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HYPER- AND SUSPENDED-ACCRETION STATES OF ROTATING BLACK HOLES AND THE DURATIONS OF GAMMA-RAY BURSTS

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ABSTRACT

We analyze the temporal evolution of accretion onto rotating black holes subject to large-scale magnetic torques. Wind torques alone drive a disk toward collapse in a finite time $\sim t_{\rm ff} E_k/E_B$, where $t_{\rm ff}$ is the initial free-fall time and E_k/E_B is the ratio of kinetic energy to poloidal magnetic energy. Additional spin-up torques from a rapidly rotating black hole can arrest the disk's inflow. We associate short/long gamma-ray bursts with hyperaccretion/suspended-accretion onto slowly/rapidly spinning black holes. This model predicts afterglow emission from short bursts and may be tested by *HETE-2*.

Subject headings: black hole physics — gamma rays: bursts — gamma rays: theory

1. INTRODUCTION

Active black holes are believed to be surrounded by magnetized disks or tori. In particular, such systems are the most favorable inner engines of cosmological gamma-ray bursts (GRBs). The BATSE catalog shows a broad, bimodal distribution of GRB durations (T_{90}) centered on short bursts of ~0.3 s and long bursts of ~30 s (Kouveliotou et al. 1993; Paciesas et al. 1999). The cosmological origin of GRBs and their rapid variability indicate an association with stellar-mass progenitors (Woosley 1993; Paczyński 1991, 1997, 1998; Fryer, Woosley, & Hartman 1999). Proposed scenarios include the coalescence of compact object binaries (Paczyński 1991) and the core collapse of massive stars (Woosley 1993). Both scenarios involve a black hole disk or torus system. This disk will be magnetized because of seeding from the progenitor's magnetic flux and disk dynamo amplification (Narayan, Paczyński, & Piran 1992; Paczyński 1997; Brandenburg et al. 1995; Hawley, Gammie, & Balbus 1996; Hawley 2000). Various lines of observational evidence support the collapsar scenario, at least for long bursts (e.g., van Putten 2001); here we shall suggest that both short and long bursts can be explained in a unified model with durations extended by coupling to a high-spin black hole.

Accretion provides a definite power source for GRBs, but the black hole's spin energy, e.g., 2 $M_\odot c^2$ for a 7 $M_\odot$ black hole, can be at least as important. Efficient spin energy extraction requires an external magnetic field supported by surrounding matter. Magnetically mediated energy extraction mechanisms include radiation in DC Alfvén waves (Blandford & Znajek 1977; Thorne, Price, & Macdonald 1986, p. 145; see also, e.g., Lee, Wijers, & Brown 2000; Beskin & Kuznetsova 2000; Krolik 2001) or, as a potential source of intermittency, superradiant scattering of fast magnetosonic waves (Uchida 1997a, 1997b; van Putten 1999).

The evolution of the disk or torus surrounding a black hole depends crucially on magnetic torques. Negative torques are expected from magnetic winds, which can dramatically stimulate accretion rates (e.g., Blandford & Payne 1982; Lee & Kim 2000). Positive torques $T = -\dot{J}_{\rm BH}$ may derive from coupling to the black hole's angular momentum $J_{\rm BH}$ by equivalence in poloidal topology to pulsar magnetospheres (van Putten 1999; Brown, Bethe, & Lee 2001) and may be key to driving hypernovae (Brown et al. 2000). The disk may also redistribute its mass and angular momentum via internal, turbulent, magnetic, and Reynolds stresses (Balbus & Hawley 1998). The relative contributions from *external* (wind/black hole) stresses versus *internal* stresses is presently unknown and may depend crucially on the net poloidal flux introduced to the disk (see, e.g., Matsumoto 1999).

