10^{44} \epsilon_{0.1} f_{\mathrm{B}}(\lambda_\gamma/0.272)\,\mathrm{erg\,s^{-1}}$.

The observed spectrum (νF_ν) from the nucleus of M87 peaks somewhere in the vicinity of 10 μm (e.g. Di Matteo et al. 2003), but with large uncertainties. We can estimate the bolometric luminosity as $L_{\mathrm{bol}} \simeq 4\pi d^2 \nu F_\nu \sim 1.4 \times 10^{41}\,\mathrm{erg\,s^{-1}}$, using the observed 10.8-μm flux of $F_\nu = 16.7 \pm 0.9$ mJy (Perlman et al. 2001b) and assuming a distance of 16 Mpc (Tonry et al. 2001). Biretta, Stern & Harris (1991) estimated a total core luminosity from the radio to the X-ray of $9.1 \times 10^{41}\,\mathrm{erg\,s^{-1}}$. Note, however, that this estimate is mainly based on an interpolation of flux from the radio to the optical bands, and is thus quite uncertain. Applying a similar interpolation for several knots of emission in the jet, Biretta et al. (1991) estimate that it is radiating at $\sim 2 \times 10^{42}\,\mathrm{erg\,s^{-1}}$. Note that the directly observed luminosities of the core and jet are only $\sim 10^{41}\,\mathrm{erg\,s^{-1}}$ in each of the optical and X-ray bands (Biretta et al. 1991; Perlman et al. 2001a; Marshall et al. 2002).

The above luminosity estimates ignore kinetic energy that is escaping from the inner (~20 arcsec) region. Including this contribution, Reynolds et al. (1996b) estimate that the total power of the nucleus may be about an order of magnitude greater, i.e. $2 \times 10^{43}\,\mathrm{erg\,s^{-1}}$, although this is sensitive to the 10^6 yr time-scale adopted for the 5 kpc radio halo.

Thus, although the total luminosity is quite uncertain, both because of the limited spectral coverage for the radiated part, and the difficulty of estimating the kinetic part, it appears that the nucleus is underluminous by about an order of magnitude with respect to models of Bondi accretion with $\epsilon = 0.1$.

Note that it is also possible that the nuclear activity is varying on time-scales shorter than the dynamical time-scale at the Bondi radius,

$$t_{\mathrm{dyn}} = r_{\mathrm{B}}/c_{\mathrm{s},\infty} = 6.6 \times 10^4\, \mu_{0.6,\infty}^{3/2} M_9 T_{\mathrm{keV},\infty}^{-3/2}\ \ \mathrm{yr},$$

which is $\sim 3 \times 10^5$ yr for M87. See Waters & Zepf (2005) for a discussion on the flaring of one of the knots of the jet on time-scales of just a few years. M87 is the only one of the ellipticals studied by Di Matteo et al. (2000, 2003) that is observed to have a powerful jet at the present time.

In the next section we propose a new mechanism that helps to explain why the central black holes of massive ellipticals are underluminous: they are underfed compared to their Bondi rates because much of the accretion flow condenses into stars. Our mechanism is complementary to (not contradictory to) any additional low-luminosity accretion mode or outflow mechanism that is operating in the immediate vicinity of the black hole.

3 THE ACCRETION DISC

The accreting material is likely to have an overall net angular momentum with respect to the central black hole. While the inflow is subsonic, it is relatively easy for turbulent motions, which themselves are limited to about the sound speed, to transport angular momentum outwards. This has been observed in the numerical simulations of Abel, Bryan & Norman (2002) in the context of primordial star formation from gas that is only able to cool quite slowly. Following the approach of Tan & McKee (2004), we assume that angular momentum is conserved inside the sonic point $r_{\mathrm{sp}} = r_{\mathrm{B}}(5 - 3\gamma)/(4\gamma)$ for $\gamma \leqslant 5/3$, and that the material passing through this point has a mean circular velocity $v_{\mathrm{circ}}(r_{\mathrm{sp}}) = f_{\mathrm{Kep}} v_{\mathrm{Kep}}(r_{\mathrm{sp}})$. The density at the sonic point is $\rho_{\mathrm{sp}} = \rho_\infty[2/(5 - 3\gamma)]^{1/(\gamma-1)}$. The simulations of Abel et al. (2002) suggest that the circular velocities are approximately equal to the sound speed, and adopting this as a fiducial value we have $f_{\rm Kep} = (2\gamma)^{-1/2}$. If this is a typical value for the infalling gas, then the median radius at which material joins a rotationally supported disc is

$$r_{\rm d} = f_{\rm Kep}^2 r_{\rm sp} = \frac{r_{\rm sp}}{2\gamma} = r_{\rm B}\frac{5-3\gamma}{8\gamma^2}. \tag{8}$$

Depending on the details of the angular momentum distribution, material will be joining the disc over a range of radii. One possible description of this assumes uniform rotation at the outer boundary, which was adopted by Terebey, Shu & Cassen (1984) and Tan & McKee (2004). The estimate implied by equation (8) is crude: the actual median radius depends on the details of the angular momentum distribution of the infalling gas. The outer radius of the disc may be several times larger than the median, possibly as large as the Bondi radius.

3.1 Gravitational instability and the inevitability of star formation

What is the fate of material that becomes centrifugally supported at $\sim r_{\rm d}$? For typical conditions, the densities are high enough that cooling is sufficient to allow the material to settle into a thin disc. Local viscous processes and/or global instabilities can then act to transport angular momentum outwards and matter inwards. Using the 'alpha' formalism (Shakura & Sunyaev 1973) for viscosity $\nu = \alpha c_{\rm s,d} h = \alpha c_{\rm s,d}^2/\Omega$, where $h = c_{\rm s,d}/\Omega$ is the disc half-thickness and $\Omega = (GM/r^3)^{1/2}$ is the orbital angular velocity, the conservation of mass and angular momentum in a locally steady, thin disc imply the following relation between the mass accretion rate, viscosity and disc surface mass density, $\Sigma_{\rm d} = 2h\rho_{\rm d}$:

$$\dot{M}_{\rm d} = 3\pi\nu\Sigma_{\rm d} = 3\pi\alpha\beta_{\rm d}^b c_{\rm s,d}^2\Omega^{-1}\Sigma, \tag{9}$$

where subscript 'd' refers to conditions in the disc and b is a parameter that allows viscosity to be proportional to gas pressure ($b = 1$) or total pressure ($b = 0$) (Goodman 2003). Note that use of the relation $h = c_{\rm s,d}/\Omega$ does not account for possible pressure support from radiation pressure from a young, massive stellar population in the disc that could couple to the gas via dust. The magnetorotational instability can yield values of $\alpha \sim 10^{-3}$ to 10^{-1} (Balbus & Hawley 1998), while local self-gravity can give $\alpha \lesssim 0.3$ (Gammie 2001). If these processes are relatively inefficient, then for a given accretion rate the disc surface density builds up to such a high level that the disc becomes gravitationally unstable and fragments into bound structures that eventually lead to the formation of stars. This condition can be expressed in terms of the dimensionless Toomre parameter (Toomre 1964)

$$Q = \frac{c_{\rm s,d}\Omega}{\pi G\Sigma_{\rm d}} \simeq \frac{\Omega^2}{2\pi G\rho_{\rm d}}, \tag{10}$$

where gravitational instability occurs for $Q < 1$.

