CAN THE FORMATION OF X-RAY–OBSCURING TORI AND JETS IN ACTIVE GALAXIES BE DETERMINED BY ONE PARAMETER?

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ABSTRACT

A torus of reduced differential rotation can form in the inner \( \lesssim 10 \) pc core of active galactic nuclei (AGNs), incurring a density enhancement that can account for obscuration of X-rays in Seyfert galaxies when the initial inner core–to–black hole mass ratio is \( \gtrsim 50 \). The same density enhancement can also lead to dynamo growth of significant poloidal magnetic fields when accreted onto the central engine. Since radio jet models often employ poloidal fields as agents in extracting power for the jet luminosity, we suggest that jetted AGNs might require this poloidal field production. Although the poloidal field would originally be produced in the obscuring torus, jetted objects could exist without obscuring tori: The poloidal field would only aid in powering jet emission after it accretes with the torus matter to the central engine. Thus, only during the relatively short torus accretion timescale could there be both a jet and a torus.

Subject headings: galaxies: active — galaxies: jets — galaxies: magnetic fields — radio continuum: galaxies

1. INTRODUCTION

At the centers of galaxies, and particularly in active galactic nuclei (AGNs), are likely black holes (BHs) that power the central luminosity by accretion (e.g., Rees 1984). Time variability and the estimated accretion efficiency seem to require black hole masses \( \gtrsim 10^7 M_\odot \) in many jetted and jet-free sources. Little is known about the detailed mass distribution in the “inner core” (IC) \( \lesssim 10 \) pc regions of all galaxy types, but in our own Galaxy the rotation curve seems to incur a change from Keplerian to flat (Genzel & Townes 1987). The transition region is quasi-rigidly rotating. This leads to a density enhancement over a purely Keplerian curve, as we show below, and may explain the presence of the circumnuclear ring (CNR) of the Galaxy and X-ray–obscuring Seyfert tori (Yi, Field, & Blackman 1994; Duschl 1987), in accordance with unified models of AGNs (e.g., Antonucci 1993).

We show that such a region also favors the dynamo production of the poloidal magnetic field (PMF) to a magnitude that, when accreted onto AGN central engine gas, can approach \( \sim 10^4 \) G—in agreement with that inferred by other equipartition estimates (Begelman, Blandford, & Rees 1984). Such a PMF is often required in jet models (Blandford & Znajek 1987; Lovelace, Wang, & Sulkanen 1987; AppI & Camenzind 1993; Lynden-Bell 1995). The PMF dynamo growth timescale is much smaller than the torus depletion time, so whether a significant PMF is produced depends only on the initial IC-to-BH mass ratio. Although the jet PMF originates in the torus, the field can only play a role in jets after accreting to the central engine. Only during the time when the torus is depleting could there be both a jet and a torus in one object. This is consistent with the fact that direct evidence for tori comes mainly from radio-quiet objects (Urry & Padovani 1995), but more data are needed.

2. ESTIMATION OF TIMESCALES AND ADIABATIC BLACK HOLE GROWTH

Rotation curves of spiral galaxies show quasi-rigid rotation in the inner few kiloparsecs, and flat rotation curves outside a few kiloparsecs (Oort 1978; Binney & Tremaine 1987). Models that account for observed galaxy gas rotation curves seem to require (Binney & Tremaine 1987) (1) a central BH, (2) a stellar IC within a few to \( 10 \) pc, (3) a more diffuse nuclear bulge of several \( 100 \) pc, and (4) an isothermal sphere of dark matter on kiloparsec scales. Here we are interested in the first, the IC region substructure to the overall rotation curves, which is likely similar for all galaxy types.

We first show that a typical BH grows by accretion slowly compared with orbital timescales but rapidly compared with the IC relaxation time, so the hole’s growth is nearly adiabatic. At any time, the accreting region will be in an approximately steady state if the viscous timescale \( \tau_{\text{vis}} \ll \tau_\beta \), where \( \tau_\beta \) is the BH growth timescale. An estimate for \( \tau_{\text{vis}} \) is

\[
\tau_{\text{vis}} \sim \frac{R^3}{\nu} \sim \frac{R_d}{V_r} \sim \frac{R_d}{10^{-3} V_\phi},
\]

\[
\sim 5 \times 10^6 \text{ yr} \left( \frac{R_d}{3 \times 10^7 \text{ cm}} \right)^{3/2} \left( \frac{M_\text{H}}{10^7 M_\odot} \right)^{-1/2},
\]

