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The end of the black hole dark ages and the origin of warm absorbers

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ABSTRACT

We consider how the radiation pressure of an accreting supermassive black hole (SMBH) affects the interstellar medium around it. Much of the gas originally surrounding the hole is swept into a shell with a characteristic radius somewhat larger than the black hole's radius of influence (∼1–100 pc). The shell has a mass directly comparable to the $(M-\sigma)$ mass that the hole will eventually reach, and may have a complex topology. We suggest that outflows from the central SMBHs are halted by collisions with the shell, and that this is the origin of the warm absorber components frequently seen in active galactic nucleus (AGN) spectra. The shell may absorb and reradiate some of the black hole accretion luminosity at long wavelengths, implying both that the bolometric luminosities of some known AGN may have been underestimated, and that some accreting SMBH may have escaped detection entirely.

Key words: black hole physics – galaxies: active – quasars: general – X-rays: galaxies.

1 INTRODUCTION

Astronomers now generally agree that the centre of most galaxies contains a supermassive black hole (SMBH). Active galactic nuclei (AGN) correspond to phases when the hole is growing in its mass by accreting gas from a very small scale disc around it. But is it not immediately obvious why these phases are in practice directly observable, as there are good reasons to expect that a significant mass of gas largely surrounds the hole at such epochs. First, a majority of AGN show significant signs of obscuration (cf. the discussion in Elvis 2000). A large gas mass is needed close enough to the hole to grow it in a reasonable time, and this must cover a significant solid angle since the total mass of a geometrically thin disc is severely limited by self-gravity. Finally, simple estimates of the column densities of matter subject to a galactic potential confirm the impression that most AGN are likely to be at least formed in fairly dense gas environments.

This suggests that when we do see black holes accreting, they may have perturbed the gravitational equilibrium of the matter which would otherwise block our view. An obvious way of doing this is to push it away, spreading it over a larger area and reducing its column density. We investigate this idea here.

2 PUSHING FOR TRANSPARENCY

In luminous AGN, by far the strongest energy supply potentially pushing matter away from a black hole is its accretion luminosity L. (This is of course not true for low-luminosity radio galaxies, where jets may interact with the surroundings.) By contrast, accretion disc *winds* are generally limited to mechanical luminosities $\eta L/2 \simeq 0.05L$, where $\eta \sim 0.1$ is the accretion efficiency (e.g. King & Pounds 2003; King 2003), and much of this is likely to be lost in shocks (King 2003, 2010).

If the surroundings have high optical depth to scattering (i.e. are strongly obscuring), photons scatter many times, and so radiation pressure must become significant. In scattering slightly inelastically, the luminosity L does work pushing against the gravitational force on the surrounding gas in the central spheroid of the galaxy. If the gas is not in large-scale dynamical motion, we assume that it is distributed isothermally, i.e. with density

$$\rho(r) = \frac{f_{\rm g}\sigma^2}{2\pi G r^2}, \tag{1}$$

so that the gas mass within radius R is

$$M_{\rm g}(R) = \frac{2 f_{\rm g}\sigma^2 R}{G}, \tag{2}$$

and the total mass (including stars, and any dark matter) is

$$M(R) = \frac{2\sigma^2 R}{G}. \tag{3}$$

Here, σ is the velocity dispersion and $f_{\rm g}$ is the gas fraction relative to all matter (e.g. dark matter and stars) which has cosmic value 0.16. We assume that $f_{\rm g}$ does not vary strongly across the central region of the galaxy.

The pressure of trapped radiation sweeps the gas up progressively into a shell of inner radius R and mass $M_{\rm g}(R)$. If the shell is