Combining equations (9) and (10) (Goodman 2003) we have

$$G\dot{M}_{\rm d}Q = 3\alpha\beta_{\rm d}^b c_{\rm s,d}^3 \tag{11}$$

for a Keplerian disc. The numerical factor 3 becomes $2\sqrt{2}$ for a flat rotation curve. Although radiation pressure should be small compared to gas pressure at $r_{\rm d}$, justifying the assumption that $\beta_{\rm d} \simeq 1$, stellar feedback on dusty gas could provide partial pressure support analogous to that of radiation pressure. Here we do not consider this further and set $\beta_{\rm d} = 1$.

First consider a disc that is heated purely by viscous dissipation and is cooled by radiation from an optically thick surface. The effective temperature is

$$\begin{aligned} T_{\rm eff,d} &= \left(\frac{3}{8\pi\sigma}\frac{GM\dot{M}}{r^3}\right)^{1/4} \\ &= 4.63\left[\frac{f_{\rm B}^2 n_\infty^2 T_{\rm keV,\infty}^3}{\mu_{0.6,\infty}}\left(\frac{\lambda_\gamma}{0.272}\right)^2\left(\frac{8\gamma^2}{5-3\gamma}\right)^6\right]^{1/8} \\ &\quad\times\left(\frac{r}{r_{\rm d}}\right)^{-3/4}\ {\rm K}, \end{aligned} \tag{12}$$

where σ is the Stefan–Boltzmann constant and we have used equation (2). If the disc is optically thick (an assumption that we check later), then the midplane temperature is given by

$$T_{\rm d}^4 \simeq \frac{3\kappa_{\rm d}\Sigma_{\rm d}}{8}T_{\rm eff,d}^4, \tag{13}$$

assuming the heating rate per unit mass and the opacity per unit mass are constant. As seen from equation (12) we are likely to be in a regime where dust grains can form and survive (i.e. $T_{\rm d} \lesssim 2000$ K). The Rosseland mean opacity per unit gas mass then depends on $T_{\rm d}$ as well as on the gas-to-dust ratio and the properties of the grains. For example, at 100 K and for a gas-to-dust ratio similar to that of the Milky Way, the opacity is about $1\,{\rm cm^2\,g^{-1}}$ (Li & Draine 2001).

The value of the disc surface mass density is

$$\begin{aligned} \Sigma_{\rm d} &= \frac{\dot{M}_{\rm d}\Omega}{3\pi\alpha c_{\rm s,d}^2} \\ &= 153\alpha_{0.3}^{-4/5}\left(f_{\rm B}n_\infty\frac{\lambda_\gamma}{0.272}\right)^{3/5}\mu_{0.6,\infty}^{9/10}M_9^{4/5}T_{\rm keV,\infty}^{-3/20} \\ &\quad\times\kappa_{\rm d}^{-1/5}\mu_{\rm d}^{4/5}\left(\frac{8\gamma^2}{5-3\gamma}\right)^{3/5}\left(\frac{r}{r_{\rm d}}\right)^{-3/5}\ {\rm g\,cm^{-2}}, \end{aligned} \tag{14}$$

where $\mu_{\rm d}$ is the mean particle mass in the disc in units of $m_{\rm H}$. The midplane temperature is

$$\begin{aligned} T_{\rm d} &= 12.7\left(\frac{\kappa_{\rm d}\mu_{\rm d}M_9}{\alpha_{0.3}}\right)^{1/5}\left(f_{\rm B}n_\infty\frac{\lambda_\gamma}{0.272}\right)^{2/5}\mu_{0.6,\infty}^{1/10} \\ &\quad\times T_{\rm keV,\infty}^{27/80}\left(\frac{8\gamma^2}{5-3\gamma}\right)^{9/10}\left(\frac{r}{r_{\rm d}}\right)^{-9/10}\ {\rm K}. \end{aligned} \tag{15}$$

Note that this is the minimum temperature that an optically thick disc can have, since the calculation assumes only viscous heating. Including the factors that depend on γ, the fiducial temperatures are 321 and 78 K for $\gamma = 3/2$ and 1. Such a disc is indeed optically thick ($\tau_{\rm d} = \Sigma_{\rm d}\kappa_{\rm d} > 1$), if dust can form that has properties similar to typical Milky Way dust grains.

What is the value of Q for the optically thick disc? Expressing equation (11) in terms of the fiducial parameters of Bondi accretion in a typical giant elliptical galaxy, we have

$$\begin{aligned} Q &= 4.86\times10^{-4}\,\mu_{\rm d}^{6/5}\kappa_{\rm d}^{13/40}\alpha_{0.3}^{7/10}f_{\rm B}^{-2/5} \\ &\quad\times\left(\frac{\lambda_\gamma}{0.272}\right)^{-2/5}n_\infty^{-2/5}\mu_{0.6,\infty}^{-47/20}M_9^{-17/10}T_{\rm keV,\infty}^{291/160} \\ &\quad\times\left(\frac{8\gamma^2}{5-3\gamma}\right)^{27/20}\left(\frac{r}{r_{\rm d}}\right)^{-27/20}. \end{aligned} \tag{16}$$

The disc is likely to be unstable with respect to its self-gravity: for example, applying the parameters appropriate to M87 yields

$$Q = (0.600, 0.0175)\alpha_{0.3}^{7/10}f_{\rm B}^{-2/5}\kappa_{\rm d}^{13/40}$$

for $\gamma = 3/2$, 1. Note that the relatively large value of Q in the $\gamma = 3/2$ case is a consequence of the small value of r_d. If a radius closer to the observed disc size in M87 (∼100 pc) is adopted, then $Q \ll 1$.

In summary, the accretion of gas to supermassive black holes in the centres of elliptical galaxies is likely to lead to the formation of a gravitationally unstable gas disc. The natural expectation is that this disc will fragment into bound objects: stars. We consider this process and its consequences in the next section.

4 STAR-FORMING DISCS AROUND GALACTIC BLACK HOLES

We expect that gravitational instability in the disc leads to fragmentation and star formation. This process removes mass from the gaseous disc, reducing Σ_d. The newly formed stars heat, and thus stabilize, the remaining gas.

First consider the reduction in the surface density of gas due to star formation. If the stars remain in a thin disc, the system will still be prone to gravitational instabilities. These will heat the stellar population relative to the gas since the latter can cool. As the scaleheight of the stellar distribution becomes larger than that of the gas, the stars have less and less influence on the gravitational stability of the gas. Furthermore, for continued star formation, the gas itself must be self-gravitating, not just the star–gas system. This effect can be seen in the star-forming discs of typical disc galaxies: the threshold for star formation is still adequately described by the Toomre parameter, Q, which depends on the total surface density of gas (Martin & Kennicutt 2001). The presence of the more massive stellar disc raises the particular value of Q below which star formation occurs, but only by factors of a few or less (Jog & Solomon 1984). One way that this can occur is via the local concentration of gas in stellar spiral arms (e.g. Kuno et al. 1995). In summary, as Σ_d is reduced via star formation, the gas disc becomes more and more stable with respect to its self-gravity.

An additional effect of the reduction of Σ_d is a smaller optical depth for the disc's cooling radiation. Cooling becomes most effective when $\tau_d \sim 1$. Other things being equal, we would then expect a somewhat cooler disc midplane temperature than in the optically thick limit, given above. However, now the disc is being heated by stellar photospheres that are much hotter than T_d. For a uniform distribution of gas over the face of the disc, the optical depth through the disc to this radiation from dust absorption is still large, even if the optical depth to thermal radiation at T_d is ∼1.

