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 disk to directly impact (let alone blow out of the galaxy) dense molecular clouds at ∼kpc. But if such feedback can drive a weak wind or outflow in the hot, diffuse ISM (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% of the radiated energy (∼100% momentum) needed if the initial feedback must directly heat or “blow out” the galactic gas, if only ∼0.5% of the luminosity (∼10% 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 disk of only ∼5−10% of the BH accretion rate.

Key words: quasars: general — galaxies: active — galaxies: evolution — 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; Hopkins et al. 2007a; Aller & Richstone 2007). Together with constraints indicating that most of the BH mass is assembled in optically bright quasar phases (Soltan 1982; Salucci et al. 1999; Yu & Tremaine 2002; Hopkins et al. 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 et al. 2005; Murray et al. 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,b). These scenarios have been able to explain many quasar observables, including luminosity functions, lifetimes, and BH mass functions (Hopkins et al. 2005b, 2006a, 2006b, 2008a; Volonteri et al. 2006; Menci et al. 2008; Somerville et al. 2008; Lapi et al. 2006; 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; Hopkins et al. 2008a; Antonuccio-Delogu & Silk 2008). The argument in the models is that, under various simple assumptions, if sufficient energy or momentum is injected into the ISM near the BH on a timescale 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 & Fontana 2005; Hopkins et al. 2006a; Shin et al. 2009].

But 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 et al. 2005); these range from intense winds (\(v \sim 10^4 \text{ km s}^{-1}\)) associated with the central engine seen in the broad emission line regions and broad absorption line quasars (e.g. Weymann et al. 1981) to more moderate outflows (\(v \sim 10^2 \sim 10^3 \text{ km s}^{-1}\)) 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; Turner et al. 2008; Miller et al. 2008). Indeed, high-velocity winds driven near the accretion disk 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 \(\lesssim M_{\text{BH}}\) (although in at least some cases, these outflows are extremely dense, and might have much higher mass-loading factors; see e.g. Hall & Hutsemekers 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 (GMCs), 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; Holt et al. 2006; Gabel et al. 2006; Krongold et al. 2007a; Batcheldor et al. 2007; Tremonti et al. 2007; McKernan et al. 2007; Ganguly et al. 2007; Prochaska & Hennawi 2009).

In this paper, we argue that it is not necessary that the small-scale, high-velocity AGN outflows directly entrain any cold gas at scales \(\sim R_e\). 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 timescale, 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 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
cloud of initial characteristic (quasi-spherical) radius $R_0$ and density contrast $\chi$ (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” timescale $t_{cc} = v_s/\chi^{1/2} R_0/v_s$. In the simple case of a pure hydrodynamic strong shock, if $t_{cc}$ is much less than the characteristic timescales 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 timescale $\sim a few t_{cc}$ (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 $v_s/\chi^{1/2}$, $t_{cc}$ is much less than the dynamical timescales of interest at all the spatial scales of interest ($t_{cc} \ll 10^2$ 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 $\sim \chi^{1/2}$ in radius in the perpendicular shock direction. Eventually, despite the initial compression, reflection shocks lead to an expansion by a factor $\sim 2$ 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 $\pi R_0^2$, post-shock, the effective cross section owing to this expansion and equilibration is $\pi R_{eq}^2 \sim \chi \pi R_0^2$.

We illustrate this behavior with a simple toy model system in Figure 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 $\sim 2$ wind. Color 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 behavior 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 behavior.

The qualitative behavior 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 timescale 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 ($\sim$ 10 yr).

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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 & Alexandel 2007; Antonuccio-Delouei & Silk 2008). A cloud in principle can be stabilized against such instabilities by being strongly magnetically dominated (Mac Low et al. 1994; Jun & Jones 1999; Fragile et al. 2005; Shin et al. 2008). However, in this limit, as the hot outflow sweeps up 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 over-pressurized, it will expand isothermally as the exterior post-shock pressure drops (the free expansion/equilibrium time of the cloud being short compared to the other timescales of interest). Because this stops when the system is equilibrated, the “effective” net expansion of $R_0$ 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 outflows until they reach equilibrium with the ISM and re-cool into new clouds.

