ASYMMETRIC SUPERNOVAE FROM MAGNETOCENTRIFUGAL JETS

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ABSTRACT

Strong toroidal magnetic fields generated in stellar collapse can generate magnetocentrifugal jets in analogy to those found in simulations of black hole accretion. Magnetocentrifugal jets may explain why all core collapse supernovae are found to be substantially asymmetric and predominantly bipolar. We describe two phases: the initial LeBlanc-Wilson jet and a subsequent protopulsar or toroidal jet that propagates at about the core escape velocity. The prompt LeBlanc-Wilson jets will produce an excess of neutron-rich matter and hence cannot be the common origin of supernova explosions; similar but less severe problems arise with the protopulsar jet that may be alleviated by partial evacuation along the axis by rotation. The jets will produce bow shocks that tend to expel matter, including iron and silicon, into equatorial tori. This may help to account for observations of the element distribution in Cas A. A magnetic “switch” mechanism may apply in rare instances (when there is low density and large magnetic field), with subsequent increase in the speed and collimation of the toroidal jet. The conditions that turn the magnetic switch “on” would yield a jet that propagates rapidly and with small opening angle through the star, depositing relatively little momentum. The result could be enough infall to form a black hole. A third, highly relativistic jet from the rotating black hole could catch up to the protopulsar jet after it has emerged from the star. The interaction of these two jets plausibly could be the origin of the internal shocks thought to produce γ-ray bursts and could explain the presence of iron lines in the afterglow. Recent estimates that typical γ-ray burst energy is \( \sim 3 \times 10^{50} \) ergs imply either a very low efficiency for conversion of rotation into jets by the Blandford-Znajek mechanism or a rather rapid turnoff of the jet process even though the black hole still rotates rapidly. Magnetars and “hypernovae” might arise in an intermediate parameter regime of energetic jets that yield larger magnetic fields and provide more energy than the routine case, but that are not so tightly collimated that they yield failed supernova.

Subject headings: gamma rays: bursts — ISM: jets and outflows — pulsars: general — supernovae: general

1. INTRODUCTION

The problem of core-collapse has been with us for over 40 yr (Hoyle & Fowler 1960). Immediately after the discovery of pulsars, it was reasonable to explore the issue of whether or not the rotation and magnetic fields associated with pulsars could be a significant factor in the explosion mechanism (Ostriker & Gunn 1971; Bisnovatyi-Kogan 1971; Bisnovaty-i-Kogan & Ruzmaikin 1976; Kundt 1976). With typical dipole fields of \( 10^{12} \) G and rotation periods of several to several tens of milliseconds, implying eddydynamic power of only \( 10^{44} - 10^{45} \) ergs s\(^{-1}\), a strong robust explosion seemed unlikely. Several factors have led to a need to reexamine this conclusion. The principle one is the accumulating evidence that core collapse supernovae are distinctly asymmetric. Aside from its famous rings, Hubble Space Telescope observations of SN 1987A resolving the debris show that the ejecta are asymmetric with an axis that roughly aligns with the small axis of the rings (Pun et al. 2001; L. Wang et al., in preparation). Chandra X-Ray Observatory (CXO) observations of Cas A show that the jet and counterjet and associated structure are observable in the X-ray (Hughes et al. 2000; Hwang, Holt, & Petre 2000) as well as the optical (Fesen & Gunderson 1996 and references therein). The most direct evidence bearing on this topic is from supernova spectropolarimetry, which shows that substantial asymmetry is ubiquitous in core-collapse supernovae and that a significant portion show strong evidence for a single, wavelength-independent axis of symmetry (Wang et al. 1996, 2001). Many, even perhaps most, core-collapse supernovae are bipolar (Wang et al. 2001, and L. Wang, in preparation).

The strength of the asymmetry observed with polarimetry is higher (several percent) in supernovae of Type Ib and Ic that represent exploding bare nondegenerate cores (Wang et al. 2001). The degree of asymmetry also rises as a function of time for Type II supernovae that have retained their hydrogen envelopes (from \( \lesssim 1\% \) to \( \gtrsim 1\% \)) as the ejecta expand and one looks more deeply into the core material (Wang et al. 2001; Leonard et al. 2000, 2001). Both of these trends suggest that it is the inner machine, the core collapse mechanism itself, that is responsible for the asymmetry. The observed polarization requires significant asymmetry—axis ratios exceeding 2:1. To impose the observed strong asymmetry in the final homologous expansion it is plausible that an axial flow must be established and maintained for at least several dynamical timescales of the outer mantle/envelope. This is the operational definition of a jet.

