

This is the accepted manuscript made available via CHORUS. The article has been published as:

Black holes, disks, and jets following binary mergers and stellar collapse: The narrow range of electromagnetic luminosities and accretion rates

Stuart L. Shapiro

Phys. Rev. D **95**, 101303 — Published 30 May 2017

DOI: [10.1103/PhysRevD.95.101303](https://doi.org/10.1103/PhysRevD.95.101303)

Black Holes, Disks and Jets Following Binary Mergers and Stellar Collapse: The Narrow Range of EM Luminosities and Accretion Rates

Stuart L. Shapiro^{1,*}

¹*Department of Physics, University of Illinois at Urbana-Champaign, Urbana, IL 61801*

We have performed magnetohydrodynamic simulations in general relativity of binary neutron star and binary black hole-neutron star mergers, as well as the magnetorotational collapse of supermassive stars. In many cases the outcome is a spinning black hole (BH) immersed in a magnetized disk, with a jet emanating from the poles of the BH. While their formation scenarios differ and their BH masses, as well as their disk masses, densities, and magnetic field strengths, vary by orders of magnitude, these features conspire to generate jet Poynting luminosities that all lie in the same, narrow range of $\sim 10^{52\pm1}$ erg s⁻¹. A similar result applies to their BH accretion rates upon jet launch, which is $\sim 0.1 - 10 M_{\odot}$ s⁻¹. We provide a simple model that explains these unanticipated findings. Interestingly, these luminosities reside in the same narrow range characterizing the observed luminosity distributions of over 400 short and long GRBs with distances inferred from spectroscopic redshifts or host galaxies. This result, together with the GRB lifetimes predicted by the model, supports the belief that a compact binary merger is the progenitor of an SGRB, while a massive, stellar magnetorotational collapse is the progenitor of an LGRB.

PACS numbers: 04.25.D-, 04.25.dg, 98.70.Rz, 47.75.+f

I. INTRODUCTION AND MOTIVATION

We have recently performed simulations of merging binary neutron stars (NSNSs) [1] and binary black hole-neutron stars (BHNSs) [2] in full general relativistic magnetohydrodynamics (GRMHD). The neutron stars (NSs) were threaded by dynamically weak dipole magnetic fields initially. These simulations revealed the launching of jets from the poles of the BH-disk remnants following the mergers. Our simulations began from the late binary inspiral phase, continued through the NS tidal disruption and merger phases and did not terminate until the magnetized disk of NS debris reached a state of quasistationary accretion onto the remnant BH. These are the first GRMHD simulations that demonstrated the launching of *bonafide* incipient jets following such mergers, i.e., outward streams of plasma from the poles of the spinning BH remnants, with flows confined and driven outward in a narrow beam by a collimated, tightly wound, helical magnetic field (see [1–3] for summaries of and references to this and earlier work). The plasma jet is accompanied by beam of electromagnetic (EM) Poynting radiation. This demonstration lends support to the suggestion [4–6] and widely-held notion that these mergers are the engines that power short gamma-ray bursts (SGRBs). If true, then observable EM radiation typical of SGRBs could accompany the gravitational waves (GWs) detected from such events.

We also have performed GRMHD simulations of the magnetorotational collapse of a supermassive star (SMS) [7] triggered by a relativistic dynamical instability at the endpoint of stellar evolution [8–11]. Such