In this work, we focus on the duration of magnetized accretion regulated by external torques. We show that wind torques alone drive the disk to a finite-time collapse singularity. The additional positive torque from the black hole, when its angular velocity exceeds a critical value, may arrest accretion onto the black hole for the duration of its spin-down lifetime. This introduces two cases: a short-lived state of "hyperaccretion" for low-spin black holes and a long-lived state of "suspended accretion" for high-spin black holes. The respective timescales are consistent with the bimodal distribution in GRB durations for plausible values of two presently uncertain ratios: the kinetic–to–poloidal magnetic energy ratio in the disk and the black hole–to–disk mass ratio. Numerical simulations that will soon be within reach may establish the likely ranges of these ratios.

Previously, bimodality in GRB durations has been attributed to, e.g., hydrodynamic timescales in Newtonian versus relativistic reverse shocks in shells decelerating in the interstellar medium (Sari & Piran 1995) or to the formation of millisecond pulsars spinning above or below a gravitationally unstable limit (Yi & Blackman 1998). Previous studies of Poynting-dominated winds in GRBs have focused on plasma processes and possible connections to gamma-ray emission (e.g., Usov 1992, 1994; Mészáros & Rees 1997; Blackman & Yi 1998; Lyutikov & Blackman 2001).

2. ORBITAL EVOLUTION OF ROTATING, MAGNETIZED RINGS SUBJECT TO WIND TORQUES

Magnetic disk winds transport angular momentum by Reynolds and Maxwell stresses, which, for outgoing (ingoing) flows, have the same (opposite) sign. A force-free limit is defined by neglecting Reynolds stresses. Applied to outgoing winds, Maxwell stresses $-RB_rB_\phi/4\pi$ provide a lower limit on the disk's accretion rate, where B denotes the magnetic field with components in spherical coordinates (r, θ, ϕ) and $R = r\sin\theta$. The force-free limit represents pure "magnetic braking."

A magnetized disk ring, with mass δM and radius $\mathcal{R}$ in quasi-Keplerian rotation at a rate of $\Omega = (GM_{\rm BH}/\mathcal{R}^3)^{1/2}$, evolves as it sheds angular momentum in a wind. Assuming a radially asymptotic flow, the magnetic torque $\delta\tau_B$ on the ring equals (Goldreich & Julian 1969)

$$\delta\tau_B = \frac{\Omega}{c}\sin^2\theta(\partial_\mu A_\phi)\delta A_\phi. \quad (1)$$

Here $\mu = \cos\theta$, the poloidal magnetic field is $B_r = -r^{-2}\partial_\mu A_\phi$, and $2\pi\delta A_\phi$ is the magnetic flux through the ring. The toroidal field is $B_\phi = -\Omega R B_r/c$ (see the wind asymptotics of Goldreich & Julian 1969 and Michel 1969 for aligned rotators; for disk winds, see Blandford 1976 and Lee & Kim 2000). The ring evolution equation, $\delta\tau_B = \frac{1}{2}\Omega\mathcal{R}\dot{\mathcal{R}}\delta M$, yields

$$\mathcal{R}\dot{\mathcal{R}} = \frac{2}{c}\sin^2\theta(\partial_\mu A_\phi)\left(\frac{\delta A_\phi}{\delta M}\right). \quad (2)$$

The evolution of a collapsing ring depends on its degree of magnetization and the magnetic wind's geometry. A split-monopole geometry (SMG) for the asymptotic poloidal field is a good approximation when $|B_r|/|B_\phi| \gg 1$, which is true at the onset of accretion for modest rotation rates. As the disk collapses and spins up, however, magnetic field line winding produces large toroidal magnetic fields, and a force-free toroidal-field geometry (TFG) becomes more appropriate. SMG/TFG represent two limiting cases, for which simple solutions to equation (2) may be obtained.[1]

