Constraints on Gamma-ray Burst Inner Engines in a Blandford-Znajek Framework

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ABSTRACT
Under the assumption that a Gamma-ray Burst (GRB) is ultimately produced by a Blandford-Znajek (BZ) jet from a highly spinning black hole (BH), we put limits on the magnetic field and BH mass needed to power observed long and short GRBs. For BHs in the range of $2 - 10 M_\odot$ (for long GRBs) and $0.5 - 4 M_\odot$ (for short GRBs), we find magnetic fields in the range of $5 \times 10^{14} \lesssim B \lesssim 10^{17}$ G are needed to power the observed GRBs. Under the simple assumption of flux conservation, we estimate the magnetic fields of the progenitor systems for both long and short GRBs, finding that single massive star progenitors require fields $\sim 10^6$ G and NS merger systems require fields $\sim 10^{15}$ G. We also discuss the implications and consequences of high magnetic fields in GRB BH-disk systems, in terms of MRI field growth and magnetically arrested disks. Finally, we examine the conditions under which the progenitor systems can retain enough angular momentum to create BHs spinning rapidly enough to power BZ jets.

Key words: stars: GRBs

1 INTRODUCTION
The progenitor system that gives rise to a gamma-ray burst (GRB) has long been an unresolved problem, but a general framework has emerged. Long GRBs appear to be associated with massive stars and short GRBs with the merger of two compact objects (reviews that compile and discuss the evidence behind these associations for both long and short bursts include Piran (2004); Zhang & Mészáros (2004); Mészáros (2006); Gehrels et al. (2009); Kumar & Zhang (2015); Levan et al. (2016); for short bursts specifically, see Lee & Ramirez-Ruiz (2007); Berger (2014); D’Avanzo (2015)).

A black-hole (BH) accretion disk system is a plausible outcome of these cataclysmic events (Woosley 1993; MacFadyen & Woosley 1999; Heger et al. 2003; Fernández & Metzger 2013; Fryer et al. 2015), although it is possible that these progenitors produce a highly magnetized ($B \gtrsim 10^{14}$ Gauss) and potentially hypermassive ($M \gtrsim 2M_\odot$) neutron star (NS) after the collapse/merger (Usov 1992; Duncan & Thompson 1992; Thompson 1994; Zhang & Mészáros 2001; Mazzali et al. 2014; Rowlinson et al. 2014; Metzger et al. 2015; Rea et al. 2015; Stratta et al. 2018). The BH-disk system has been shown to be a viable engine for the GRB jet, from the standpoint of the timescales and energetics involved (e.g., consider the energy available from a $20 M_\odot$ star accreting $5M_\odot$ at a rate of $0.1 - 0.1 M_\odot s^{-1}$; for more discussion on this, see, for example, Popham et al. (1999)). Observationally, we see BH-accretion disks in active galactic nuclei (which can, in many ways, be considered scaled cousins of GRBs; Nemmen et al. (2012); Zhang et al. (2013); Wu et al. (2016); Lloyd-Ronning et al. (2018)), and infer they are behind the relativistic jets observed in these objects.

It was suggested in the 1970’s (Blandford & Znajek 1977) that a magnetized BH-accretion disk system is capable of launching a powerful jet through the so-called Blandford-Znajek (BZ) process. The energy reservoir for the BZ process is the rotational energy of the central BH, extracted as the BH rotates in an external magnetic field. The process can be envisioned in a number of ways - for example, in terms of the torque exerted by the BH on magnetic flux tubes threading the horizon, and/or in terms of a dc-circuit analogy (a current is induced by changing magnetic flux near the horizon, and both the BH and force-free region around it have some

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Our paper is organized as follows: In §2, we discuss the theoretical framework of a BZ jet and the assumptions we make in our calculations. In §3, we present our results, exploring constraints put on the magnetic field and BH mass of long and short GRBs. We also make an estimate of the theoretical framework and assumptions.

