CHARACTERIZING THE TIME VARIABILITY IN MAGNETIZED NEUTRINO-COOLED ACCRETION DISKS: SIGNATURES OF THE GAMMA-RAY BURST CENTRAL ENGINE

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ABSTRACT

The central engine of gamma-ray bursts (GRBs) is hidden from direct probing with photons mainly due to the high densities involved. Inferences on their properties are thus made from their cosmological setting, energetics, low-energy counterparts, and variability. If GRBs are powered by hypercritical accretion onto compact objects, on small spatial scales the flow will exhibit fluctuations, which could in principle be reflected in the power output of the central engine and ultimately in the high-energy prompt emission. Here, we address this issue by characterizing the variability in neutrino-cooled accretion flows through local shearing box simulations with magnetic fields, and then convolving them on a global scale with large-scale dynamical simulations of accretion disks. The resulting signature is characteristic and sensitive to the details of the cooling mechanism, providing in principle a discriminant for GRB central engine properties.

Key words: accretion, accretion disks – gamma-ray burst: general – hydrodynamics – magnetic fields

Online-only material: color figures

1. INTRODUCTION

Gamma-ray bursts (GRBs) are thought to be driven by mass accretion onto compact objects (Woosley & Bloom 2006; Nakar 2007; Lee & Ramirez-Ruiz 2007; Gehrels et al. 2009). The enabling cooling mechanism allowing accretion is neutrino emission, raising the usual Eddington rate for photons by 16 orders of magnitude (Popham et al. 1999; Narayan et al. 2001; Kohri & Mineshige 2002; Beloborodov 2003; Di Matteo et al. 2002; Setiawan et al. 2004; Lee et al. 2005; Chen & Beloborodov 2007). This applies both to long and short GRBs, perhaps due to the collapse of massive rotating stars (Woosley 1993; MacFadyen & Woosley 1999) and compact binary mergers (Eichler et al. 1989; Paczyński 1991; Narayan et al. 1992), respectively, or magnetized neutron stars (Usov 1992). The prompt gamma-ray emission originates at $\sim 10^{14}$–$10^{16}$ cm from this source, possibly from internal shocks in the relativistic outflow generated by the engine (Zhang & Mészáros 2004). The engine itself is hidden from view due to the high opacities and can be potentially probed directly only through gravitational waves and neutrinos.

The observed time series in GRB prompt emission show diversity between events (see, e.g., Norris et al. 1996): some have a single peak, others have multiple emission episodes, with correlations between the fluence of the active period and the length of the quiescent interval preceding it (Ramirez-Ruiz & Merloni 2001). On top of this, rapid (ms) fluctuations are routinely observed. Fourier analysis of the high-energy light curves of the prompt emission in GRBs in the source frame reveals power-law spectra (Beloborodov et al. 1998, 2000; Ryde et al. 2003), with index $\sim -5/3$ and a break at $\sim 1$–2 Hz. Variability is likely due to a combination of several effects, allowing in principle an additional way to probe the central engine indirectly. Some are probably intrinsic to the progenitor: the distribution of angular momentum with radius inside the star in the case of a collapsar may lead to distinct episodes of energy release (Kumar et al. 2008; Perna & MacFadyen 2010; Lopez-Camara et al. 2010); the fall back at late times of material stripped from a tidally disrupted neutron star is capable of powering secondary accretion episodes (Rosswog 2007; Lee et al. 2009); hydrodynamical or magnetic instabilities in the accretion disk may result in intermittent accretion (Perna et al. 2006; Proga & Zhang 2006; Taylor et al. 2010). Others can come from the relativistic outflow: the interaction of a jet with high Lorentz factor with the stellar envelope before breakout can lead to irregularities and shocking (Morsony et al. 2010); the outcome of internal shocks between shells in the flow depends on the variation in mass and energy upon ejection (Panaitescu et al. 1999; Ramirez-Ruiz et al. 2001; Bosnjak et al. 2009; Mendoza et al. 2009). An additional factor, upon which we focus here, is related to the variability present in the accretion disk as a result of turbulent motions. The dissipation is related to the local hydrodynamical variables, and as these vary in time, so will the energy output.