where \( V_\phi \) is the rotational velocity, \( V_r \) is the radial velocity, and \( R_d \) is the radial extent of the accretion flow. We estimate the BH growth time by assuming that the objects radiate at Eddington luminosity \( L_{\text{Edd}} \). For a central BH of mass \( M_{\text{BH}} \), \( L_{\text{Edd}} = 4 \pi G M_{\text{BH}} c / k \sim 1.3 \times 10^{38} (M_{\text{BH}}/M_\odot) \) ergs s\(^{-1}\), where \( k \) is the Thomson opacity, \( \sim 0.4 \) cm\(^2\) g\(^{-1}\). As the AGN radiates at \( \sim L_{\text{Edd}} \), we have \( M_{\text{Edd}} = 2.2 \times 10^{-7} \) \( M_\odot \) yr\(^{-1}\), where \( \epsilon \ll 1 \) is the efficiency factor for radiation by accretion. Then

\[
\tau_{\text{vis}} \ll \tau_\beta \approx M_{\text{BH}}/M_{\text{Edd}} = 4.5 \epsilon \times 10^6 \text{ yr}.
\]
We must also compare $\tau_\sigma$, the gravitational relaxation time of the IC region, with $\tau_\Omega$, the orbital timescale. When the BH is small, its radius of influence $R_{\text{in}} \sim GM_{\odot}/\sigma^2$ satisfies $R_{\text{in}} < R_{\text{ic}}$, where $R_{\text{ic}}$ is the IC radius and $\sigma$ is the stellar velocity dispersion, so that $\tau_\sigma$ is independent of $M_{\odot}$. When $R_{\text{in}} \gg R_{\text{ic}}$, the relaxation time depends on $M_{\odot}$. In the Fokker-Planck approximation, for the case $R_{\text{in}} < R_{\text{ic}}$ we have (Binney & Tremaine 1987)

$$\tau_\sigma \sim 10^{11} \text{ yr} \left(\sigma/100 \text{ km s}^{-1}\right)^2 \left(10^5 M_{\odot} \text{ pc}^{-3}/\rho_{\odot}\right)^{-1}$$

(3)

where $\rho_{\odot}$ is the IC mass density. For $R_{\text{in}} \gg R_{\text{ic}}$ the relaxation time is

$$\tau_\sigma \sim \sigma^3/(GM_{\odot} \rho_{\odot}) = \sigma^3 \left(4\pi R_{\text{in}}/3\right)/(G^2 M_{\odot}^2 N)$$

$$\sim 10^{13} \text{ yr} \left(M_{\odot}/10^8 M_{\odot}\right)^3 \left(\sigma/100 \text{ km s}^{-1}\right)^{-3} (N/10^6)^{-1}$$

(4)

where $N$ is the number of IC stars.

Because of equations (2), (3), and (4), $\tau_\sigma \ll \tau_{\text{H}} \ll \tau_{\text{n}}$, where $\tau_{\text{H}}$ is the Hubble time, $\sim 10^{10}$ yr, and $\tau_{\text{n}} \approx 10^8$ yr ($R_{\text{ic}}/10$ pc) ($\sigma/100$ km s$^{-1}$)$^{-1}$, the adiabatic approach is appropriate.

Young (1980) considers the adiabatic evolution of an isothermal sphere with a growing BH. Quinlan, Hernquist, & Sigurdsson (1995) confirm the results of Young (1980) and extend them to nonisothermal spheres. As applied to the IC, these are reasonably consistent with a total mass dependence on the radius given by

$$M_{\text{tot}}(r) = M_{\text{h}} \left[1 + m(r/R_{\text{ic}})^{n/3}\right]$$

$$= M_{\text{h}} \left[1 + m_{\text{ic}} r/R_{\text{ic}} \right], \quad r \leq R_{\text{ic}},$$

$$= M_{\text{h}} + m_{\text{ic}} r/R_{\text{ic}}, \quad r > R_{\text{ic}},$$

where $m_{\text{ic}} = M_{\text{ic}}/M_{\text{h}}$ is the ratio of BH to stellar core mass, and $n \sim 3$.