The amount of heating per mass of new stars depends on their initial mass function (IMF). The upper part of the IMF appears to follow approximately a Salpeter (1955) distribution, $d\mathcal{N}/dm \propto m^{-2.35}$, for $1\,M_\odot \lesssim m \lesssim 100\,M_\odot$ with a dearth of stars beyond the upper mass limit. The IMFs in the Galaxy (Salpeter 1955; Kroupa 2002) and in metal-rich starbursts (Schaerer et al. 2000) appear to be similar in this regime. Below a solar mass the Galactic IMF flattens somewhat (e.g. Muench et al. 2002), but much of the total mass is still contained in subsolar-mass stars. There are no observational constraints on this part of the IMF in starburst galaxies.

For the above IMF, while most of the mass is in low-mass stars, most of the energy injection into the interstellar medium (ISM) comes from high-mass stars. However, the value of the high-mass cutoff in the IMF does not significantly affect the overall feedback properties of the population, so long as it is greater than $\sim 60\,M_\odot$ or so. The uncertainty in the low-mass end introduces an overall uncertainty in the normalization of feedback per unit stellar mass of factors of a few: for example, if the Salpeter IMF extends down to $0.1\,M_\odot$, then the feedback per stellar mass is reduced by a factor of 0.39 from the case with a lower-mass cutoff of $1\,M_\odot$.

Goodman & Tan (2004) have argued that gravitational instability in the inner regions of quasar accretion discs (at radii of a fraction of a parsec) may lead to very massive stars. Conditions there are very different from those at a typical value of r_d of the Bondi-fed discs considered in the present paper. In particular, the inner quasar discs are closer to being dominated by radiation pressure as opposed to gas pressure; temperatures are too high to allow the presence of dust grains; and these regions are only marginally self-gravitating (by definition). We expect that these differences are important enough to lead to very different IMFs. In short, given existing data we feel that adopting a Salpeter IMF is the most reasonable approximation to make for star formation in Bondi-fed gas discs of ellipticals.

For a Salpeter IMF from 0.1 to $100\,M_\odot$, the luminosity associated with a given star formation rate, $\dot{M}_*$, is

$$\begin{aligned} L_* &\simeq 3.9 \times 10^9 \left(\dot{M}_*/\mathrm{M_\odot\,yr^{-1}}\right)\ \mathrm{L_\odot} \\ &= 1.5 \times 10^{43} \left(\dot{M}_*/\mathrm{M_\odot\,yr^{-1}}\right)\ \mathrm{erg\,s^{-1}} \end{aligned}$$

(Leitherer et al. 1999).[1] After about 30 Myr, this stellar population produces a supernova at a rate of

$$7.8 \times 10^{-3} \left(\dot{M}_*/\mathrm{M_\odot\,yr^{-1}}\right)\ \mathrm{yr^{-1}},$$

causing a mean injection rate of mechanical energy of

$$\sim 2.3 \times 10^{41} \left(\dot{M}_*/\mathrm{M_\odot\,yr^{-1}}\right)\ \mathrm{erg\,s^{-1}}.$$

The mechanical luminosity of stellar winds is estimated to be about a factor of 6 smaller. Thus the maximum energy input from stars is about

$$\epsilon_* \equiv E_*/\left(M_* c^2\right) = 2.6 \times 10^{-4}.$$

The number of H-ionizing photons produced is

$$(9.2, 7.9, 7.2) \times 10^{52} \left(\dot{M}_*/\mathrm{M_\odot\,yr^{-1}}\right)\ \mathrm{s^{-1}}$$

for metallicity $Z = 1.0, 2.0, 3.0\ \mathrm{Z_\odot}$ (Smith, Norris & Crowther 2002), where the highest metallicity value is based on a power-law extrapolation from the lower two values. Assuming case B recombination in the ionized regions with no escape of ionizing photons yields an Hα luminosity of (Kennicutt, Tamblyn & Congdon 1994)

$$L_{\mathrm{H}\alpha} = 0.99 \times 10^{41} f_{\mathrm{trap,i}} \left(\dot{M}_*/\mathrm{M_\odot\,yr^{-1}}\right)\ \mathrm{erg\,s^{-1}}, \tag{17}$$

with the normalization being for the $Z = 3\,\mathrm{Z_\odot}$ case and where $f_{\mathrm{trap,i}}$ is the fraction of ionizing photons that are trapped in the disc's H II region and do not escape into the hot interstellar medium of the galaxy. This estimate ignores dust absorption in the disc, which would lower the numerical factor in the above equation. Note that a recombination flow of one solar mass per year would produce an Hα luminosity that is a factor of 4.4×10^{-4} times smaller than the amount produced by one solar mass per year of star formation (with the above Salpeter IMF).

The Hα flux from the circumnuclear disc of M87 is $(4.4 \pm 1.5) \times 10^{-14}\ \mathrm{erg\,cm^{-2}\,s^{-1}}$ (Ford et al. 1994) (see Fig. 2), corresponding to a luminosity of $1.35 \times 10^{39}\ \mathrm{erg\,s^{-1}}$ for an assumed distance of 16 Mpc (Tonry et al. 2001). This implies a star formation rate of

$$1.4 \times 10^{-2} f_{\mathrm{trap,i}}^{-1}\ \mathrm{M_\odot\,yr^{-1}}$$

[1] This is actually the luminosity of a stellar population that has been continuously forming stars at $1.0\,\mathrm{M_\odot\,yr^{-1}}$ for the past 10^7 yr, which is approximately the age at which ionizing feedback reaches a constant level.

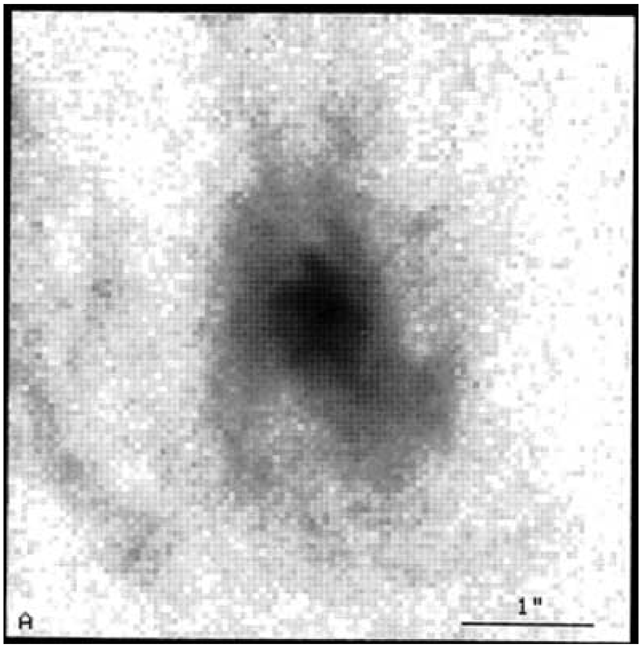

Figure 2. Adapted from Ford et al. (1994), and reproduced by permission of the AAS: *HST* F658N on-band image minus *HST* F547M off-band image of the M87 nucleus designed to highlight Hα emission. North is towards the top and east to the left; the famous synchrotron jet (not visible here) extends to the north-west. The black hole is at the centre of the image and is surrounded by a disc of material: the velocity difference from 1 arcsec to the north-east to 1 arcsec to the south-west is 1000 km s^{-1} (Harms et al. 1994; Macchetto et al. 1997). The presence of this disc, including its spiral structure and dust lanes, as well as its total Hα luminosity, is evidence in support of a model of circumnuclear star formation.

using the normalization of equation (17). Note that this estimate is a lower limit because it does not allow for any internal extinction in the star-forming disc that would attenuate the observed Hα flux.

