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POYNTING JETS FROM BLACK HOLES AND COSMOLOGICAL GAMMA-RAY BURSTS

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ABSTRACT

We discuss the properties of magnetically dominated jetlike outflows from stellar mass black holes surrounded by debris tori resulting from neutron star disruption. These jets may have narrow cores (along the rotation axis) that are almost free of baryons and attain very high bulk Lorentz factors $\gtrsim 10^6$. The jets give rise to a characteristic MeV to TeV emission as well as to relativistic shocks producing the usual MeV bursts. Because the outflow is highly directional, the properties of the observed gamma rays would depend on the viewing angle relative to the rotation axis. Even for the most intense bursts, which under the assumption of isotropic emission and substantial redshifts would be inferred to emit 10^{52}–10^{53} ergs, the efficiencies required are only 10^{-2} to 10^{-4}.
Subject headings: black hole physics — gamma rays: bursts — magnetic fields

1. INTRODUCTION

Gamma-ray bursts (GRBs) pose two sets of problems. The first is to account for the required large and sudden energy release. The prime candidate here is the formation of a compact object or the merger of a compact binary: this can trigger the requisite energy release (few 10^{51} ergs s^{-1} for an isotropic burst at cosmological distances), with a characteristic dynamical timescales as short as milliseconds. The second problem is how this energy is transformed into a relativistically outflowing plasma able to emit intense gamma rays with a nonthermal spectrum.

The literature on the second of these problems is already extensive. There have, in particular, been detailed calculations on the behavior of relativistic winds and fireballs. We have ourselves, in earlier papers (see, e.g., Rees & Mészáros 1992, 1994; Mészáros, Laguna, & Rees 1993; Mészáros, Rees, & Papathanassiou 1994, hereafter MRP94; Papathanassiou & Mészáros 1996; see also Paczyński 1990; Katz 1994; Sari, Narayan, & Piran 1996) addressed the physical processes in relativistic winds and fireballs. Motivating such work is the belief that compact objects can indeed generate such outflows. There have, however, been relatively few attempts to relate the physics of the outflow to a realistic model of what might energize it. (Among these are early suggestions by Paczyński 1991; Mészáros & Rees 1992; Narayan, Paczyński, & Piran 1992; Thompson 1994; and Usov 1994). These have involved either the reconversion of a burst of neutrino energy into pairs and gamma rays or else strong magnetic fields. Although the former model cannot be ruled out, because the beamed neutrino annihilation luminosity is marginal in black holes with a disrupted neutron star torus (Jaroszyński 1996), initial calculations for neutron star mergers are discouraging (see, e.g., Ruffert et al. 1997); we here focus attention on magnetic mechanisms. In this Letter, we try to incorporate our earlier idealized models into a more realistic context and consider some of the distinctive consequences of an outflow that is directional rather than isotropic.

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2. MAGNETICALLY DRIVEN OUTFLOWS

Magnetic fields must be exceedingly high in order to transform rotational energy quickly enough into Poynting flux. At the "base" of the flow, at a radius $r_l \sim 10^6$–10^7 cm, where the rotation speeds may be of order c, the field strength must be at least $\sim 10^{14}$ G to transmit 10^{51} ergs in a few seconds. These fields are of course higher than those in typical pulsars. However, as several authors have noted, the field could be amplified by differential rotation, or even by dynamo action (see, e.g., Thompson & Duncan 1993). If, for instance, a single fast-rotating star collapses or two neutron stars spiral together producing (even very transiently) a differentially rotating disklike structure, it would take only a few dynamical timescales for the field to amplify to the requisite level (which is, incidentally, at least 2 orders of magnitude below the limit set by the virial theorem).