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geometrically thin, its electron scattering optical depth at radius R is

$$\tau_{\rm sh}(R) = \frac{\kappa M_{\rm g}(R)}{4\pi R^2} = \frac{\kappa f_{\rm g}\sigma^2}{2\pi G R}, \tag{4}$$

where $\kappa \simeq 0.34\,{\rm cm}^2\,{\rm g}^{-1}$ is the electron scattering opacity. If the shell is geometrically thicker, $\tau_{\rm sh}(R)$ is an upper limit to its optical depth, as on average the gas is more spread out (i.e. at larger radii). The undisturbed gas outside R has optical depth

$$\tau(R) = \int_R^\infty \kappa\rho(r){\rm d}r = \frac{\kappa f_{\rm g}\sigma^2}{2\pi G R} = \tau_{\rm sh}(R), \tag{5}$$

most of which is concentrated near the inner radius R. The radiation thus encounters total optical depth

$$\tau_{\rm tot}(R) = \tau(R) + \tau_{\rm sh}(R) \simeq \frac{\kappa f_{\rm g}\sigma^2}{\pi G R} \tag{6}$$

whatever the thickness of the shell be. Gas distributed in this way is optically very thick near the black hole when its inner edge R is small (cf. equation 8 below). Then the accretion luminosity L of the AGN is initially largely trapped and isotropized by scattering, increasing the interior radiation pressure P. This growing pressure pushes against the weight

$$W(R) = \frac{GM(R)M_{\rm g}(R)}{R^2} = \frac{4 f_{\rm g}\sigma^4}{G} \tag{7}$$

of the swept-up gas shell at radius R (which is constant with R since $GM(R)M_{\rm g}(R)/R^2 \propto R.R/R^2 = \text{constant}$).

The appendix discusses in detail the shell's equation of motion as it expands. But it is already clear that the effectiveness of radiation pressure is eventually limited because the shell's optical depth falls off like $1/R$ as it expands. The force exerted by the radiation drops as it begins to leak out of the cavity, until for $\tau_{\rm tot}(R) \sim 1$ it is unable to drive the shell further.

This shows that the sweeping up of gas by radiation pressure must stop at a 'transparency radius'

$$R_{\rm tr} \sim \frac{\kappa f_{\rm g}\sigma^2}{\pi G} \simeq 50\left(\frac{f_{\rm g}}{0.16}\right)\sigma_{200}^2\,{\rm pc}, \tag{8}$$

where (up to a logarithmic factor) the optical depth $\tau_{\rm tot}$ is of the order of 1, so that the radiation just escapes, acting as a safety valve for the otherwise growing radiation pressure. Here, $\sigma_{200} = \sigma/200\,{\rm km\,s}^{-1}$.

The appendix shows that the total gravitational potential energy which the accretion luminosity must supply to push the galactic gas to this radius is

$$E_{\rm tr} \simeq 3WR_{\rm tr} = \frac{12\kappa f_{\rm g}^2\sigma^6}{\pi G^2}, \tag{9}$$

so that the central black hole must accrete a mass

$$\Delta M \gtrsim \frac{E_{\rm tr}}{\eta c^2} \sim 3\times 10^3\sigma_{200}^6\,{\rm M}_\odot, \tag{10}$$

where $\eta \simeq 0.1$ is the accretion efficiency. This is much smaller than the black hole mass itself, so we expect transparency to be achieved early in the life of the central SMBH, and easily maintained after this. In addition, the radiation field of the accreting SMBH ionizes many of the photoelectric absorbing species outside this radius, affecting the photoelectric absorption column.

Our discussion so far assumes that the swept-up shell remains spherical, whereas in reality it is likely to fragment to some degree. We consider the effects of this further in Section 5 below.

3 THE SIGNIFICANCE OF THE TRANSPARENCY RADIUS

We can rewrite equation (8) as

$$R_{\rm tr} = \frac{M_\sigma}{M}R_{\rm inf}(M) = R_{\rm inf}(M_\sigma), \tag{11}$$

where

$$M_\sigma = \frac{f_{\rm g}\kappa}{\pi G^2}\sigma^4 \tag{12}$$

is the $M-\sigma$ mass (King 2003, 2005) and

$$R_{\rm inf}(M) = \frac{GM}{\sigma^2} \tag{13}$$

is the gravitational influence radius of a hole of mass M. From equation (3) and the first form of equation (8), we also have

$$M(R_{\rm tr}) = 2f_{\rm g}\frac{f_{\rm g}\kappa\sigma^4}{\pi G^2} = 2f_{\rm g}M_\sigma \sim M_\sigma. \tag{14}$$

