Quasar feedback: more bang for your buck

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ABSTRACT

We propose a 'two-stage' model for the effects of feedback from a bright quasar on the cold gas in a galaxy. It is difficult for winds or other forms of feedback from near the accretion disc to directly impact (let alone blow out of the galaxy) dense molecular clouds at ∼kpc. However, if such feedback can drive a weak wind or outflow in the hot, diffuse interstellar medium (a relatively 'easy' task), then in the wake of such an outflow passing over a cold cloud, a combination of instabilities and simple pressure gradients will drive the cloud material to effectively expand in the direction perpendicular to the incident outflow. This shredding/expansion (and the corresponding decrease in density) may alone be enough to substantially suppress star formation in the host. Moreover, such expansion, by even a relatively small factor, dramatically increases the effective cross-section of the cloud material and makes it much more susceptible to both ionization and momentum coupling from absorption of the incident quasar radiation field. We show that even a moderate effect of this nature can dramatically alter the ability of clouds at large radii to be fully ionized and driven into a secondary outflow by radiation pressure. Since the amount of momentum and volume which can be ionized by observed quasar radiation field is more than sufficient to affect the entire cold gas supply once it has been altered in this manner (and the 'initial' feedback need only initiate a moderate wind in the low-density hot gas), this reduces by an order of magnitude the required energy budget for feedback to affect a host galaxy. Instead of ∼5 per cent of the radiated energy (∼100 per cent momentum) needed if the initial feedback must directly heat or 'blow out' the galactic gas, if only ∼0.5 per cent of the luminosity (∼10 per cent momentum) can couple to drive the initial hot outflow, this mechanism could be efficient. This amounts to hot gas outflow rates from near the accretion disc of only ∼5–10 per cent of the black hole accretion rate.

Key words: galaxies: active – galaxies: evolution – quasars: general – cosmology: theory.

1 INTRODUCTION

Observations have established that the masses of supermassive black holes (BHs) are tightly correlated with various host galaxy properties (Magorrian et al. 1998; Ferrarese & Merritt 2000; Gebhardt et al. 2000; Aller & Richstone 2007; Hopkins et al. 2007b). Together with constraints indicating that most of the BH mass is assembled in optically bright quasar[1] phases (Soltan 1982; Salucci et al. 1999; Yu & Tremaine 2002; Hopkins, Narayan & Hernquist 2006b), this has led to the development of models where feedback processes from accretion self-regulated BH growth at a critical mass (Silk & Rees

1998; Di Matteo, Springel & Hernquist 2005; Murray, Quataert & Thompson 2005). Gas inflows triggered by some process fuel rapid BH growth, until feedback begins to expel nearby gas and dust. This 'blowout' results in a short-lived, bright optical quasar that, having expelled its fuel supply, fades and leaves a remnant on the observed BH–host correlations (Hopkins et al. 2005a,c). These scenarios have been able to explain many quasar observables, including luminosity functions, lifetimes and BH mass functions (Menci et al. 2003; Hopkins et al. 2005b, 2006c, 2008a, 2009; Lapi et al. 2006; Volonteri, Salvaterra & Haardt 2006; Somerville et al. 2008; Tortora et al. 2009).

It is much less clear, however, what the impact of whatever feedback regulates BH growth will be on the host galaxy. In models, such feedback is invoked to explain the rapid 'quenching' of star formation and sustained lack of cooling in massive galaxies (Granato et al. 2004; Scannapieco & Oh 2004; Croton et al. 2006; Antonuccio-Delogu & Silk 2008; Hopkins et al. 2008b). The argument in the

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[1] In this paper, we use the term 'quasar' loosely as a proxy for high-Eddington ratio activity, rather than as a reference to specific optical properties.

models is that, under various simple assumptions, if sufficient energy or momentum is injected into the interstellar medium (ISM) near the BH on a time-scale short enough to halt accretion, then it will yield a supersonic pressure or momentum-driven outflow that propagates to large scales (see e.g. Monaco & Fontanot 2005; Hopkins et al. 2006a; Shin, Ostriker & Ciotti 2009).

However, the actual mechanisms of feedback and physics of the ISM relevant for this remain highly uncertain. Highly energetic outflows are associated with bright quasars (for a review, see Veilleux, Cecil & Bland-Hawthorn 2005); these range from intense winds ($v \sim 10^4 \, \mathrm{km \, s^{-1}}$) associated with the central engine seen in the broad emission-line regions and broad absorption line (BAL) quasars (e.g. Weymann, Carswell & Smith 1981) to more moderate outflows ($v \sim 10^2$–$10^3 \, \mathrm{km \, s^{-1}}$) associated with the narrow-line region and the 'warm absorber' (Laor et al. 1997; Crenshaw et al. 2000) as well as with small-scale quasar absorption and occultation systems (e.g. McKernan & Yaqoob 1998; Miller, Turner & Reeves 2008; Turner et al. 2008). Indeed, high-velocity winds driven near the accretion disc are theoretically hard to avoid (see e.g. Blandford & Payne 1982; Begelman 1985; Konigl & Kartje 1994; Elvis 2000; Proga 2000, 2007). However, these are probably tenuous, with an initial mass loading $\lesssim \dot{M}_{\mathrm{BH}}$ (although in at least some cases, these outflows are extremely dense and might have much higher mass-loading factors; see e.g. Hall & Hutsemèkers 2003; Hall et al. 2007). It is not clear whether such 'hot' outflows could efficiently entrain gas at larger radii. If most of the gas mass of the galaxy is at some appreciable fraction of the galaxy effective radius R_{e} and in the form of cold, dense giant molecular clouds, then it is difficult to imagine such a diffuse wind directly 'launching' the clouds out of the galaxy. It remains unclear whether, in fact, the momentum associated with the winds that *are* known to emanate from the central engine of a quasar is sufficient to unbind the cold gas in the host (see e.g. Baum & McCarthy 2000; de Kool et al. 2001, 2002; Steenbrugge et al. 2005; Gabel, Arav & Kim 2006; Holt et al. 2006; Batcheldor et al. 2007; Ganguly et al. 2007; Krongold et al. 2007; McKernan, Yaqoob & Reynolds 2007; Tremonti, Moustakas & Diamond-Stanic 2007; Prochaska & Hennawi 2009).