The most severe constraint on any acceptable model for gamma-ray bursts is that the baryon density in the outflow must be low enough to permit attainment of the requisite high Lorentz factors: the absolute minimum, for wind-type models invoking internal dissipation, is $\Gamma \sim 10^2$; impulsive models depending on interaction with an external medium require $\Gamma \sim 10^3$. Since the overall efficiency is unlikely to exceed 10^{-1}, this requires any entrained baryons to acquire at least 10^3 times their "pro rata share" of the released energy. There are of course other astrophysical situations in which an almost baryon-free outflow occurs—for instance, the wind from the Crab pulsar, which may contain essentially no baryons. However, this is not too surprising because the internal dissipation in pulsars is far too low to generate the Eddington luminosity. On the other hand, in GRBs, the overall luminosities are $\gtrsim 10^{13} L_{\rm Edd}$. It is hardly conceivable that the fraction channeled into thermal radiation is so low that radiation pressure does not drive a baryon outflow at some level. The issue is whether this level can be low enough to avoid excess "baryon poisoning."

When two neutron stars coalesce, some radiation-driven outflow will be induced by tidal dissipation before coalescence (Mészáros & Rees 1992). When the neutron stars have been disrupted, bulk differential rotation is likely to lead to more violent internal dissipation and a stronger radiation-driven outflow. Almost certainly, therefore, some parts of the outflow must, for some reason, be less accessible to the baryons.

3. AXISYMMETRIC DEBRIS AND JETS AROUND BLACK HOLES

One such reason might be that the bursts come from a black hole orbited by a disk annulus or torus (see, e.g., Paczyński 1991; Levinson & Eichler 1993). This is of course what happens when the central part of the rotating gaseous configuration has collapsed within its gravitational horizon; otherwise, there is no reason that material should avoid the center—indeed, there is more likely to be a central density peak in any noncollapsed configuration supported largely by rotation (either a single star or a compact binary after disruption).

Such a configuration could come about in two ways:

1. The spinning disk that forms when two neutron stars merge (see, e.g., Davies et al. 1994) probably exceeds the maximum permitted mass for a single neutron star; after viscosity had redistributed its angular momentum, it would evolve into a black hole (of 2–3 $M_\odot$) surrounded by a torus of mass about 0.1 $M_\odot$ (Ruffert et al. 1996).
2. The system may result from coalescence of a neutron star and black hole binary. If the hole mass is $\lesssim 5\ M_\odot$, the neutron star would be tidally disrupted before being swallowed, leading to a system resembling that described in scenario (1) but with a characteristic radius larger by a factor of 2 and a torus mass of ~ 1 instead of 0.1 $M_\odot$.

Numerical simulations yield important insights into the formation of such configurations (and the relative masses of the hole and the torus), but the Lorentz factors of the outflow are sensitive to much smaller mass fractions than they can yet resolve.

It is, however, a general feature of axisymmetric flows around black holes that the region near the axis tends to be empty. This is because the hole can swallow any material with angular momentum below some specific value: within a roughly paraboloidal "vortex" region around the symmetry axis (Fishbone & Moncrief 1976), infall or outflow are the only options. Loops of magnetic field anchored in the torus can enter this region, owing to "buoyancy" effects operating against the effective gravity owing to centrifugal effects at the vortex walls, just as they can rise from a flat gravitating disk. These would flow out along the axis. In addition, there can be an ordered poloidal field threading the hole, associated with a current ring in the torus. This ordered field (which would need to be the outcome of dynamo amplification rather than just differential shearing) can extract energy via the Blandford-Znajek (1977) effect. In the latter, the role of the torus is mainly to anchor the field: the power comes from the hole itself, whose spin energy can amount to $\sim 10^{53}$ ergs. Irrespective of the detailed field structure, there is good reason to expect any magnetically driven outflow to be less loaded with baryons along the rotation axis than in other directions. Field lines that thread the hole may be completely free of baryons.

4. BARYON-FREE OUTFLOWS

As a preliminary, we consider a magnetically driven outflow in which baryonic contamination can be neglected. In the context of many models, this may seem an unphysical limiting case: the difficulty is generally to understand how the baryonic contamination stays below the requisite threshold. But the comments in the previous section suggest that it is not impossible for the part of the flow that emanates from near a black hole and is channelled along directions aligned with the rotation axis.