If we assume that the bolometric luminosity of the circumnuclear disc in M87 is similar to the observed peak in νL_ν, i.e. $\sim 1.4 \times 10^{41}$ erg s^{-1} (Perlman et al. 2001b); see Section 2), and is due entirely to star formation, then the above relations imply a star formation rate of 9.5×10^{-3} M$_\odot$ yr^{-1}. Fig. 3, described in more detail in the next section, shows how the bolometric luminosity due to star formation compares to the observed spectrum of the M87 nucleus.

Despite the uncertainties in the IMF, internal extinction and equation of state of the accreting gas, the key result from the above star formation rate estimates is that they are both roughly consistent with the inferred Bondi accretion rate. This agreement suggests that the star-forming accretion flow model provides a plausible mechanism for preventing significant gas accretion on to the central black hole, at least in M87. This is in addition to the fact that a relatively thin disc is observed to be present on scales just inside the Bondi radius (see further discussion in Section 5.1).

4.1 Estimate of the star formation efficiency and residual accretion on to the black hole

Here we present a simple one-zone model that describes the star formation efficiency of a $Q = 1$ disc as a function of the coupling of star formation feedback to the gas. From equation (11) we have

$$Q = \frac{3\alpha}{G\dot{M}_{\rm d}} \left(\frac{kT_{\rm d}}{\mu_{\rm d}} \right)^{3/2}. \tag{18}$$

For an optically thick disc, $T_{\rm d}^4 \simeq 3\kappa_{\rm d}\Sigma T_{\rm eff}^4/8$ and its luminosity is

$$L \simeq 2\pi r_{\rm d}^2 \sigma T_{\rm eff}^4 = f_{\rm trap,L} L_*$$
$$\simeq f_{\rm trap,L} \times 3.9 \times 10^9\,(\dot{M}_*/{\rm M_\odot\ yr^{-1}})\ {\rm L_\odot},$$

assuming a negligible contribution from viscous heating, valid at $r \sim r_{\rm d}$ if we are to achieve a quasi-stable disc with $Q \sim 1$. In the last step above, we have used the IMF-dependent relation between luminosity and star formation rate discussed previously. Now by using the above relations, including equation (10), assuming that $Q = 1$ and that most of the Bondi accretion rate goes into stars, i.e. $\dot{M}_* \simeq \dot{M}$, we derive the mass accretion rate in the disc that can be supported against self-gravity by stellar feedback:

$$\begin{aligned}\dot{M}_{\rm d} &= 2.27 \times 10^{-6}\,\alpha_{0.3} \left(\frac{\kappa_{\rm d}^2 f_{\rm trap,L}^2 L_1^2 M_9}{\mu_{\rm d}^8} \right)^{3/14} \\ &\quad \times \left(\frac{\dot{M}_*}{M_\odot\,{\rm yr}^{-1}} \right)^{3/7} \left(\frac{r_{\rm d}}{100\,{\rm pc}} \right)^{-3/2} \ {\rm M_\odot\ yr^{-1}} \\ &= 2.686 \times 10^{-6}\,\alpha_{0.3} \left(\frac{\kappa_{\rm d} f_{\rm trap,L} L_1 T_{\rm keV,\infty}^2 f_{\rm B} n_\infty}{\mu_{\rm d}^4 M_9 \mu_{0.6,\infty}} \frac{\lambda_\gamma}{0.272} \right)^{3/7} \\ &\quad \times \left(\frac{8\gamma^2}{5-3\gamma} \right)^{3/2} \ {\rm M_\odot\ yr^{-1}},\end{aligned} \tag{19}$$

where $L_1 = L_*/(3.9 \times 10^9\,{\rm L_\odot})$ is the luminosity generated by a star formation rate of 1 M$_\odot$ yr^{-1} for our adopted IMF. This is the residual amount of gas accretion that can be stabilized with respect to its own self-gravity by transforming most of the initial mass flux into stars. We have written equation (19) in two stages so that different constraints can be used: either the observed disc size and the Bondi accretion rate to the disc (or any general mass feeding rate to the disc), or the properties of the hot gas in the elliptical (in this case the theoretical estimate for the disc size is used; equation 8). For M87, in the first case the accretion rate in the disc is

$$(7.0 \times 10^{-7}, 1.3 \times 10^{-6})\alpha_{0.3} \left(\kappa_{\rm d} f_{\rm trap,L} f_{\rm B} \mu_{\rm d}^{-4} \right)^{3/7}\ {\rm M_\odot\ yr^{-1}}$$

for $\gamma = 3/2, 1$ with $r_{\rm d} = 100$ pc; while in the second case it is

$$(1.36 \times 10^{-4}, 9.25 \times 10^{-6})\alpha_{0.3} \left(\kappa_{\rm d} f_{\rm trap,L} f_{\rm B} \mu_{\rm d}^{-4} \right)^{3/7}\ {\rm M_\odot\ yr^{-1}}$$

for $\gamma = 3/2, 1$. These are much lower than the initial Bondi accretion rates of (0.036, 0.148) $f_{\rm B}$ M$_\odot$ yr^{-1}.

If the disc is optically thin to its cooling radiation, then

$$T_{\rm d}^4 \simeq T_{\rm eff,d}^4 \big/ \tau_{\rm d} = T_{\rm eff,d}^4 \big/ (\kappa_{\rm d}\Sigma_{\rm d}),$$

and the equation for the mass accretion rate in the disc becomes

$$\begin{aligned}\dot{M}_{\rm d} &= 8.57 \times 10^{-5}\,\alpha_{0.3} \left(\frac{f_{\rm trap,L}^2 L_1^2}{\mu_{\rm d}^8 \kappa_{\rm d}^2 M_9} \right)^{1/6} \\ &\quad \times \left(\frac{\dot{M}_*}{M_\odot\,{\rm yr}^{-1}} \right)^{1/3} \left(\frac{r_{\rm d}}{100\,{\rm pc}} \right)^{-1/6} \ {\rm M_\odot\ yr^{-1}} \\ &= 2.63 \times 10^{-5}\,\alpha_{0.3} \left(\frac{f_{\rm trap,L} L_1 f_{\rm B} \mu_{0.6,\infty}^2 n_\infty M_9}{\mu_{\rm d}^4 \kappa_{\rm d} T_{\rm keV,\infty}} \frac{\lambda_\gamma}{0.272} \right)^{1/3} \\ &\quad \times \left(\frac{8\gamma^2}{5-3\gamma} \right)^{1/6} \ {\rm M_\odot\ yr^{-1}}.\end{aligned} \tag{20}$$

Now, for M87, in the first formula of (20), the accretion rate in the disc is

$$(2.3 \times 10^{-5}, 3.7 \times 10^{-5})\alpha_{0.3} \left(\kappa_{\rm d}^{-1} f_{\rm trap,L} f_{\rm B} \mu_{\rm d}^{-4} \right)^{1/3}\ {\rm M_\odot\ yr^{-1}}$$

for $\gamma = 3/2, 1$ with $r_{\rm d} = 100$ pc; while in the second formula it is

$$(4.2 \times 10^{-5}, 4.7 \times 10^{-5})\alpha_{0.3} \left(\kappa_{\rm d} f_{\rm trap,L} f_{\rm B} \mu_{\rm d}^{-4} \right)^{3/7}\ {\rm M_\odot\ yr^{-1}}$$

for $\gamma = 3/2, 1$.