The magnetorotational instability (MRI; Balbus & Hawley 1998) is a possible mechanism that will allow for the transport of angular momentum in accretion disks with differential rotation. Its behavior under the physical conditions in neutrino-cooled disks, which are similar to those occurring in supernovae, but different from the usual ones present in X-ray binaries and active galactic nuclei, has come under scrutiny more recently (Thompson et al. 2005; Masada et al. 2007; Rossi et al. 2008; Obergaulinger et al. 2009). One such difference lies in the sensitivity of the cooling rate to temperature and is at the heart of this work. Whereas, for example, the photon bremsstrahlung emissivity scales as $\dot{q} \propto T^{3/2}$ in the optically thin limit, for neutrinos $\dot{q} \propto T^{\beta}$, where $\beta \simeq 6$–9 depending on the cooling process. Further, while photon-cooled disks are typically optically thick, $\tau_{\nu} \gg 1$, for a wide range of relevant parameters their neutrino-cooled counterparts are optically thin, $\tau_{\nu} \lesssim 1$, tightly coupling the local conditions to the emitted luminosity.

In this Letter, we characterize the local variability in neutrino-cooled disks through shearing box MHD simulations (Section 2), and then use the results of global disk simulations to scale the results for the central engine as a whole (Section 3). Prospects for placing constraints on GRBs are discussed in Section 4.
2. THE SMALL SCALE: NUMERICS AND PHYSICS IN THE SHEARING BOX CALCULATIONS

The local flow in the disk is modeled by the shearing box approximation (Hawley et al. 1995): a rectangular Cartesian coordinate system represents a local neighborhood inside the disk, at an arbitrary orbital radius $R_0$ with dimensions which are much smaller than $R_0$. The $x$, $y$, and $z$-axes represent the radial, azimuthal, and vertical directions in the disk, respectively. The radial component of the central object’s gravity is included, and differential rotation is replaced by a Keplerian shear flow along the $y$-direction. Periodic boundary conditions are set in $y$ and $z$, while in $x$ these are “shearing periodic”: upon crossing a radial boundary, a fluid element is displaced along $y$ by an amount given by the shear value. The Eulerian ZEUS-3D code (Stone & Norman 1992a, 1992b) is used to solve the equations of ideal MHD. These are

$$\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0,$$  

(1)

$$\frac{\partial \mathbf{v}}{\partial t} + \mathbf{v} \cdot \nabla \mathbf{v} = -\frac{1}{\rho} \nabla \left( P + \frac{B^2}{8\pi} \right) + \frac{(\mathbf{B} \cdot \nabla) \mathbf{B}}{4\pi \rho} - 2\rho \mathbf{v} \times \frac{3}{2} \mathbf{e} \times \mathbf{e},$$

(2)

$$\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{v} \times \mathbf{B}),$$

(3)

$$\frac{\partial e}{\partial t} = -P \nabla \cdot \mathbf{v} - C\rho T^\beta,$$

(4)

and an ideal gas equation of state with $P = (\gamma - 1) e$, where $\mathbf{v}$, $\rho$, and $P$ are the gas velocity, density, and pressure, respectively, $\mathbf{B}$ is the magnetic field, $\mathbf{e} = (0, 0, \Omega_0)$ is the disk angular frequency, $e$ is the internal energy density, and $\mathbf{e}$ is a unitary radial vector.

The second term on the right-hand side of Equation (4) is added to model neutrino cooling processes, with the nature of the particular emission mechanism fixing the value of $\beta$. In GRB central engines, the dominant processes are $e^\pm$ pair annihilation (Itoh et al. 1996), with $q_{\text{ann}} \propto T^6$, and $e^\pm$ capture onto free neutrons and protons (Langanke & Martinez-Pinedo 2001), with $q_{\text{cap}} \propto T^9$. A fraction $\zeta$ of the initial internal energy of the fluid is removed over 100 orbital periods at $R_0$, the duration of the simulation. We used $\zeta = 0.1$, thus the cooling time is much longer than the dynamical time.

Initially the density, internal energy, and pressure in the box are uniform, and the magnetic field is vertical with non-zero net magnetic flux across the box. The ratio of gas to magnetic pressure is $P_{\text{gas}}/P_{\text{mag}} = 400$. The velocity field is randomly perturbed with amplitude $v_{\text{pert}} = 10^{-3} c_s$, where $c_s$ is the sound speed. We have performed simulations at varying resolution to test for convergence (the highest having $128 \times 256 \times 128$ zones in $x$, $y$, $z$, respectively). From the power spectrum of the velocity field $\tilde{v}(\mathbf{k})$, we have verified that our results are converged and unaffected by resolution changes at the level reported. Turbulent motions induced by the MRI develop in the simulation on a timescale $\Omega_{\text{MRI}} \approx 1/\Omega_0$, leading to turbulent transport of angular momentum. The fundamental output of the simulations is the luminosity as a function of time of the shearing box, which we denote as $s(t)$.