The analysis below is carried out in the context that the GRB jet is a result of the BZ process. We can distill the details of the BZ process into a relatively simple expression for the luminosity from the jet (see, e.g., Tchekhovskoy et al. (2010); Tchekhovskoy & Giannios (2015)):

\[ L_{\text{BZ}} = (k f c^5 / 64 \pi G^2) u^2 \phi_{BH}^2 M_{BH}^2 \]

where \( k \) is a geometrical factor related to the magnetic field geometry (of order \( \sim 0.05 \)), and \( f \) is a correction factor related to the BH spin \( a = Jc / G M_{BH}^2 \) (\( J \) being the BH angular momentum). For \( a \lesssim 0.95 \), \( f \approx 1 \). The parameter \( \phi_{BH} \) is the magnetic flux near the horizon, and \( M_{BH} \) is the mass of the BH. Again, although this equation was originally derived for slowly spinning (\( a \sim 0.1 \)) BHs, it has been shown to describe the BZ luminosity from more rapidly spinning BHs (up to \( a \sim 0.95 \)) within the correction factor \( f \) (for additional discussion of the validity of equation 1, see McKinney (2005); Tchekhovskoy et al. (2010); Tchekhovskoy & Giannios (2015)).

If the GRB ultimately derives its power from the rotational energy of the central BH via the BZ jet, the observed

\[ \approx 10^{53} \text{erg} \left( a / 0.9 \right)^2 (\phi_{BH} / 10^{28} \text{Gcm}^2)^2 (5 M_{\odot} / M_{BH})^2 \]

1 The BZ process can also be modeled as a Penrose process. See Brito et al. (2015) for a review and introductory discussion of Penrose processes and their wave analogues.
where than that of short GRBs, although a value of $\theta$ opening angle for long GRBs seems to be somewhat larger so (Wang et al. 2018). We point out the inferred average the inferred opening angle can be less by a factor of 2 or

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AGNs (green circles), using data from the previous two ref-

rameters from X-ray binary observations (blue circles) and

(2014) and Reynolds (2014)). Figure 1 shows BH spin pa-

(2014)). A larger jet opening angle will require more power

0.1 - i.e. 10 percent of of the Blanford-Znajek jet power goes in to producing the gamma-ray luminosity. The effi-

ciency of conversion of jet energy to GRB radiation has been discussed extensively in the literature (see, e.g., Guetta et al. (2001); Lloyd-Ronning & Zhang (2004); Fan & Piran (2006)). Although we do not have a definitive handle on this number $\eta$ either observationally or theoretically, our value of 0.1 is consistent with (and even conservative compared to) estimates from GRB afterglow observations (Lloyd-Ronning & Zhang 2004; Fan & Piran 2006; Ryde & Pe'er 2009). Note that our results below for the values of the magnetic field and BH masses will scale accordingly, changing by a factor of about 3 for every order of magnitude change in the efficiency $\eta$.

We assume a jet opening angle of $\theta \approx 10^\circ$, so that the luminosity in the GRB jet is $L_{\text{GRB}} = \frac{1}{2}(1 - \cos(\theta))L_{\text{iso}}$, where $L_{\text{iso}}$ is the isotropic equivalent luminosity. Although - again - we don’t have extremely strong constraints on the jet opening angle from observations of long or short GRBs, in a number a cases a clear break in the afterglow light curve suggests the presence of jet with opening angles of $\sim 5$ to 10 degrees (for example, see results reported in Ghirlanda et al. (2004); Berger (2014); Fong et al. (2015); Dainotti et al. (2017); Wang et al. (2018)). In some cases, the inferred opening angle can be less by a factor of 2 or so (Wang et al. 2018). We point out the inferred average opening angle for long GRBs seems to be somewhat larger than that of short GRBs, although a value of $\theta = 10^\circ$ lies in the range for both populations (see, e.g. Figure 18 of Berger (2014)). A larger jet opening angle will require more power from our BZ jet; hence, we consider $\theta \sim 10^\circ$, a conservative estimate in the sense that it requires more extreme values of magnetic fields and BH masses compared to smaller jet opening angles.

We also assume a roughly poloidal magnetic field configuration on the BH, which produces the strongest possible jet (the details of the field structure are contained in the factor $k$ in equation 1 above). The magnetic field is estimated from the magnetic flux through the relationship $\phi \approx B\Omega R^2$, where $B$ is the magnetic field, $\Omega$ is the solid angle of the jet, and $R$ is the Kerr radius given by $R = GM/c^2 + \sqrt{(GM/c^2)^2 - a^2}$. Note that expressing magnetic flux in terms of magnetic field and relating the Kerr radius to the BH mass, we see that:

$$L_{\text{GRB}} \approx 10^{50} \text{erg} (\eta/0.1)(a/0.9)^2 (B/10^{16} \text{G})^2 (M_{\text{BH}}/5M_\odot)^2$$

