# THEORY OF THE INTERMEDIATE ZONE

JULIAN H. KROLIK
*Department of Physics and Astronomy*
*Johns Hopkins University*
*Baltimore Maryland 21218*
*U.S.A.*

ABSTRACT: The "intermediate zone" is the several decades in radius between the broad and narrow emission line regions in active galactic nuclei. As Antonucci and Miller (1985) first discovered, it comprises two basic structures: an obscuring torus, and a reflection region which fills the hole of the torus and extends some ways above its top. Recent theoretical work, reviewed here, has concentrated on understanding both the thermal/ionization state, and the dynamical state of both regions. The torus is probably composed of a large number of dusty, molecular clouds of large column density and unusually high temperature and ionization fraction, while the reflection region is hot and highly ionized. Spectroscopic diagnostics exist for both. The key problem, still essentially unsolved, in the dynamics of the torus is its geometrical thickness. The reflection region is probably a pressure-driven transonic wind whose source is material evaporated by photoionization from the inner surface of the torus, but a small fraction of it may actually be accreting onto the nucleus.

## 1. Introduction

At the time of the last Santa Cruz Workshop on Active Galactic Nuclei (1984), there was still argument whether type 1 and type 2 Seyfert galaxies had the same sort of central engine, or whether they were more than fortuitously related at all (see, *e.g.,* Wilson and Heckman 1985). As Miller discusses elsewhere in this volume, at that time he already had the key data in hand which would indicate the answer to that question: type 2 Seyfert galaxies differ from type 1 Seyfert galaxies (or possibly a subset of type 1 Seyfert galaxies) only in viewing angle. The reason we see strong narrow emission lines from both types, but broad lines and a strong UV/X-ray continuum only from type 1's, is that a very optically thick torus surrounds the nucleus, occupying some portion of the several decades in radius outside the broad line region and inside the narrow line region. This blocking of the nonstellar continuum also largely explains the discrepancy in the ratio of radio luminosity to nonstellar optical continuum luminosity between types 1 and 2 which was a subject of much interest four years ago. Viewed along the polar axis, such a system would appear to be a type 1 Seyfert galaxy; viewed in the equatorial plane, the only way the nucleus can be revealed is through polarizing reflection from the material filling the hole of the torus.

The data obtained four years ago (and published the following year: Antonucci and Miller 1985) was high quality spectropolarimetry of the archetypical type 2 Seyfert galaxy, NGC 1068. Since that time, considerable additional evidence has been found in support of this picture (see Miller's article in this volume for a more detailed review of observational work). Miller and co-workers have shown that other type 2 Seyfert galaxies also reveal type 1 nuclei in spectropolarimetry, and that



*D. E. Osterbrock and J. S. Miller (eds.), Active Galactic Nuclei, 285–292.*

in NGC 1068 collimated nuclear light can be seen escaping through the torus's hole *via* reflection from dust clouds hundreds of parsecs out. X-ray observations detecting a weak, but hard, X-ray continuum from NGC 1068 (Elvis and Lawrence 1988, Koyama this volume), as well as an Fe K$\alpha$ line of very large equivalent width (Koyama, this volume) are also in line with this view.

In addition to the intrinsic interest accruing to observations which explain the distinction between two types of active galaxies, these observations have the additional significance of opening up an entirely new field of inquiry. Because no photons had ever been previously identified as coming from this region, it had been impossible to study. Now, given such a strong observational impetus, theoretical efforts have become both possible and attractive.

The theoretical problems taken up so far divide naturally into four groups, divided both by location (obscuring torus vs. reflection region) and by subject (thermal and ionization state of the material vs. dynamical state and evolution). This review will be organized by that subdivision. Although considerable progress has been made, one of the particular goals of this review is to highlight areas requiring more work.