It also is significant that CXO observations have determined that jets are routinely associated with pulsars, not only in binary systems like SS 433, but also in isolated objects like the Crab and Vela pulsars (Weisskopf et al. 2000; Helfand, Gotthelf, & Halpern 2001). One interpreta-
tion of these observations is that the present-day jets in these young objects are vestiges of a much more powerful MHD jet era that occurred in the first few seconds of the protopulsar phase, when the compact objects were still inside their progenitor stars as they first attained nuclear densities. During this period the transient values of the magnetic field and rotation could have greatly exceeded those observed today, and indeed, as we suggest here, the reduction of the rotation rate and magnetic field to their present-day values could have been the process of energy release that powered the initial explosion.

Asymmetries associated with neutrino emission may produce dynamical asymmetries, but it is not clear that they can account for the polarization observations. Neutrino asymmetries may yield a short-lived, essentially impulsive effect (Shimizu, Yamada, & Sato 1994; Burrows & Hayes 1996; Fryer & Heger 2000; Lai et al. 2001). Expansion and transverse pressure gradients can wipe out transient asymmetries before homologous expansion is achieved (Chevalier & Soker 1989). “Finger” asymmetries might be preserved, but it is unclear that they can reproduce the common feature of a single symmetry axis that is substantially independent of space and time (Wang et al. 2001). Sufficient neutrino impulse might be delivered to the neutron star to yield a substantial runaway velocity (Burrows & Hayes 1996; Spruit & Phinney 1998), but it is difficult to see how this impulse can be communicated in a substantial and permanent way to the final ejecta trajectories.

Highly resolved, fully three-dimensional, adaptive grid numerical calculations (Khokhlov 1998) have, however, established that nonrelativistic axial jets of energy of order $10^{51}$ ergs originating within the collapsed core can initiate a bipolar asymmetric supernova explosion that is consistent with the spectropolarimetry (Khokhlov et al. 1999; Khokhlov & Höflich 2001; Höflich, Khokhlov, & Wang 2001). Some imbalance in axial jets can also account for pulsar runaway velocities, specifically velocities that are parallel to the spin axis (Helfand et al. 2001 and references therein). While a combination of neutrino-induced and jet-induced explosion may prove necessary for complete understanding of core-collapse explosions, jets as computed by Khokhlov et al. are sufficient. In this paper, we will ignore the possibility of neutrino-induced explosions as we focus on the possibilities of jet formation and jet-induced supernovae. For purposes of discussion, we will consider the sort of structure that forms a stalled shock after core bounce.

Further work on jet-induced supernovae has been done by Khokhlov & Höflich (2001) and Höflich et al. (2001). These papers present well-resolved, three-dimensional jets that propagate from within the region of the original iron core until breakout from the star or stoppage in a red giant envelope in a single calculation that does not require artificial halting, rescaling, and restarting the jet. A bow shock forms at the head of the jet and spreads roughly cylindrically around each jet. The stellar matter is shocked by the bow shock and acts as a high-pressure confining medium by forming a cocoon around the jet. The jets become long bullets of high-density material moving through the background low-density material almost ballistically. Spreading is limited by a secondary shock that forms around each jet between the jet and the material already shocked by the bow shock. The laterally expanding bow shocks generated by the jets move toward the equator, where they collide with each other. The result is that the material in the equatorial plane is compressed and accelerated more than material in other directions (excluding the jet material). The result is that heavy elements (e.g., O, Ca) are characteristically ejected in tori along the equator. Iron, silicon, and other heavy elements in Cas A are distributed in this way (Hwang et al. 2000), and there is some evidence for this distribution in SN 1987A (Wang et al. 2001b).