stars could provide the seeds of the supermassive black holes (SMBHs) that reside in most galaxies, including the Milky Way, and are believed to power AGNs and quasars [12]. These simulations, which can be scaled to stars of arbitrary mass, may also furnish crude models of the collapse of very massive Pop III stars, alternative sources of SMBH seeds [13]. Scaling the mass downwards to “collapsars”, they may provide simplified models of long GRBs (LGRBs) [14]. We shall collectively refer to these simulations as the SMS scenario. The simulations, performed in full 3+1 dimensions, showed that, upon arriving at the onset of instability, an SMS threaded by a dynamically weak magnetic field and spinning uniformly at the mass-shedding limit collapses to a spinning BH remnant immersed in a magnetized accretion disk. The results found for the BH and disk masses and BH spins were in good agreement with all previous GR simulations that started with the same uniformly rotating, unstable star and adopted the same radiation pressure-dominated EOS. But all such previous simulations were performed in axisymmetry (see, e.g., [15–17]), and only [17] incorporated magnetic fields (stellar interior only). These latest 3D simulations also agreed with analytic predictions [18] for the key remnant parameters. However, by threading the star with a dynamically weak dipolar magnetic field and evolving it for many dynamical timescales following collapse ($\Delta t \gtrsim 30,000M$) these 3D simulations launch a jet from the poles of the BH once the disk has settled down to a state of quasistationary accretion. Once again, the GW burst is accompanied by an appreciable EM Poynting flux.

In both our compact binary merger and SMS collapse scenarios we showed that the EM power generated by the jets is likely the result of the Blandford-Znajek (BZ) mechanism [19]. In such a case the Poynting luminosity

*Also Department of Astronomy and NCSA, University of Illinois at Urbana-Champaign, Urbana, IL 61801

is given roughly by

$$L_{\text{BZ}} \sim B_p^2 M_{\text{BH}}^2 \left(\frac{a}{M_{\text{BH}}} \right)^2 [\mathcal{L}_0] \sim 10^{52 \pm 1} \text{ erg s}^{-1} \quad (1)$$

where B_p is the poloidal B -field above the BH poles, M_{BH} is the BH mass, and a/M_{BH} is the BH spin parameter, all in gravitational units with $G = c = 1$. The corresponding quasistationary rest-mass accretion rate is given roughly by

$$\dot{M}_{\text{BH}} \sim 4\pi\rho M_{\text{BH}}^2 [\dot{\mathcal{M}}_0] \sim 0.1 - 10 M_\odot \text{ s}^{-1}, \quad (2)$$

where ρ is the rest-mass density near the BH horizon. The quantities \mathcal{L}_0 and $\dot{\mathcal{M}}_0$ appearing in Eqs. (1) and (2) restore the factors of G and c and with them the physical dimensions for L_{BZ} and \dot{M}_{BH} :

$$\begin{aligned} \mathcal{L}_0 &\equiv c^5/G = 3.6 \times 10^{59} \text{ erg s}^{-1}, \\ \dot{\mathcal{M}}_0 &\equiv c^3/G = 2.0 \times 10^5 M_\odot \text{ s}^{-1}. \end{aligned} \quad (3)$$

The collective numerical results of our simulations pose the following puzzle: *why do the physical magnitudes of L_{BZ} and \dot{M}_{BH} found for the different remnant BH-disk-jet systems in the NSNS, BHNS and SMS cases we have simulated all turn out to be comparable?* The fact is that their physical values are *all* within the rather narrow ranges indicated on the right-hand sides of Eqs. (1) and (2). This finding is most unexpected, as the scenarios leading the formation of the BH-disk-jet systems are very different and, more significantly, the masses, length and time scales, and the B -fields and densities characterizing these systems can differ by *many* orders of magnitude, particularly when we consider SMS progenitors. For example, the compact binary mergers form stellar mass BHs, while the SMS collapse scenario applies to the formation of arbitrarily large SMBH masses. How do we explain why their values of L_{BZ} and \dot{M}_{BH} are all within an order of magnitude or two of each other in light of the middle expressions in Eqs. (1) and (2), which exhibit quadratic scaling with M_{BH} , as well as quadratic scaling with B_p and linear scaling with ρ ? The remnant SMBH masses can vary over many decades for different SMS mass progenitors, and they can differ by orders of magnitude from the masses in the compact binary scenarios. Similar differences apply to B_p and ρ . How can these parameters conspire to generate similar luminosities and similar rates of accretion for the initial data we evolved?

We will resolve this puzzle below and discuss some of its implications. Unless explicitly noted otherwise, we adopt geometrized units with $G = 1 = c$.