We can judge which of the optically thin or thick limits is appropriate by the value of

$$\Sigma_{\rm d} = \dot{M}_{\rm d}\Omega\big/\left(3\pi\alpha c_{\rm s,d}^2\right) = [\dot{M}_{\rm d}/(3\alpha)]^{1/3}\Omega\big/[\pi(GQ)^{2/3}]$$

and the disc temperature, which sets $\kappa_{\rm d}$. For the optically thick case the surface density is

$$\Sigma_{\rm d} = 7.05\times10^{-3}\left(\frac{\kappa_{\rm d} f_{\rm trap,L} L_1 M_9^4}{\mu_{\rm d}^4}\frac{\dot{M}_*}{\rm M_\odot\ yr^{-1}}\right)^{1/7} \times\left(\frac{r_{\rm d}}{100\,\rm pc}\right)^{-2}\ \rm g\ cm^{-2}, \tag{21}$$

while in the optically thin case it is

$$\Sigma_{\rm d} = 0.0236\left(\frac{f_{\rm trap,L} L_1 M_9^4}{\kappa_{\rm d}\mu_{\rm d}^4}\frac{\dot{M}_*}{\rm M_\odot\ yr^{-1}}\right)^{1/9} \times\left(\frac{r_{\rm d}}{100\,\rm pc}\right)^{-14/9}\ \rm g\ cm^{-2}. \tag{22}$$

For a dust-to-gas ratio three times that of the Milky Way, $A_V \simeq 600$ Σ, so that the above two numerical factors correspond to $A_V = 4.2$, 14, respectively. The effective temperature at the disc surface is

$$T_{\rm eff,d} = \left(\frac{f_{\rm trap,L} L_*}{2\pi r_{\rm d}^2\sigma}\right)^{1/4} = 25.8\left(\frac{f_{\rm trap,L} L_1 \dot{M}_*}{\rm M_\odot\ yr^{-1}}\right)^{1/4}\left(\frac{r_{\rm d}}{100\,\rm pc}\right)^{1/2}\ \rm K \tag{23}$$

For most elliptical galaxies, the disc will be optically thin to its cooling radiation in the outer parts, and will make a transition to the optically thick regime in the inner region. Note that the discs are expected to be quite optically thick to their stellar radiation field, although clustering of young stars and ISM inhomogeneities will act to reduce this. Because of these complications, we have left the factor $f_{\rm trap,L}$ as a free parameter, but expect it to have a value of the order of unity.

Despite the uncertainties in the parameters and the dependence on the adopted disc size, the estimates suggest that only a small fraction of the initial mass accretion flux through the disc would be supported against star formation and accreted by the central black hole. The amount of the reduction implied by equations (19) and (20) is more than sufficient to explain the low luminosity of the nucleus of M87. However, the residual accretion rates predicted by the above simple model are so small that other processes that have not been considered, such as mass supply from stellar winds, may now be relatively important. Another caveat is that the above estimates only consider thermal pressure support in the gas disc and ignore the additional stabilizing effect of magnetic fields.

Note that some material will be expelled from stars as they undergo and complete their stellar evolution. For Salpeter-like IMFs, most of this matter will be returned after a relatively long time-scale: for our adopted Salpeter IMF from 0.1 to 100 $\rm M_\odot$, only 14 per cent of the mass is in stars with masses greater than 8 $\rm M_\odot$. Irrespective of the time-scale of mass return, most of this material would be recycled back into the gas disc rather than accreted directly by the black hole. This is because most stars are forming and evolving at distances from the black hole that are still quite large compared to the sizes of their stellar wind bubbles or planetary nebulae, and the specific angular momentum of the returned gas is the same as that of the stars.

Other authors have highlighted the issue of self-gravity in the accretion discs that feed supermassive black holes, mainly in the context of quasars that are accreting at much higher rates, approaching their Eddington limits (Schlosman & Begelman 1989; Goodman 2003; Sirko & Goodman 2003; Thompson, Quataert & Murray 2005). Goodman (2003) pointed out that stellar heating of the disc was unlikely to be able to stabilize the disc, because the energy release per mass of stars formed is relatively low. We have extended this line of reasoning to conclude that most of the mass flux goes into stars, which are then able to stabilize a small residual fraction of the gas in an accretion disc to the black hole. The implications of this model for the fuelling of quasars are discussed in Section 5.4.

4.2 Estimate of the total mass of gas in the disc

One prediction of our model of (Bondi-fed) star-forming accretion flows is the presence of cold, presumably molecular gas. To estimate the amount, we assume that a steady state has been achieved so that the star formation rate in the disc balances the Bondi accretion rate to the disc. The disc mass is $M_{\rm d} \simeq \pi r_{\rm d}^2\Sigma_{\rm d}$. In the optically thick case this is

$$M_{\rm d} = 1.06\times10^6\left(\frac{\kappa_{\rm d} f_{\rm trap,L} L_1 M_9^4}{\mu_{\rm d}^4}\frac{\dot{M}_*}{\rm M_\odot\ yr^{-1}}\right)^{1/7}\ \rm M_\odot, \tag{24}$$

while in the optically thin case it is

$$M_{\rm d} = 3.55\times10^6\left(\frac{f_{\rm trap,L} L_1 M_9^4}{\kappa_{\rm d}\mu_{\rm d}^4}\frac{\dot{M}_*}{\rm M_\odot\ yr^{-1}}\right)^{1/9} \times\left(\frac{r_{\rm d}}{100\,\rm pc}\right)^{4/9}\ \rm M_\odot. \tag{25}$$

For M87, the disc mass is then

$$(1.3,\,1.6)\times10^6\left(\kappa_{\rm d} f_{\rm trap,L} L_1 \mu_{\rm d}^{-4}\right)^{1/7}\ \rm M_\odot$$

for $\gamma = 3/2$, 1 in the optically thick case, and

$$(4.1,\,4.8)\times10^6\left(\kappa_{\rm d}^{-1} f_{\rm trap,L} L_1 \mu_{\rm d}^{-4}\right)^{1/9}\ \rm M_\odot$$

for $\gamma = 3/2$, 1 in the optically thin case.

As an alternative method, we use an empirical relation between star formation rate and disc gas mass to estimate the properties of the gas disc. For a sample of about a hundred galactic and circumnuclear discs, Kennicutt (1998) found an empirical linear relation between the surface density of star formation and the product of the total mean gas surface density and the mean orbital angular velocity:

$$\Sigma_{\rm SFR} \simeq 0.017\Sigma_{\rm d}\Omega, \tag{26}$$

where $\Sigma_{\rm SFR} \equiv \dot{M}_*/(\pi r_{\rm d}^2)$. It is not clear if this relation applies to Bondi-fed circumnuclear discs in ellipticals, as the orbital time-scales are much shorter than in Kennicutt's sample and the rotation curves may be closer to Keplerian (e.g. M87, Macchetto et al. 1997), rather than being flat or rising. This may lead to somewhat different star formation efficiencies, at least in the context of shear-dependent models of star formation (Tan 2000).