However, we are interested in all of this occurring in the background of a luminous AGN, which will both ionize and ex-
ert a radiation pressure force. The cloud – especially a realistic cloud with a large dust mass and corresponding opacity – is optically thick 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 $p_{\text{rad}} = L_{\text{abs}}/c$ where $L_{\text{abs}} = L_{\text{QSO}} \Delta \Omega$. Comparing this to the gravitational force $F_{\text{grav}} = -M_c \partial \phi/\partial r$ 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 $\sim M_c v_{\text{esc}}$). This limit is when the two are equal: equivalently

$$\frac{L_{\text{QSO}}}{c} \frac{\pi R_c^2}{r^2} = M_c \frac{G M_{\text{BH}}}{(r + R_{\text{eff}})^2}. \quad (5)$$

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

$$\left(\frac{R_c}{R_0}\right) \gtrsim 2.7 \frac{m_{\text{abs}}}{r + R_{\text{eff}}} \frac{\rho_i}{\rho_i + \rho_{\text{gas}}}. \quad (6)$$

In other words, a cloud on the “normal” size mass relation at large radii $r \sim R_0$ is sufficiently dense and sufficiently high column-density to avoid being unbound by radiation pressure. But 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.

Figure 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 $\gtrsim 3$ times the initial background density but with no cell at $x < x_0$ 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 $t < t_f$ 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 $v_c = V_f$), $\sim 0.01 M_c v_c (R_c/R_f)/v_f$. 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 Figure 1). Ignoring geometric effects of photon diffusion, the volume of a cloud of mean density $n_i$ ionized is $V_{\text{ion}} = N/n_i^2 \beta$, where $\beta \approx 2 \times 10^{-11} \text{ cm}^2 \text{ s}^{-1}$ is the recombination coefficient for gas at the temperature for hydrogen ionization ($T_i \sim 10^4 \text{ K}$) and $N$ is the rate at which ionizing photons hit the cloud. The total rate of production of ionizing photons from the quasar is $N_{\gamma} = \lambda L_{\gamma}/h_0, 12$ (a $\sim 0.07$ comes from a proper integration over this quasar spectrum; here from [Hopkins et al. 2007]), 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

$$h_{\text{ionized}} \approx \frac{h_{\text{ionized}}}{R_c} \approx 10^{-3} m_{\text{ion}} \left(\frac{r + R_{\text{eff}}}{r}\right)^2 \left(\frac{R_c}{R_0}\right)^5. \quad (7)$$

Give 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 $r \gtrsim R_{\text{eff}}$. 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 $\propto R_c^5$ owing to the decreased density lowering the recombination rate.

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

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as $R_c \propto t^\alpha$, where $\alpha$ 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 (Figure 2, bottom left). For typical galaxy density profiles, $\alpha = 5/3$ 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 onto a BH scales roughly $\propto r_p$ (in Bondi-Hoyle accretion; or similarly $\propto \Sigma$ for viscous accretion from a disk), the accretion rate, and hence luminosity $L = 0.1 M_{BH} c^2$, will decay as well (Figure 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 a double power-law.

The shock therefore crosses a radius $r_t$ (when $R_s = r$). In the wake of the shock, the post-shock ambient pressure $P_{\text{ext}}$ will drop, reflecting the density decline from material being blown out (Figure 2, top center). 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 $P \propto [(t - t_f)/t_f]^{-\beta}$, where $t$ is the crossing time of the shock relative to the cloud ($\sim R_c/v_c$) 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 $P \propto [(t - t_f)/t_f]^{-\beta'}$ and $\beta' \sim 2 \leq \beta$ (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 behavior of each cloud — at least the key parameter of interest, the effective radius of the cloud in the direction perpendicular to the shock ($R_c$) 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) cloud would expand isothermally (e.g. conserving total magnetic energy) as it is now over-pressurized, such that pressure equilibrium would be conserved. For a background pressure declining with some power law $P \propto (t/t_c)^{-\beta}$, this implies expansion of the cloud with $R_c/R_0 = (t/t_c)^{3/5}$ (generally, $R_c/R_0 = (P/P_0)^{-1/3}$, for isothermal expansion). If the cloud were supported by thermal energy with no new inputs this index would be slightly modified (for adiabatic expansion, $R_c/R_0 = (P/P_0)^{1/7}$) but the behavior is qualitatively similar. Technically this approximation assumes that the time for the cloud to equilibrate is short relative to the timescale 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 $\sim R_c/c_{s,\text{eff}}$; for quasi-virial clouds this is simply the dynamical time $1/\sqrt{G\rho}$ where the effective density $\rho$ follows from the observed size-mass relation (Equation 3). Using the observed values, this gives a timescale of $\sim 0.1 - 1 \times 10^7$ yr (for cloud sizes $\sim 0.1 - 10^6$ pc). Compare this to the characteristic timescale for the evolution of the background, a few $t_e$ (itself of order $t_{\text{dyn}}$, the galaxy dynamical time at $R$). For a typical $\sigma \sim 200$ km s$^{-1}$ spheroid, $t_e \gtrsim R_c/c_{s,\text{eff}}$ For all radii $\gtrsim 100$ pc – in other words, this condition is easily satisfied at the radii $\sim R_c$, which contain most of the mass of the galaxy. Figure 2 (bottom center) shows how the clouds will expand in effective cross section ($\sim R_e$), 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 timescales $t_{cc} = \chi^{1/2}R_0/v_c$. But since the shock velocity is of order the galaxy escape velocity for the “interesting” diffuse outflows considered here ($v_c \sim$ a few $\sigma$), and typical $\sigma \sim 200$ km s$^{-1}$, compared with a typical effective sound speed of a virialized cloud $\sim 1 - 10$ 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 $\sim 1/2$ in the perpendicular direction (and a small, $\sim$ constant factor in the parallel direction), where $\chi$ is the initial density contrast; in other words, until the effective density and pressure drop to approximate equilibrium with the background. Thus, for timescales $\sim t_{cc}$ 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 timescales, as in the isothermal expansion case).