II. EXPLANATION

Recall that the BZ luminosity measures a Poynting flux $\vec{S} = (\vec{E} \times \vec{B})/4\pi$ propagating out from above the

BH poles and is given, very roughly, by

$$L_{\text{BZ}} \sim S\pi r_{\text{H}}^2 \sim \frac{B_p^2}{4\pi} \pi r_{\text{H}}^2 \sim B_p^2 M_{\text{BH}}^2, \quad (4)$$

where r_{H} is the BH horizon radius. Since no flux is generated for nonspinning BHs, there must also be a dependence on the BH spin parameter. So more precisely, combining \vec{S} with the MHD relation $\vec{E} = -\vec{v} \times \vec{B}$ for a rotating cylindrical region of gas above a BH pole gives a Poynting luminosity $L_{\text{Poynt}} \sim (B_p B_\phi / 4\pi) (r_{\text{H}} \Omega_{\text{H}}) (\pi r_{\text{H}}^2)$. Here $\Omega_{\text{H}} = a/(2M_{\text{BH}} r_{\text{H}})$ is the angular velocity of the horizon. The maximum rate at which this flux can be generated occurs for the maximal value of B_ϕ , $B_\phi(\text{max}) \sim (r_{\text{H}} \Omega_{\text{H}}) B_p$ (see, e.g., [20], Eq. 2) which gives the BZ luminosity quoted in [21], Eq. (4.50),

$$L_{\text{BZ}} \sim \frac{1}{128} B_p^2 r_{\text{H}}^2 \left(\frac{a}{M_{\text{BH}}} \right)^2, \quad (5)$$

or, scaled to the BHNS remnant parameters in the simulations of [2],

$$L_{\text{BZ}} \sim 10^{51} \left(\frac{B_p}{10^{15} \text{ G}} \right)^2 \left(\frac{M_{\text{BH}}}{5.6 M_\odot} \right)^2 \left(\frac{a}{M_{\text{BH}}} \right)^2 \text{ erg s}^{-1}. \quad (6)$$

Now an outward jet can form only when the magnetic field is sufficiently strong to confine the plasma, reverse the inflow above the BH poles and drive an outflow. This ability requires the polar field to be force-free, $B_p^2/8\pi\rho_p \gtrsim 1$, which when satisfied is found to trigger the launching of a jet within a tightly wound, collimated, helical magnetic funnel above the poles. The force-free condition is satisfied once the plasma density above the poles has been sufficiently reduced and the initial magnetic field has been sufficiently amplified. Magnetic winding, MRI and the Kelvin-Helmholtz instability all contribute [1, 2, 7, 22]. The magnetic field exhibits a coherent polar component with a net vertical flux threading the BH horizon. The characteristic magnitude of the plasma density is established very roughly by the mean density in the disk,

$$\begin{aligned} \rho &\sim \frac{M_{\text{disk}}}{2H_{\text{disk}}\pi R_{\text{disk}}^2} \\ &\sim \left(\frac{1}{2\pi} \right) \left(\frac{M_{\text{disk}}}{M_{\text{BH}}} \right) \left(\frac{M_{\text{BH}}}{R_{\text{disk}}} \right)^2 \left(\frac{M_{\text{BH}}}{H_{\text{disk}}} \right) M_{\text{BH}}^{-2}. \end{aligned} \quad (7)$$

The disk is a bloated toroid of half-thickness H_{disk} and outer radius R_{disk} that engulfs the BH. The radius is determined by the angular momentum of the gas, which is set by the outermost layers of the NS following tidal break-up for binary mergers and by the angular momentum in the spinning outer envelope of the SMS progenitor for magnetorotational collapse. In all cases the disks are geometrically thick. For typical cases, whenever a jet is launched, ρ_p has fallen to values less than $\sim 0.1 - 0.01$ times the mean density ρ in the disk. At the same time the ratio $B_p^2/8\pi\rho_p$ grows to over $10 - 100$ in the jet once

it is fully established. So very crudely these two factors cancel and we can write

$$B_p^2 \sim 8\pi\rho \quad (8)$$

to relate B_p near the BH poles to ρ in the disk. In this spirit we will neglect the detailed toroidal geometry of the disk and simply set $2H_{\text{disk}} \sim R_{\text{disk}}$ in Eq. (7) to get

$$\rho \sim \frac{M_{\text{disk}}}{\pi R_{\text{disk}}^3} \sim \frac{1}{\pi} \left(\frac{M_{\text{disk}}}{M_{\text{BH}}} \right) \left(\frac{M_{\text{BH}}}{R_{\text{disk}}} \right)^3 M_{\text{BH}}^{-2}. \quad (9)$$