## 2. Thermal and Ionization State of the Obscuring Torus

The weakness of even hard X-rays from type 2 Seyfert galaxies (Mushotzky 1987) immediately tells us that the mean column density in the equatorial plane of the torus must be at least a few $\times 10^{24}$ cm$^{-2}$. The minimum, *i.e.*, unclumped, emission measure if the torus material were ionized, $\sim 10^{68} r_{pc}^3$ cm$^{-3}$, is so large that the luminosity it would produce—$\sim 10^{46\pm1}$ ergs s$^{-1}$—would certainly have been observed long ago. We estimate the scale of the torus as $\sim$ 1 pc because the spectropolarimetry results show that it must be outside the broad line region ($\sim 0.1 L_{44}^{1/2}$pc) and inside the narrow line region ($\sim 10 - 100 L_{44}^{1/2}$pc). Here $L_{44}$ is the nuclear luminosity in units of $10^{44}$erg s$^{-1}$. The torus must, then, be mostly neutral. The only other broad band opacity source besides electron scattering capable of blocking the optical light from the nucleus is dust. Such a large column implies, for a conventional dust/gas ratio, an enormous extinction: $A_V \sim 10^3$ mag.

At the same time, however, the statistics of Seyfert galaxy types (as garnered from, *e.g.*, the CfA redshift survey: Burg 1987; Edelson, *et al.* 1987) suggest that the torus is geometrically as well as optically thick. For that to be the case, the material in the torus must possess a velocity dispersion in the vertical direction which is a significant fraction of its orbital velocity. Even without any stellar contribution to the potential, the orbital velocity must be at least $50 L_{44}(L_E/L)$ km s$^{-1}$, so these random motions, if thermal, would correspond to temperatures easily high enough to destroy dust grains ($L_E$ is the Eddington luminosity). Consequently, we expect that the material in the obscuring torus is clumped, with the required large random velocities arising from the bulk motion of the clumps.

Although there may be additional sources of heating and ionization within the torus, the nuclear continuum is certain to be present at some level. However, at a depth of $\sim 10^{24}$ cm$^{-2}$, only the hardest X-rays penetrate. Dust eliminates every nuclear photon from the infrared through the extreme UV, while K-shell opacity absorbs photons from 1 Ryd through $\simeq$ 10 keV. The Compton opacity seen by harder photons is independent of whether the electrons are bound in atoms or not, so they, too, are screened, but the effectiveness of the shadowing depends on what fraction of the solid angle around the nucleus is covered by the torus. In the limit as the torus approaches a complete spherical shell, the only diminution with radius suffered by the X-ray flux $F_x$ would be the Compton recoil losses of photons above $\sim$ 100 keV; in the opposite limit, when the difference between the outer and inner radii of the torus is very large compared to its vertical thickness, the flux would drop nearly exponentially. No one has yet done a transfer solution for the realistic geometry which lies between these two extremes, so we are forced, as often happens, to parameterize our ignorance.

https://doi.org/10.1017/S0074180900141063 Published online by Cambridge University Press

$F_x$ may be bounded between its value at the inner edge of the torus, $\sim 8 \times 10^5 L_{44} r_{pc}^{-2}$ ergs $\mathrm{cm^{-2}s^{-1}}$, and its minimum possible value at the outer edge, several e-foldings smaller.

The pressure in the clumps is even less certain, but a plausible estimate comes from the minimum pressure which, at its inner edge, sustains a cool photoionized equilibrium against the full nuclear ionizing continuum: $p \sim 2 \times 10^{10} L_{ion44} r_{pc}^{-2}$ K $\mathrm{cm^{-3}}$. At such pressures, dusty clump material is primarily molecular, but comparatively warm and ionized: $T \sim 10^3$K, and $x_e \sim 10^{-3}$ (Krolik and Lepp 1988). The column density of CO is so large that even the rarer isotopes are optically thick, and $H_2$ vibrational transitions dominate the cooling. Few $2\mu$ line photons escape, though, for the extinction is so great that the torus is only transparent longward of $\sim 60\mu$. Most of the energy absorbed from the nuclear continuum is transformed into thermal radiation of characteristic temperature $\sim 350 L_{44}^{1/4} r_{pc}^{-1/2}$ K. At wavelengths longer than $\sim 60\mu$, where the torus is optically thin, the spectrum reverts to whatever the nucleus supplies. Again, no one has yet performed the detailed transfer calculation necessary to predict the infrared spectrum as a function of viewing angle.