There have been earlier discussions (dating back to Phinney 1982) of relativistic MHD flows from black holes in AGNs. In our present context, the values of L/M are larger by at least 10 orders of magnitude. This means that the effects of radiation pressure and drag are potentially much stronger relative to gravity; also, pair production due to photon-photon encounters is vastly more important. For an outflow of magnetic luminosity L and $e^{\pm}\gamma$ luminosity $L_w \lesssim L$ channeled into jets of opening angle θ at a lower radius $r_l = 10^6 r_6$ cm, the initial bulk Lorentz factor is $\Gamma_l \sim L/L_w$, and the comoving magnetic field, temperature, and pair density are $B_l' \sim 2.5 \times 10^{14} L_{51}^{1/2} \Gamma_l^{-1} r_6^{-1} \theta^{-1}$ G, $T_l' \sim 2.5 \times 10^{10} L_{51}^{1/4} \Gamma_l^{-3/4} r_6^{-1/2} \theta^{-1/2}$ K, and $n_l' \sim 4 \times 10^{32} L_{51}^{3/4} \Gamma_l^{-9/4} r_6^{-3/2} \theta^{-3/2}$ cm^{-3} (primed quantities are comoving). Unless $\Gamma_l \gg 1$, the jet will be loaded with pairs, very opaque and in local thermal equilibrium. It behaves as a relativistic gas that is "frozen in" to the magnetic field and that expands with $T' \propto r^{-1}$. The lab-frame transverse field $B \propto r^{-1}$, and the comoving $B' \sim B/\Gamma$. The comoving energy density (predominantly magnetic) is $\epsilon' \propto r^{-2}\Gamma^{-2}$, and the pair density is $n' \propto T'^3 \propto r^{-3}$ so $\Gamma \propto n'/\epsilon' \propto r$, or $\Gamma \sim \Gamma_1 (r/r_l)$.

When the comoving temperature approaches $m_e c^2/k$, the pairs start to annihilate, and their density drops exponentially, but as long as the scattering depth $\tau_T' > 1$, the annihilation photons remain trapped and continue to provide inertia, so $T' \propto r^{-1}$, $\Gamma \propto r$ persists until $T_a' \sim 0.04 m_e c^2 \simeq 17$ keV at a radius r_a, where $\tau_T'(r_a) \sim 1$. Between r_l and r_a this leads to $(r_a/r_l) \sim (T_l'/T_a') \simeq 10^2 L_{51}^{1/4} \Gamma_l^{-3/4} r_6^{-1/2} \theta^{-1/2}$, with $\Gamma_a \sim \Gamma_l (r_a/r_l) \simeq 10^2 L_{51}^{1/4} \Gamma_l^{1/4} r_6^{-1/2} \theta^{-1/2}$. At r_a, the adiabatic density $n'_{a,\mathrm{ad}} \sim n_l'(r_a/r_l)^3 \simeq 4 \times 10^{26}$ cm^{-3} is mostly photons, while the pair density from Saha's law is $n_a' \sim 1.5 \times 10^{18} \Gamma_l r_6^{-1}$, the photon to pair ratio being $\sim 10^8$. The annihilation photons streaming out from r_a appear, in the source frame, as photons of energy around $\Gamma_a 3kT_a' \sim 5 L_{51}^{1/4} \Gamma_l^{1/4} r_6^{-1/2} \theta^{-1/2}$ MeV.

Beyond r_a, the lab-frame pressure $B^2 \propto r^{-2}$, but the inertia is drastically reduced, since it is provided only by the surviving pairs, of which there are many fewer than the free-streaming photons. In the absence of any restraining force, $\Gamma \propto n'/\epsilon' \propto n'/B'^2 \propto B^2/n'$, and the gas would accelerate much faster than the previous $\Gamma \propto r$. (The pair density is still above that needed to carry the currents associated with the field, analogous to the Goldreich-Julian density, so MHD remains valid.) However, the Compton drag time remains very short, since even after $\tau_T' < 1$ when most photons are free-streaming, the pairs experience multiple scatterings with a small fraction of the (much more numerous) photons for some distance beyond r_a.