Combining Eqs. (8) and (9) yields

$$B_p^2 M_{\text{BH}}^2 \sim 8\pi\rho M_{\text{BH}}^2 \sim 8 \left(\frac{M_{\text{disk}}}{M_{\text{BH}}} \right) \left(\frac{M_{\text{BH}}}{R_{\text{disk}}} \right)^3. \quad (10)$$

Inserting Eq. (10) into Eq. (5) now yields

$$L_{\text{BZ}} \sim \frac{1}{10} \left(\frac{M_{\text{disk}}}{M_{\text{BH}}} \right) \left(\frac{M_{\text{BH}}}{R_{\text{disk}}} \right)^3 \left(\frac{a}{M_{\text{BH}}} \right)^2 [\mathcal{L}_0], \quad (11)$$

where employing gravitational units makes each factor in parenthesis nondimensional, so that inserting \mathcal{L}_0 again restores physical dimensions to L_{BZ} . Similarly inserting Eq. (10) into Eq. (2) gives

$$\dot{M}_{\text{BH}} \sim 4 \left(\frac{M_{\text{disk}}}{M_{\text{BH}}} \right) \left(\frac{M_{\text{BH}}}{R_{\text{disk}}} \right)^3 [\dot{\mathcal{M}}_0] \quad (12)$$

Eq. (10) provides the crucial relations: this equation reveals that the products $B_p^2 M_{\text{BH}}^2$ and ρM_{BH}^2 appearing in Eqs. (1) and (2) are comparable in magnitude and determined by the disk mass and size normalized to the corresponding remnant BH parameters. For all the different compact binary merger and SMS collapse simulations that we have found to launch *bonafide* incipient jets, these nondimensional, *normalized* disk mass and size ratios fall within a narrow range. For this reason it is no longer so surprising why the values of \mathcal{L}_{BZ} and $\dot{\mathcal{M}}_{\text{BH}}$, given by Eqs. (11) and (12), also fall in a narrow range. We give examples in the next section.

Combining Eqs. (11) and (12) yields the Poynting radiation efficiency ϵ :

$$\epsilon = \frac{L_{\text{BZ}}}{\dot{M}_{\text{BH}}} \sim \frac{1}{40} \left(\frac{a}{M_{\text{BH}}} \right)^2. \quad (13)$$

The Poynting efficiency for the jet is less than the EM efficiency for a corotating, *thin* Keplerian disk accreting onto a rapidly spinning BH and is zero for a nonspinning Schwarzschild BH. However, the EM emission from the *thick* toroids found here may be less efficient and will likely appear with a different spectrum, particularly if the Poynting flux from the jets is ultimately pumped into prompt γ -rays.

TABLE I: NSNS remnant after jet launch

Source	ρ (g/cm ³) ^a	B_p (G) ^a	\dot{M}_{BH} (M _⊙ /s)	L_{BZ} (erg/s)
model ^b	10 ⁹	10 ¹⁵	0.1	10 ⁵¹
simulation ^c	10 ⁹	10 ¹⁶	0.1	10 ⁵¹

^a ρ and B_p^2 scale as $(M_{\text{BH}}/2.85M_{\odot})^{-2}$

^bEqs. (10)-(12,14)

^c[1]