We then solve for the behavior of each cloud in the wake of this hot outflow, according to the approximations in [Klein et al. (1994)] and § 2 (Figure 2, top right). In particular, given the time-evolution in $R_c/R_0$ shown in Figure 2, we use the scalings derived in § 2 to estimate the fraction of the cloud (at some initial radius $r$) which will be ionized (i.e. $f_{\text{ion}} = \rho \sim 1/2$ of $R_{cr}/R$ at Equation 7) where the AGN accretion rate $m$ 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 $R_{0,pc} = 1$ 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), $P_{\text{rad}}/Edd = 0.14 m (R_c/R_0)^2 [(r + R_c)/r]^3$ (re-arranging Equations 5 & 6, where again $m$, $R_c/R_0$, and $r$ for the clouds is given). Since these both scale steeply with the increasing effective cloud size ($\sim R_c^2$ and $R_c^2$, respectively), both increase rapidly in time. At early times, only the clouds within a narrow region $\sim 100$ 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 $P_{\text{rad}}/Edd > 1$).
crossing time at all radii (as we discuss above) are generically much larger than the cloud hot/diffuse outflow as a function of time (compare the galaxyeffective radius times of all but the most massive molecular cloud complexes.

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% 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 – 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 scales (~a few 10^8 yr), and the AGN luminosity has correspondingly decayed to ~ 1% 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 lightcurve 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 black hole) 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 e.g. 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 $f_{\text{hot}} \sim 0.1$ (the mass fraction in the hot diffuse ISM); i.e., rather than the canonical $\sim 5\%$ of the radiant energy ($\sim 100\%$ momentum) needed in the initial outflow if it were to entrain the entire gas supply directly, only $\sim 0.5\%$ ($\sim 10\%$ momentum) is sufficient to drive to the hot gas and enter the regime of interest here. Another way of stating this is, for accretion dominated galaxies, gas near the radiant energy ($\sim \eta \mathcal{L}$). 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 ($R_0 \geq \text{pc}$), observed to be in rough pressure equilibrium with the ambient medium and containing most of the ISM mass. Dense cores ($R_0 \ll \text{pc}$) are observed to be in self-gravitating collapse; these will continue to collapse and form stars on a very short timescale despite a diffuse outflow. The important thing is that 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. ULIRGs, the column densities are so high ($N_H \gtrsim 10^{23} \text{cm}^{-2}$, see e.g. Komossa et al. 2003, Li et al. 2007) that the quasar can do little until star formation exhausts more of the gas supply. In disk-dominated galaxies, gas near $R_{\text{eff}}$ can be similarly self-shielded. If the radiation is isotropic, then for some disk 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 disks, or the central regions of disks (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 e.g. 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 “multi-stage” 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 $\sim v_{\text{esc}}$ at the radii of launching ($\sim 10^7 \text{km} \text{s}^{-1}$), 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 timescales are $\sim$ a few $t_{\text{dyn}}$ 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$^9$ yr, at Eddington ratios $\sim 0.01$ 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 $M_{\text{BH}} - \sigma$ 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 $\sim 100\%$ of the radiative energy/momentum to couple to cold gas. If, for example, quasar accretion-disk (or broad-line) winds (with characteristic velocities $v \sim 10^4 \text{km} \text{s}^{-1}$) 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 $\sim 0.1 M_{\text{BH}}$, a fraction of the BH accretion rate.

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