(2)

Throughout this paper, we assume an $a = 0.9$, which is consistent (and in some cases comfortably above) the value needed to launch a GRB jet according to the simulations referenced in the introduction (e.g. McKinney et al. (2012); Ruiz et al. (2018)). Such high values of BH spin have indeed been inferred through observations of stellar mass BHs in X-ray binaries, as well as in supermassive BHs powering AGN (for reviews on these measurements, see McClintock et al. (2014) and Reynolds (2014)). Figure 1 shows BH spin parameters from X-ray binary observations (blue circles) and AGNs (green circles), using data from the previous two references. The data indicate that the majority of these objects have high values BH spins $a \gtrsim 0.6$. On the other hand, estimates of the spins of remnant BHs from BH-BH merger seen by LIGO indicate very low ($a \sim 0.05$) inferred spin values. Hence, the distribution of BH spin parameter and its dependence on the progenitor system is still an open question. We discuss the implications of our relatively high spin constraint - particularly in terms of the requirements it puts on GRB progenitor angular momentum - in §5.2 below.

3 RESULTS

3.1 Long GRBs

Given the assumptions above, we can explore the range of magnetic fields and BH masses required to produce observed GRB luminosities. For a representative set of long GRB luminosities, we use those in Table 1 from Wu et al. (2016). The Wu et al. (2016) sample is originally from Sonbas et al. (2015) who analyze a sample of Swift and Fermi GRBs with measurable prompt (gamma-ray) temporal statistics. For our purposes, we consider this sample representative of the long GRB population.

To estimate a reasonable range of BH masses, we assume the GRB comes from either a massive star collapse or a Helium merger (and indeed produces a BH rather than a NS inner engine). In this scenario, we can estimate the expected BH mass from the duration of the long GRB. For example, for typical long GRB durations between $20 - 100$ s, accreting at 0.1$M_\odot$/s (Popham et al. 1999), we need at least 2 to 10$M_\odot$ of material in the disk. For the progenitors mentioned above, this leaves anywhere from a few to tens of solar masses left over to form the BH (ignoring any additional mass expelled in the process; note the black hole seed mass from the iron core is $\sim 2M_\odot$ for a 10–15$M_\odot$ star before collapse; Woosley & Heger (2006)).

Within this mass range, we can then compute the magnetic field needed to satisfy equation 1, given our observed GRB luminosities. The left panel of Figure 2 shows the magnetic field at the horizon as a function of BH mass, for 22 values of observed long GRB luminosities. As one can see in this figure, for BH masses between 2 and 10 solar masses, we require field strengths from $\sim 5\times10^{14}G$ to $\sim 10^{17}G$ to power the GRB jet. We discuss both the generation of these magnetic fields as well as the implications of such high magnetic fields in §4 below.

3.1.1 Progenitor Magnetic Field

If we naively assume all of our BH flux originated from the progenitor star of the GRB, we can make an estimate of the nascent magnetic field needed to provide our required flux on the BH. In other words, we can get a crude estimate of the progenitor ("p") magnetic field under the simple assumption of magnetic flux conservation $B_p R_p^2 \approx B_{BH} R_{BH}^2$ (where we have assumed flux conservation over the same solid angle between the progenitor and BH system).

The left panel of Figure 3 shows a histogram of magnetic fields for a single progenitor of radius $R_\star = 10^{12}$cm and a remnant black hole mass of $5M_\odot$, under the assumption of magnetic flux conservation. The inferred fields are

$$L_{\text{GRB}} = \eta L_{\text{BZ}}.$$
roughly $\sim 10^6$ G. Typical surface magnetic fields of massive stars have been observed to be anywhere from 300G to 30kG (Walder et al. 2012; Hubrig et al. 2015). However, we are currently unable to probe the magnetic fields in the interiors of these stars, which in principle may be much larger. Recent astroseismology analysis of intermediate mass stars ($\sim 2M_\odot$) using Kepler data have shown that some stars exhibit evidence of interior fields $\sim 10^6$ to $10^7$ G (Stello et al. 2016). Hence, a dynamo in the convective cores of massive stars may indeed be able to generate such fields. Even if this is a rare phenomenon, it is possible that GRBs come from the subset of those stars able to produce such fields.