Particularly interesting potential diagnostics of conditions inside the torus are the $H_2O$ mega-masers which are observed to lie so close to the nucleus in NGC 1068 (Claussen and Lo 1986), and possibly other Seyfert galaxies as well (Baan 1985). The intensity of infrared radiation is so great that photon-mediated transitions completely dominate the population equilibrium of molecular excited states. Therefore, the most likely pumping mechanism for these masers is some deviation from exact equilibrium in the spectrum. More work is required to follow up this suggestion, or invent others. It is also possible, particularly if the X-ray intensity is diminished a factor $\sim 10^{-4}$ by scattering and Compton losses, that the hot electron/cold neutrals pumping scheme of Kylafis and Norman (1986, 1987) operates.

## 3. Torus Dynamics

Three geometric properties of the torus can be measured: the mean equatorial column density $C\langle N_{cl}\rangle$, where $C$ is the mean covering factor of clumps and $\langle N_{cl}\rangle$ is the mean clump column density; its geometrical thickness $h/r$, where $h$ is its vertical scale height and $r$ is the cylindrical radius; and its inner radius $r_t$. In their 1988 paper, Krolik and Begelman attempted to explain all of these.

They suggested that the product $C(h/r)$ is fixed by a competition between cloud-cloud collisions, which can both merge clumps and generate smaller fragments, and tidal shearing, which breaks them apart. It happens that the density inferred for the clumps ($\sim 10^7 \mathrm{cm^{-3}}$) is quite near the Roche limit for an Eddington-limited central object, so the tidal destruction time even for self-gravitating clouds is close to an orbital period. In fact, self-gravity is quite likely to be significant, for the Jeans column density is $\simeq 6 \times 10^{24} (p/10^{10}$ K $\mathrm{cm^{-3}})^{1/2}$. In order for collisional merger to balance tidal shearing, the rate of collisions must then be nearly one per orbital period, which, in turn, sets the vertical covering factor $C(h/r) \sim 1$. The similarity of a number of different rates: orbital frequency, self-gravitational free-fall, sound-wave crossing, and collisions means that these clouds will never be able to attain an equilibrium structure.

These collisions transport angular momentum and dissipate cloud orbital energy, so there is a net mass inflow. In a steady state, this inflow matches whatever is being fed into the torus from the outside. Because the local rate of mass transport depends on the mass per cloud (given the orbital and random speeds, and the cloud-cloud collision rate), matching $\dot{M}_{in}$ to the outer boundary condition gives the mean cloud thickness:

$$\langle N_{cl}\rangle \simeq 7 \times 10^{23} \left(\frac{v_{orb}}{200\mathrm{km\ s^{-1}}}\right)^{-1} \left(\frac{\Delta v}{v_{orb}}\right)^{-2} r_{pc} \left(\frac{\dot{M}_{in}}{M_\odot \mathrm{yr^{-1}}}\right) \mathrm{cm^{-2}}, \tag{1}$$

where $v_{orb}$ is the speed required for a circular orbit at radius $r$, $\Delta v$ is the 3-d velocity dispersion of the clouds, and we have normalized $\dot{M}_{in}$ to $1 M_\odot$ yr$^{-1}$ because that seems to be the approximate mass outflow rate in the wind (see §4). This estimate justifies the assertion in the preceding paragraph of the relevance of self-gravity.

The hardest dynamical problem having to do with the torus is explaining the large $\Delta v$. Stellar stirring is almost certainly inadequate. Even if new stars were created inside the torus at a rate equal to $\dot{M}_{in}$, and all of them blew up as type 2 supernovae, the energy input would still only be marginally sufficient to compete with the losses in cloud-cloud collisions.

