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 towards collapse in a finite time \( \sim t_{ff} E_k / E_B \), where \( t_{ff} \) is the initial free-fall time and \( E_k / E_B \) is the ratio of kinetic-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-II.

Subject headings: black hole physics — gamma-rays: bursts and 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. The BATSE catalogue shows a broad, bimodal distribution of GRB durations \( T_{90} \) centered on short bursts of \( \sim 0.3 \) s and long bursts of \( \sim 30 \) 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 et al. 1999). Proposed scenarios include coalescence of compact object binaries (Paczynski 1991), and core-collapse of massive stars (Woosley 1993). Both scenarios involve a black hole-disk or -torus system. This disk will be magnetized, due to 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., \( 2M_\odot c^2 \) for a \( 7M_\odot \) 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 et al. 1986); see also e.g. Lee et al. (1999); Beskin & Kuznetsova (2000); Krolik (2000), or superradiant scattering of fast magnetosonic waves (Uchida 1997; 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 \( T = -\dot{J}_H \) may derive from coupling to the black hole’s angular momentum \( \dot{J}_H \) by equivalence in poloidal topology to pulsar magnetospheres (van Putten 1999); Brown et al. (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 vs. 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 time-scales are consistent with the bimodal distribution in GRB durations for plausible values of two presently-uncertain ratios: the kinetic-to-poloidal-magnetic energy in the disk and the black hole-to-disk mass. Numerical simulations soon within reach may establish the likely ranges of these ratios.}

Previously, bimodality in GRB durations has been attributed to, e.g., hydrodynamic time-scales in Newtonian vs. relativistic reverse shocks in shells decelerating in the ISM (Sari & Piran 1995), or 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 (in-
going) flows have the same (opposite) sign. A force-free limit is defined by neglecting Reynolds stresses. Applied to outgoing winds, Maxwell stresses \(-RB_0B_\phi/4\pi\) provide a lower limit on the disk’s accretion rate, where \(B\) denotes the magnetic field with components in spherical coordinates \((r, \theta, \phi)\), and \(R = r \sin \theta\). The force-free limit represents pure “magnetic braking.”

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

\[
\delta \tau_B = \frac{\Omega}{c} \sin^2 \theta \left( \partial_\phi A_\phi \right) \delta A_\phi. \tag{1}
\]

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

\[
\dot{\Omega} R = \frac{2}{c} \sin^2 \theta \left( \partial_\phi A_\phi \right) \left( \frac{\delta A_\phi}{\delta M} \right). \tag{2}
\]

The evolution of a collapsing ring depends on its degree of magnetization and the ring’s degree of geometry. A split-monopole geometry (SMG) for the asymptotic poloidal field is a good approximation when \(|B_\phi|/|B_\theta| >> 1\), 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 split-monopole geometry** — In SMG, we have \(A_\phi = A(1 - \mu/\mu_0)\), with \(\theta_0 = \arccos(\mu_0)\) the angle of the innermost streamline with respect to the pole; the outermost streamline lies on the equator. With \(\varpi = \Omega R/\delta M\), for \(\delta M_0\) the the initial ring radius, the solution of equation (2) describes a finite-time singularity:

\[
\varpi(t; A_\phi) = \left( 1 - t/t_{ff}^{mg} \right)^{1/2}, \tag{3}
\]

where \(t_{ff}^{mg} = (1/4)(\delta J_0/|\delta E_0^B|)\), and \(\delta J_0 = \delta M\Omega_0 R_0^2\) and \(\delta E_0^B\) denote the initial ring angular momentum and torque. The initial ring kinetic energy is \(\delta E_0^B = \Omega_0^2 \delta J_0^0/2\). The energy in the magnetic field of a (double) wedge-shaped portion \(A_\phi\) rooted at \(\delta M_0\) is \(\delta E_0^B = (1/2) \int B^2 \sin^2 \delta \Omega_0 R_0^2 \delta \Omega_0 \delta A_\phi A_\phi(\Omega_0 R_0)/(2M_0) = (\delta \tau_B^0/c) R_0^2 (\varpi^2 \sin^2 \theta)^{-2}\). For split-monopole geometry, \(\delta E_0^B \rightarrow \delta E_{B,s} = A(\partial_\phi A_\phi)/(2R_0\delta M_0)\). The initial freefall time is \(t_{ff} = \pi/(\Omega_0\varpi^2)\). For \(R_\delta = 2GM/H/c^2\) the Schwarzschild radius of the black hole, the collapse time in SMG becomes

\[
t_{ff}^{mg} = \frac{1}{2cR_\delta \sin^2 \theta \delta E_{B,s}} \left( \frac{R_0}{R_\delta} \right)^{1/2} \frac{\delta E_k}{\delta E_{B,s}}. \tag{4}
\]

In the above, superscripts on the energies are dropped since \(\delta E_k^B/\delta E_{B,s} = (\delta M/\delta A_\phi)(GM/H/A)\) 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 toroidal field geometry** — 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 \(B_\phi \propto 1/R^4\) (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, \(\partial_\phi A_\phi = -r^2 B_\phi = 2I_p[\Omega sin^2 \theta]^{-1}\) so that \(\delta \tau_B = 2I_p \delta A_\phi/c\). Here, \(I_p\) relates to the total normalized flux \(A \equiv \int d\mu \partial_\phi A_\phi\) by \(I_p \approx -A \Omega (\varpi_0)/[2(\varpi_0^2)]\), where \(\varpi_0\) is the innermost disk radius.