3.2 Short GRBs

As mentioned above, we assume the progenitor system for short GRBs is the merger of two NSs. We take our short GRB luminosities from Fong et al. (2015), where we define GRB prompt isotropic luminosity as $E_{iso}/T_{90}$. The Fong et al. (2015) sample includes all short GRBs with measured redshifts, and we consider this sample representative of the short GRB population.

Because NSs fall in a fairly narrow mass range, we have a limited mass range for a resultant BH, which we take between 0.5 and 4.0 $M_\odot$. The right panel of Figure 2 shows the magnetic field at the horizon as a function of BH mass, for 22 representative values of observed short GRB luminosities. For our range of BH masses, we require field strengths from $\sim 10^{15}$G to $\sim 10^{17}$G to power the short GRB jet.

3.2.1 Progenitor Magnetic Field

We can again simply naively assume that all of the magnetic flux in the BH-accretion disk system came from the progenitor system and was conserved. Under this (admittedly overly-simplistic) assumption, we can estimate the magnetic fields of the two NSs before the merger. The right panel of Figure 3 shows a histogram of the inferred magnetic fields of a DNS system in which the radius of each NS is $R_* = 10^6$ cm. These fields span the range from $3 \times 10^{14}$ G to $5 \times 10^{15}$ G. This is in the range inferred for so-called magnetars (Usov 1992; Thompson 1994), and NSs with magnetic fields in these ranges have been invoked as the engine behind a number of astrophysical phenomena, including long and short GRBs, soft gamma-ray repeaters, fast radio bursts, and superluminous supernovae (see, for example, the discussion in Margalit et al. (2018)). For larger assumed NS radius (or alternatively a harder equation of state), these field strengths will decrease.

4 ON HIGH MAGNETIC FIELDS

The magnetic fields inferred from our analysis are extreme, and it is worth discussing how such fields in practice can be generated and sustained. A number of studies have looked at the growth of magnetic fields, particularly in the case of double NS mergers. In these systems, Rosswog & Davies (2002) showed that massive fields $\sim 10^{17}$G can be produced through differential rotation of the central object in a DNS merger. Both Price & Rosswog (2006) and Obergaulinger et al. (2010) showed that Kelvin-Helmholtz instabilities can amplify fields to $\sim 10^{16}$G on very short timescales ($\sim ms$), although Obergaulinger et al. (2010) argue these fields are not long lasting (only a few milliseconds). Zrake & MacFadyen (2013) examined turbulent amplification of magnetic fields in DNS systems and showed fields up to $\sim 10^{16}$G can be produced.

For the more general case of a BH-accretion disk system, it is well known that a small extant magnetic field in the disk can be amplified by the magneto-rotational instability (MRI) (Velikhov 1959; Chandrasekhar 1960; Acheson & Hide 1973; Balbus & Hawley 1991).
4.1 MRI Growth

Simulations that examine the MRI in accretion disks\(^2\) generally start with some seed magnetic field \(B_0\), which grows at a rate \(\dot{B} = B_0 \nu e^{\omega t}\); the parameter \(\omega\) is the maximum growth rate of the MRI, \(\omega = (1/2)\nu d\sigma/dr\), where \(\sigma\) is the rotational velocity in the disk (Balbus & Hawley 1998).

From this equation for the growth rate, we can estimate the magnetic field growth from the MRI in the context of GRB disks. Because \(\sigma\) (which for our purposes here we simply take as the Keplerian velocity in the disk) depends on radius, we evaluate the maximum growth rate at some fiducial radius. One possibility for this radius is the innermost stable circular orbit (ISCO), \(R_{\text{ISCO}} = 6GM_{\text{BH}}/c^2\) for a non-rotating BH. For prograde rotation, which is a reasonable assumption for the BH and the disk, since they arose from the same (rotating) progenitor, the ISCO will be smaller than this and is given by:

\[
R_{\text{ISCO}} = \frac{6GM_{\text{BH}}}{c^2} \sqrt{3 + \frac{Z_2}{Z_1} + \frac{1 - x^2}{(1 + x)^{1/3}}} \frac{1}{(1 - \frac{x}{\sqrt{3}})^{1/3}}
\]

(3)

and \(x = a/R_{\text{Sch}} = a c^2 / 2M_{\text{BH}}\).