Solution for SMG.—In SMG, we have $A_\phi = A(1 - \mu/\mu_0)$, with $\theta_0 = \arccos(\mu_0)$ being the angle of the innermost streamline with respect to the pole; the outermost streamline lies on the equator. With $\varpi \equiv \mathcal{R}/\mathcal{R}_0$, where $\mathcal{R}_0$ is the initial ring radius, the solution of equation (2) describes a *finite-time singularity:*

$$\varpi(t; A_\phi) = (1 - t/t_f^{\rm smg})^{1/2}, \quad (3)$$

where $t_f^{\rm smg} = \frac{1}{4}(\delta J^0/|\delta\tau_B^0|)$, and $\delta J^0 = \delta M\Omega_0\mathcal{R}_0^2$ and $\delta\tau_B^0$ denote the initial ring angular momentum and torque.

The initial ring kinetic energy is $\delta E_k^0 = \Omega_0\delta J^0/2$. The energy in the magnetic field of a (double) wedge-shaped portion δA_ϕ rooted at $\mathcal{R}_0$ is $\delta E_B^0 = \frac{1}{2}\int B_r^2\, d\mu r^2\, dr = \delta A_\phi|\partial_\mu A_\phi|/(2\mathcal{R}_0) = |\delta\tau_B|c(2\mathcal{R}_0\Omega_0)^{-1}(\sin\theta)^{-2}$. For SMG, $\delta E_B^0 \to \delta E_{B,s}^0 \equiv A\delta A_\phi/(2\mathcal{R}_0)$. The initial free-fall time is $t_{\rm ff} = \pi/(\Omega_0 2\sqrt{2})$. For $R_S \equiv 2GM_{\rm BH}/c^2$, the Schwarzschild radius of the black hole, the collapse time in SMG becomes

$$t_f^{\rm smg} = \frac{1}{2cR_S}\frac{\mathcal{R}_0^2}{\sin^2\theta}\frac{\delta E_k}{\delta E_{B,s}} = \frac{t_{\rm ff}/\pi}{\sin^2\theta}\left(\frac{\mathcal{R}_0}{R_S}\right)^{1/2}\frac{\delta E_k}{\delta E_{B,s}}. \quad (4)$$

In the above, superscripts on the energies are dropped since $\delta E_k^0/\delta E_{B,s}^0 = (\delta M/\delta A_\phi)(GM_{\rm BH}/A)$ is an evolutionary invariant. Equation (4) shows that accretion is more efficient with increasing $\sin\theta$; outer-disk rings have magnetic field lines closer to the equator than inner-disk rings, providing greater moment arms for the magnetic torque.

Solution for TFG.—Under large magnetic hoop stresses, magnetic flux surfaces adjust their latitudes such that the *toroidal* magnetic field in the wind approaches a force-free configuration with $B_\phi \propto R^{-1}$ (see, e.g., the protostellar disk wind in Shu et al. 1995). The magnetosphere is then effectively current-free except along the pole and the equator, with the poloidal pole current I_p producing $B_\phi \approx -2I_p/(cR)$ elsewhere. Hence, $\partial_\mu A_\phi = -r^2B_r = -2I_p(\Omega\sin^2\theta)^{-1}$ so that $|\delta\tau_B| = 2I_p\delta A_\phi/c$. Here I_p relates to the total normalized flux $A \equiv \int d\mu\partial_\mu A_\phi$ by $I_p \approx -A\Omega(\mathcal{R}_{\rm in})/[2|\ln(\theta_0/2)|]$, where $\mathcal{R}_{\rm in}$ is the innermost disk radius.