The orbital motion of the clouds is potentially a much larger energy reservoir; the problem is to tap it. Krolik and Begelman (1988) pointed out that, to the degree that the collisions have any elasticity, orbital energy is transferred into cloud random motions (*i.e.*, , cloud "heat") by the same collisions whose inelasticity drains energy from their random motions. In fact, the two processes scale identically with $\Delta v$, and the maximum kinematic viscosity possible for orbiting particles is achieved when, just as in this case, $C(h/r) \sim 1$ (Goldreich and Tremaine 1978). Quantitatively, in order for heating to win out over cooling, it is only necessary that

$$(1 - \epsilon_{diss}) > \epsilon_{diss} \left( \frac{d \ln \Omega}{d \ln r} \right)^{-2}, \tag{2}$$

where the inequality is in fact *modulo* an unknown dimensionless factor of order unity, $\epsilon_{diss}$ is the fractional inelasticity per collision, and $\Omega$ is the orbital frequency. In order to accomplish this, they speculated that the clouds were sufficiently strongly magnetized (plasma $\beta$, or ratio of gas to magnetic pressure, $\lesssim 0.1$) that the shocks created in cloud-cloud collisions do most of their work on the (elastic) field rather than the (inelastic) gas pressure.

At the inside edge of the torus, exposed to the full blast of the nuclear ionizing continuum, the temperature must be a great deal higher than farther out, where dust and photoionization opacity shelter the gas. In fact, if we define the inner edge of the torus as that place where the local ionization parameter $\Xi$ is exactly the critical one $\Xi_c^*$ above which no cool equilibrium exists, we can immediately estimate the ablation rate:

$$\dot{M}_{abl} \sim \frac{h}{r} L_{44} \left( \frac{c_{s*}}{30 \mathrm{km\ s}^{-1}} \right)^{-1} M_\odot \mathrm{yr}^{-1}, \tag{3}$$

where $c_{s*}$ is the sound speed at the sonic point of the evaporating flow. If clouds evaporate independently, $c_{s*}$ is determined by atomic physics; if the streams merge before passing through the sonic point, then it is determined by a balance between Compton heating and expansion cooling (Balbus and McKee 1982; Begelman, *et al.* 1983; Krolik and Begelman 1986).

The form of equation 3 suggests that the inner ablation rate and the outer feeding rate of the torus are entirely independent. If that were the case, any balance would be completely accidental, and the torus would either extend all the way to the center (if $\dot{M}_{in} > \dot{M}_{abl}$) or be rapidly burned away (if $\dot{M}_{abl} > \dot{M}_{in}$). Even if $c_{s*}$ is related to $r_t$ (the merged stream case), the equilibrium is unstable. The only way to resolve these paradoxes is to couple $\dot{M}_{abl}$ to $L$.

In fact, such a coupling is quite plausible. Suppose, for example, that the luminosity is derived from the accretion of a fraction of the ablated material, and that fraction is a function of $r_t$. If the accretion efficiency is as high as $\sim 10\%$, only a few percent of the mass efflux in the wind suffices. Given that the distribution of vertical cloud velocities is quite broad, it would seem reasonable that the distribution of azimuthal cloud velocities, *i.e.*, specific angular momenta, is similarly broad. Then some fraction of the clouds would have angular momenta very much smaller than required to support a circular orbit at $r_t$. The material from these clouds could go very close to the central axis, and would be much more likely to be accreted than gas with higher specific angular momentum. By the same argument, the accreted fraction would decrease as $r_t$ increases. For stability in the $\dot{M}_{abl} = \dot{M}_{in}$ equilibrium, all that is needed is for the logarithmic derivative of the capture fraction with respect to $r_t$ to be less than $-1/2$. Numerical hydrodynamics simulations are required to make these guesses quantitative.

If the assumption is made that a portion of the evaporated material is captured, then a special lengthscale is also picked out, the accretion radius at the characteristic temperature set by the balance between Compton heating and adiabatic cooling:

$$S = 0.26 L_{44} \left(\frac{L}{L_E}\right)^{-3} \text{pc}. \tag{4}$$

Thus, the position of the inner edge is fixed by balancing evaporation with inflow, provided that the toroidal inflow also contains the accretion flow. It is also of the right order of magnitude to be consistent with the observationally-based estimates of $r_t$, provided $L/L_E$ is not too small. This argument, taken as a lower bound on $L/L_E$, also places an upper limit on the time over which the active nucleus has been accreting (Krolik and Begelman 1986).