III. APPLICATIONS

A. NSNS Mergers

Consider the merger of identical, magnetized, irrotational NSNSs in circular orbit and obeying a stiff Γ -law EOS with $\Gamma = 2$. Take each star to have a fraction 0.81 of the rest-mass of the maximum-mass TOV star constructed from the same EOS and thread them with a dynamically weak, interior, dipole B -field (i.e, pressure ratio $P_{\text{mag}}/P_{\text{gas}} \ll 1$) and a strong, exterior dipole field (i.e., a pulsar-like magnetosphere, $P_{\text{mag}}/P_{\text{gas}} \gg 1$) [1]. The merger results in the formation of a hypermassive neutron star (HMNS) that undergoes delayed collapse to a spinning black hole immersed in a magnetized disk. The delayed collapse of the HMNS is crucial for the buildup of the magnetic field and jet launch, which do not occur for prompt collapse characterizing more massive NSNSs [23]. We extract the following results from the simulation for the nondimensional ratios appearing in Eqs. (10)-(12) after jet launch:

$$\frac{M_{\text{disk}}}{M_{\text{BH}}} \sim 1.1 \times 10^{-2}, \quad \frac{R_{\text{disk}}}{M_{\text{H}}} \sim 56, \quad \frac{a}{M_{\text{BH}}} \sim 0.74. \quad (14)$$

Most of the rest-mass goes into the remnant BH following collapse of the HMNS. Inserting the above ratios into Eqs. (10)-(12) yields the model values listed in the first row of Table I. The second row displays the values extracted from the simulation data following jet launch. Given the approximate nature of the above model equations, where the numerical coefficients are crude estimates at best, and given the raw extraction of time-varying simulation data from representative snapshots, we quote only the orders of magnitude of the tabulated quantities. Allowing for the rough nature of our simple analysis, the model expectations are consistent with the simulation data.

B. BHNS Mergers

Now consider the merger of a magnetized, irrotational NS in circular orbit about a spinning BH three times its mass [2]. The NS, which again obeys a $\Gamma = 2$ EOS, is a fraction 0.83 of the rest-mass of the maximum-mass TOV star and is again threaded by a dynamically weak, interior, dipole B -field and a strong, pulsar-like, exterior dipole field. The black hole has an initial spin

TABLE II: BHNS remnant after jet launch

Source	ρ (g/cm ³) ^a	B_p (G) ^a	\dot{M}_{BH} (M _⊙ /s)	L_{BZ} (erg/s)
model ^b	10 ¹⁰	10 ¹⁶	1	10 ⁵²
simulation ^c	10 ¹⁰	10 ¹⁵	0.1	10 ⁵¹

^a ρ and B_p^2 scale as $(M_{\text{BH}}/5.6M_{\odot})^{-2}$ ^bEqs. (10)-(12,15)^c[2]

TABLE III: SMS remnant after jet launch

Source	ρ (g/cm ³) ^a	B_p (G) ^a	\dot{M}_{BH} (M _⊙ /s)	L_{BZ} (erg/s)
model ^b	10 ⁻²	10 ¹⁰	0.1	10 ⁵¹
simulation ^c	0.1	10 ¹¹	1	10 ⁵²

^a ρ and B_p^2 scale as $(M_{\text{BH}}/10^6 M_{\odot})^{-2}$ ^bEqs. (10)-(12),16)^c[7]

$a/M_{\text{BH}} = 0.75$ aligned with orbital angular momentum. We extract the following results from the simulation for the nondimensional ratios appearing in Eqs. (10)-(12) after jet launch:

$$\frac{M_{\text{disk}}}{M_{\text{BH}}} \sim 2.5 \times 10^{-2}, \quad \frac{R_{\text{disk}}}{M_{\text{BH}}} \sim 30, \quad \frac{a}{M_{\text{BH}}} \sim 0.85. \quad (15)$$

When we insert the above ratios into Eqs. (10)-(12) and obtain the values recorded in Table II we find that the model is again consistent with the simulation data.