In this paper, we argue that it is not necessary that the small-scale, high-velocity active galactic nucleus (AGN) outflows directly entrain any cold gas at scales $\sim R_{\mathrm{e}}$. Rather, so long as these are sufficient to drive a significant wind in the 'hot' diffuse ISM, then clouds will be effectively destroyed or deformed and 'secondary' feedback mechanisms – namely the radiative effects of dust absorption and ionization – will be able to act efficiently on the cold gas at large scales. This will effectively terminate star formation on a short time-scale, with greatly reduced energy/momentum requirements for the 'initial' outflow drivers.

2 RADIATIVE FEEDBACK IN THE PRESENCE OF HOT OUTFLOWS

Consider a typical galaxy, where the ISM gas is composed of a mix of diffuse warm/hot gas and cold clouds.[2] Radiation will always

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[2] By 'diffuse' ISM, we refer to the warm/hot phases of the ISM, with effective temperatures $T \sim 10^4$–10^6 K and densities $n_{\mathrm{H}} \sim 10^{-2}$–$10^0 \, \mathrm{cm^{-3}}$. The 'clouds', on the other hand, represent the cold phase (but *not* collapsing molecular cloud cores) with $T \lesssim 100$ K molecular gas and $n_{\mathrm{H}} \sim 10^2$–$10^4 \, \mathrm{cm^{-3}}$. The volume filling factor of clouds is small, but by mass, they dominate the total gas mass; the hot gas with filling factor near unity has a typical mass fraction $f_{\mathrm{hot}} \lesssim 0.1$ (see e.g. McKee & Ostriker 1977).

act on the cold clouds (in the form of ionization and momentum injection from absorbed photons), but they may be too dense and self-shielding to be significantly affected. In such a case, one could invoke a blastwave or cold shell, driven by AGN feedback on small scales, to entrain this material. Various models and simulations have shown that if feedback needs to directly launch a blastwave in *both* the hot and cold gas together (sufficient to entrain most of the galactic gas), then an efficiency $\eta_{\mathrm{E}} \sim 0.05$ is the relevant value (for energy injection where the feedback $\dot{E} = \eta_{\mathrm{E}} L = \eta_{\mathrm{E}} \epsilon_{\mathrm{r}} \dot{M}_{\mathrm{BH}} c^2$; $\epsilon_{\mathrm{r}} \sim 0.1$ is the radiative efficiency). If the outflow is instead momentum driven, the relevant value for driving hot+cold phases is $\eta_{\mathrm{p}} \sim 1 (\dot{p} = \eta_{\mathrm{p}} L/c)$.

If, however, the 'initial' feedback from the central source need only drive a wind in the low-density hot gas, and does not necessarily directly entrain the cold clouds, then the energy required will be much less. In both cases above, the efficiency needed simply scales linearly with the mass of material to be driven (scaling with its binding energy or momentum, respectively). So, if a fraction f_{hot} of the gas is in the hot, diffuse ISM, and only that needs to be initially driven, we obtain

$$\eta_{\mathrm{E}} = \frac{\dot{E}_{\mathrm{fb}}}{L} \sim 0.05 \, f_{\mathrm{hot}} \tag{1}$$

for energy-driven and

$$\eta_{\mathrm{p}} = \frac{\dot{p}_{\mathrm{fb}}}{L/c} \sim f_{\mathrm{hot}} \tag{2}$$

for momentum-driven outflows. For typical $f_{\mathrm{hot}} \lesssim 0.1$ (e.g. Blitz 1993), this implies an order-of-magnitude reduction in the necessary feedback input for 'interesting'' behaviour. And in detail, since the hot-phase gas is already virialized (rather than in e.g. a cold disc), the efficiency gains may be even higher. The question is then, if most of the mass is in cold clouds and only the hot gas is affected by this initial outflow, will any interesting behaviour result? Consider a cold molecular cloud in a galaxy. The cloud has a mass M_{c} and an initial ('equilibrium') effective spherical radius R_0. In the observed typical ISM, these are related by

$$M_{\mathrm{c}} \sim 300 \, \mathrm{M_\odot} \, R_{0,\mathrm{pc}}^2 \tag{3}$$

(where $R_{0,\mathrm{pc}} \equiv R_0/1 \, \mathrm{pc}$; Blitz & Rosolowsky 2006). Since we will later allow the cloud to stretch or deform, define the instantaneous radius of the cloud as R_{c}. The cloud resides at a spherical distance r from the centre of a Hernquist (1990) profile bulge of total mass M_{bul} and characteristic scalelength R_{eff}, which hosts a BH on the characteristic BH–host relations, $M_{\mathrm{BH}} = \mu_{\mathrm{BH}} M_{\mathrm{bul}} (\mu_{\mathrm{BH}} \approx 0.0014$; Häring & Rix 2004). The BH radiates with a luminosity

$$L_{\mathrm{QSO}} = \dot{m} \, L_{\mathrm{Edd}}, \tag{4}$$

where $\dot{m}$ is the dimensionless Eddington ratio and $L_{\mathrm{Edd}} = 1.3 \, M_{\mathrm{BH,8}} \times 10^{46} \, \mathrm{erg \, s^{-1}}$ is the Eddington luminosity ($M_{\mathrm{BH,8}} = M_{\mathrm{BH}}/10^8 \, \mathrm{M_\odot}$).

We are interested in the case where the cold clouds are embedded in some kind of hot outflow generated by the 'primary' AGN feedback. Specifically, assume that the quasar somehow succeeds in driving an outflow through the diffuse warm/hot ISM; to be conservative, the outflow can be tenuous and we will assume that the outflow 'impacting' the cloud carries negligible momentum compared to the binding momentum of the cloud. In other words, assume that at some smaller scale, a wind or outflow is generated sufficient to sweep up the tenuous, diffuse ISM at large radii, but insufficient to affect cold dense clouds, which contain most of the ISM gas mass.