Luminosity variations for $s(t)$ are shown in Figure 1 for $\beta = 6, 9$. Rapid variability is present, reflecting the turbulent motions in the box. These changes can be characterized by their Fourier transform, $\tilde{s}(\nu)$, where $\nu$ is the signal frequency (Figure 1). Roughly, the spectrum is a power law, $\tilde{s}(\nu) \propto \nu^{-2}$, over a wide range of frequencies in the case of $\beta = 6$, with more complicated behavior superimposed in the case of $\beta = 9$. It is clear that the variations in energy release have a characteristic behavior, with the details varying according to the specific cooling mechanism.

3. THE LARGE SCALE: EXTENSION TO GLOBAL DISK VARIABILITY

To find the global variability we require the superposition of a very large number of smaller regions represented by the
shearing box calculations, and approximations are necessary to proceed further. We assume equatorial and azimuthal symmetry, and also take the disk to be in a stationary state. Obtaining global quantities thus requires integrating over the polar angle \( \phi \in [0, 2\pi] \) and radius \( r \in [r_{\text{in}}, \infty) \). We are still faced with the superposition of a potentially large number of zones, each uncorrelated with the other, and which could destroy any coherent signature of the processes occurring on small scales. A natural quantity in the disk, however, enforces causality for a given frequency: the local sound speed, \( c_s \).

Consider a disk annulus at radius \( R \), with sound speed \( c_s \). The correlation length over which signals with frequency \( \nu \) propagate is \( l_\nu = c_s/\nu \). The azimuthal velocity is \( v_{\phi} = R \Omega \), and from vertical hydrostatic balance, the pressure scale height \( H \) satisfies \( H/R \approx c_s/v_{\phi} \). Thus, \( c_s = 2\pi R f (H/R) \), with \( f \) being the orbital frequency at radius \( R \). In neutrino-cooled accretion disks, \( A = H/R \approx \text{cst} \) (Lee et al. 2005; Lee & Ramirez-Ruiz 2007), so that \( l_\nu = 2\pi R A f / v \). The number of distinct azimuthal zones within the ring for frequency \( \nu \) is \( N_{\phi,\nu} = 2\pi R / l_\nu = \nu / (A f) \).

Dynamic simulations of such disks show that \( A \approx 1/10 \). If \( N_{\phi,\nu} \gg 1 \), the signal at frequency \( \nu \) is thus diminished with respect to the idealized case of entirely in-phase contributions by a factor \( N_{\phi,\nu} \), due to the random nature of the phase decorrelation. If, on the contrary, \( N_{\phi,\nu} \ll 1 \), the signal at frequency \( \nu \) is added constructively. This acts as a broad low-pass filter by slightly suppressing the signal at high frequencies, when \( \nu \gg A f \). Carrying out this filtering procedure also requires scaling the fiducial model to different radii, which we now address.

The emissivity is not a function of azimuthal position, but it is one of radius. Given the extreme sensitivity of the cooling rate to temperature, only the innermost regions of the disk contribute to the luminosity, and thus to its variability properties. The argument given above for the decoupling of azimuthal zones in the disk for a particular frequency is also applicable to the radial coordinate. Advection may modify this slightly, but for our present purpose it is valid enough to say that as the flow is split up into \( N_{\phi,\nu} \) zones in azimuth, it is divided into rings of characteristic size \( l_\nu \). The number of radial zones in the disk at frequency \( \nu \) is \( N_{R,\nu} = r_{\text{out}}/l_\nu \), where \( r_{\text{out}} \) is the outer radius of the disk. Using the previous results, \( N_{R,\nu} = r_{\text{out}}/(2\pi R A f) \).

A second low-pass filter needs to be considered in analogy with the azimuthal one, and the signal is now suppressed by a factor \( N_{R,\nu}^{1/2} \) if \( N_{R,\nu} \gg 1 \).

In order to carry out this procedure, we have used the results of global, two-dimensional simulations of neutrino-cooled accretion disks (Lee et al. 2005; Lee & Ramirez-Ruiz 2007) orbiting a stellar mass black hole (BH), evolved in the pseudorelativistic potential of Paczyński & Wiita (1980). For consistency purposes, we have re-computed their evolution for this work for two test cases, with \( M_{\text{BH}} = 4 M_\odot \) and \( q = M_{\text{disk}}/M_{\text{BH}} = 7.5 \times 10^{-3} \). In the first model, T6, we have eliminated all sources of cooling except from pair captures, where \( \dot{q} \propto T^6 \), while in the second, T9, we have only kept that due to pair annihilation, with \( \dot{q} \propto T^7 \). Each model can then be more faithfully compared with the two fiducial runs for the shearing box with \( \beta = 6, 9 \), respectively, using a near-stationary state, where the accretion rate is \( M = 0.05 M_\odot \text{ s}^{-1} \).