C. SMS Collapse

Consider finally the magnetorotational collapse of an SMS rotating uniformly at the mass-shedding limit and marginally unstable to a general relativistic dynamical (radial) instability [7]. The star is governed by a Γ -law, radiation pressure-dominated EOS with $\Gamma = 4/3$ and is again threaded by a dynamically weak, interior, dipolar B -field and a strong, pulsar-like, exterior field, all aligned with the spin axis. The mass of the star is arbitrary. About 92% of the mass goes into forming a spinning black hole, with most of the rest comprising a bloated, toroidal disk. Eventually jets of gas confined by collimated, helical magnetic fields emerge from both poles of the BH. A simulation yields the following nondimensional ratios after jet launch:

$$\frac{M_{\text{disk}}}{M_{\text{BH}}} \sim 0.1, \quad \frac{R_{\text{disk}}}{M_{\text{BH}}} \sim 100, \quad \frac{a}{M_{\text{BH}}} \sim 0.68. \quad (16)$$

Using the data in Eq. (16), Eqs. (10)-(12) yield the results recorded in Table III, which are again compared with the data extracted from our SMS simulation. Once again the model expectations are roughly consistent with the numerical results. This SMS scenario for BH-disk-jet formation affords a strong test of our model, as the mass,

length, timescales and magnetic field strengths all differ by orders of magnitude from those characterizing the compact binary merger scenarios. A comparison of the values tabulated for ρ and B bear this out. In fact, these parameters can vary by orders of magnitude even within the SMS scenario itself by allowing the SMS progenitor mass to vary by orders of magnitude. Yet the predicted values of L_{BZ} and \dot{M}_{BH} for the compact binary merger and SMS collapse scenarios are within an order of magnitude of each other. This establishes the near-universality of these quantities, at least for the BH-disk-jet formation mechanisms we have simulated. For our SMS collapse scenario alone, L_{BZ} and \dot{M}_{BH} are *strictly* independent of the progenitor mass, assuming a fixed initial magnetic field topology and ratio of magnetic to gravitational potential energy.

IV. IMPLICATIONS

The restriction of jet Poynting luminosities and disk rest-mass accretion rates to a narrow range in our simulations was unanticipated, but is no longer a mystery. Even within our model we still expect some variation in these quantities for different merger or collapse cases, as the disk and black hole spin ratios appearing on right-hand sides of Eqs. (11) and (12) will never be identical for different scenarios leading to BH-disk-jet systems. But each of these ratios can never exceed unity, so that our model predicts that the extreme upper limit, never approached in our simulations to date, is $\sim \mathcal{L}_0$ for the Poynting luminosity and $\sim \mathcal{M}_0$ for the accretion rate. The fact that the luminosities and accretion rates in our simulations are universally 6-7 orders of magnitude smaller simply reflects the fact that in all cases the disk mass is typically $\lesssim 0.1$ times smaller than the BH mass and the disk radius is larger than the BH horizon radius by a factor $\gtrsim 50$.

There are several distinguishing, observable features that are expected to be quite different for the different cases. One of these is the disk lifetime, which we can estimate using Eq. (12):

$$t_{\text{disk}}/M_{\text{BH}} \sim \frac{(M_{\text{disk}}/M_{\text{BH}})}{\dot{M}_{\text{BH}}} \sim \frac{1}{4} \left(\frac{R_{\text{disk}}}{M_{\text{BH}}} \right)^3 \left[\frac{1}{\mathcal{M}_0} \right]. \quad (17)$$

Substituting Eqs. (14)-(16) into Eq. (17) then gives

$$\begin{aligned} t_{\text{disk}} &\sim 0.6 (M_{\text{BH}}/2.85M_{\odot}) \text{ s} & (\text{NSNS}) \\ &\sim 0.2 (M_{\text{BH}}/5.6M_{\odot}) \text{ s} & (\text{BHNS}) \\ &\sim 1 \times 10^6 (M_{\text{BH}}/10^6 M_{\odot}) \text{ s} & (\text{SMS}). \end{aligned} \quad (18)$$