From equation (2), the TFG limit gives \(\Omega R \dot{\varpi} = -(4I_p/c)(\delta A_\phi/\delta M)\). Using \(\Omega(\varpi) = (\Omega(0) \varpi^{-3/2})\) we find the finite-time singularity solution:

\[
\varpi(t; A_\phi) = \left( 1 - t/t_{ff}^{Tfg} \right)^{1/2}, \tag{5}
\]

where

\[
t_{ff}^{Tfg} = \frac{\delta J_0}{|\delta E_0^B|} = \frac{2 \Omega_0}{\pi \delta E_{B,s}}. \tag{6}
\]

Here, \(\varpi_0 \equiv (c/\sqrt{2}) R_0^1 \int d\theta (\varpi \sin \theta)^{-1} \sim (R_\delta^3 / R_\phi^3)^{1/2} \left| \ln(\theta_0/2) \right|\). Going from \(\varpi_0\) to the outside of the disk at \(\varpi_d\), the radial factor in \(\varpi_0\) decreases from \(\varpi_0\) to \(\varpi_d\); which is order-unity to \((R_\delta^3 / R_\phi^3)^{1/2}\) (which is of order \(\varpi_0\) to \(\varpi_d\)); the product \(t_{ff}^{Tfg}\) varies \(\propto \varpi_0^1\). Note that \(\delta E_{B,s}\) in TFG is larger than \(\delta E_{B,s}\) by a factor \(\partial_\phi A_\phi/\partial \theta \sim (\Omega_0 \varpi_0 \varpi_0^3)^{1/2}\), which is of order unity or less unless the effective beaming angle \(\theta_0\) 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

\[
t_{in} = 0.011 s \left( \frac{R_{in}}{100 \text{ km}} \right)^2 \left( \frac{M_H}{10 M_\odot} \right)^{-1} \frac{\delta E_k}{\delta E_{B,s}}, \tag{7}
\]

where \(g(\theta_0) = (1 + \ln(\theta_0/10))/\ln(0.009)\) and \(\theta_0\) is the angle of the innermost wind field line. Correspondingly, the maximum outer-disk accretion time is

\[
t_{dd} = 0.057 s \left( \frac{R_d}{1000 \text{ km}} \right)^2 \left( \frac{M_H}{10 M_\odot} \right)^{-1} \frac{\delta E_k}{\delta E_{B,s}}, \tag{8}
\]

assuming \(R_d >> R_{in}\) with SMG; otherwise, if \(R_d \sim R_{in}\), we would use the TFG result, which increases the time by a factor \(\sim 4 \left| \ln(\theta_0/2) \right|\) consistent with equation (7). These limits will bracket the interval over which disk accretion takes place.

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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).
3. DISK ACCRETION AND BLACK HOLE SPINDOWN TIMES COMPARED

A ring spends the greatest portion of its lifetime near its initial radius; during the subsequent rapid infall phase, magnetic flux is approximately conserved. For fiducial estimates, we shall normalize to an initial kinetic-to-magnetic flux ratio of about 100 for all rings, as found from Hawley (2000)’s 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); net poloidal flux, as in Hawley, Gammie, & Balbus (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 time classes of GRBs.

**Time-scale 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 \( \lesssim \) a few seconds—comparable to that of the 2s event GRB 000310C (Jensen et al. 2001). In this scenario, short bursts represent hyperaccretion onto a slowly-spinning black hole.

**Time-scale 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 \( t_H \) required for angular momentum extraction by the Blandford-Znajek (1977) and related processes. Given a horizon flux \( 2\pi A_H \), the black hole torque for open fields is \( \mathcal{L}_H \sim \Omega H^2 \) (see Thorne et al. (1986)). To estimate \( A_H \), assume that the black hole field strength is comparable to that from the inner part of the disk. In TFG, we find \( A_H \approx A(2|\ln(\theta_0/2)|)^{-1} \) so that \( \mathcal{L}_H \sim \Omega H^2 (4c|\ln(\theta_0/2)|)^{-1} \). For comparison, the total disk wind angular momentum luminosity is equal to \( \tau \equiv \int dt\mathcal{L}_B \); i.e. \( \mathcal{L}_B = \Omega(R_{in})^2 c|\ln(\theta_0/2)|^{-1} \). The spin-down time \( t_H = 2GM^2_H/3\tau \). Here, \( \tau = R_H c^2(\lambda/2) \) is the radius of the black hole horizon; \( \lambda = a/M \) is the ratio of specific angular momentum to the maximal Kerr value.