For the range of BH masses and spins we consider above for both long and short GRBs, we find the maximum growth rate \(\omega \sim 10^3 \text{s}^{-1}\). This means for a seed field \(B_0 \sim \mu\text{G}\), the MRI will take only \(t \sim 5 \times 10^{-5} \text{s}\) to reach fields of \(B \sim 10^{15}\text{G}\) at the maximum growth rate. This timescale is much shorter than the relevant timescales in both long and short GRBs, \(M/\dot{M} \sim \text{seconds} \) (where, again, \(M\) is the mass in the disk and \(\dot{M}\) is the accretion rate). In other words, if the MRI operates in a GRB disk in the linear regime, it is an efficient way to grow the magnetic field to the values we require to power the GRB-BZ jet.

This is of course a simple analytic estimate and ignores the non-linearities that come into play as the field develops. In particular, we’ve not considered at what point the magnetic field growth may saturate in the GRB disk. This is not a well understood problem and depends on many factors related to the microphysics as well as global structure of the disk. Popham et al. (1999) estimated GRB magnetic fields in a BZ context, assuming that the magnetic field energy density reaches 1% of the accretion disk kinetic energy density: \(B^2/8\pi \sim 0.01 \rho v^2\), where \(\rho\) is the density in the disk and \(v\) is the velocity. Using an \(a-\) disk prescription (Shakura & Sunyaev 1973), they estimated magnetic fields in the range of \(10^{14} - 10^{16}\text{G}\) (see their Table 4). Allowing the magnetic field energy to reach closer to equipartition with the disk energy can allow fields up to \(10^{19}\text{G}\), under these arguments.

We are also neglecting that the field grows at a different rate in different parts of the disk (the growth rate is fastest at the ISCO), and that the orbital velocity is not Keplerian near the horizon. Numerical simulations (see, e.g., figure 7 in Shiokawa et al. (2012)) indicate that magnetically-driven turbulence becomes dominant in the disk, and thus that the magnetic field is large, after about 500 \(GM_{\text{BH}}/c^2\), or \(t \sim 5 \times 10^{-3} \text{seconds}\). This is an order of magnitude longer than our rough growth time estimate above, but still far shorter than the characteristic time of the GRB. For further discussions of these issues, we refer the reader to Pessah & Psaltis (2005); Begelman & Pringle (2007); Bai & Stone (2013) and references therein.

4.2 MAD Disks

Given the possibility of growth of a large amount of magnetic flux near the ISCO of our BH-accretion disk, our system may very well be in a magnetically arrested disk (MAD) configuration (Narayan et al. 2003). In a MAD state, the magnetic field pressure is large enough to balance the ram

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\(^2\) See Abramowicz & Fragile (2013) and references therein for a review.
pressure of the infalling matter, and the accretion can be halted or arrested. Although it is an open question whether the disk can actually maintain such a high flux near the BH, simulations seem to indicate it may be viable (see, e.g., Tchekhovskoy et al. (2011) who examined MAD disks for similar values of magnetic flux and spin that we consider here). Studies by Liska et al. (2018a,b) and Morales Teixeira et al. (2018) found that even thin disks can maintain large magnetic flux on rotating BHs and a powerful outflow is launched (note these simulations were done in the context of AGN disks).

Recently Lloyd-Ronning et al. (2016, 2018) examined GRB properties in the context of the MAD model and found that the GRB variability timescale, as well as observed corre-

la) that the GRB variability timescale, as well as observed corre-

lating the bulk Lorentz factor \( \Gamma \) with minimum variability timescale \( t_{\text{min}} \) in the gamma-ray light curve \( (t_{\text{min}} \propto \Gamma^{-4.8 \pm 1.5}) \), and b) gamma-ray luminosity with the minimum variability timescale in the prompt light curve \( (t_{\text{min}} \propto L^{-1.1 \pm 0.1}) \) are naturally explained in a MAD model with a Blanford-Znajek jet.

In Lloyd-Ronning et al. (2018), accretion rates of long GRBs were estimated in the context of the MAD model using

\[
M = (0.1 M_\odot s^{-1}) \eta^{-2} \epsilon_{\text{MAD}} (L_{\text{GRB}}/10^{52} \text{ergs}^{-1})
\]

where \( \epsilon_{\text{MAD}} \sim 0.001 \) is a parameter estimating the degree of “arrestedness” of the accretion flow in the disc. Using isotropic equivalent luminosities, they found accretion rates for long GRBs \( \sim 0.5 M_\odot/s \), although ranging from \( \sim 0.005 \) to \( 10 M_\odot \) in the most extreme cases (see their Table 1). Because the average luminosity of our sample of short GRBs is lower than that of the long GRBs by roughly an order of magnitude, the accretion rates for short GRBs will be correspondingly lower by the same factor (i.e. \( M \sim 0.05 M_\odot/s \)). Correcting for the opening angle of the jet, these rates would drop by a factor of \( \sim 0.1 \) (our assumed jet opening angle in this work).