If the BH-disk-jet remnants in these systems are the engines that power GRBs, then the disk lifetimes computed above (and not simply the compact binary merger or stellar collapse timescales) might be directly related to the observed prompt γ -ray burst (rest-frame) lifetimes. Eqn. (18) is then consistent with the widely-held view

that NSNS and/or BHNS mergers power SGRBs, with burst timescales $\lesssim 2$ s. The widely-held view that magnetorotational, massive stellar collapse may give rise to LGRBs, with burst timescales $\gtrsim 2$ s, is also consistent with Eq. (18), provided our crude magnetorotational collapse model can be extended downward in mass to stars that collapse to black holes with $M_{\text{BH}} \gtrsim 10 M_{\odot}$ (“collapsars” [14]). Our model suggests that such collapse will give rise to Poynting luminosities very comparable to those for compact binary mergers, but with longer burst timescales because M_{BH} and $(R_{\text{disk}}/M_{\text{BH}})^3$ are both larger. As the disk is consumed and the BH accretion rate falls, the jet luminosity presumably will decline rapidly, as in observed GRBs and in some MAD accretion models (see, e.g., [24] and references therein).

A recent survey [25] of 407 GRBs, which includes 375 GRBs with spectroscopic redshift measurements and 32 others with host galaxy information, is most revealing. Table 1 and the panels in Fig. 1 of this survey show a fairly narrow width of a couple of orders of magnitude centered at $\gtrsim 10^{52}$ erg/s in the distributions of the isotropic peak (prompt) γ -ray luminosities of both SGRBs and LGRBs. The narrow range of luminosities in our simulations centered near this value and explained with our simple model is consistent with this data. (The agreement might be improved further by correcting for beaming.) The magnitude of this luminosity enables these sources to be observed out to high redshift ($z_{\text{max}} = 8.23$) by present γ -ray satellites. This property bodes well for future instruments designed to use GRBs at high redshift as a tool to probe the early Universe [26].

Simulations describing alternative scenarios for SMS collapse also have been performed in GR. Consider a rapidly *differentially* rotating star whose collapse is triggered by lowering Γ to 1.33 to account for e^+e^- pair production [27]. When the star is seeded by a small initial $m = 2$ density perturbation it forms a binary SMBH during the collapse. The binary ultimately merges, forming a single BH-disk system. No magnetic fields are incorporated, so there is no jet. Had the star been threaded by a dynamically weak magnetic field and a

jet formed, the expected Poynting luminosity can be inferred from Eqn. (13), using the reported accretion rate ($\dot{M} = 6.7 \times 10^{-5} \dot{M}_0 = 13 M_{\odot}/\text{s}$) and remnant BH spin ($a/M_{\text{BH}} = 0.9$): $L_{\text{BZ}} \sim 10^{53}$ erg/s. This estimate falls in the narrow range of Poynting luminosities predicted by our model for BH-disk-jets formed following stellar collapse. It suggests that the GW signals calculated in [27] could be accompanied by counterpart EM bursts, were a magnetic field present. Detailed simulations in GRMHD are required to confirm these expectations.

While the *Fermi* detection [28] of a weak SGRB during the LIGO binary black hole (BHBH) event GW150914 [29] may be a chance coincidence, there are scenarios whereby a BHBH merger could occur within an ambient disk. If any scenario results in a BH-disk-jet remnant similar to those found in our simulations, then our model can estimate the disk parameters from the observed γ -ray luminosity ($L \sim 1.8 \times 10^{49}$ erg/s) and lifetime ($t \sim 1$ s). Eqs. (11), (12) and (17) yield

$$\frac{M_{\text{disk}}}{M_{\text{BH}}} \sim 1 \times 10^{-5}, \quad \frac{R_{\text{disk}}}{M_{\text{BH}}} \sim 20, \quad \dot{M}_{\text{BH}} \sim 0.9 \times 10^{-3} M_{\odot} \text{ s}^{-1}. \quad (19)$$

Further simulations and future detections will be required to assess this possibility.

Acknowledgments: It is a pleasure to thank E. Berger, C. Gammie, V. Paschalidis, M. Ruiz and L. Sun for useful discussions. This paper was supported in part by NSF Grants PHY-1300903 and PHY-1602536 and NASA Grant NN13AH44G at the University of Illinois at Urbana-Champaign. This work used the Extreme Science and Engineering Discovery Environment (XSEDE), which is supported by NSF grant number OCI-1053575. This research is part of the Blue Waters sustained-petascale computing project, which is supported by the National Science Foundation (award number OCI 07-25070) and the state of Illinois. Blue Waters is a joint effort of the University of Illinois at Urbana-Champaign and its National Center for Supercomputing Applications.