We first compute the equatorial radial profile of emissivity in simulations T6 and T9, \( \dot{q}(R) \), shown in Figure 2. At small radii the dominant contribution comes from pair capture, while annihilation dominates for \( R \gg R^* = 30 R_{\text{Sch}} \). The combination is a strongly decreasing function of \( R \), so the radial convolution need not be extended to infinity, and pair capture dominates the total emission as well as its general behavior. The inner radius of the disk is \( r_{\text{in}} = r_{\text{ms}} = 3 R_{\text{Sch}} \), where the effects of general relativity truncate the flow. We have used \( r_{\text{out}} = 85 R_{\text{Sch}} \), and checked that the final result is insensitive to this choice, as long as \( r_{\text{out}} \gg R^* \).

To perform the radial integral, the time series \( \dot{q}(t) \) and \( \dot{\tau}(t) \) need to be scaled to the disk at different radii \( R \). The characteristic break frequency in the fiducial spectrum for the volume-averaged total (Maxwell and Reynolds) stress, measured as an equivalent \( \alpha \) viscosity parameter, is roughly at the orbital frequency, \( f \). This is related to the nature of the MRL, where, \( \Omega_{\text{MRI}} \approx 1/\sqrt{f} \). The orbital angular frequency is known as a function of radius, \( \Omega = (G M_{\text{BH}}/R)^{1/2}/(R/(R - R_{\text{Sch}})) \), so we can convert from frequency to radius, \( \dot{\tau}(v) \) to \( \dot{\tau}(R) \) and integrate over different annuli (the expression for \( \Omega \) is appropriate for the pseudorelativistic potential used).

The power associated with the variability in the neutrino luminosity can now be computed as

\[
P(v) = \int_{r_{\text{in}}}^{r_{\text{out}}} \dot{q}(R) \Phi(N_{\phi,\nu}) \rho(N_{R,\nu}) \dot{\tau}(R) 2\pi R dR, \tag{5}
\]

where \( \dot{q}(R) \) corresponds to model T6 or T9, and \( \Phi(N_{\phi,\nu}) \) and \( \rho(N_{R,\nu}) \) account for the filtering at high frequencies. Figure 3 shows the resulting spectra, where several features clearly stand out. At low frequencies, \( v \lesssim v_0 \approx 0.1 \text{ Hz} \), the spectrum is the same in both cases, \( P(v) \propto v^{-2} \). For \( v \gtrsim v_0 \) the results diverge: model T6 approximately maintains the power-law decay, while model T9 shows an excess, and in general more complex behavior, with broad features at \( \approx 2 \) and \( 40 \text{ Hz} \) overlaid on a decay with lower index, \( P(v) \propto v^{-1.7} \). At \( v = v_1 \approx 100 \text{ Hz} \), a break to a more rapidly decaying power law terminates both spectra. The errors are greater for model T9, due to greater

\[
\text{Figure 2. Equatorial (z = 0) neutrino emissivities, } \dot{q}(R/R_{\text{Sch}}), \text{ for two-dimensional, azimuthally symmetric global dynamical simulations of accretion disks around a BH with } M = 4 M_\odot. \text{ The blue (red) line corresponds to simulation } T6 \text{ (T9), where only pair captures (annihilation) were used to compute the cooling rate. The solid black line is the result of a simulation including both processes. (A color version of this figure is available in the online journal.)}
\]
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The reason for the difference when modifying the cooling mechanism is fundamentally related to the local energy balance. For \( \dot{q} \propto T^9 \) (model T9) and a given internal energy supply, \( e \), the local cooling time \( t_{\text{cool}} = e/\dot{q} \) is shorter under a temperature perturbation \( \Delta T \) than when \( \dot{q} \propto T^6 \) (model T6). The power is thus higher at such frequencies, leading to the observed displacement in Figure 3. The argument holds when comparing model T6 with the test run at \( \beta = 1 \), and gives a way to characterize the accretion flow and discriminate between competing mechanisms, if the variations in the neutrino luminosity are directly reflected in the accretion power output which drives a relativistic flow.