One can define an "isotropic" frame moving at $\Gamma_i \propto r$, in which the annihilation photons are isotropic. In the absence of the magnetic pressure, the drag would cause the electrons to continue to move with the radiation at $\Gamma_i \propto r$. The magnetic pressure, however, acting against a much reduced inertia, will tend to accelerate the electrons faster than this, and as soon as $\Gamma \gtrsim \Gamma_i$, aberration causes the photons to be blueshifted and incident on the jet from the forward direction, so the drag acts now as a brake. In the isotropic frame, the jet electron energy is $\gamma = \Gamma/\Gamma_i$, and its drag timescale is that needed for it to encounter a number of photons whose cumulative energy after scattering equals the energy per electron, $t_{\mathrm{dr},i} \sim m_e c^2/(u_{\mathrm{ph},i}\sigma c\gamma) = (m_e c^2 4\pi r^2 \Gamma_i^3/L\sigma_T\Gamma)$. In the lab frame this is Γ_i times longer, and the ratio of the drag time to the expansion time r/c must equal the ratio of the kinetic flux to the Poynting flux, $n_j' m_e c^2 \Gamma^2/[(B_l^2/4\pi)(r_l/r)^2]$, where σ is the scattering cross section and n_j' is the comoving pair density in the jet. This is satisfied for $\Gamma \sim \Gamma_a$ at r_a, and since the drag time is much shorter than the annihilation time, the pair number is approximately frozen, and $\Gamma \propto r^{5/2}$ for $r > r_a$. The upscattered photons will, in the observer frame, appear as a power-law extension of the annihilation spectrum, with photon number index -2, extending from $\sim 0.12\Gamma_a m_e c^2$ to $\lesssim \Gamma_p m_e c^2$.

The acceleration $\Gamma \propto r^{5/2}$ abates when the annihilation photons, of isotropic frame energy $0.12 m_e c^2 (r_a/r)$, are blueshifted in the jet frame to energies $\gtrsim m_e c^2$. Their directions are randomized by scattering, and collisions at large angles above threshold lead to runaway $\gamma\gamma \rightarrow e^{\pm}$ (the compactness parameter still being large). This occurs when $\Gamma_p \sim 10^7 L_{51}^{1/4}\Gamma_l^{1/4} r_6^{-1/2}\theta^{-1/2}$, at $r_p \sim 10^{10} L_{51}^{1/4}\Gamma_l^{-3/4} r_6^{1/2}\theta^{-1/2}$ cm. Thereafter, the threshold condition implies $\Gamma \propto r^2$, until the inertial limit $\Gamma_{\mathrm{in}} \sim 10^9 L_{51}^{1/4}\Gamma_l^{1/4} r_6^{1/2}\theta^{-1/2}$ is reached at $r_{\mathrm{in}} \sim 10^{11} L_{51}^{1/4}\Gamma_l^{-3/4} r_6 \theta^{-1/2}$ cm. Besides going into pairs, a reduced fraction of the Poynting energy may continue going into a scattered spectrum of number slope -2 up to $\Gamma_{\mathrm{in}} m_e c^2$.