The presence of a red giant envelope can stop the jet and dissipate the effects of the asymmetry in the outer envelope, but the asymmetry will still be manifest in the core. Radiative matter ejected in the jets can alter the ionization structure and hence the shape of the photosphere of the envelope even if the density structure is spherically symmetric (Höflich et al. 2001). This will generate a finite polarization, even though the density distribution is spherical and the jets are stopped deep within the star and may account for the early polarization observed in Type II supernovae (Leonard et al. 2000; Wang et al. 2001; L. Wang, in preparation). The jets are unstable to Kelvin-Helmholtz instabilities, but the growth time is long compared to the jet propagation time and no significant modulation of the jet is observed.

Faster jets tend to make weaker explosions by propagating so rapidly and narrowly through the star that there is relatively little time to generate the transverse shocks that cause the explosion. The opening half angle of the jet is approximately (Wheeler et al. 2000)

$$\theta \approx \frac{v_{\text{env}}}{v_{\text{bow}}} \approx 0.1 \text{ rad} \frac{v_{\text{env}, 8}}{v_{\text{bow}, 9}} \approx 5 \frac{v_{\text{env}, 8}}{v_{\text{bow}, 9}},$$

where $v_{\text{env}}$ is the speed of sound of the envelope and $v_{\text{bow}}$ is the speed of the bow shock, which is less than that of the inflow velocity of the jet material. As $v_{\text{bow}}$ approaches the speed of light, the jet will be very narrow and affect a very small volume of the star as it propagates out. This feature may be relevant to the outcome of the explosion process and to the possibility of making $\gamma$-ray bursts, as we will discuss in § 4.3. Radiative transfer effects within optically thin portions of the jet may lead to significant instabilities in faster, lower density, jets (P. Höflich 2001, private communication). (For related work in the context of black hole formation, see MacFadyen & Woosley 1999; Aloy, et al. 2000; MacFadyen, Woosley, & Heger 2001; Zhang & Woosley 2001.)

The question then arises as to the mechanism of the production of the jets in routine core collapse events. A significant role for rotation and magnetic fields is an obvious candidate. Any mechanism that purports to account for routine pulsar formation may involve large transient magnetic fields, but must ultimately be consistent with the distribution of deduced dipole strengths of “normal” pulsars of $\sim 10^{12}$–$10^{13}$ G and with the distribution of initial rotation periods. There is growing evidence for a considerable spread in these quantities (e.g., Kaspi et al. 2001 and references therein). In a minority of cases, the final dipole field might be consistent with a value of $\sim 10^{15}$ G, yielding a “magnetar” (Duncan & Thompson 1992). Growing evidence for neutron stars of that field strength has been obtained by RXTE and other facilities (Kouveliotou et al. 1999; Ibrhim et al. 2001). The nature of the birth event of a magnetar is a separate problem that is significant in its own right. The ensemble of normal pulsar and magnetar births must also be consistent with observed nucleosynthetic abundances, a
potentially crucial constraint on jet-induced supernova models that we discuss in § 4.

Possible physical mechanisms for inducing axial jets, asymmetric supernovae, and related phenomena driven by magnetorotational effects were considered by Wheeler et al. (2000). The means of amplifying magnetic fields by linear wrapping associated with differential rotation in the neutron star and possibly by $\alpha-\Omega$ dynamos was discussed. Attention was focused on the effect of the resulting net dipole field on the torquing of the infalling plasma and newly formed neutron star and on the creation of strong Poynting flux (in the form, initially, of ultrarelativistic MHD waves) at the time-variable speed of light circle in analogy to pulsar radiation mechanisms, albeit buried deeply in the core-collapse supernova ambience. Allusion was made to the existence and possible role of the primary toroidal field that is expected to form in routine collapse where proto–neutron-star rotation rates are not expected to trigger $\alpha-\Omega$ dynamos and exponential field growth.

In this paper, we will explore the capacity of the predominantly toroidal field generated near the core-mantle boundary to directly generate axial jets within core-collapse conditions by analogy with magnetocentrifugal models of jets in active galactic nuclei (AGNs) (Romanova et al. 1998; Koide, Shibata, & Kudoh 1999; Meier et al. 1997; Ustyugova et al. 1999; Meier 1999; Koide et al. 2000; and references therein). The magnetocentrifugal models in the literature usually focus on black hole conditions where the field is anchored in the disk and at infinity. The physical picture discussed in this paper may prove especially robust since the field is anchored in substantial, corporeal objects: the outer, still-collapsing layers of the progenitor star and the collapsed object—a neutron star. We find that the production of a strong toroidal field, substantially stronger than the $10^{12}$ G dipole field of a pulsar, is nearly inevitable, and strong axial jets driven by that field equally so. The mechanisms described here also may prove to have a close astrophysical analog in the production of supersonic magnetocentrifugal jets in young stellar objects where there is also an outer envelope of infalling matter and an inner, differentially rotating object (Konigl & Pudritz 2000; Lery & Frank 2000; Ouyed, Pudritz, & Stone 1997).