From equation (2), the TFG limit gives $\Omega\mathcal{R}\dot{\mathcal{R}} = -(4I_p/c) \times (\delta A_\phi/\delta M)$. Using $\Omega(\mathcal{R}) = \Omega(\mathcal{R}_0)\varpi^{-3/2}$, we find the *finite-time singularity* solution:

$$\varpi(t; A_\phi) = (1 - t/t_f^{\rm tfg})^2, \quad (5)$$

where

$$t_f^{\rm tfg} = \frac{\delta J^0}{|\delta\tau_B^0|} = t_{\rm ff}\frac{4\mathcal{I}_0}{\pi}\frac{\delta E_k}{\delta E_{B,s}}. \quad (6)$$

Here $\mathcal{I}_0 \equiv (c/\sqrt{2})\mathcal{R}_0^{-1}\int d\theta(\Omega\sin\theta)^{-1}$ is $\sim(\mathcal{R}_{\rm in}^3/R_S\mathcal{R}_0^2)^{1/2} \times |\ln(\theta_0/2)|$. Going from $\mathcal{R}_{\rm in}$ to the outside of the disk at $\mathcal{R}_d$, the radial factor in $\mathcal{I}_0$ decreases from $(\mathcal{R}_{\rm in}/R_S)^{1/2}$ (which is order unity) to $(\mathcal{R}_{\rm in}^3/R_S\mathcal{R}_d^2)^{1/2}$ (which is of order $\mathcal{R}_{\rm in}/\mathcal{R}_d$); the product $t_{\rm ff}\mathcal{I}_0$ varies proportional to $\mathcal{R}_0^{1/2}$. Note that δE_B^0 in TFG is larger than $\delta E_{B,s}^0$ by a factor $|\partial_\mu A_\phi|/A \sim \Omega(\mathcal{R}_{\rm in})[\Omega(\mathcal{R}_0)\sin^2\theta \times |\ln(\theta_0/2)|]^{-1}$.

For the inner disk, $t_f^{\rm tfg} \ll t_f^{\rm smg}$ because of the geometrical factor $4\theta_0^2|\ln(\theta_0/2)| \ll 1$; the enhanced efficiency arises from the squeezing of the poloidal field lines toward the pole by toroidal hoop stresses. For the outer disk, $t_f^{\rm tfg} : t_f^{\rm smg} = 4(\mathcal{R}_{\rm in}/\mathcal{R}_d)^{3/2}|\ln(\theta_0/2)|$, which is of order unity or less unless the effective beaming angle θ_0 is extremely small. Thus, although inner-disk accretion is greatly enhanced when the magnetic wind becomes more stratified toward the poles, the overall lifetime set by outer-disk accretion is not as sensitive to these geometric changes.

Dimensional estimates of accretion times.—Dimensionally, the minimum inner-disk accretion time, from TFG, is

$$t_{\rm in} = 0.011\ {\rm s}\left(\frac{\mathcal{R}_{\rm in}}{100\ {\rm km}}\right)^2\left(\frac{M_{\rm BH}}{10\ M_\odot}\right)^{-1}\frac{\delta E_k}{\delta E_{B,s}}g(\theta_0), \quad (7)$$

where $g(\theta_0) = [1 + \ln(\theta_0/1°)/\ln 0.009]$ and θ_0 is the angle of the innermost wind field line. Correspondingly, the maximum outer-disk accretion time is

$$t_d = 0.057\ {\rm s}\left(\frac{\mathcal{R}_d}{1000\ {\rm km}}\right)^2\left(\frac{M_{\rm BH}}{10\ M_\odot}\right)^{-1}\frac{\delta E_k}{\delta E_{B,s}}, \quad (8)$$

assuming $\mathcal{R}_d \gg \mathcal{R}_{\rm in}$ with SMG; otherwise, if $\mathcal{R}_d \sim \mathcal{R}_{\rm in}$, we would use the TFG result, which increases the time by a factor of $\sim 4|\ln(\theta_0/2)|$ consistent with equation (7). These limits will bracket the interval over which disk accretion takes place.

3. DISK ACCRETION AND BLACK HOLE SPIN-DOWN TIMES COMPARED

A ring spends the greatest portion of its lifetime near its initial radius; during the subsequent rapid infall phase, magnetic

[1] Although, as shown by Blandford & Payne (1982) for (baryonic) MHD disk winds, field line focusing may produce nonradial asymptotic poloidal fields, cylindrically collimated MHD winds have been found to have relatively low asymptotic flow speeds (Ostriker 1997).