## 4. Thermal and Ionization State of the Reflection Region

The job of describing the thermal and ionization state of an optically thin cosmic abundance plasma largely boils down to finding just two numbers: its ionization parameter $\Xi$ and its temperature $T$. These can, of course, be discussed in terms of location in the $\Xi - T$ plane (Figure 1). If it is in thermal balance, $\Xi$ determines $T$ (though occasionally multiple solutions are possible). Work in this area has concentrated on finding observational ways to bound $\Xi$ and $T$, and then to explain those numbers from dynamical considerations (§5).

We begin by estimating $\Xi$:

$$\begin{aligned}\Xi &\equiv \frac{L_{ion}}{4\pi r^2 c n_H kT} \\ &= \frac{L_{ion}\sigma_T \Delta r/r}{4\pi r c \tau kT} = \frac{L_{ion}\sigma_T}{4\pi r c k T f_{refl}} \frac{\Delta r}{r} \frac{\Delta\Omega_{refl}}{4\pi} \\ &\simeq 10 L_{ion,44} r_{pc}^{-1} T_6^{-1} \left(\frac{f_{refl}}{0.01}\right)^{-1} \left(\frac{\Delta r}{r}\right) \left(\frac{\Delta\Omega_{refl}/4\pi}{0.2}\right),\end{aligned} \tag{5}$$

where $L_{ion}$ is the portion of the luminosity between 1 and 1000 Ryd, $n_H$ is the density of hydrogen nuclei, $f_{refl}$ is the fraction of the nuclear photons reflected, $\tau$ is the electron scattering optical depth across the reflection region, $\Delta r$ is its geometrical thickness, and $\Delta\Omega_{refl}$ is the solid angle subtended by it. $f_{refl} = (\Delta\Omega/4\pi)\min(1,\tau)$ is found by comparing the reflected broad H$\beta$ or hard X-ray luminosity in type 2 Seyfert galaxies to the bolometric, and supposing that the intrinsic ratio is the same as in type 1 galaxies; the estimate of 0.01 comes from applying both methods to NGC 1068 (Antonucci and Miller 1985, Elvis and Lawrence 1988). $\Delta\Omega_{refl}$ must certainly be less than the solid angle $\Delta\Omega_{open}$ about the nucleus which is not obscured by the torus; the fraction of Seyfert galaxies with strong, thermal infrared continua which are type 1 (see, *e.g.*, Edelson, *et al.* 1987) gives the latter quantity.

A second way to bound $\Xi$ and $T$ is by X-ray spectroscopy. Krolik and Kallman (1987) predicted that type 2 Seyfert galaxies should display Fe K$\alpha$ lines of very large equivalent width ($\simeq 0.5[Fe/H]/[Fe/H]_\odot$ keV) because we are able to see the full fluorescent line luminosity, but

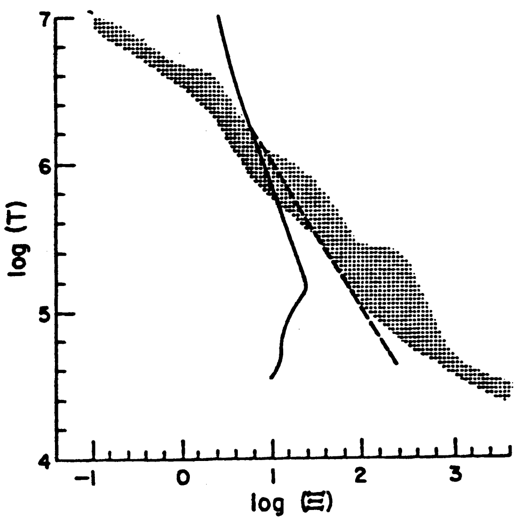


Fig. 1 The $\Xi - T$ plane. The solid line is the radiative equilibrium curve; the broken line is the line $\Xi = 10(T/10^6\mathrm{K})^{-1}$ (*cf.* equation 5); within the stippled region the mix of Fe ionization states is consistent with the energy of the K$\alpha$ line seen in NGC 1068.

only $f_{refl}$ times the hard X-ray continuum which is responsible for generating it. As Koyama reports in this volume, the *Ginga* satellite has confirmed this prediction. Moreover, the energy of the K$\alpha$ line places a limit on the range of ionization states contributing to it: FeXIX – FeXXIII. That range of ionization states corresponds to a band in the $\Xi - T$ plane which largely overlies the line $\Xi = 10/T_6$, confirming the estimate of equation 5.