Finally, we briefly comment on QED processes that arise in the presence of strong magnetic fields. When the energy associated with the gyro-frequency of an electron (\( \omega_g \)) is equal to twice its rest mass (\( m_e \)) energy, pair production can occur in the presence of a magnetic field. This requires \( h \omega_g \gtrsim 2m_e c^2 \), where \( \omega_g = eB/m_e c \). Therefore in fields greater than \( B \gtrsim 2m_e c^2/h \approx 4 \times 10^{13} G \), electron-positron pairs can be produced (Dalgerty & Harding 1983). Some amount of pair production is necessary to keep the region around the BH force-free which allows the BH layer to operate with maximum power output (MacDonald & Thorne 1982). However, a third order QED process - photon splitting - can suppress pair production in high magnetic field (this has long been suggested as a mechanism to explain radio quiet pulsars Baring & Harding (1995, 2001)), if all three channels of photon splitting operate (Baring & Harding 2001).

In this work, we assume that neither of these processes (pair production or photon splitting) measurably compare to the energy in our magnetic field, nor do they affect the operation of the BH itself, so that they can effectively be ignored. A detailed examination of these effects in very strong magnetic fields is the subject of a future investigation.

5 PROGENITOR ANGULAR MOMENTUM

Throughout this paper, we have assumed a BH spin parameter of \( a \sim 0.9 \). As discussed above and displayed in Figure 1, there is ample observational evidence for BH spin parameters of this magnitude, in both stellar mass and supermassive BHs. Theoretically, it is still a question how to achieve these spin parameters in the process of forming the BH. Below, we discuss the conditions for which that assumption is reasonable, given our long and short gamma-ray progenitors.

5.1 Long GRBs

Figure 4 shows the spin parameter as a function of BH mass for a variety of initial stellar masses and metallicities, using a suite of models from Belczynski et al. (2017), which employs the Geneva stellar evolutionary code (Georgy et al. 2013). These models assume there are no dynamo-generated magnetic fields driving coupling between layers in the stars, and give a maximum to the rotation rate of any single star. Material is accreted onto the BH layer by layer and the angular momentum added to the BH is set by the angular momentum of the accreting material: \( j = \min (j_{\text{infall}}, j_{\text{SCCO}}) \) where \( j_{\text{infall}} \) is the specific angular momentum of the accreting material and \( j_{\text{SCCO}} \) is the angular momentum of a disk at the innermost stable circular orbit. This equation assumes that, if there is sufficient angular momentum to form a disk, only the angular momentum at the innermost stable circular orbit is accreted onto the BH. From this process, we can estimate the spin of the BH.

A wide range of angular momenta can be produced in stellar models, depending both on mass and metallicity. Not all of these systems will actually form the BH accretion disks needed to form GRBs. The size of the disk produced in these collapsing stars is determined by the radius at which the centrifugal force of the rotating, infalling material equals the gravitational force (based on the enclosed mass). For many of these systems, there is not enough angular mo-
momentum and this radius falls below the ISCO (so that no disk forms). Some systems with enough angular momentum (corresponding to a BH spin parameter of \( \sim 0.8 \)) do form a disk, depending on the initial properties of the progenitor and the BH mass. Figure 5 illustrates this point. This plot finds the radius where the centrifugal force equals gravity: \( r_{\text{disk}} = j^2 / G M_{\text{BH}} \) where \( j \) is the specific angular momentum of the infalling material, \( G \) is the gravitational constant and \( M_{\text{BH}} \) is the mass of the BH set by the mass of the material enclosed at each Lagrangian mass point.

A number of mechanisms have been studied showing that angular momentum is transported out of stellar cores, producing even more slowly rotating cores (see, e.g., the recent review by Aerts et al. (2018)); these stars will not have enough angular momentum to produce GRBs (Janiuk et al. 2018). Hence, GRBs may be the special subset of collapsing stars for which the remnant BH spin parameter is high (above 0.9).