However, in an outflow of $L \sim 10^{51} L_{51}$ ergs s^{-1} maintained for a time $t_w \sim$ few seconds, the relativistic jet would have to push its way out through an external medium, with consequent dissipation at the front end, as in double radio sources. Except for a brief initial transient ($\ll 1$ s in observer frame), the shock will be far outside the characteristic radii discussed so far. The external shock will be ultrarelativistic and slows down as $r^{-1/2}$. The jet material, moving with $\Gamma \gg 1$, therefore passes through a (reverse) shock, inside the contact discontinuity, which strengthens as $r^{1/2}$ as the external shock slows down. Since it is highly relativistic and the medium is highly magnetized, this reverse shock can emit much of the overall burst energy on a time $\sim t_w$. When $t \sim t_w$, the external shock has reached $r_d \sim 10^{16} L_{51}^{1/4} n_o^{-1/4} t_w^{1/2}\theta^{-1/2}$ cm, where n_o is the external density, and the Lorentz factor has dropped to $\sim 10^3 L_{51}^{1/8} n_o^{-1/8} t_w^{-1/4}\theta^{-1/4}$. This, as well as the expected radiation of the shocks at (and after) this stage is rather insensitive to whether the initial Lorentz factor is indeed $\sim 10^7$–10^9 or whether baryon loading has reduced it to $\sim 10^3$. After this, the flow is in the impulsive regime and produces an "external" shock GRB on an observer timescale $r_d/c\Gamma^2 \simeq t_w$, as in, e.g., Mészáros & Rees (1993) and MRP94.

5. RADIATION FROM HIGH-Γ MAGNETICALLY DRIVEN JETS

If the jet is indeed baryon free and therefore has a Lorentz factor $\sim 10^7$–10^9, an extra mechanism can tap the Poynting energy along its length (before it runs into external matter)—namely, interaction of pairs in the jet and annihilation photons along the jet with an ambient radiation field. In our context, this would be radiation emitted by the torus or baryon-loaded wind that is expected to flow outward, mainly in directions away from the rotation axis, forming a funnel that surrounds the jet.

The ambient radiation causes an additional drag, which limits the terminal Lorentz factor of the jet below the values calculated in § 4 in those regions where this is important. As a corollary, the Poynting flux is converted into directed high-energy radiation, after pair creation in the jet by interaction with the annihilation radiation. We cannot generally assume that the ambient radiation field is uniform across the whole jet, because it may not be able to penetrate to the axis, but the boost in photon energy can in principle be $\lesssim \Gamma^2$. This is similar to what is discussed in AGNs but with more extreme Lorentz factors and radiation densities. Since the ambient radiation has a luminosity $\gtrsim L_{\mathrm{Edd}}$ and the burst luminosity is larger by 10^{12}–10^{13}, this mechanism is significant only when the jet Lorentz factor has the very high values $\gtrsim 10^6$ characteristic of these baryon-free outflows.

The photons from the sides of the funnel emitted into the jet have energies $x_f = E_\gamma/m_e c^2 \sim 1/20$. When they pass transversely into the jet and are scattered by "cold" pairs, the scattering will be in the K-N regime for any $\Gamma > 20$. The electron will therefore recoil relativistically, acquiring a Lorentz factor $\sim\Gamma/20$ (in the comoving frame), and will then cool by synchrotron emission, isotropic in the jet frame. This radiation will, in the observer frame, be beamed along the jet and blueshifted by Γ. One readily sees that the net amplification (energy/incident energy) is $\sim\Gamma^2$. If the external photons instead interact with one of the beamed gamma rays of energy E in the source frame (typically near threshold for pair production) the resultant pair will have a Lorentz factor of order $\Gamma/(E/m_e c2)$ in the comoving frame and will again emit synchrotron, yielding the same amplification factor as before.

The interaction of the ambient photons with the beamed annihilation photons in the jet leads to a complicated cascade, where the jet Lorentz factor must be calculated self-consistently with the drag exerted by the ambient photons. A schematic cascade at $r \sim r_p$, where $\Gamma_p \sim 10^7$, would be as follows:

1. Interactions near r_p between funnel photons of lab frame energy $x_f \sim 10^{-1}$ and beamed annihilation jet photons of lab frame energy $x_1 \sim x_a \sim 10$ (produced by pair recombination in the outflow) lead to $e^{\pm}$ of energy $\gamma_2 \sim x_1 \sim 10$ (lab), or $\gamma_2' \sim \Gamma_p/\gamma_2 \sim 10^6$ (comoving jet frame). In the comoving magnetic field $B' \sim 6 \times 10^3$ G, these pairs produce synchrotron photons of energy $x_2' \sim 10^1$ (comoving), or $x_2 \sim x_2'\Gamma_p 2 \times 10^8$ (lab).
2. Photons x_2 interact with ambient photons x_f to produce pairs of energy $\gamma_3 \sim x_2 \sim 2 \times 10^8$ (lab), or $\gamma_3'/\Gamma_p \sim 2 \times 10^1$ (comoving). In the same comoving field, these produce synchrotron photons $x_3' \sim 10^{-8}$ (comoving) or $x_3 \sim 10^{-1}$ (lab).
3. Photons $x_3 \sim 10^{-1}$ are below threshold for pair production with funnel photons $x_f \sim 10^{-1}$, ending the cascade. The resulting photons have lab energies $E_3 \sim x_3 m_e c^2 \sim 50 L_{51}^{1/4}\Gamma_l^{1/4} r_6^{-1/2}\theta^{-1/2}$ keV. Of course, in a more realistic calculation, the self-consistent jet Lorentz factor may vary across the jet, and one needs to integrate over height. However, this simple example illustrates the types of processes involved.

6. DISCUSSION

From the scenario above, it follows that Poynting-dominated (or magnetically driven) outflows from collapsed stellar objects would lead, if viewed along the baryon-free jet core, to a γ-ray component resembling the GRB from the external shocks in "impulsive fireball" models. However, the reverse shock is here relativistic and may play a larger role than in impulsive models. Viewed at larger angles, the GRB component would resemble that from internal shocks (see below). The characteristic duration t_w and the (possibly complicated) light curve are controlled by the details of the Poynting luminosity production as a function of time, e.g., the cooling or accretion time of the debris around a centrally formed black hole.

Besides the "standard" GRB emission, Poynting-dominated outflows viewed along the jet core would be characterized by additional radiation components from the jet itself (§§ 4 and 5). The annihilation component peaks at $E_1 \sim 5L_{51}^{1/4}\Gamma_l^{1/2}r_6^{-1/2}\theta^{-1/2}$ MeV (or 50 MeV if $\theta \sim 0.1$, $\Gamma_l \sim 10$), with a power-law extension of photon number slope -2 going out to $\lesssim \Gamma_p m_e c^2 \sim 5L_{51}^{1/4}\Gamma_l^{1/2}r_6^{-1/2}\theta^{-1/2}$ TeV, if ambient photon drag limits Γ to $\lesssim \Gamma_p$ (otherwise, it could extend to $\Gamma_{\rm in} m_e c^2 \sim 500$ TeV. If, as argued by Illarionov & Krolik (1996) for AGNs, an outer skin of optical depth unity protects the core of the jet from ambient photons, this annihilation component could have a luminosity not far below the GRB MeV emission, and the Γ of the jet core would not be limited by ambient drag, only that of the skin. However, the skin depth is hard to estimate without taking the drag self-consistently into account. The cascade from the interaction of ambient (funnel) photons with jet photons leads to another jet radiation component. The simplified estimate of § 5 gives $E_3 \sim 50L_{51}^{1/4}\Gamma_l^{1/4}r_6^{-1/2}\theta^{-1/2}$ keV as a characteristic energy, with a power-law extension. The amplification factor of the cascade is $A_c \lesssim \Gamma^2 \lesssim 10^{14}L_{51}^{1/2}\Gamma_l^{1/2}r_6^{-1}\theta^{-1}$, and since the funnel emits $\sim L_{\rm Edd}$, the luminosity could be a (possibly small) fraction of L. The duration of these components is $\sim t_w$, preceding the normal MeV burst by $\sim t_w$, and it may be more narrowly beamed, arising from regions with $\Gamma_p \sim 10^7$, as opposed to $\Gamma \lesssim 10^3$ for the MeV components, and $\Gamma_a \sim 10^2$ for the peak annihilation component.