This problem of the survival of cold clouds in a post-shock hot medium is well studied in the context of star formation and supernova feedback (see e.g. Klein, McKee & Colella 1994, and references therein). In general, the collision of a shock or wind with velocity v_s with a cloud of initial characteristic (quasi-spherical) radius R_0 and density contrast χ (ratio of cloud density to external medium density $\chi \equiv n_c/n_0$) will launch secondary shocks within the cloud with velocity $v_s/\chi^{1/2}$. This defines a 'cloud-crushing' time-scale $t_{cc} = \chi^{1/2} R_0/v_s$. In the simple case of a pure hydrodynamic strong shock, if t_{cc} is much less than the characteristic time-scales for the density to change behind the shock and $v_s/\chi^{1/2}$ is comparable to or larger than the characteristic internal velocities of the cloud, then the cloud will be stretched and 'shredded' by a combination of Rayleigh–Taylor and Kelvin–Helmholtz instabilities on a time-scale approximately a few t_{cc} (see e.g. Klein et al. 1994; Xu & Stone 1995; Fragile et al. 2004; Orlando et al. 2005; Nakamura et al. 2006). Given the definitions above, for $v_s \sim v_{esc}$, t_{cc} is much less than the dynamical time-scales of interest at all the spatial scales of interest ($t_{cc} \ll 10^7$ yr for all $r \gtrsim \chi^{1/2} R_0$ – i.e. for clouds not in the very nuclear regions).

The material being mixed off of the surface of the cloud from these instabilities expands into and mixes with the low-pressure zones created by the passage of the shock on the sides of the cloud. This leads to an effective net expansion of the cloud by a factor of $\sim \chi^{1/2}$ in radius in the perpendicular shock direction. Eventually, despite the initial compression, reflection shocks lead to a expansion by a factor of ~ 2 in the parallel shock direction, bringing the original cloud material into an effective density and pressure equilibrium with the external medium. The surface area of the cloud can increase dramatically; for our purposes, we are interested in the effective cross-section the cloud presents to the perpendicular shock direction. The radii defined above should be thought of in this manner: the initial cloud cross-section to the hot shock is πR_0^2; post-shock, the effective cross-section owing to this expansion and equilibration is $\pi R_c^2 \sim \chi \pi R_0^2$.

We illustrate this behaviour with a simple toy model system in Fig. 1. Specifically, we show an example of a hydrodynamic simulation of an idealized system, using the ZEUS code (Stone & Norman 1992a,b). The initial conditions consist of a Plummer sphere cloud embedded in a uniform background, with density contrast of $\chi = 100$ (peak density of the cloud relative to background), with the initial system in pressure equilibrium (uniform pressure) and periodic

boundary conditions in a large grid. At time $t = t_S$, the low-density material is rapidly accelerated into a mach ~ 2 wind. Colour encodes the gas density, from black (the arbitrary background density) to red (the initial maximum). Note that the example shown is purely for illustrative purposes – we do not include many possible complexities, such as gas cooling, star formation or magnetic fields. The behaviour of clouds in response to outflows with such sophistications has been extensively studied in the references above, and more detailed extensions of such simulations to the regime of interest here will be the subject of future work. Nevertheless, this simple experiment illustrates much of the important qualitative behaviour.

The qualitative behaviour we care about – the mixing/ stretching/deformation of the cloud in the perpendicular shock direction leading to an increase in the effective cross-section of the cloud – is in fact quite general. Simulations have shown that the same instabilities operate regardless of whether the 'hot outflow' is a strong shock, weak shock or wind (since we assume the hot material is being unbound in this wind, it cannot be substantially subsonic). The time-scale of cloud expansion increases by a factor of a few in the weaker wind case, but it is still much less than the relevant local galactic dynamical times (Klein et al. 1994; Jones, Ryu & Tregillis 1996). The process is also similar in the case of a cloud being impacted by AGN jets, despite the different densities, temperatures and magnetic field states associated with jets and 'bubbles' (see Krause & Alexander 2007; Antonuccio-Delogu & Silk 2008). A cloud could in principle be stabilized against such instabilities by being strongly magnetically dominated (Mac Low et al. 1994; Jun & Jones 1999; Fragile et al. 2005; Shin, Stone & Snyder 2008). However, in this limit, as the hot outflow sweeps up the material, the pressure of the diffuse ISM trailing the outflow will decline as a steep function of time t/t_{cc} (Ostriker & McKee 1988). Since, in this limit, the cloud is then overpressurized, it will expand isothermally as the exterior post-shock pressure drops (the free expansion/equilibration time of the cloud being short compared to the other time-scales of interest). Because this stops when the system is equilibrated, the 'effective' net expansion of R_c is the same as in the hydrodynamic shock case, even though the details are quite different.

If nothing more were to happen to the cloud, this would only suppress star formation for a short time. The cooling instabilities that produced the cloud in the first place would operate. In the 'typical' ISM, clouds mix in the wake of stellar or supernovae-driven
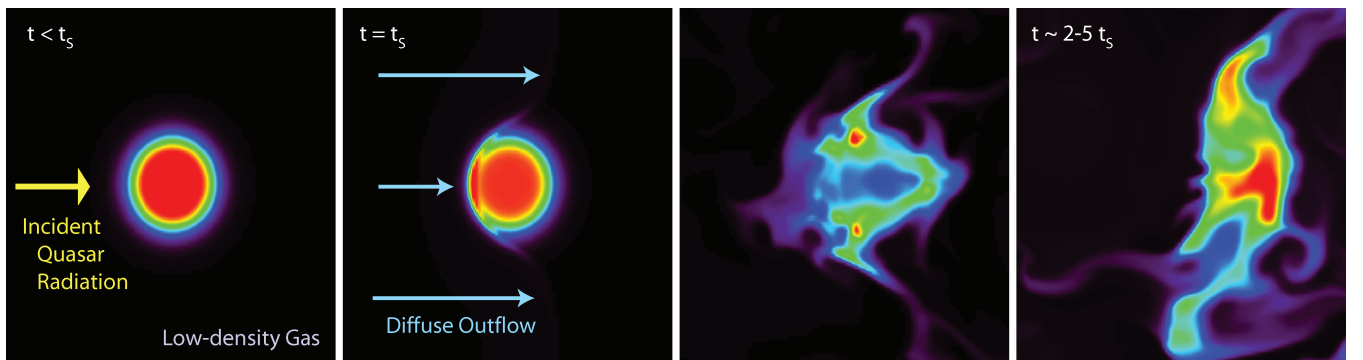