So, we can think of the transparency radius $R_{\rm tr}$ as roughly the radius initially containing a gas mass comparable to the final mass of the black hole. The pressure of the trapped radiation rearranges much of this gas into a shell at $R_{\rm tr}$. While this is now transparent to the accretion luminosity of the central SMBH, its enormous mass makes it a severe obstruction to mechanical outflows. These must shock against it and effectively stop completely for any SMBH mass below the critical $M-\sigma$ value. (The significance of the M_σ mass is that at this point, winds carrying the Eddington thrust of the SMBH are finally able to drive the gas to large radii, where the wind shocks no longer cool. The outflow makes a rapid transition to energy driving, which largely clears the gas from the galaxy spheroid, simultaneously halting the SMBH growth – cf. King 2003, 2005.)

The most important outflows are black hole winds driven by radiation pressure. These have velocities $v \simeq 0.1c$ and momentum scalars $\dot{M}v \simeq L_{\rm Edd}/c$, where $L_{\rm Edd}$ is the Eddington luminosity (King & Pounds 2003). The impact of these winds on the interstellar gas is what ultimately fixes the $M-\sigma$ relation (King 2003, 2005). Many of these impacts are likely to occur close to the transparency radius $R_{\rm tr}$, which we can also write (using equation 11) as

$$R_{\rm tr} = 10^6\frac{M_\sigma}{M\sigma_{200}^2}R_{\rm s}(M), \tag{15}$$

where $R_{\rm s} = 2GM/c^2$ is the black hole gravitational radius.

4 OBSERVATIONAL CONSTRAINTS

We argued above that black hole winds are halted by collisions near $R_{\rm tr}$. The shocked winds must rapidly cool, slow and recombine, and mix with swept-up interstellar medium (ISM). So we expect this gas to have modest ionization, and much slower velocities than the winds themselves. These properties are very similar to those inferred for the so-called warm absorber (WA) components in AGN spectra, and we suggest that WA result from these wind impacts. Here, we look for observational tests of this idea.

Powerful highly ionized winds in AGN have been widely observed in X-ray spectra over the past decade (Pounds et al. 2003; Reeves, O'Brien & Ward 2003; Tombesi et al. 2010). The radial location R of an AGN outflow component is notoriously difficult to determine from the quantities usually measured – ionization parameter ξ and equivalent hydrogen column density $N_{\rm H}$. These both involve the electron density, and the column density, an unknown filling factor as well. To date this degeneracy has been resolved in

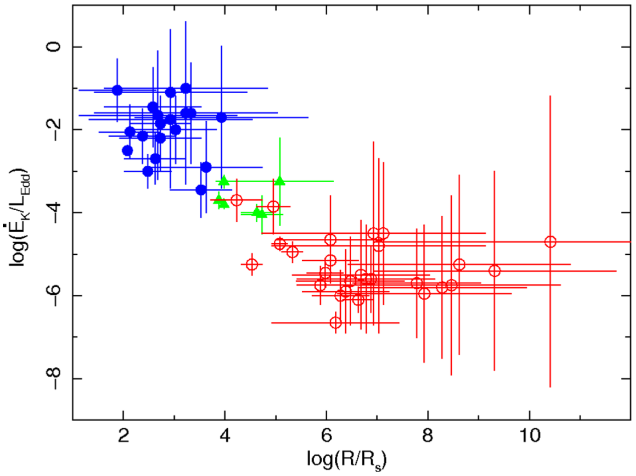

Figure 1. Comparison between the radial distance log (R/Rs) and the estimated outflow kinetic energy rate for an overlapping sample of WAs (red) and UFOs (blue) in nearby, bright AGN (from Tombesi et al. 2013). For a black hole mass of $10^8\ \mathrm{M}_\odot$, the distance scale may be converted to parsecs by noting that $10^6 R_\mathrm{s}$ is $\sim$10 pc.

a very few cases, including the fast outflow of NGC 4051 (Pounds & King 2013).