Two possible scenarios in which the difference between the cooling regimes studied here may be of relevance are worthy of note. First, the mass of the BH introduces a scaling into the problem when combined with neutrino cooling. If \( M_{\text{BH}} \) is too large, the density and temperature in the accretion flow can be too low for neutrino cooling to operate. As the BH mass is reduced, pair annihilation, with \( \dot{q} \propto T^9 \), first becomes effective as an energy sink, followed by pair captures, with \( \dot{q} \propto T^6 \). A signature of the BH mass is thus in principle available in the variability of the flow. The second case is related to the time evolution of the flow, assuming that a certain amount of mass \( M_{\text{disk}} \) is initially available for accretion, and no further feeding of the central engine takes place. Global disk simulations show that the density and temperature drop as the disk drains into the BH on the viscous timescale. Given enough time, the whole disk will lie on the branch cooled by pair annihilation where \( \alpha \) particles have formed, below \( \log(\rho/(\text{g cm}^{-3})) \simeq 6.5 \) and \( \log [T/\text{(K)}] \simeq 10 \) for the adopted BH mass (see also Figure 2 in Lee et al. 2009). Thus, the variability may initially behave as in case T6, and end as in case T9 (the luminosity will have decreased substantially by then, along with the accretion rate). The associated GRB, if one occurs, need not necessarily be powered by neutrinos in order for these effects to be apparent. The fact that neutrinos are responsible for the cooling process allowing accretion makes them relevant in this context.

A number of limitations apply to this study and can be matters for further study. First, the scaling we have used to infer the variability at different radii is strictly valid if the flow is adiabatic. Explicit cooling thus violates this assumption. However, given the choice of \( \xi \) and \( t_{\text{cool}} \), the associated timescales are such that \( t_{\text{cool}} \gg t_{\text{dyn}} \), making this a reasonable approximation. Second, we have used the neutrino luminosity as a proxy for the manifestation of central engine activity, which could have a neutrino component, but need not be restricted to it. For example, magnetic fields may power relativistic outflows leading to the observed high-energy emission. In this sense, an alternative characterization based on the mass accretion rate \( \dot{M} \) through the disk may be of use as well, which can then be associated with the energy output of the relativistic outflow as \( L_{\text{rel}} \propto \dot{M} c^2 \). Third, this is only the first filter any variability originating in a disk needs to go through before emerging as high-energy photons. A variable mass accretion rate and energy conversion efficiency, propagation through the stellar envelope (for a massive star progenitor), external shocks, and general relativistic effects (Birkl et al. 2007), among others, can each potentially leave their own fingerprint on the power spectrum. The background upon which they will do so, however, must ultimately be related to the flow in the disk itself if that is what is driving the energy release, and is what we focus on in this work. Finally, other mechanisms can potentially be responsible for angular momentum transport.

scatter in the Fourier transform. The thin black line in Figure 3 is from a shearing box test calculation with \( \dot{q} \propto T^6 \), i.e., with \( \beta = 1 \) (the smaller error bars are not plotted for clarity). As the sensitivity of the cooling term to the temperature rises, the trend is for an excess in power at higher frequencies and a more moderate background power-law decay.

4. SUMMARY, DISCUSSION, AND PROSPECTS FOR OBSERVABILITY IN GRBs

Through shearing box MHD simulations we have characterized the time variability of the energy release at small scales in a neutrino-cooled accretion disk around a BH, where energy losses primordially come from \( e^+ \) pair capture onto free nucleons and protons and pair annihilation. With the use of cooling profiles from large-scale, two-dimensional simulations of full disks, we have convolved this local variability to obtain a global signature of time variations in the power output, through the power density spectrum of the neutrino luminosity. Since accretion is enabled by the cooling through neutrinos, we take this as an indicator of central engine variability which will be reflected in the relativistic outflow eventually giving rise to a GRB.

The power spectrum exhibits characteristic features related to the general nature of the neutrino cooling in the optically thin regime, and particularly to its temperature dependence. A background power-law decay, with index \( -1.7 \text{--} 2 \), extends approximately from 0.1 to 100 Hz. For emissivities \( \dot{q} \propto T^6 \), additional power appears at high frequencies as \( \beta \) rises, with the particular values of the slopes and cutoffs scaling with the cooling mechanism (Figure 3). This result thus provides in principle a discriminant for GRB central engines powered by neutrino-cooled accretion flows and illustrates how any local mechanism can be used in the same way to test its viability.
and variability in the disk itself, such as thermal, viscous, or gravitational instabilities, none of which we have considered here.

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