Figure 1. An illustration from a simple simulation of the effects of a hot outflow on radiation feedback. Left: initial dense cloud in pressure equilibrium with the diffuse ISM. Ionization and momentum flux from the quasar is negligible, because the effective surface area for absorption is small. Centre: at $t = t_S$, an outflow generated in the hot/diffuse ISM hits the cloud. Ambient pressure drops in the wake (low-pressure regions form on the perpendicular sides of the cloud). A mix of instabilities shortly causes the cloud to mix in the perpendicular direction. Right: after a few cloud crossing times, the mixing has increased the effective cross-section by a factor $(R_c/R_0)^2 \gtrsim 10$ in the perpendicular direction. The cloud may now (with lower density and higher cross-section to absorb quasar photon momentum) be vulnerable to radiative feedback, and will be accelerated to $v \sim v_{esc}$.

outflows until they reach equilibrium with the ISM and recool into new clouds.

However, we are interested in all of this occurring in the background of a luminous AGN, which will both ionize and exert a radiation pressure force. The cloud – especially a realistic cloud with a large dust mass and corresponding opacity – is optically thick[3] to the quasar radiation, with an effective cross-section $\Delta\Omega \sim (\pi R_c^2)/(4\pi r^2)$. There is therefore an inescapable deposition of photon momentum from the radiation field with a deposition rate $\dot{p}_{rad} = L_{abs}/c$, where $L_{abs} = L_{QSO}\,\Delta\Omega$. Comparing this to the gravitational force $F_{grav} = -M_c\,\partial\phi/\partial r$ defines an effective Eddington limit for the cloud: if the absorbed flux exceeds some limit, the cloud will be unbound (equivalently, the absorbed momentum in a single dynamical time, over which the cloud could redistribute that momentum, will exceed the cloud binding momentum $\sim M_c\,v_{esc}$). This limit is when the two are equal; equivalently

$$\frac{L_{QSO}}{c}\,\frac{\pi R_c^2}{4\pi r^2} = M_c\,\frac{G\,M_{bul}}{(r + R_{eff})^2}\,. \tag{5}$$

Assuming that the cloud lies on the observed size–mass relation (equation 3) and that the galaxy lies on the M_{BH}–M_{bul} relation, this reduces to the criterion for unbinding the cloud:

$$\left(\frac{R_c}{R_0}\right) \gtrsim 2.7\,\dot{m}^{-1/2}\left(\frac{r}{r + R_{eff}}\right). \tag{6}$$

In other words, a cloud on the 'normal' size–mass relation at large radii $r \sim R_{eff}$ is sufficiently dense and sufficiently high column density to avoid being unbound by radiation pressure. However, if the effective size of the cloud (the effective coupling surface area) could be increased by a factor of a couple, or the effective column lowered, the cloud would rapidly be unbound by the incident radiation field momentum. This condition is easily satisfied in post-shock clouds.

Fig. 1 also includes and illustrates this effect. Specifically, we include a very simple, time-independent momentum deposition rate in the 'facing' cells to the incident radiation field, which we simply approximate as all cells with a density above $\gtrsim$ three times the initial background density but with no cell at $x < x_i$ above this density (i.e. implicitly assuming that such a cell would shield the cells 'behind' it with respect to the incident radiation). The magnitude of this is initialized such that at $t < t_S$ the 'total' deposition rate over the surface of the cloud is equal to a small fraction of the 'binding' momentum over the dynamical time (assuming $v_s \sim V_c$), $\sim 0.01\,M_c\,v_s/(R_c/v_s)$. But the details make little qualitative difference to the global acceleration.

A similar effect pertains to the ionization of the cloud (although this is not explicitly included in Fig. 1). Ignoring geometric effects of photon diffusion, the volume of a cloud of mean density n_c ionized is $V_{ion} = \dot{N}/n_c^2\,\beta$, where $\beta \approx 2 \times 10^{-13}$ cm^3 s^{-1} is the recombination coefficient for gas at the temperature for hydrogen ionization ($T_e \sim 10^4$ K) and $\dot{N}$ is the rate at which ionizing photons hit the cloud.

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[3] For any incident spectrum where a significant fraction of the incident energy is in the optical/ultraviolet (UV) or higher wavelengths, the effective optical depth from dust within the cloud will be $\tau \sim 1$ (Murray, Quataert & Thompson 2005; Thompson, Quataert & Murray 2005) – this is simply a statement that the clouds will be optically thick to some portion of that spectral energy distribution [SED; whether ionizing photons in the far-UV, or, if the cloud is not yet ionized, then there is dust which will absorb in the optical and infrared (IR)] and reradiate that energy. The cloud can effectively be thought of as a single absorbing 'mega-grain' with effective cross-section $\sim\pi R_c^2$.

The total rate of production of ionizing photons from the quasar is $\dot{N}_Q = \lambda L/h\nu_{912}$ ($\lambda \approx 0.07$ comes from a proper integration over the quasar spectrum; here from Hopkins et al. 2007a), and a fraction $\Delta\Omega$ are incident on the cloud. Together with the typical values above, this implies that clouds will be ionized to a depth $h_{ionized}$:

$$\frac{h_{ionized}}{R_c} \approx 10^{-3}\,\dot{m}\,R_{0,pc}\left(\frac{r + R_{eff}}{r}\right)^2\left(\frac{R_c}{R_0}\right)^5. \tag{7}$$

Given the cloud size–mass relation, this is equivalent to the statement that all clouds below a mass $M_c \lesssim 10^8\,M_\odot\,(R_c/R_0)^{-10}$ will be self-shielded at $r \gtrsim R_{eff}$. For typical clouds, the depth ionized is clearly quite small; but there is a steep dependence on cloud radius. As R_c increases, the ionized depth increases by a factor $\propto R_c^2$ owing to the increased photon capture cross-section and a factor $\propto R_c^3$ owing to the decreased density lowering the recombination rate.