3. GAS DENSITY

The surface gas density is given by $\Sigma(r) = 2H(r)\rho(r)$, where $H(r)$ is the half–scale height of the gas and $\rho(r)$ is the gas density. Following the standard treatment (e.g., Pringle 1981), we take the viscosity coefficient to be $\nu = \gamma v_T H$, where we assume $\gamma \leq 1$ is the viscosity parameter and $v_T$ is the turbulent energy speed associated with an external source of turbulent energy. Combining this with the conservation of gas mass and angular momentum, the surface density satisfies (Pringle 1981; Yi et al. 1994)

$$d\Sigma/dr + f(r)\Sigma = g(r),$$

(6)

where $f(r) = (r^3 \Omega')/(r^3 \Omega^3)$, $g(r) = -[\dot{M}/(2\pi \nu \Omega^3)] [2\Omega(\Omega') + 1]$, and $\Omega$ is the angular velocity determined by the potential. The prime indicates $d/dr$. The solution of equation (6) is given by

$$\Sigma(r) = \exp \left[ - \int_{r_0}^{r} ds f(s) \right] \left[ \Sigma_0 + \int_{r_0}^{r} ds g(s) \exp \left[ \int_{r_0}^{s} d\lambda f(\lambda) \right] \right],$$

where $\Sigma(r_0) = \Sigma_0$. To find the relationship between the surface density and the stellar mass density, we note that, for circular orbits, $\Omega = (GM/r^3)^{1/2}$. Plugging in the expressions for $f(r)$ and $g(r)$, using equation (5), $S = r/R_{\text{ic}}$, and $\lambda = m^{1/2}$, we have

$$\Sigma(r) = \frac{1 + m S^n_{\text{h}}}{R_{\text{ic}}^2 \left(n - 3\right) m S^n_{\text{h}} - 3 S^n_p}$$

$$\times \frac{R_{\text{ic}}^n [m - 3 m S^n_p^{1/2} - S^n_p]}{1 + m S^n_{\text{h}}}
\int_{m S^n_{\text{h}}}^{m S^n_{\text{h}} + 1} \frac{d\alpha}{(\lambda + \alpha^{n+1})^{1/2}}.$$  

(7)

where $K = \dot{M}m^{(1 - 1/n)2} \left(G m_{\text{ic}} R_{\text{ic}}^{1/2}/(2\pi \nu \Omega^3)\right)$.

The third term in equation (7) is negative, so the second term gives the correct order of magnitude. Figure 1 shows the surface density at $r = R_{\text{ic}} \sim R_{\text{c}}$, the edge of the torus (outside of which the potential drops), for four values of $n$ as a function of $m$, and vice versa. For large $m$ and $n \sim 3$ the density can be 100 times the Keplerian (i.e., when $m \ll 1$) value. That $n \sim 3$ is consistent with Young (1980) and Quinlan et al. (1995). This large density enhancement can easily account for X-ray obscuration in Seyfert galaxies (Yi et al. 1994).

4. RIGID ROTATION AND POLOIDAL FIELD GROWTH

As shown above, the enhanced density torus results from reduced differential rotation. We can calculate the reduction in $\Omega'$ by setting $\Omega = [GM_{\text{ic}}(r) r^{3/2}]^{1/2}$. Using equation (5) with $n = 3$ gives $\Omega'(R_{\text{ic}}) = -(3/2) G^2 R_{\text{ic}}^{1/2} M_{\text{h}}^{1/2}/(1 + m)^{1/2}$. For $m > 50$ this gives more than a factor of 7 reduction in $\Omega'$ from the Keplerian value. This region may be important for the dynamo generation of the PMF for radio sources, as we now describe.

The mean field dynamo theory (Moffatt 1978; Parker 1979) splits the velocity and the magnetic field into mean ($\bar{V}, \bar{B}$)
and fluctuating \((r', b')\) components. The time evolution of mean magnetic field is given by (Moffatt 1978; Parker 1979)

\[
\dot{B} = \nabla \times [\dot{V} \times \vec{B} + \alpha \vec{B} - (\lambda_m + \beta) \nabla \times \vec{B}],
\]

where \(\alpha\) and \(\beta\) are the helicity and diffusion dynamo coefficients, respectively, and are functions of the turbulent velocity. A nonvanishing \(\alpha\) is the result of buoyant eddies rising in an upwardly decreasing density gradient while conserving their angular momenta. The magnetic viscosity, \(\lambda_m\), satisfies \(\lambda_m \ll \beta\).

For small \(\Omega\), the “\(\alpha^2\) dynamo” (Moffatt 1978) is favored because the maximum growth rate depends on \(\alpha^2\), as we will see. Sufficient reduction of \(\Omega\) means that the radial PMF produced by the \(\alpha\) effect is not sheared into the toroidal field. Simulations (e.g., Donner & Brandenburg 1990) show that dipole modes are favored in such a dynamo, in contrast to the \(\alpha - \Omega\) dynamo often applied to disks (Parker 1979).