The helium merger model (produced when a NS or BH spirals into the helium core of a massive star) produces even faster rotation rates (Fryer & Woosley 1998; Zhang & Fryer 2001). The disks formed tend to be large (above 10,000 km) and the subsequent accretion rates are lower. The BH spins for all of these systems will quickly spin up to values above 0.8-0.9 unless the initial compact remnant mass is large compared to the helium core.

5.2 Short GRBs

The spin period of the remnant of a double NS merger is very much an open question and depends on a number of things, including the masses of the NSs, their equations of state, the initial spins of each NS, etc. (for a recent discussion of these issues, see Radice et al. (2016); Piro et al. (2017)). Using results of Newtonian merger calculations (Korobkin et al. 2012) and various equilibrium states, Fryer et al. (2015) examined a suite of models to determine the fate of the compact remnant of a DNS merger (see, e.g., their Table 1). For core masses above \( 2M_\odot \), they find angular momenta (column 5 of their Table 1) corresponding to remnant core spin parameters between \( \sim 0.35 \) and \( \sim 0.65 \). Zappa et al. (2018) looked at the spin parameter of the BH remnant of a DNS merger using a large sample of numerical relativity simulations with different binary parameters and input physics. They find BH remnants with spin parameters between 0.6 and 0.9 for a wide range of models (finding an empirical relation between the the total gravitational radiation and the angular momentum of the remnant; see their Figure 4). Finally, Ruiz et al. (2016) also simulated DNS mergers for equal mass, magnetized NSs with an \( n=1 \) polytrope equation of state. They find a BH remnant with \( a \sim 0.74 \). These initial spins from the BH remnant of a DNS merger approach the high value of the spin parameter we assumed above to explain short GRB luminosities.

### 6 CONCLUSIONS

We have examined the constraints that observed GRB luminosities place on GRB inner engine properties in the context of a BZ jet powering the GRB. Our main results are as follows:

- For long GRBs with BH masses in the range of \( 2 - 10M_\odot \), and a BH spin parameter of \( a \sim 0.9 \), magnetic fields from \( \sim 10^{17} G \) (for the less massive BHs) down to \( \sim 5 \times 10^{15} G \) (for more massive BHs) are needed to explain the range of observed GRB luminosities, assuming 10% of the BZ power goes into GRB emission.
- For short GRBs with BH masses in the range of \( 0.5 - 4M_\odot \), and a BH spin parameter of \( a \sim 0.9 \), magnetic fields from \( \sim 10^{17} G \) (for the less massive BHs) down to \( 10^{15} G \) (for more massive BHs) are needed to explain the observed sGRB luminosity assuming 10% efficiency.
- The inferred fields of the progenitor systems under the simple assumption of magnetic flux conservation are \( \sim 10^6 G \) for single star/collapsar systems powering long GRBs, and \( \sim 10^{15} G \) for double NS systems powering short GRBs.
- These magnetic field values can be reached through MRI growth on timescales much shorter than the duration or variability timescales in GRBs. The consequences of such high magnetic flux include magnetically arrested disks (MADs), which have been shown to explain a number of observed phenomena in GRB light curves (Lloyd-Ronning et al. 2016, 2018).

We have assumed a relatively conservative jet opening angle of 10° throughout. For more narrowly beamed systems (and therefore less emitted luminosity), the constraints on our BZ jets are lessened.

The detailed properties of GRB disks and how they ultimately relate to the BZ jet power need to be explored further. For example, the estimates in Popham et al. (1999) indicate that for a constant accretion rate, the disk density will scale inversely with BH mass; therefore - under the assumption that the magnetic field is some fraction of the disk kinetic energy - higher mass BHs will result in smaller magnetic fields. This simple scaling argument, however, should be examined in detail. In particular, the physics of the accre-
tion process and how the rate of accretion scales with BH mass (and varies throughout a GRB) are important processes that will affect this picture. We note that in this work, we have neglected any contribution from a jet powered by neutrino annihilation (Eichler et al. 1989). For some disk models, this process could enhance jet power, possibly substantially.

Finally, simulations exploring the generation of high magnetic flux at the BH horizon and its consequences, as well as a more detailed look at the connection between progenitor angular momentum and the spin parameter of the BH remnant could help validate some of the assumptions in this work, and — more importantly — uncover the nature of the GRB inner engine.

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