3 IMPLICATIONS IN A GLOBAL FEEDBACK SCENARIO

There are many caveats to the simplified derivations above: clouds have some size and mass spectrum, and are distributed at various radii, with the background quasar changing in time. Nevertheless, embedding this in more detailed models for AGN feedback, the results are interesting.

Consider an $\sim L_*$ bulge with $M_{bul} = 10^{11}\,M_\odot$ and $M_{BH} = 1.4 \times 10^8\,M_\odot$, with a Hernquist (1990) density profile and scale radius $R_{eff} = 4$ kpc. Assume gas traces stars (initially), with mass fraction $f_{gas} = 0.1$, and that 90 per cent of the gas is in cold clouds while 10 per cent is in a hot diffuse phase (which we assume is in hydrostatic equilibrium). This yields a density profile of hot or cold gas of

$$\rho_i = f_i\,f_{gas}\,\frac{M_{bul}}{2\pi}\,\frac{R_e}{R\,(R + R_e)^3}, \tag{8}$$

where f_i represents the fraction in either the hot or the cold phase, ($f_{hot} = 0.1$, $f_{cold} = 0.9$).

At a given radius (within some small radial annulus), the cold gas (with mean volume density given above) is assumed to be locked into cold clouds with a small volume filling factor. The clouds are initially placed on the observed size–mass relation (equation 3; determining an initial radius R_0 for each cloud of mass M_c), and are distributed in mass according to the observed mass spectrum $dN/dM_c \propto M_c^{-1.8}$ (up to a maximum $M_c = 10^7\,M_\odot$; see Rosolowsky 2005, and references therein). Given some total cold gas mass ρ_{cold} per unit volume in an annulus, this mass spectrum, integrated from arbitrarily small cloud mass (it makes no difference if we adopt some lower mass cut-off) to the maximum M_c, must integrate to ρ_{cold} per unit density; i.e. this determines the number density n_c of clouds at each galaxy radius R and mass interval $M_c \to M_c + dM_c$ [and corresponding initial cloud radius $R_0(M_c)$].

At a time $t = 0$, we assume that the BH 'turns on', radiating (initially) at the Eddington limit. We allow it to drive a shock/wind through the diffuse ISM according to the analytic solutions derived in Hopkins & Hernquist (2006). In these models, the AGN is assumed to drive a simple, Sedov–Taylor-type outflow via any 'small-scale' feedback channel. The detailed behaviour is derived and compared with hydrodynamic simulations in Hopkins & Hernquist (2006), but can be simply summarized as follows: the BH, once on, couples a fraction $\eta \approx 0.05\,f_{hot}$ of its luminous energy to the diffuse-phase gas in its vicinity. Because the spatial and time-scales in the vicinity of the BH are small compared to the rest of the galaxy, this appears to the galaxy as a point-like energy injection in a hot medium. The result is therefore a roughly self-similar (power law) Sedov–Taylor-type outflow.
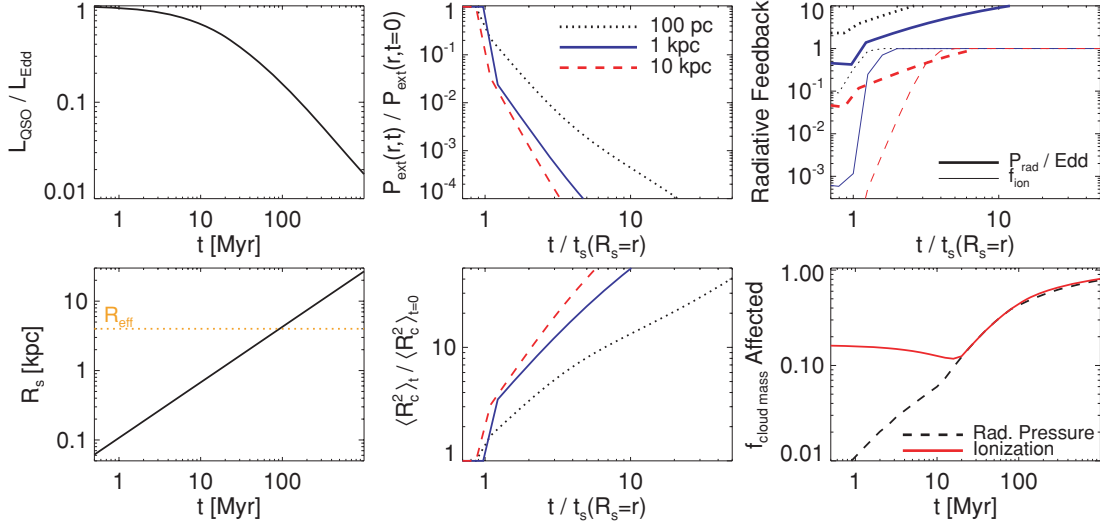
Figure 2. Effects of a two-stage feedback model in the wake of a quasar-induced hot outflow, with parameters of a typical ∼L_* quasar. Top left: quasar Eddington ratio as a function of time, after feedback from the quasar begins to drive an outflow in the diffuse ISM (time $t = 0$). Bottom left: radius of the hot/diffuse outflow as a function of time (compare the galaxy effective radius $R_{\rm eff}$). Top centre: change in pressure of the (post-shock) diffuse ISM at a given radius r from the BH, at time $t/t_{\rm s}$ ($t_{\rm s}$ is the time when the hot outflow first reaches r). We show three different radii: $r = 100$ pc, 1 and 10 kpc. Bottom centre: change in effective cross-section of a typical ($R_0 \sim 1$ pc) cloud at each r versus time $t/t_{\rm s}$. Top right: ratio of the radiation pressure force to the Eddington limit for expulsion of each cloud (thick; $P_{\rm rad}/{\rm Edd}$) and fraction of the cloud ionized (thin; $f_{\rm ion}$). Bottom right: total fraction of cloud mass that can be expelled by radiation pressure or ionized (we consider each separately), as a function of time. The curves reflect the integral over the observed mass spectrum of clouds at each radius, integrated over all radii following the galaxy density profile.