Nevertheless, using equation (26) we can estimate the surface density of the gas disc:

$$\Sigma_{\rm d} = 0.184\frac{\dot{M}_*}{\rm M_\odot\ yr^{-1}} M_9^{-1/2}\left(\frac{r_{\rm d}}{100\,\rm pc}\right)^{-1/2}\ \rm g\ cm^{-2}. \tag{27}$$

In the case of M87 this takes values

$$(3.7,\,15.3)\times10^{-3}\ \rm g\ cm^{-2} \equiv (18,\,73)\ M_\odot\ pc^{-2}$$

for $\gamma = 3/2$, 1, which is in the middle of the range of values covered by Kennicutt's sample. The total gas mass in this disc is

$$M_{\rm d} = \pi r_{\rm d}^2 \Sigma_{\rm d} = \frac{\dot{M}_*}{0.017\Omega}$$
$$= 2.77 \times 10^7 \frac{\dot{M}_*}{{\rm M}_\odot\,{\rm yr}^{-1}} M_9^{-1/2} \left(\frac{r_{\rm d}}{100\,{\rm pc}}\right)^{3/2} \;{\rm M}_\odot, \qquad (28)$$

which takes values (0.6, 2.3 $\times$ 10^6) $\mathrm{M}_\odot$ in the case of M87 with $\gamma = 3/2$, 1. The coincidence of the two different methods of estimating the disc mass suggests that this feedback model may have relevance to setting the normalization of the empirical Kennicutt (1998) relation.

Note that, if the disc were disrupted by a particular event, such as an outburst from the AGN, then it would take $\sim 10^7$–10^8 yr to reform, if supplied at the Bondi rate.

5 DISCUSSION

We have presented a model of star-forming accretion flows with particular attention to the flow inside the Bondi radius of the central black hole of an elliptical galaxy. Our approach was as follows: First we calculated the accretion rate, showing the explicit dependence on the effective equation of state of the gas. We then estimated the heating and cooling rates of the gas in the accretion flow near the Bondi radius, showing that idealized accretion solutions $\gamma \simeq 1$–1.5 are probably quite a reasonable description of the actual accretion flow, at least in the case of M87. We then argued that angular momentum conservation inside the sonic point of the flow leads to formation of a disc on scales that are a factor of a few or so smaller than the Bondi radius. The disc was shown to be gravitationally unstable if the heating is due to viscous processes in the disc. We found that self-gravity is so strong that most of the mass flux will be converted into stars. The energy input from these stars can stabilize only a small residual fraction of the initial mass flux, i.e. the accretion rate to the black hole is substantially reduced below the Bondi accretion rate. We argue that this is a major cause of the 'low luminosity' of black holes observed in giant elliptical galaxies, where 'low luminosity' is in reference to a model of thin disc accretion at the Bondi rate as determined from the conditions in the kiloparsec-scale, X-ray-emitting gas. A radiatively inefficient accretion flow may exist at the centre of the disc, much closer to the black hole. However, its inefficiency must be argued in the context of an accretion rate and a luminosity determined at a much smaller scale, for example where the disc is no longer self-gravitating.

The predictions of this model include the presence of a star-forming gas disc (i.e. containing relatively cold molecular gas) around the black holes of ellipticals and the presence of hot, young stars in these regions. These may reveal themselves by spectral features (e.g. clustered Hα emission from gas ionized by massive stars; stellar atmosphere/wind features, e.g. from Wolf–Rayet stars; or stellar ultraviolet continuum, although this may be swamped by AGN emission) or by the occurrence of Type II supernovae. For example, a typical star formation rate at the centre of a giant elliptical galaxy is expected to be $\sim 0.1\,\mathrm{M}_\odot\,\mathrm{yr}^{-1}$, about a factor of 30 less than the star formation rate of a Milky Way-like disc galaxy. For the standard Salpeter IMF adopted here, the Type II supernova rate from such an elliptical galaxy is $\sim 8 \times 10^{-4}\,\mathrm{yr}^{-1}$.

Nayakshin (2004) considered how the presence of many young stars, left over from a previous quasar phase, can affect the accretion and appearance of low-luminosity AGN. However, he did not consider the possibility that the stars are currently forming in these systems.

We now summarize the potential application of the star-forming accretion disc model to M87, other giant ellipticals, the Galactic Centre, quasars and other AGN.

5.1 Application to M87

We have used M87 as an example to illustrate the model throughout the text. Here we summarize this application, discuss the existing evidence that supports the model, and suggest future observations that can confirm or refute its validity.

The central black hole of M87, with mass $(3.2 \pm 0.9) \times 10^9\,\mathrm{M}_\odot$ (Macchetto et al. 1997), should have a Bondi accretion rate of 0.036–0.15 $\mathrm{M}_\odot\,\mathrm{yr}^{-1}$ (this range corresponds to uncertainties in $\gamma = 3/2$–1) from the hot (0.8 keV), low-density ($n_{\rm e} \simeq 0.17\,\mathrm{cm}^{-3}$) gas (Di Matteo et al. 2003) at the centre of the giant elliptical (uncertainties in the black hole mass and gas temperature could also change the accretion rate by factors of 3 or so). The Bondi radius is about 110 pc, but could be twice as large if the high end of the black hole mass determination is valid and the temperature in this inner region is actually 0.5 keV (see discussion in Section 2). The expected *median* disc radius is several times smaller than the Bondi radius: our fiducial case has it smaller by factors of 36 and 4 for $\gamma = 3/2$ and 1. The actual outer disc radius would be larger than the median by factors of a few, and in any case the exact size depends on the details of the initial angular momentum distribution of the hot accreting gas, which is not well constrained. Since the Hα disc seen by *HST* is $\simeq$100 pc in extent (Ford et al. 1994; see Fig. 2), this is probably the best estimate of the disc scale in this system. Note also that the *HST* observations show some filamentary structure out to several hundred parsecs, so that the accretion process may not be quite a simple as given by the purely Bondi solution.

Despite the caveats, the coincidence of the observed disc size with the estimated Bondi radius supports a model in which that disc is forming from a Bondi accretion flow with initial circular motions of the order of the hot gas sound speed.

The Hα disc was described by Ford et al. (1994) as having 'spiral structure' and 'dust lanes' (see Fig. 2), which supports the theoretical conclusion that it is gravitationally unstable (Section 3) and contains cool gas that is liable to undergo star formation (Sections 3 and 4). The total Hα luminosity implies a star formation rate (assuming a standard Salpeter IMF) within factors of a few of the Bondi accretion rate at which the disc is being fed. This is consistent with the implied star formation rate being equal to the Bondi accretion rate given the uncertainties.