LIGO-P010055-00-R

flux is approximately conserved. For fiducial estimates, we shall normalize to an initial kinetic–to–poloidal magnetic energy ratio of ~100 for all rings, as found from Hawley's (2000) numerical simulations of global, hot, MHD accretion disks. This energy ratio certainly varies in real systems; estimates based on simulations with different physics or initial conditions (e.g., heating from resistive dissipation, as in Stone & Pringle 2000, and net poloidal flux, as in Hawley et al. 1996) may decrease this ratio toward unity. The natural diversity in mean magnetizations may in part be responsible for the broad distributions of the two duration classes of GRBs.

Timescale for short GRBs.—Applying $\delta E_k/\delta E_{B,s} = 10$–$100$ in equation (7) or (8) (for a disk of size $\lesssim 1000$ km), the accretion time will be a few seconds or less—comparable to that of the 2 s event GRB 000310C (Jensen et al. 2001). In this scenario, short bursts represent hyperaccretion onto a slowly spinning black hole.

Timescale for long GRBs.—What accounts for the long-duration bursts? If the black hole is initially rapidly spinning, then the burst may persist for the time $t_{\rm BH}$ required for angular momentum extraction by the Blandford-Znajek (1977) and related processes. Given a horizon flux $2\pi A_{\rm BH}$, the black hole torque for open fields is $L_{\rm BH} \sim \Omega_{\rm BH} A_{\rm BH}^2/c$ (see Thorne et al. 1986). To estimate $A_{\rm BH}$, We assume that the black hole field strength is comparable to that from the inner part of the disk. In TFG, we find that $A_{\rm BH} \approx A\,[2|\ln(\theta_0/2)|]^{-1}$, so that $L_{\rm BH} \sim \Omega_{\rm BH} A^2\,[4c|\ln(\theta_0/2)|^2]^{-1}$. For comparison, the total disk-wind angular momentum luminosity is equal to $\tau_- \equiv \int d\tau_B$; i.e., $L_d^{\rm tfg} = \Omega(\mathcal{R}_{\rm in})A^2\,[c|\ln(\theta_0/2)|]^{-1}$. The spin-down time $t_{\rm BH} = 2GM_{\rm BH}^2 r_{\rm BH}\Omega_{\rm BH}c^{-2}L_{\rm BH}^{-1}$. Here $r_{\rm BH} = R_{\rm S}\cos^2(\lambda/2)$ is the radius of the black hole horizon; $\sin\lambda \equiv a/M_{\rm BH}$ is the ratio of specific angular momentum to the maximal Kerr value.

Taking $\delta E_k/\delta E_{B,s} \sim GM_{\rm BH}M_d/A^2$, we obtain $t_{\rm BH} = (8r_{\rm BH}/c) \times (\delta E_k/\delta E_{B,s})(M_{\rm BH}/M_d)|\ln(\theta_0/2)|^2$. Assuming $a \sim M$ so that $r_{\rm BH} \sim R_{\rm S}/2$, we find that

$$t_{\rm BH} = 88\ {\rm s}\left(\frac{M_{\rm BH}}{10\ M_\odot}\right)\left(\frac{M_{\rm BH}/M_d}{100}\right)\left(\frac{\delta E_k/\delta E_{B,s}}{100}\right) g^2(\theta_0), \quad (9)$$

consistent with a bimodal offset of long- from short-duration bursts for $M_d \ll M_{\rm BH}$. Note that $t_{\rm BH}/t_{\rm in} \sim M_{\rm BH}/M_d$, as magnetic torques on the disk and black hole are comparable, but $J_{\rm BH}$ exceeds J of the inner disk by a factor of the mass ratio, assuming comparable specific angular momenta, as for a maximal-Kerr hole.