Fig. 2 illustrates some of the basic behaviours of this outflow. The energy injection leads to a shock that expands outwards with radius $R_{\rm s}(t)$, in an approximately power-law fashion as $R_{\rm s} \propto t^\alpha$, where α is a function of the local density profile, gas equation of state, and is coupled (weakly) to the declining energy injection rate of the BH (Fig. 2, bottom left). For typical galaxy density profiles, a $\gamma = 5/3$ gas, and the conditions assumed here, Hopkins & Hernquist (2006) show $\alpha \approx 4/5$. In the wake of the expanding bubble/shock, the post-shock density drops in a related power law manner. Since the spherical accretion rate on to a BH scales roughly $\propto \rho$ (in Bondi–Hoyle accretion; or similarly $\propto \Sigma$ for viscous accretion from a disc), the accretion rate, and hence luminosity $L = 0.1 \, \dot{M}_{\rm BH} c^2$, will decay as well (Fig. 2, top left). Again, we refer to the full derivation in Hopkins & Hernquist (2006) for details, but the self-consistent solution derived therein can be approximated as $L \propto [1 + t/t_{\rm Q}]^{-\eta_L}$, with $\eta_L \sim 1.5$ and $t_{\rm Q} \sim 1 - 5 \times 10^7$ yr for the parameters here (consistent with various observational constraints; see Martini 2004; Yu, Lu & Kauffmann 2005; Hopkins & Hernquist 2009). The value of η_L follows from α and generic behaviour of Sedov–Taylor post-shock gas and $t_{\rm Q}$ is simply related, modulo appropriate numerical coefficients, to the local dynamical time near the BH radius of influence.

The shock therefore crosses a radius r at a time $t_{\rm s}$ (when $R_{\rm s} = r$). In the wake of the shock, the post-shock ambient pressure $P_{\rm ext}$ will drop, reflecting the density decline from material being blown out (Fig. 2, top centre). Again, this approximately follows a standard decline in the wake of a Sedov–Taylor blastwave; the exact solution for the pressure internal to the blastwave under the conditions here must be obtained numerically, but Ostriker & McKee (1988) show that it can be approximated as a double power law. Roughly speaking, there is a rapid drop in pressure in the immediate post-shock region, as the thin shell at the front of the blastwave clears the diffuse material away from the region, and the pressure declines as a steep power $P \propto [(t - t_{\rm s})/t_{\rm c}]^{-\beta}$, where $t_{\rm c}$ is the crossing time

of the shock relative to the cloud ($\sim R_{\rm c}/v_{\rm s}$) and $\beta \sim 3$–5 (the exact index depends on the local density profile slope and rate of decay of the driving source, so is not the same at all radii). This is followed by a more gradual decline, as the diffuse medium internal to the shock relaxes, is heated, and expands, with $P \propto [(t - t_{\rm s})/t_{\rm s}]^{-\beta'}$ and $\beta' \sim 2$ (again, for the detailed derivation of the double power-law structure in the wake of the blastwave, for the conditions of the feedback-driven blastwaves considered here, we refer to Hopkins & Hernquist 2006).

In the wake of the shock, with a rapidly declining background pressure and density, the cloud will be mixed and effectively increase its surface area. We can solve for the behaviour of each cloud – at least the key parameter of interest, the effective radius of the cloud in the direction perpendicular to the shock ($R_{\rm c}$) as a function of time, according to the approximations in Klein et al. (1994) in the wake of the shock defined above. If the cloud could somehow resist shredding (say, via sufficient magnetic field or turbulent support), this would be trivial: the cloud (initial pressure equilibrium) would expand isothermally (e.g. conserving total magnetic energy) as it is now overpressurized, such that pressure equilibrium would be conserved. For a background pressure declining with some power law $P \propto (t/t_{\rm s})^{-\beta}$, this implies expansion of the cloud with $R_{\rm c}/R_0 = (t/t_{\rm s})^{\beta/3}$ [generally, $R_{\rm c}/R_0 = (P_{\rm c}/P_0)^{-1/3}$, for isothermal expansion]. If the cloud were supported by thermal energy with no new inputs, this index would be slightly modified [for adiabatic expansion, $R_{\rm c}/R_0 = (P_{\rm c}/P_0)^{1/5}$] but the behaviour is qualitatively similar. Technically, this approximation assumes that the time for the cloud to equilibrate is short relative to the time-scale on which the background is changing, but this is easily satisfied. The cloud expands/equilibrates on its internal crossing time, given by the effective sound speed as $\sim R_{\rm c}/c_{\rm s,eff}$; for quasi-virial clouds this is simply the dynamical time $1/\sqrt{G\rho}$, where the effective density ρ follows from the observed size–mass relation (equation 3). Using the observed values, this gives a time-scale of ~ 0.1–1×10^6 yr (for

cloud sizes ~ 0.1–$10\,\mathrm{pc}$). Compare this to the characteristic time-scale for the evolution of the background, a few t_s (itself of order t_{dyn}, the galaxy dynamical time at R). For a typical $\sigma \sim 200\,\mathrm{km\,s^{-1}}$ spheroid, $t_s \gg R_c/c_{s,\mathrm{eff}}$ for all radii $\gtrsim 100\,\mathrm{pc}$ – in other words, this condition is easily satisfied at the radii $\sim R_e$, which contain most of the mass of the galaxy. Fig. 2 (bottom centre) shows how the clouds will expand in effective cross-section ($\sim R_c^2$), relative to their initial sizes, given this declining background pressure for the simple isothermal case.