As Miller discusses in his article, a comparison of the Balmer line profiles in the nuclear reflected spectrum with the profiles in the external dust-reflected spectrum suggests $T \sim 10^5$K. If that is so, the energy of the X-ray line requires $\Xi \sim 100$. Two arguments suggest that this is unlikely. The first is that the heating time in that region of the diagram is very short compared to the dynamical time: $t_{heat}/t_{sound} \simeq 5 \times 10^{-5} T_5^3 (\Xi/10) L_{ion44}^{-1} r_{pc}$. The second is that if the clouds' total flux albedo is $g$, then their ionization parameter will always be less than $g^{-1}$. In order for the reflection region $\Xi$ to be as high as $\sim 100$, the gas pressure must have dropped by at least $100g$ between the evaporation surface and the reflection region. Since $g$ is likely to be not too much less than unity, this seems a bit extreme.

## 5. Dynamics of the Reflection Region

All of the estimates of the preceding section agree in one crucial respect: they place the reflection region either near the radiative equilibrium curve, or on the side of net heating. Because the heating time is so short, the gas cannot be stationary; in fact, supersonic motion is called for. If the Mach number is sufficiently high and the streamlines diverge, adiabatic cooling can ultimately

https://doi.org/10.1017/S0074180900141063 Published online by Cambridge University Press

be competitive with radiative heating. When this is the case (Krolik and Begelman 1986), the temperature tops out at

$$T_{char} \simeq 3 \times 10^6 L_{44}^{2/3} r_{pc}^{-2/3} K, \tag{6}$$

and conservation of momentum flux from the subsonic base of the wind to its supersonic exterior gives a mass efflux of:

$$\dot{M}_{wind} \simeq 0.5 L_{44}^{2/3} r_{t,pc}^{1/3} M_\odot \mathrm{yr}^{-1}. \tag{7}$$

In such a wind $\tau \simeq 0.05$, in line with the previous estimates.

However, the argument of the previous paragraph assumed that the trajectory the gas follows in the $\Xi - T$ plane is far from the equilibrium curve. This is not necessarily the case. If the maximum $\Xi$ at which *molecular* gas can maintain an equilibrium is less than that at which *photoionized* gas can (as suggested by the calculations of Lepp, *et al.* 1985), then in the subsonic portion of the wind, where the gas is still not far from hydrostatic, its vertically rising track in the $\Xi - T$ plane will intersect the photoionized equilibrium curve below $\Xi_c^*$. From there on, the temperature rises at a rate which is controlled by the rate of change of the pressure, and the slope of the equilibrium curve (which, of course, is very steep near $\Xi_c^*$). If this case applies, heating is retarded for as long as the pressure stays relatively high. Because the curves of constant effective (*i.e.,* including conserved specific angular momentum) potential probably run more or less perpendicular to the torus's inner edge initially, a slow fall in pressure along streamlines will only occur if the rate of evaporation at high latitude is relatively great. Sooner or later, the streamlines must bend sharply outward and move more nearly perpendicular to the equipotential surfaces. In fact, a review of equation 5 shows that the smaller $\Xi$ and $T$ suggested in this picture also require $(\Delta r/r)(\Delta\Omega_{refl}/4\pi)$ smaller than 0.2, as would be produced by a sharp outward bend in the streamlines and a concentration of evaporation at high latitudes. Once again, a proper hydrodynamical calculation (*i.e.,* numerical simulation, given the difficulty of doing this analytically) is necessary to determine which picture is the physical one.