To estimate when the \(\alpha^2\) dynamo is favored, we work in cylindrical coordinates \((r, \phi, z)\) and write \(\vec{V} = r \Omega (r, z) \hat{k}_\phi, \vec{B} = \vec{B}_r(r, z) \hat{e}_r + \vec{B}_\phi,\) where the subscript \(P\) indicates the poloidal \((r, \phi)\)-component and the subscript \(\phi\) indicates the toroidal component. \(\vec{B}_r = \vec{V} \times \vec{A} \hat{r}(r, z) \hat{e}_\phi.\) Assuming \(\alpha\) and \(\beta\) are constant, equation (8) can be written (Moffatt 1978)

\[
\dot{B}_r = \alpha \vec{B}_r - \beta (\nabla \times \vec{B}) \hat{r},
\]

and

\[
\dot{A}_\phi = \alpha B_r + \beta (\nabla \times \vec{B}) \hat{\phi}.
\]

An \(\alpha^2\) dynamo will dominate the \(\alpha - \Omega\) dynamo when the second source term on the right of equation (12) dominates the first, that is, when \(\alpha r >> |\Omega|\). From Parker (1979), we have \(\alpha \sim 0.4 \Omega r\). For a turbulently supported dust torus, observations require \(v_r / V_\phi \sim 0.5\), where \(V_\phi = r \Omega\) (Krolik & Begelman 1988; Urry & Padovani 1995). Using these and equation (5), the requirement near \(r = R_{\text{depl}} \sim R_r\) becomes \(0.2 \alpha^2 M_{\text{disk}}^{1/2} (1 + m)^{1/2} > (3/2) M_{\text{torus}}^{1/2} (1 + m)^{1/2}\), or

\[
m \gg 6.5.\]

When equation (11) is satisfied, we can ignore the first term on the right in equation (8) near \(r \sim R_{\text{depl}}\). We capture the essence of an \(\alpha^2\) dynamo by considering the case when the \(z\)-variation dominates and by assuming solutions of the form \(\vec{B}_{\phi, r} \propto \exp (\gamma t + k_z z)\). Plugging these into equations (9) and (10) gives

\[
\gamma \vec{B}_r = \alpha \vec{A} k_z^2 - \beta \vec{B}_r k_z^2,
\]

and

\[
\gamma \vec{A} = \alpha \vec{B}_r - \beta \vec{A} k_z^2,
\]

so that the growth rate is given by

\[
\text{Re} [\gamma] = -\beta k_z^2/2 + 3k_z \alpha/2.
\]

The growth rate is positive if \(k_z < 3 \alpha/\beta\). The maximum growth rate is \(\text{Re} [\gamma]_{\text{max}} = (9/8) \alpha^2/\beta\), showing the \(\alpha^2\) dependence. The second term on the right of equation (14) provides a more conservative estimate.

Let us see why an \(\alpha^2\) dynamo favors the PMF. For an AGN torus, the fluctuation scale is determined by the size of dust-containing clouds. The dust must be in clouds because it could not survive if the random velocities of \(\sim 100 \text{ km s}^{-1}\) were thermal. Thus, the cloud size \(r_c\) satisfies \(r_c < R_r\), where the torus radius \(R_r\) is the scale for the variation of the mean quantities. We can estimate the dust cloud size from observations of the CNR of our Galaxy (Genzel & Townes 1987), which is thought to be similar to the AGN dust tori (Krolik & Begelman 1988). These observations (Genzel & Townes 1987) show clouds with \(0.1 \lesssim r_c \lesssim 0.25\) pc. Now we estimate the PMF produced: Setting \(\text{Re} [\gamma] \sim k \alpha \text{ and using } k_r \sim B,\) with \(\alpha \sim 0.4 \Omega r\) and \(\beta \sim (1/3) v_r / r,\) in equations (12) and (13), gives \(B_r \sim B,\) for the \(\alpha^2\) dynamo. The equations analogous to equations (12) and (13) for the \(\alpha - \Omega\) dynamo, derived by keeping the first term on the right of equation (8) and dropping the term linear in \(\alpha\), give \(B_r / B_\phi \sim r / r_c\). Thus, \(B_r / B_\phi \sim 300,\) showing that the \(\alpha^2\) dynamo favors the PMF in comparison to the \(\alpha - \Omega\) dynamo. This ratio is important for jet models, particularly when the resulting luminosity depends on \(B^2\) (e.g., Blandford & Znajek 1977). Thus, a factor of \(50 - 100\) in the field corresponds to a factor of \(2.5 \times 10^{-10}\) in the jet luminosity.