We remark that a black hole mass $M_{\rm BH}$ of about 10 $M_\odot$, formed promptly in the hypernova scenario, is expected in view of the Kerr constraint $J_{\rm BH} \leq GM_{\rm BH}^2/c$. For prompt collapse of a He core of radius $R_{\rm He}$, stripped of its hydrogen envelope and tidally locked with a binary companion of mass m (Woosley 1993; Paczyński 1997; Iben & Tutukov 1996) with an orbital separation $\xi R_{\rm He}$,

$$M > M_{\rm BH} \geq M\left(\frac{2}{5}\right)^3\left(\frac{c^2 R_{\rm He}}{GM\xi^3}\right)^{3/2}\left(1+\frac{m}{M}\right)^{3/2}\left(\frac{\bar\rho}{\rho_c}\right)^2. \quad (10)$$

Here the central density ρ_c is that extrapolated from the He mantle, with $\bar\rho$ denoting the mean density. For the canonical values of $m = M = 20\ M_\odot$, $R_{\rm He} = R_\odot$, and $\xi = 4$, a black hole may form promptly provided the inequality on the left-hand side of equation (10) is satisfied; i.e., $1 > 1.3 \times 10^3\,(\bar\rho/\rho_c)^2$. This would be true if, e.g., the He core is approximately fitted by a polytrope with index $n = 3$, for which the Lane-Emden relationship gives $\rho_c \sim 50\bar\rho$ (Kippenhahn & Weigert 1990, p. 178). Thus, rapidly rotating black holes formed in prompt collapse tend to be massive.

4. SUSPENDED ACCRETION AND LONG GRBs

Long GRBs rely on the continuing presence of the torus. Hyperaccretion times (§ 2), however, are far shorter. We speculate that the resolution of this paradox depends on the ability of the black hole to spin up the disk over interconnecting magnetic field lines equivalent in poloidal topology to pulsar magnetospheres (see Fig. 2 in van Putten 1999). The black hole torque on the torus is (adapted from Thorne et al. 1986)

$$\delta\tau_+ \approx \frac{1}{c}(\Omega_{\rm BH} - \Omega)\sin^2\theta(-\partial_\mu A_\phi)\delta A_\phi \quad (11)$$

for a flux tube of field lines between them (counting both above and below the midplane). The torus now has an inner face interacting with the black hole and an outer face interacting with infinity. The primary role of the inner interaction is to extend the lifetime of the torus by preventing its accretion onto the black hole.

The total black hole torque on the torus, τ_+, depends on the associated horizon flux $2\pi A_{\rm BH}$. Approach to the horizon tends to increase $A_{\rm BH}$; if $2\pi A_{\rm BH}$ becomes large compared with the open torus flux, an equilibrium state $\tau_+ \sim \tau_-$ may be reached, and accretion onto the black hole stalls. This state is expected to be stable since τ_+ decreases with increasing $\mathcal{R}$: if the torus gains (loses) angular momentum, it moves outward (inward) and reduces (increases) τ_+. Such oscillations may give rise to intermittency. The black hole spin-down proceeds on essentially the timescale in equation (9); the difference is that now the torus—with its two faces—mediates the interaction.

More quantitatively, integration of equation (11) gives the net torque on the torus by the black hole as $\tau_+ \approx (\Omega_{\rm BH} - \Omega)A^2 f_{\rm BH}^2/c$, while the wind torque is now $\tau_- \approx \Omega A^2 f_w^2 \times [c|\ln(\theta_0/2)|]^{-1}$. Here A is now $(2\pi)^{-1}$ times the total torus magnetic flux, and $f_{\rm BH}$ and f_w are the fractions of A passing through the inner and outer light surfaces, respectively, so that $f_{\rm BH} + f_w \sim 1/2 - 1$. The torus, "sandwiched" between the black hole and infinity, obeys an evolution equation

$$\mathcal{R}\dot{\mathcal{R}} \approx \frac{2A^2}{M_d c}\left[f_{\rm BH}^2\left(\frac{\Omega_{\rm BH}}{\Omega} - 1\right) - \frac{f_w^2}{|\ln(\theta_0/2)|}\right]. \quad (12)$$