We can model the bolometric luminosity associated with this star formation rate. Fig. 3 shows the observed infrared and radio spectrum of the M87 nucleus, much of which results from non-thermal emission from the AGN (see Section 2). Care must be taken since the fluxes at different wavelengths often correspond to regions of different scales. Also shown in Fig. 3 is a simple blackbody spectrum of the gas and dust disc heated by star formation at rates 0.036–0.15 $\mathrm{M}_\odot\,\mathrm{yr}^{-1}$, corresponding to luminosities (1.6–5.9) $\times$ $10^8\,\mathrm{L}_\odot$. The temperature is set by assuming that this luminosity is emitted over the two faces of a disc of radius 100 pc. This model is highly idealized, since it assumes a single temperature, ignores additional heating from the interstellar radiation field of the galaxy, ignores deviations from a pure blackbody spectrum, such as PAH (polycyclic aromatic hydrocarbon) emission features in the range 1–20 μm, and assumes that all the light from stars is reprocessed by dust. The overall luminosity normalization is also IMF-dependent. Nevertheless, it is interesting that these simple models are now in a position to be tested by high-resolution infrared and submillimetre

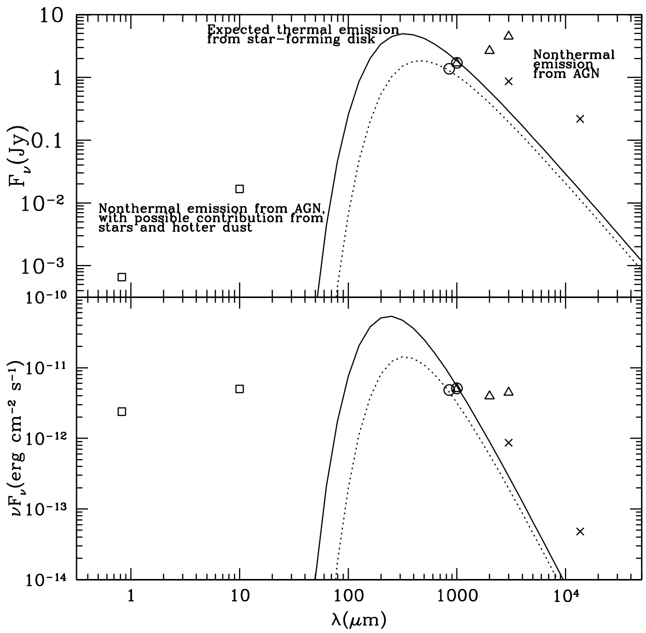

Figure 3. Continuum spectrum of the nucleus of M87 shown as F_ν (top) and νF_ν (bottom). Data in order of increasing wavelength: squares – *HST* optical (Ho 1999), and *Gemini* 10 μm (Perlman et al. 2001b); circles – SMA calibration data (sensitive to flux on scales of the order of 2 arcsec); triangles – IRAM 30-m data (sensitive to flux on scales of the order of 11 and 28 arcsec at 1.3 and 3 mm, respectively (data also extracted from the SMA calibrator list); crosses – VLBI data at 100 GHz (Baath et al. 1992) and 22 GHz (Spencer & Junor 1986) sensitive to flux on scales of 0.2 and ∼0.00015 arcsec. The solid curve is a blackbody spectrum from a star-forming disc with the star formation rate equal to the upper limit of the Bondi feeding rate of the disc (0.15 $M_\odot$ yr^{-1}) and a standard Salpeter mass function down to 0.1 $M_\odot$. The temperature used here is 16 K set by having the emitting area be equal to that of the observed Hα disc with radius $r_d \simeq 100$ pc. This is a lower limit, as it neglects additional heating from the evolved stellar population of the galaxy. The dotted curve is the equivalent model for the lower limit of the Bondi accretion rate (0.036 $M_\odot$ yr^{-1}).

observations that are now feasible with the *Spitzer Space Telescope* and the Sub-Millimeter Array (SMA).

The presence of cold, presumably molecular, gas in the disc with total mass $\sim 10^6\,M_\odot$ could be searched for with millimeter interferometers. Such a search is particularly well suited to the capabilities of the SMA and the Atacama Large Millimeter Array (ALMA).

To estimate the efficiency of the central engine, we need to know the accretion rate of the black hole. As noted in Section 2, to explain the observations requires a reduction in the accretion rate by a factor of ∼10 from the Bondi value (this factor is itself uncertain by about an order of magnitude), or a similar reduction in the efficiency of energy liberation from the fiducial value of $\epsilon = 0.1$. Note that most of the energy is thought to be released as non-radiating kinetic energy of the jet/outflow. The reduction in black hole accretion rate from the Bondi rate due to star formation is also quite uncertain. We estimated (equation 20) a possible reduction by factors of up to ∼1000. However, as discussed in Section 4, this simple estimate is based on a simple one-zone model for the disc and does not include effects such as accretion of some very low angular momentum gas from the larger-scale galaxy and the return of material evolved stars. In summary, the observations of the M87 nucleus can be explained by models of efficient ($\epsilon \sim 0.1$) accretion if a large fraction (∼90 per cent) of the initial mass flux to the disc forms stars. While such a high fraction is certainly allowed by the simple estimates of self-regulated star formation we have presented in Section 4, quantitative estimates of the residual accretion rate are difficult.

5.2 Other elliptical galaxies

Optical imaging with *HST* has been used to detect dusty discs in elliptical galaxies. Jaffe et al. (1993) reported a disc of cool gas and dust with radial extent of 65 pc around the nucleus of NGC 4261 (3C 270), another elliptical galaxy in the Virgo cluster (see also Martel et al. 2000; van Bemmel et al. 2005). Bower et al. (1997) reported a disc, with radius 82 pc, of ionized gas in M84 (NGC 4374 = 3C 272.1). Jaffe et al. (1994) surveyed 14 Virgo ellipticals, finding clear evidence for dust discs (seen in absorption) in several. Note that this method of detection could be sensitive to the orientation of the disc and to other properties of the nucleus: for example, in this study no dust disc was seen in absorption in the case of M87. They note that essentially all their systems exhibit a change in morphology on scales of ∼100–200 pc from the nucleus: either in the extent of an absorbing dust disc or in a change in the brightness profile. We identify this scale with the Bondi radius, inside which a cool, star-forming disc can form. Van Dokkum & Franx (1995) found that dust ‘lanes’ were present in the majority (∼80 per cent) of a sample of 64 early-type galaxies. These were usually perpendicular to the axis of any radio jet present in the system. Carollo et al. (1997) and Tomita et al. (2000) also found evidence for nuclear dust absorption in the majority of the galaxies they analysed. Martel et al. (2004) found nuclear discs of dust and ionized gas in a large fraction of the nine early-type galaxies they studied with *HST*-ACS. Again the typical scales of these features are ∼100 pc. They estimated star formation rates of (0.6–6) $\times 10^{-3}\,M_\odot$ yr^{-1} assuming all the Hα emission is produced by star formation and using the normalization of Kennicutt (1998). Lauer et al. (2005) found dust in the centres of about half their sample of 77 galaxies. Only one of their ‘dusty’ galaxies had dust that was more prominent on scales greater than 4 arcsec from the centre than closer to the centre; most of the dust is highly concentrated, and in all the systems where the dust is clearly in a disc or a ring, no dust clouds are visible external to the outer edge of the disc or ring. Significantly, no dust was seen in relatively low-luminosity, low-mass ellipticals ($M_V > -21$), which might be expected since the Bondi accretion rate increases with galaxy mass, assuming the black hole mass is a fixed fraction of the galaxy mass.

Jaffe & McNamara (1994) detected radio absorption in H I and CO from neutral gas in the disc of NGC 4261 and estimated a mass $\sim 2 \times 10^5\,M_\odot$, within an order of magnitude of the theoretical mass estimates we have made for the disc in M87. Knapp & Rupen (1996) surveyed 42 elliptical galaxies for CO emission, concluding that the detection rate was just under 50 per cent. However, since the half-power beam of these observations was about 30 arcsec, these detections most likely do not relate to molecular gas forming inside the Bondi radius, but to material formed from evolved stars or associated with the infall of gas-rich satellite galaxies. Young (2002) detected CO emission in five out of seven ellipticals, but again, because of the resolution of the observations, the focus was on roughly kiloparsec scales. Vila-Vilaró, Cepa & Butner (2003) detected CO(3–2) emission from six out of 10 early-type galaxies that already showed CO(2–1) emission, with a beam of about 20 arcsec. They inferred typical star formation rates (from Hα) of a few tenths of a solar mass per year, but again this is mostly on scales larger than the Bondi radius (nuclear Hα emission was not considered in this study).