5. POLOIDAL FIELD IN AGNS

For the PMF to be produced by a working torus dynamo, equation (11) must hold. In addition, the dynamo growth time must be shorter than the torus lifetime. That is,

\[
\tau_d < \tau_{\text{jet}} / \tau_{\text{torus}}.
\]

Note that \(\tau_d\) for the torus must be less than the field diffusion time: \(\tau_{\text{diff}} = (10 \text{ pc})^2/\beta \sim 100 \text{ kpc}^2/(100 \text{ km s}^{-1} 10 \text{ kpc}) \sim 10^7 \text{ yr}\). For a density \(\rho_r \sim 5 \times 10^{-8} \text{ g cm}^{-3}\) (Krolik & Begelman 1988) and radius \(5 \text{ pc}\) with height \(2.5 \text{ pc}\), \(M_{\text{torus}} \sim 10^7 M_{\odot}\). From equation (2) with \(\epsilon \sim 0.1, \tau_d \sim 5 \times 10^7 \text{ yr}\). Thus, violating equation (15) requires the extreme case of \(M_{\text{torus}} \sim 10^4 M_{\odot}\), so that equation (8) is more stringent.

The larger \(m\) is, the greater the density enhancement and produced PMF. Equipartition between turbulent and magnetic energy densities gives an upper limit to the field. For \(\rho_r \sim 10^{-17} \text{ g cm}^{-3}\) (Krolik & Begelman 1988), corresponding to \(m \sim 50\) and \(v_r \sim 300 \text{ km s}^{-1}\), the turbulent energy is \((1/2) \rho v_r^2 \sim 4.5 \times 10^{-16} \text{ ergs cm}^{-3}\). Setting this equal to \(B_{r, \text{torus}}^2 / 8\pi\), we have \(B_{r, \text{torus}} \sim 0.3 \text{ G}\). The field is accreted to the central engine as the torus depletes. The radial component of the PMF is then sheared, but only that fraction of the field at a much lower scale height than that of the \(\sim 5\)–\(10\) pc torus. The \(z\)-component is unaffected by the shear. An estimate of the accreted PMF can then be made from flux freezing. For an ion-electron accretion disk with height \(H_{\text{disk}} \sim 10^{13} \text{ cm}\) and density \(\rho_{\text{disk}} \sim 10^{-6} \text{ g cm}^{-3}\), flux conservation implies that the \(z\)-PMF will accrete to

\[
B_{r, \text{disk}} = 5 \times 10^3 G(B_{r, \tau} / 0.5 \text{ G}) \times (\rho_{\text{disk}} H_{\text{disk}} / 10^5 \text{ cm}^{-2})(\rho_{\tau} H_{\tau} / 10 \text{ cm}^{-2})^{-1},
\]

and the field accreted to the BH can be much larger.

Only when the torus depletes by accretion can the field produced there move to the central engine and play its role in jets. Since any jet formation time in the engine is likely much shorter than the torus depletion time, only during the latter timescale can a jetted object show both a jet and a torus. For a \(10^7 M_{\odot}\) hole accreting at the Eddington rate with \(\epsilon \sim 0.1\) and a torus mass of \(10^7 M_{\odot}\), this accretion period lasts \(\sim 10^8\) yr. Note, also, that a torus scale dynamo need not determine the final scale of the magnetic field but would mediate the initial energy extraction from the rotating central
source. The initial jet flow could drag the resulting field to kpc–Mpc scales as in Blandford & Rees (1974).

6. CONCLUSIONS

Reduced $\Omega'$ tori in the central $\leq 10$ pc regions of AGNs can obscure X-rays and incur dynamo production of the PMF likely required for jets, when $m \approx 50$ initially. PMF growth in the torus allows angular momentum transport, and the torus will accrete to the central engine carrying its field. PMF transport and torus depletion are associated processes, and objects with jets that last $10^8$ yr or more would be less likely to show obscuring tori. Jet-free AGNs would either have an initial $m < 50$ or have their jets beamed away from us. There are a few radio-loud quasars or Seyfert galaxies with inferred BH masses greater than $10^6 M_\odot$ (e.g., NGC 5548; Krolik et al. 1991). The absence of direct evidence for dust tori in the latter may be consistent with our paradigm but may require non-adiabatic analysis. This picture would suggest the following evolutionary phases for an AGN: (1) quiescent galaxy, (2) jet-free AGN with torus, (3) jetted AGN with torus, (4) jetted AGN without torus, and (5) quiescent galaxy. Differences between elliptical and spiral AGNs with respect to jets could be the result of the formation timescale for ellipticals, and differences in the ambient gas environments.

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