Accretion may therefore be halted—or reversed—if $\Omega_{\rm BH} \gtrsim \Omega[1 + (f_w/f_{\rm BH})^2|\ln(\theta_0/2)|^{-1}]$. Early on, $f_{\rm BH} \ll 1$, and this equation will not be satisfied. Beginning near the midpoint of accretion when $f_{\rm BH} \gtrsim f_w$, and provided that $\Omega_{\rm BH}$ is initiated close to its maximal value, accretion may be suspended for an interval $\sim t_{\rm BH}$, until the value of $\Omega_{\rm BH}$ drops below the critical value. Since, in the "suspended" state, the torus gains and loses almost equal quantities of angular momentum, significant baryonic matter may be carried off in the wind.

5. DISCUSSION AND COMPARISON WITH OBSERVATIONS

We have revisited accretion timescales in magnetized black hole plus disk/torus systems, analyzing evolution when external torques rather than internal "turbulent" torques dominate. We propose two dynamical states—hyperaccretion and suspended accretion—for systems with low- or high-spin black holes. Around a slowly spinning black hole, wind torques drive the disk to a finite-time singularity, producing short bursts. Rapidly spinning black holes transfer sufficient angular momentum to the disk to arrest matter inflow until the black hole spins down, producing long bursts. A bimodal duration distribution occurs, provided that $M_d \ll M_{\rm BH}$.

Our principal assumption is that external ordered torques dominate internal disordered ones; this requires significant poloidal magnetic flux. Determining how poloidal flux is created and when external torques prevail is a major unsolved problem. Numerical studies to date have shown the development of magnetic power spectra dominated by the largest scales permitted, in local disk simulations (e.g., Hawley, Gammie, & Balbus 1995), and that magnetically driven winds may develop from both cold and hot disks (e.g., Shibata & Uchida 1986; Stone & Norman 1994; Stone & Pringle 2000), with dependence of the angular momentum loss rate on E_B similar to analytic predictions (Kudoh, Matsumoto, & Shibata 1998). Because the choice of numerical boundary conditions (particularly for B_ϕ) may affect the solution at the largest scale (see, e.g., Ustyugova et al. 1999), a very large dynamic range would be needed to assess the external magnetic torque.

Our model differs from expectations for a turbulence-dominated disk in several ways. First, a wind-dominated system converts gravitational energy into Poynting flux rather than thermalizing it. Such a "clean" system could have less radiative contamination/spectral degradation from a low-energy thermal component. Second, a magnetized wind is naturally collimated, unlike the fireball in the alternative scenario. Finally, only if significant poloidal flux is present can field geometry become unfavorable for matter inflow, allowing an extended disk lifetime. In turbulent disks, magnetic fields lie primarily parallel to the midplane, so flow onto the black hole proceeds even as magnetic stresses transfer spin energy and angular momentum outward (Krolik 1999; Gammie 1999).

Because most of the energy is liberated in the innermost regions, our hyperaccretion/short-GRB and suspended-accretion/long-GRB proposal implies similar intrinsic luminosities, but larger fluences in the second case. We thus expect short GRBs to feature afterglows similar to those of long GRBs, unless the environments of high and low angular momentum progenitors differ appreciably. This may be tested by the *HETE-2* mission and is already supported by the afterglow detection to the 2 s event GRB 000301C (Jensen et al. 2001).

The iron emission lines from GRB 970508 (Piro et al. 1999) may reflect the composition of progenitor matter. In the suspended-accretion (long-GRB) state, some torus matter is expected to be blown to infinity by the powerful black hole torque. We associate this baryonic wind with line-emitting regions for long-burst systems; they are expected to be less powerful in short GRBs.

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