For the more complex case of cloud shredding, it turns out that, in aggregate, a similar scaling obtains. The characteristic time for the cloud to effectively be mixed via instabilities and so effectively increase its cross-section is a few cloud-crushing time-scales $t_{cc} = \chi^{1/2}\,R_0/v_s$. However, since the shock velocity is of order the galaxy escape velocity for the 'interesting' diffuse outflows considered here ($v_s \sim$ a few σ), and typical $\sigma \sim 200\,\mathrm{km\,s^{-1}}$, compared with a typical effective sound speed of a virialized cloud ~ 1–$10\,\mathrm{km\,s^{-1}}$, this time is almost always much shorter than (or at least comparable to) the cloud dynamical time. Following Klein et al. (1994), cloud shredding will equilibrate when the system expands by a factor $\sim \chi^{1/2}$ in the perpendicular direction (and a small, $\sim$constant factor in the parallel direction), where χ is the initial density contrast; in other words, until the *effective* density and pressure drop to approximate equilibrium with the background. Thus, for time-scales $\sim t_s$ over which the background is evolving, long compared to the cloud-crushing time, we can consider the systems to effectively expand with an average effective radius scaling in the same way in equilibrium with the external/hot medium background pressure (i.e. similar effective net expansion, averaged over these time-scales, as in the isothermal expansion case).

We then solve for the behaviour of each cloud in the wake of this hot outflow, according to the approximations in Klein et al. (1994) and Section 2 (Fig. 2, top right). In particular, given the time-evolution in R_c/R_0 shown in Fig. 2, we use the scalings derived in Section 2 to estimate the fraction of the cloud (at some initial radius r) which will be ionized (i.e. $f_{\mathrm{ion}} = h_{\mathrm{ionized}}/R_c$ from equation 7; where the AGN accretion rate $\dot{m}$ and cloud expansion R_c/R_0 are given as a function of time above, the initial cloud radius r is one of those specified in the figure, and we chose a representative initial cloud with radius $R_{0,\mathrm{pc}} = 1$ for illustrative purposes). We also show the relative strength of radiation pressure on the cloud, i.e. the radiation pressure relative to the local Eddington limit (that which would unbind the cloud), $P_{\mathrm{rad}}/\mathrm{Edd} = 0.14\,\dot{m}\,(R_c/R_0)^2\,[(r+R_e)/r]^2$ (re-arranging equations 5 and 6, where again $\dot{m}$, R_c/R_0 and r for the clouds is given). Since these both scale steeply with the increasing effective cloud size ($\propto R_c^5$ and R_c^2, respectively), both increase rapidly in time. At early times, only the clouds within a narrow region $\sim 100\,\mathrm{pc}$ around the quasar are efficiently ionized, and the effects of radiation pressure are weak – similar to what is observed in the narrow-line regions of AGN (Crenshaw et al. 2000; Rice et al. 2006). This changes rapidly in the wake of the outflow at large radii – the deformation induced makes the clouds vulnerable to ionization and radiative momentum driving.

Integrating over the entire cloud population and galaxy mass, i.e. integrating over the initial cloud mass or size (R_0) spectrum at each galactic radius r and then over the total cold gas density at each radius r, we obtain the *total* fraction of cloud mass that can be ionized or effectively accelerated by radiation pressure (we define the latter as the integral of mass in clouds where $P_{\mathrm{rad}}/\mathrm{Edd} > 1$). Fig. 2 (bottom right) shows this as a function of time, given the above outflow properties. Most of the cold mass can be effectively accelerated by radiation pressure at large radii in a time-scale $\sim$a few

galaxy dynamical times, despite the declining AGN luminosity. We have experimented with varying the exact parameters adopted here, and find that the qualitative results are robust. In a couple of dynamical times, ~ 90 per cent of the original cold cloud mass becomes vulnerable to secondary radiative feedback; i.e. these numbers – the fraction that can be ionized and/or the strength of radiation pressure approach or exceed unity.

This and the previous results also provide an important check of an implicit assumption in this model: that the cloud acceleration time is long relative to the time-scale for the clouds to expand/equilibrate. If this were not true, the two behaviours could not be treated independently, and the physical consequences are unclear (it is possible, for example, that each 'parcel' of the cloud which is mixed or stripped off by instabilities would rapidly be accelerated, leading to the edges of the clouds being effectively blown away or stripped but giving little acceleration to the cloud core). As noted above, the expansion/equilibration time is given by the cloud-crushing time, comparable to the internal cloud crossing/dynamical times $< 10^6\,\mathrm{yr}$. In comparison, the acceleration times are of order a couple to a few t_s, the dynamical time at r in the galaxy, which (as we discuss above) are generically much larger than the cloud crossing time at all radii $\sim R_e$ (where most of the galaxy mass is located), indeed all radii $\gtrsim 100\,\mathrm{pc}$ (i.e. all radii which are not already affected by feedback even without a diffuse outflow). In a global sense, most of the mass is accelerated and/or ionized over a time-scale $\sim 10^7$–$10^8\,\mathrm{yr}$, much longer than the crossing/collapse times of all but the most massive molecular cloud complexes.

In fact, this acceleration time is a relatively long time, at large scales ($\sim$a few $10^8\,\mathrm{yr}$), and the AGN luminosity has correspondingly decayed to ~ 1 per cent of the Eddington limit. The model above accounts for this, but an important question remains if real AGN can sustain even this level of energetic output over these time intervals. If, for example, AGN switch to a radiatively inefficient state above or around this accretion rate, the driving will suddenly vanish. Clearly, this is an interesting regime; better knowledge of how feedback-induced hot outflows and subsequent light-curve evolution proceed will be important to understanding both how dramatic the effect on the galaxy will be and, potentially, how much variation there may be between galaxies.

4 DISCUSSION

'Feedback' from bright AGN is a topic of fundamental interest for galaxy evolution, but it remains unknown whether or not any of the obvious candidate feedback mechanisms are capable of effectively coupling to cold molecular gas, especially at kpc scales, the dominant reservoir for star formation. Here, we demonstrate that it is at least possible that the cold gas reservoir is destroyed and/or blown out of the galaxy *despite* inefficient coupling of 'initial' feedback mechanisms that originate near the BH.

If some coupling of energy or momentum near the BH – whether from e.g. Compton heating, radiation pressure, BAL winds, jets or resonant line driving – can generate a wind or shock/blastwave in the warm/hot ISM, then when the outflow passes by a cold cloud, even if it does not directly entrain the material, it will generate various instabilities that 'shred' the cloud and mix it, efficiently enhancing the cloud cross-section in the perpendicular direction. Even if the cloud is magnetically supported or extremely dense and able to resist instabilities, there is still a growing pressure imbalance that drives the cloud to expand in the same manner.