Tombesi et al. (2013, hereafter T13) provide constraints on the radial location of WAs, and set this in a wider context. They consider a sample of 35 type 1 Seyferts previously included in a study of ultrafast outflow (UFOs), many of which also show WAs. Although a direct determination is not possible, T13 constrain the radial locations of all the velocity components in their sample. They get minimum values from the assumption that the gas is moving at the local escape velocity, and maximum values from the relation

$$R = \frac{L_\mathrm{i}}{N_H \xi}. \tag{16}$$

Here, L_i is the ionizing luminosity, and T13 use values of N_H and ξ from XSTAR spectral fitting. Fig. 1 is based on a figure from T13 and compares the derived kinetic energy rates with the radial distance constraints for both UFOs and WAs in the sample. To make the plot independent of black hole mass the energy rate is in units of L_Edd, and the radius in Schwarzschild radii R_s. As expected, UFOs cluster at $\sim 10^2 R_\mathrm{s}$. Importantly, we see that WAs cluster between 10^6and10^7 Schwarzschild radii R_s.

For comparison, equation (8) predicts $R_\mathrm{tr}/R_\mathrm{s} \simeq 5 \times 10^6$, taking $\sigma = 200\ \mathrm{km\ s^{-1}}$ for an SMBH mass of $10^8\ \mathrm{M}_\odot$. Although the spread of data points is broad, with substantial uncertainty for individual radii, the coincidence of the radial distance distribution for the full sample of WAs with the transparency radius derived above is a strong indication of a physical connection.

A second feature of Fig. 1 is the much higher kinetic power carried by the UFOs as a group, underlining the importance of high-speed winds for AGN feedback. We note that T13 also show that the measured mass outflow rates are essentially constant with radius (their fig. 2). So the factor $\sim 10^4$ between the UFO and WA energy rates in Fig. 1 is consistent with the idea that these two components characterize the start and end points of the same mass-conserving outflows, with mean velocity differences of the order of $\sim$100.

UFOs have typical velocities $\sim 0.1c$ characteristic of escape from $R \sim 100 R_\mathrm{s}$ (cf. King & Pounds 2003), and must retain them until they hit an obstruction with comparable inertia. Since UFOs are also hypersonic, any deceleration must involve strong shocks. At this point they are likely to lose almost all of their kinetic energy, since the shocks are probably strongly Compton cooled by the AGN radiation field (King 2003, 2010). It is natural to assume that the obstruction is usually the surrounding ISM at R_tr. In some cases, the winds may hit previous shocked ejecta or infalling gas well within R_tr. This latter possibility seems likely for NGC 4051 (Pounds & King 2013, and references therein).

5 DISCUSSION

We have seen that the trapped radiation pressure exerted by an accreting SMBH is likely to affect the ISM in its immediate neighbourhood quite strongly. Much of the gas originally surrounding the hole is swept into a dense shell with the characteristic radius R_tr, where photons can escape and prevent a further build up of radiation pressure inside. This shell has a mass directly comparable to the final $(M-\sigma)$ mass the hole will eventually reach.

It appears that the radius R_tr is similar to the size of the region responsible for warm absorber behaviour. This is very reasonable, since the shell at R_tr is so massive that outflows from the central SMBH must be halted in shocks there if they do not collide with other structures within R_tr.

We have so far treated the swept-up shell as spherical and continuous. It is likely that in practice instabilities can fragment it before it reaches R_tr. If the topology of the fragmented shell remained simple (i.e. a punctured ball), this might relieve the excess radiation pressure and let the fragmented shell settle at a radius within R_tr. However, it is likely that the fragmented shell becomes complex, because it is effectively Rayleigh–Taylor unstable. The instabilities then produce overturning motions and hence overlapping gas, which in practice make it difficult for photons to escape without multiple scattering. Moreover, the undisturbed ISM immediately outside R_tr contributes at least as much opacity as the swept-up shell. This suggests that even in the given fragmentation, R_tr is likely to remain a characteristic radius for the central AGN. An indication of the complex topology of this region may be that the warm absorption column often has no accompanying cold absorption, as we might naively expect for a smoothly stratified shell.