In summary, there is much evidence that elliptical galaxies contain dusty gas discs on scales approximately equal to their Bondi radii, i.e. ∼100 pc. It is less clear if star formation is actively occurring in these discs. One could make a more quantitative test of whether central dust concentrations are related to the Bondi radius by comparing the size of the dust features with the r_B. If the gas temperature near the Bondi radius is proportional to the virial temperature, then $T_\infty \propto T_{vir} \propto \sigma_*^2 \propto M_{gal}^{0.32}$, where σ_* is the galaxy's stellar velocity dispersion and M_{gal} is its mass, and the last step uses the empirical relation of Faber et al. (1997). Assuming a constant black hole to galaxy mass fraction, i.e. $M \propto M_{gal}$, the Bondi radius then scales as $r_B \propto M_{gal}^{0.68}$.

Long-term star formation at the scale of the Bondi radius (∼100 pc) will have implications for the morphologies of the cores of elliptical galaxies. For example, it may account for the ∼100 pc scale discs seen in some early-type galaxies (e.g. van den Bosch, Jaffe & van der Marel 1998; Krajnović & Jaffe 2004). It may also account for the flattening in the radial profiles of 'core' elliptical galaxies (e.g. Lauer et al. 2005): this occurs on scales of ∼100 pc in galaxies that tend to be the most massive ellipticals.

5.3 The Galactic Centre

The black hole at the Galactic Centre has a mass of $\sim 3\text{–}4 \times 10^6$ M$_\odot$, appears to be embedded in hot gas (as seen in X-rays), yet is underluminous by several orders of magnitude with respect to thin disc accretion at the Bondi rate (Baganoff et al. 2003, and references therein). Some, but probably not the majority, of the diffuse X-ray emission on the arcsecond scales relevant to Bondi accretion may be the dust-scattered halo of the X-ray emission from inside the black hole's Bondi radius (Tan & Draine 2004). Basic ADAF models have been ruled out for the Galactic Centre due to polarization measurements, which indicate that the density is even lower than that predicted by ADAFs (Bower et al. 2003).

This system is probably more complicated than the supermassive black holes in elliptical galaxies because the Bondi radius is only ∼0.1 pc in size, which is relatively small compared to the feedback scales associated with stellar wind bubbles and supernova remnants from individual massive stars. For example, the dominant source of mass in the region is likely to be from stellar winds from young, massive stars. Some of these are even present inside the Bondi radius (as estimated for the diffuse X-ray-emitting gas) and because of the relatively low mass of the black hole, these winds with speeds ~ 1000 km s^{-1} may be unbound and act to prevent diffuse gas from accreting (Quataert 2004).

A further complication is that there are large amounts of neutral gas in the vicinity (Herrnstein & Ho 2005; Christopher et al. 2005). The sporadic infall of this material may well dominate the time-averaged mass accretion rate.

5.4 Quasars and AGN

The issue of self-gravity and quasar accretion discs has been discussed by Schlosman & Begelman (1989), Goodman (2003) and Thompson et al. (2005). Our model herein has been developed with low-luminosity galactic nuclei in mind, for which the central regions are being fed at the Bondi rate from their host elliptical galaxies. The accretion rates to the black holes powering quasars approach the Eddington limited value:

$$\dot{M}_{d,E} = 4\pi G M l_E/(\kappa_{es} c \epsilon) \simeq 22(l_E/\epsilon_{0.1}) M_9 \ \mathrm{M_\odot\ yr^{-1}},$$

where $\kappa_{es} \simeq 0.4\,\mathrm{cm^2\,g^{-1}}$ is the electron scattering opacity, $l_E \equiv L/L_E$ is the ratio of the luminosity to the Eddington value, and $\epsilon \equiv L/(\dot{M}_d c^2) \equiv 0.1\epsilon_{0.1}$ is the radiative efficiency of the accretion disc.

To have these accretion rates be the residual mass flux from a star-formation-supported disc on 100 pc scales requires star formation rates that are $\sim 10^6$ times larger (equation 20), which are unrealistic. However, this estimate assumes that the star formation time is short compared to the radial advection time in the disc. Thompson et al. (2005) show that the opposite may be true in quasars, in which case the initial outer-disc mass feeding rate (or equivalently the star formation rate) need only be factors of ∼100 greater.

There may be additional complications, such as the supply of some very low angular momentum material to the inner disc, which is not gravitationally unstable.

The problem of the mass supply to quasars requires further work. The possibility remains that self-gravity in a disc is so severe that black holes need to be supplied at much smaller radii, perhaps through stellar disruption events (e.g. Hills 1975; Goodman 2003).

In lower-luminosity AGN, such as Seyfert galaxies and low-ionization nuclear emission regions (LINERs), there is much evidence for nuclear star formation (e.g. Terlevich & Melnick 1985). These accretion discs are also expected to be self-gravitating, and there is some observational evidence to support this (e.g. Kondratko, Greenhill & Moran 2005). While the fuelling of these nuclei may have little to do with Bondi accretion (because of the presence of large amounts of cold gas in an extended disc), the model of an accretion disc supported by stellar feedback (Section 4) probably is relevant, predicting a general correlation of star formation and AGN activity.

6 CONCLUSIONS

Observations of the central regions of elliptical galaxies suggest that supermassive black holes that reside there are fed at the Bondi rate, of a few hundredths to a few tenths of a solar mass per year. For standard thin disc accretion, these black holes are very underluminous given this mass accretion rate. Previously, alternative accretion models invoking inefficient radiation of gravitational energy, radially dependent accretion rates, and outflows have been proposed as solutions to this problem. In this paper we have suggested an additional process, likely to be occurring, which can be of comparable or potentially greater importance: Bondi-fed star-forming discs. Accretion occurs at the Bondi rate, but because the disc that inevitably forms inside the Bondi radius is very gravitationally unstable, most of the mass flux reaching it turns into stars. Stellar feedback is relatively weak, so that only a small residual amount of gas can be stabilized to reach the black hole at the centre of the disc. Our estimates of this residual mass flux are uncertain, but appear to be potentially small enough that standard thin disc accretion models can avoid the constraints imposed by the observed low luminosities. ADAFs, CDAFs, or other models with outflows may still apply in the very inner regions, but with a smaller outer boundary value for their inward mass flux.

Our model is highly simplified: we mostly consider only a single typical scale for the disc radius, whereas in reality gas will reach the disc at a range of radii depending on the initial angular momentum distribution; our quantitative estimates depend on the assumed stellar IMF; we have adopted the standard alpha approximation for viscous stresses in the disc; and we have mostly considered only thermal pressure support in the disc, ignoring magnetic pressure.

As a result of these simplifications, it is vital to compare the model as closely as possible to real systems, the best observed being M87.

We have described how the observational data on the nucleus of this galaxy are consistent with the possibility that more than half of the Bondi accretion flow is depleted into stars: a thin disc of gas and dust, apparently self-gravitating, is seen inside the Bondi radius and the inferred star formation rate from the Hα emission is consistent with the Bondi accretion rate. This conclusion is not contradicted even when current estimates of the mechanical luminosity of the M87 jet are incorporated into the residual central engine accretion rates. Additional observations to detect molecular gas scales inside the Bondi radius are now feasible with the SMA. This instrument is also capable of detecting continuum emission from dust heated by the expected star formation.

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