This is well studied in the context of supernovae-driven winds, but there is an important difference in the presence of a bright quasar.

The effective increase in cross-section means that momentum driving and ionization heating from the quasar radiation is quickly able to act in much more dramatic fashion on clouds that were once too dense and too small (or at too large a distance from the BH) to be perturbed by the radiation field. This effect can have dramatic implications for star formation in quasar host galaxies.

Because radiation pressure always acts, this means the energy needed in 'initial' feedback from the central source to, for example, drive winds in the low-density hot gas will be much less than if it were expected to act directly on the cold clouds. We show that the energetic or momentum driving requirements for the initially driven feedback are reduced by at least factor $f_{\rm hot} \sim 0.1$ (the mass fraction in the hot diffuse ISM); i.e. rather than the canonical ~ 5 per cent of the radiant energy (~ 100 per cent momentum) needed in the initial outflow if it were to entrain the entire gas supply directly, only ~ 0.5 per cent (~ 10 per cent momentum) is sufficient to drive the hot gas and enter the regime of interest here. Another way of stating this is, for accretion with an Eddington ratio $\dot{m}$ and BH mass $M_{\rm BH}$ relative to the expectation $\langle M_{\rm BH}\rangle$ from the $M_{\rm BH}$–σ relation, the relevant outflows will be driven (and star formation suppressed) when

$$\eta\,\dot{m}\,\frac{M_{\rm BH}}{\langle M_{\rm BH}\rangle} \sim 0.05\,f_{\rm hot}, \qquad (9)$$

where η is the feedback efficiency ($\dot{E} = \eta\,L$). Given this criterion, that there is sufficient momentum in photons for the 'secondary' feedback to act is guaranteed for all but the most extremely gas-dominated systems.

Note that the derivation here pertains to large clouds ($R_0 \gtrsim {\rm pc}$), observed to be in rough pressure equilibrium with the ambient medium and containing most of the ISM mass. Dense cores ($R_0 \ll {\rm pc}$) are observed to be in self-gravitating collapse; these will continue to collapse and form stars on a very short time-scale despite a diffuse outflow. The important thing is that no *new* cold gas reservoir of large clouds will be available to form new cores.

We have also neglected the possibility that galaxies are highly self-shielding. For example, in dense nuclear star-forming regions in e.g. ultraluminous infrared galaxies (ULIRGs), the column densities are so high ($N_{\rm H} \gtrsim 10^{23}\,{\rm cm}^{-2}$; see e.g. Komossa et al. 2003; Li et al. 2008) that the quasar can do little until star formation exhausts more of the gas supply. In disc-dominated galaxies, gas near $R_{\rm eff}$ can be similarly self-shielded. If the radiation is isotropic, then for some disc mass and gas fraction, only a fraction $\sim h/R$ (the fractional scale height at R) will couple to the relevant area (and the radiation may, in fact, be preferentially polar, yielding even lower efficiency). Only the most gas-poor discs, or the central regions of discs (where systems are typically bulge dominated) will be affected by the coupling efficiencies above.

What we outline here is a simple model for the qualitative physical effects that may happen when cold clouds in the ISM encounter a hot outflow driven by an AGN. More detailed conclusions will require study in hydrodynamic simulations which incorporate gas phase structure, cooling, turbulence, self-gravity, radiation transport and possibly (if they provide significant pressure support) magnetic fields. Detailed effects which we cannot follow analytically, such as self-shielding within thin, dense fingers in Rayleigh–Taylor or Kelvin–Helmholtz instabilities may alter the effects of the radiation field on the cloud material and change our conclusions. Nevertheless, our simple calculations here demonstrate that the process of cloud deformation in the wake of a hot outflow can have dramatic implications for the susceptibility of those clouds to other modes of feedback, and should motivate further study.

According to these simple considerations, outflows driven by AGN feedback may in fact be 'multistage' or 'two-tiered', with an initial hot shockwave or strong wind driven by feedback mechanisms near the BH, which is then supplemented by a successive wind driven out as clouds in the wake of the former are deformed/mixed and increase their effective cross-section to the AGN luminosity. The characteristic velocity of this secondary outflow, which will carry most of the mass, should be $\sim v_{\rm esc}$ at the radii of launching ($\sim 10^2\,{\rm km\,s}^{-1}$), and it will behave similarly to outflows from star formation. Indeed, because the driving occurs at large radii, it is not clear whether it could be distinguished from stellar-driven outflows at all, except indirectly (e.g. in cases where the observed star formation is insufficient to power the outflow). Characteristic time-scales are $\sim$ a few $t_{\rm dyn}$ of the galaxy, so much of the outflow occurs as sub-Eddington luminosities (as the AGN fades in the wake of launching the 'primary' outflow) and the systems will not appear gas-depleted until they have evolved by a significant amount ($\sim$ few 10^8 yr, at Eddington ratios ~ 0.01 typical of 'quiescent' ellipticals). These processes should nevertheless imply effective shutdown of star formation and destruction/heating of the cold gas supply in 'massive' BH systems – bulge-dominated systems on the $M_{\rm BH}$–σ relation that have recently been excited to near-Eddington luminosities.

Most intriguing, this reduces the energetic requirements for the 'initial' feedback – whatever might drive an outflow in the hot gas from the vicinity of the BH – by an order of magnitude. Our estimates suggest that coupling only a fraction $\sim 10^{-3}$ of the luminosity of the AGN on small scales would be sufficient to drive such a hot outflow and then allow ~ 100 per cent of the radiative energy/momentum to couple to cold gas. If, for example, quasar accretion-disc (or broad-line) winds (with characteristic velocities $v \sim 10^4\,{\rm km\,s}^{-1}$) do not immediately dissipate all their energy, then the hot outflows we invoke would be generated with a mass loading in such winds of just $\sim 0.1\,\dot{M}_{\rm BH}$, a fraction of the BH accretion rate.

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