In all cases, it seems very likely that some of the AGN luminosity gets absorbed and reradiated by gas with significant optical depth situated at radii $\lesssim R_\mathrm{tr}$. If the reradiated component is roughly blackbody, we find a characteristic temperature

$$T_\mathrm{tr} = \left(\frac{l L_\mathrm{Edd}}{4\pi f R_\mathrm{tr}^2 \sigma_\mathrm{SB}}\right)^{1/4} \sim 100 \left(\frac{l M}{f M_\sigma}\right)^{1/4} \mathrm{K}, \tag{17}$$

where l is the fraction of the Eddington luminosity L_Edd, which is reradiated, and $f \times 4\pi$ the solid angle of the obscuring shell. Given this low temperature and the large photosphere, this component may have evaded detection. A completely intact shell ($f = 1$) might totally obscure an AGN. On either count, it seems possible both that the bolometric luminosities of some known AGN may have been underestimated, and that some accreting SMBH have escaped observation entirely.

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APPENDIX A: MOTION OF A GAS SHELL SWEPT UP BY TRAPPED RADIATION PRESSURE

Assuming that the swept-up optically thick gas shell is geometrically thin, its equation of motion is

$$\frac{\mathrm{d}}{\mathrm{d}t}[M_{\mathrm{g}}(R)\dot{R}] = 4\pi R^2 P - W. \tag{A1}$$

Since the radiation pressure does work on the surroundings, we also need the energy equation

$$\frac{\mathrm{d}}{\mathrm{d}t}[VU] = L - 4\pi R^2 \dot{R} P - W\dot{R}, \tag{A2}$$

where $V = 4\pi R^3/3$ is the volume interior to the shell, which is filled with radiation of energy density $U = 3P$, and is supplied with further energy at the rate L. This form is very similar to the energy equation for the case of a wind with mechanical luminosity L, assuming that none of this is lost in cooling after shocking against the surroundings ('energy-driven flow'). The equation for this case is derived in King (2005, see also King, Zubovas & Power 2011). We follow the derivation given there, for a general adiabatic relation $P = (\gamma - 1)U$, where the index $\gamma = 4/3,\ 5/3$ for the present case of radiation and the earlier case of a monatomic gas. We use equation (A1) to eliminate the pressure P from equation (A2). The result is

$$L = \frac{2f_{\mathrm{g}}\sigma^2}{3G(\gamma-1)}[R^2\dddot{R} + (3\gamma+1)R\dot{R}\ddot{R} + (3\gamma-2)\dot{R}^3] + \frac{6\gamma-5}{3\gamma-3}.\frac{4f_{\mathrm{g}}\sigma^4}{G}\dot{R}. \tag{A3}$$

This reduces to the equation for energy driving by a wind given in King (2005) and King et al. (2011) if we set $\gamma = 5/3$ (note that the mechanical luminosity L of the wind is $(\eta/2)$ times the near-Eddington radiative luminosity $\sim L_{\mathrm{Edd}}$ driving it in this case).

In the trapped radiation case of this Letter, we have $\gamma = 4/3$, giving 325

$$L = \frac{2f_{\mathrm{g}}\sigma^2}{G}[R^2\dddot{R} + 5R\dot{R}\ddot{R} + 2\dot{R}^3] + \frac{12f_{\mathrm{g}}\sigma^4}{G}\dot{R}. \tag{A4}$$

As in the wind case (see King 2005; King et al. 2011), there is a constant-velocity solution $R = v_e t$, with

$$L = \frac{4f_{\mathrm{g}}\sigma^2 v_e^3}{G} + \frac{12f_{\mathrm{g}}\sigma^4}{G}v_e, \tag{A5}$$

which is an attractor. This equation defines a unique solution v_e. We can write L as

$$L = \frac{\mathrm{d}E}{\mathrm{d}R}v_e, \tag{A6}$$

where E is the total radiation energy inside R, so that equation (A5) becomes

$$\frac{\mathrm{d}E}{\mathrm{d}R} = \left(3 + \frac{v_e^2}{\sigma^2}\right)W. \tag{A7}$$

For modest accretion luminosities L (i.e. well below the Eddington value for the final black holes mass M_σ), we must have $v_e \ll \sigma$. Then, equation (A7) implies that the total accretion energy used to push the gas to the transparency radius R_{tr} is

$$E_{\mathrm{tr}} \simeq 3WR_{\mathrm{tr}} = \frac{12\kappa f_{\mathrm{g}}^2\sigma^6}{\pi G^2}, \tag{A8}$$

(cf. equation 9) in the body of this Letter.