Taking \( \delta E_k/\delta E_{B,s} \sim GMH M_d/A^2 \), we obtain \( t_H = (8\pi/c)(\delta E_k/\delta E_{B,s})(M_H/M_d)|\ln(\theta_0/2)|/2 \). Assuming \( a \sim M \) so \( \tau \sim R_H/2 \), we find

\[
t_H = 88s \left( \frac{M_H}{10M_\odot} \right) \left( \frac{M_d}{100} \right) \left( \frac{\delta E_k/\delta E_{B,s}}{100} \right) g^2(\theta_0),
\]

consistent with a bimodal offset of long- from short-duration bursts for \( M_d << M_H \). Note that \( t_H/t_{in} \sim M/H \), as magnetic torques on the disk and black hole are comparable, but \( t_H \) decreases as \( \theta_0 \) increases from 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 \( M_H \) of about \( 10M_\odot \), formed promptly in the hypernova scenario, is expected in view of the Kerr constraint \( J_H \leq GM_H^2/c \). For prompt collapse of a He-core of radius \( R_{He} \), 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 \( \xi R_{He} \),

\[
M > M_H \geq M \left( \frac{2}{5} \right)^3 \left( \frac{2R_{He}}{GM^2} \right)^{3/2} \left( 1 + \frac{m}{M} \right)^{3/2} \left( \frac{\bar{\rho}}{\bar{\rho}_c} \right)^2.
\]

Here, the central density \( \rho_c \) is that extrapolated from the He-mantle, with \( \bar{\rho} \) denoting the mean density. For canonical values \( m = 20M_\odot, R_{He} = 100 \) and \( \xi = 4 \), a black hole may form promptly provided the inequality on the left hand-side of equation (10) is satisfied; i.e. \( 1 > 1.3 \times 10^3 (\bar{\rho}/\bar{\rho}_c)^2 \). This would be true, e.g. the He core is approximately fit by a polytrope with index \( n = 3 \), for which the Lane-Emden relationship gives \( \rho_c \sim 50\bar{\rho} \) (Kippenhahn & Weigert 1990). 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) are far shorter, however. We speculate that 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 (Fig. 2 in van Putten (1999)). The black hole-torque on the torus is (adapted from Thorne et al. (1986))

\[
\delta \tau_+ \approx \frac{1}{c}(\Omega_H - \Omega) \sin^2 \theta(\partial_{\theta} A_B) \delta A_B,
\]

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, \( \tau_+ \), depends on the associated horizon flux \( 2\pi A_H \). Approach to the horizon tends to increase \( A_H \); if \( 2\pi A_H \) becomes large compared to the open torus flux, an equilibrium state \( \tau_+ \sim \tau_\ast \) may be reached, and accretion onto the black hole stalls. This state is expected to be stable, since \( \tau_\ast \) decreases with increasing \( R \); if the torus gains (loses) angular momentum, it moves outward (inward) and reduces (increases) \( \tau_\ast \). 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 \( \tau_+ \approx (\Omega_H - \Omega) A^2 f_H^2/c \), where the wind torque is now \( \tau_\ast \sim \Omega^2 A^2 f_\ast^2 c^2|\ln(\theta_0/2)|^{-1} \). Here, \( A \) is now \((2\pi)^{-1} \) times the total torus magnetic flux, and \( f_H \) and \( f_\ast \) are the fractions of \( A \) passing through the inner and outer light surfaces, respectively, so that \( f_H + f_\ast \sim 1/2 - 1 \). The torus, “sandwiched” between the black hole and infinity, obeys an evolution equation

\[
\mathcal{R}_\mathcal{R} \approx \frac{2A^2}{M^2 c^7} \left[ f_H^2 (\Omega_H - \Omega) - f_\ast^2 \left( \frac{\Omega}{|\ln(\theta_0/2)|} \right)^2 \right].
\]

Accretion may therefore be halted - or reversed - if \( \Omega_H \lesssim \Omega(1 + f_\ast^2 f_H^2)\). Early on, \( f_H \ll 1 \), and this equation will not be satisfied. Beginning near the midpoint of accretion when \( f_H \sim f_\ast \), and provided that \( \Omega_H \) is initiated close to its maximal value, accretion may be suspended for an interval \( \sim \tau_H \), until the value of \( \Omega_H \) 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 \( M_d << M_H \).

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 \( B_0 \)) may affect the solution at the largest scale (see e.g. Ustyugova et al (1999)), 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 fluxes 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-II mission, and is already supported by the afterglow detection to the 2s 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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