Gamma rays will be detected for a time t_w (longer if external shocks are expected). There may also be prolonged after effects (see, e.g., Mészáros & Rees 1997) with X-ray and optical fluxes decreasing as a power law in time. In a rotating black hole–torus system with a super-Eddington outflow, the baryon contamination will be minimal along the rotation axis and will increase at larger angles to it. A narrow, largely baryon-free jet core such that as discussed in §§ 4 and 5 would be surrounded by a debris torus producing a baryon-loaded, slower outflow acting as a funnel, which injects X-ray seed photons into the jet. This slower outflow would carry a fraction of the GRB luminosity L in kinetic energy form. For a slow outflow $\Gamma \sim 10$, this kinetic energy can be reconverted into nonthermal X-rays after a time $t_x \sim$ day if it is shock heated at a radius $r \sim 10^{15}$ cm, either by Alfvén wave heating or by interaction with a preejected subrelativistic shell of matter of $\sim 10^{-3}\, M_\odot$ as has been detected in SN 1987A. This could lead to the substantial X-ray afterglow detected $\sim$days later in GRB 970228 (Costa et al. 1997).

The observed time structure and luminosity of the gamma-ray burst would depend on the angle between the rotation axis and the line of sight. Viewed obliquely, the outflow has $\Gamma \lesssim 10$, and only an X-ray transient with some accompanying optical emission would be seen. At smaller angles to the jet axis, outflows with $\Gamma \lesssim 10^2$ would also be seen (which might originate at $r > 10^{10}$ cm by entrainment of slower baryonic matter by the faster jet, or they might already originate closer in), and the predominant radiation would arise from internal shocks (Rees & Mészáros 1994; Papathanassiou & Mészáros 1996), which can have complex, multiple-peaked light curves. Closer to the rotation axis, outflows with $\Gamma \sim 10^3$ may dominate, either at large radii or already lower down, and radiation from external (deceleration) shocks would be prominent (see, e.g., MRP94). Nearest to the rotation axis, if there is indeed an almost baryon-free core to the jet, the annihilation radiation power law component, the relativistic reverse shock radiation, and the cascade process (§§ 4 and 5) yield significant extra contributions, which arrive ahead of the external shock burst. The luminosity function for the burst population may in part (or even entirely) be determined by the beam angle distribution, jet edge effects, and the observer angle relative to the jet axis. A time variability would be expected from intermittence in the Poynting flux extraction or wobbling of the inner torus or black hole rotation axis, on timescales down to $t_v \sim r_l/c \sim 10^{-3}$ to 10^{-4} s.

If the Poynting flux is derived from the accretion energy of a NS-NS remnant torus of 0.1 $M_\odot$, and the gamma rays are concentrated within an angle $\theta \sim 10°$, the efficiency of conversion of rest mass into magnetic energy need only be 10^{-4} to generate a burst that (if isotropic) would be inferred to have an energy of 10^{51} ergs and only 10^{-2} to simulate a 10^{53} ergs isotropic burst. For a BH-NS merger, the torus may be $\sim 1\, M_\odot$, and the corresponding magnetic efficiencies required are 10^{-5} and 10^{-3}, respectively. Thus, even bursts whose inferred isotropic luminosities are 10^{52}–10^{53} ergs, owing either to high redshifts or to exceptional observed fluxes, require only very modest efficiencies. Further reductions in the required efficiencies are possible if the Poynting flux is due to the Blandford-Znajek mechanism, which can extract $\lesssim 10^{-1}$ of a near–maximally rotating Kerr black hole rest mass (as would be expected from a NS-NS merger), leading to equivalent isotropic energies $\sim 10^{53}(4\pi/2\pi\theta^2) \gg 10^{53}$ ergs. Even without the latter, somewhat speculative, possibility, it is clear that jetlike Poynting flows from tori around black holes can produce even the most intense bursts with comfortably low efficiencies. The detailed burst properties—particularly the rapid variability and bursts seen along the jet core, the blueshifted annihilation, and high-energy cascade photons—can help to pin down the parameters of the model, while the afterglow can provide information on the dynamics and mass flux in directions away from the jet axis.

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