-
- [1] M. Ruiz, R. N. Lang, V. Paschalidis, and S. L. Shapiro, *Astrophys. J. Lett.* **824**, L6 (2016).
 - [2] V. Paschalidis, M. Ruiz, and S. L. Shapiro, *Astrophys. J. Lett.* **806**, L14 (2015).
 - [3] V. Paschalidis, *Class. Quant. Grav.* **34**, 084002 (2017).
 - [4] D. Eichler, M. Livio, T. Piran, and D. N. Schramm, *Nature* **340**, 126 (1989).
 - [5] R. Narayan, B. Paczynski, and T. Piran, *Astrophys. J. Lett.* **395**, L83 (1992).
 - [6] R. Mochkovitch, M. Hernanz, J. Isern, and X. Martin, *Nature* **361**, 236 (1993).
 - [7] L. Sun, V. Paschalidis, M. Ruiz, and S. L. Shapiro (2017), [arXiv:1704.04502](https://arxiv.org/abs/1704.04502).
 - [8] S. Chandrasekhar, *Phys. Rev. Lett.* **12**, 114 (1964a).
 - [9] S. Chandrasekhar, *Astrophys. J.* **140**, 417 (1964b).
 - [10] R. P. Feynman, unpublished work cited in W. A. Fowler, *Rev. Mod. Phys.* **36**, 545 (1964).
 - [11] T. W. Baumgarte and S. L. Shapiro, *Astrophys. J.* **526**, 941 (1999).
 - [12] M. J. Rees, *ARA&A* **22**, 471 (1984).
 - [13] P. Madau and M. J. Rees, *Astrophys. J. Lett.* **551**, L27 (2001).
 - [14] A. I. MacFadyen and S. E. Woosley, *Astrophys. J.* **524**, 262 (1999).
 - [15] M. Shibata and S. L. Shapiro, *Astrophys. J. Lett.* **572**, L39 (2002).
 - [16] M. Shibata, Y. Sekiguchi, H. Uchida, and H. Umeda, *Phys. Rev. D* **94**, 021501 (2016).

- [17] Y. T. Liu, S. L. Shapiro, and B. C. Stephens, Phys. Rev. D **76**, 084017 (2007).
- [18] S. L. Shapiro and M. Shibata, Astrophys. J. **577**, 904 (2002).
- [19] R. D. Blandford and R. L. Znajek, Mon. Not. R. Astro. Soc. **179**, 433 (1977).
- [20] M. Livio, G. I. Ogilvie, and J. E. Pringle, Astrophys. J. **512**, 100 (1999).
- [21] K. S. Thorne, R. H. Price, and D. A. MacDonald, *Black Holes: The Membrane Paradigm* (Yale Univ. Press, New Haven, 1986).
- [22] K. Kiuchi, P. Cerdá-Durán, K. Kyutoku, Y. Sekiguchi, and M. Shibata, Phys. Rev. D **92**, 124034 (2015).
- [23] M. Ruiz, V. Paschalidis, and S. L. Shapiro, in preparation.
- [24] A. Tchekhovskoy and D. Giannios, Mon. Not. R. Astro. Soc. **447**, 327 (2015).
- [25] Y. Li, B. Zhang, and H.-J. Lü, Astrophys. J. Supp. **227**, 7 (2016).
- [26] R. Salvaterra, JHEA **7**, 35 (2015).
- [27] C. Reisswig, C. D. Ott, E. Abdikamalov, R. Haas, P. Mösta, and E. Schnetter, Phys. Rev. Lett. **111**, 151101 (2013).
- [28] V. Connaughton, et al., Astrophys. J. Lett. **826**, L6 (2016).
- [29] B. P. Abbott, et al., Phys. Rev. Lett. **116**, 061102 (2016).