## 6. Summary

As we have just shown, theoretical arguments are already partners with observational results in illuminating the intermediate zone. The obscuring torus is probably composed of dusty, thick ($\langle N_{cl} \sim 10^{24}\mathrm{cm}^{-2}$) , warm ($\sim 10^3$ K), ionized ($x_e \sim 10^{-3}$) molecular clouds which maintain only a tenuous individual integrity against collisions and tidal shearing. Subjected to a stirring whose origin remains somewhat mysterious, but may be connected with orbital shear flow, they slowly drift in towards the center, carrying a total current of $\sim 1 M_\odot$ yr$^{-1}$. At about 1 pc from the nucleus itself, they are suddenly exposed to the full nuclear continuum, and evaporate. Most of the material rushes out in an ionized wind, but a small part, most likely that segment with the least specific angular momentum, is captured by the nucleus and ultimately powers the central engine.

## Acknowledgements

This work was partially supported by NASA Grant NAGW-1017. I am also grateful to Joe Miller and Ski Antonucci, and my collaborators, Mitch Begelman, Tim Kallman, and Steve Lepp, for many stimulating and enlightening conversations.

## References

Antonucci, R.R.J. and Miller, J.S., 1985, *Ap. J.* **297**, 621.

Balbus, S.A. and McKee, C.F., 1982, *Ap. J.* **252**, 529.

Begelman, M.C., McKee, C.F., and Shields, G.A., 1983, *Ap. J.* **271**, 70.

Burg, R., 1987, private communication.

Claussen, M.J. and Lo, K.-Y., 1986, *Ap. J.* **308**, 592.

Edelson, R.A., Malkan, M.A., and Rieke, G.H., 1987, *Ap. J.* **321**, 238.

Elvis, M. and Lawrence, A., 1988, *Ap. J.* in press.

Goldreich, P. and Tremaine, S., 1978, *Icarus* **34**, 227.

Krolik, J.H. and Begelman, M.C., 1986, *Ap. J. Lett.* **308**, L55.

Krolik, J.H. and Begelman, M.C., 1988, *Ap. J.* **329**, 702.

Krolik, J.H. and Kallman, T.R., 1987, *Ap. J. Lett.* **320**, L5.

Krolik, J.H. and Lepp, S., 1988, *Ap. J.* submitted.

Kylafis, N.D. and Norman, C.A., 1986, *Ap. J. Lett.* **300**, L73.

Kylafis, N.D. and Norman, C.A., 1987, *Ap. J.* **323**, 346.

Lepp, S., McCray, R.M., Shull, J.M., Woods, D.T., and Kallman, T.R., 1985, *Ap. J.* **288**, 58.

Miller, J.S., 1988, in *Active Galactic Nuclei*, eds. P.J. Wiita and H.R. Miller (Springer-Verlag: Berlin)

Mushotzky, R.F., 1987, private communication.

Wilson, A.S. and Heckman, T.M., 1985, in *Astrophysics of Active Galaxies and Quasi-Stellar Objects* (University Science Books: Mill Valley CA)

## DISCUSSION

G. BURBIDGE How much mass is contained in the dust, and where did it originate?

KROLIK The total mass of the torus is $\sim 10^5$ $M_\odot$. Its most likely source is the large reservoir of molecular gas found in the starburst region which often (as in the case of NGC 1068) exists in the central portion of the host galaxy. A mechanism (a la Shlosman's paper) which combines non-axisymmetric perturbations in the gravitational potential and dissipation in the gas is a very promising way to take this material from $\sim$ 100 pc to approximately a few pc.

WHITTLE I do not understand how you can argue on the one hand that the disk thickness requires highly supersonic cloud motion, but on the other hand that cloud collisions result in mergers, rather than shocked disruption.

KROLIK Because the cooling time of shocked material is shorter than the duration of the collision, collisions are highly inelastic, and merger of at least the physically overlapping parts of the clouds is likely. If there is substantial magnetic pressure inside the clouds, the collisions might be more elastic and result in greater fragmentation.

https://doi.org/10.1017/S0074180900141063 Published online by Cambridge University Press