Our working hypothesis is thus that jets leading to asymmetric supernova explosions are associated with routine core collapse and pulsar formation in roughly 90% of the cases. Perhaps 10% of the core collapse events would be associated with magnetars. Magnetar formation might be associated with exceptionally strong explosions and asymmetries, but this is not necessarily the case. An even smaller fraction, perhaps $10^{-2}$ (Scalo & Wheeler 2002a) might involve black hole formation and the production of $\gamma$-ray bursts.

In § 2 we outline the circumstances by which a newly formed neutron star could yield MHD jets. In § 3 we discuss the basic physics of MHD jets and their formation. Section 4 explores the application of the physical principles of jet formation to supernova conditions, how magnetocentrifugal jets could be generated in core collapse, and some of their expected properties. We sketch a scenario in which extreme values of the rotation and magnetic field could lead to the failure of the initial supernova explosion with the subsequent collapse of the core to form a black hole and the attendant possibility to form a $\gamma$-ray burst. In § 5 we summarize our conclusions.

2. THE PRODUCTION OF JETS FROM NEUTRON STAR FORMATION

To understand routine, pulsar-forming, strongly asymmetric supernovae, there must be a deeper understanding of the generation of jets during the process of core collapse to form neutron stars.

We identify three stages where magnetorotational processes may come into play in the collapse process:

1. The initial collapse or LeBlanc-Wilson phase, hereafter referred to as the “LeBlanc-Wilson jet” or the “LW jet” (LeBlanc & Wilson 1970). In this phase, the wrapping of field lines of the original, precollapse core field within the nascent neutron star can give rise to a magnetocentrifugal jet or rising magnetic “bubble” that forms along the rotation axis at $R \approx 5 \times 10^{7}$ cm and $\rho \approx 10^{9}$ g cm$^{-3}$. Parameters typical of modern presupernova calculations imply that the LeBlanc-Wilson stage is probably too weak to form a strong jet and explosion with observed supernova energies. Constraints from the production of r-process material also argue that this phase must not be robust (§ 4.1).

2. The protopulsar phase. There will be further wrapping of field lines by linear amplification of the contracting, cooling, deleptonizing proto–neutron star that is differentially rotating with respect to the infalling matter and the uncollapsed outer portions of the star. The toroidal fields can form within the newly formed, contracting, neutron star, at the shearing boundary between the neutron star and the infalling matter, and within the infalling matter. Between $10^{51}$ and $10^{52}$ ergs of rotational energy are available during the contracting deleptonization phase. A substantial portion of this energy will go into a toroidal field generated by differential rotation with a strength $\sim 10^{14}$–$10^{16}$ G. The jets are launched by toroidal magnetic “springs” that evolve into a steady state, Blandford-Payne MHD wind (§ 3.1), the energy and momentum of which eject the stellar envelope. This is essentially the earliest phase in the birth of what will be the pulsar remnant and is characterized by a period of rapid spin-down due to the production of the bipolar MHD outflow. This is a natural environment for a magnetocentrifugal jet and may be the routine mechanism for jet-induced, asymmetric supernova explosions. We will refer to this as the “protopulsar jet” or “slow toroidal jet” phase.

3. In relatively rare circumstances, depending on the distribution of rotation rates of the presupernova cores, the new neutron star may attain a rotation period of $\sim 1$ ms after deleptonization where the shear can compete with the convective motions and an $\alpha-\Omega$ dynamo becomes possible. The resulting exponential growth of radial and toroidal fields generates fields of strength $10^{16}$ G or more. The rapid growth of the field coupled with the large rotational energy reservoir associated with millisecond rotation could lead to especially large explosions. A mix of a magnetocentrifugal jet and the emission of a strong Poynting flux at the light cylinder could lead to power of $\sim 10^{52}$ ergs s$^{-1}$ in extreme cases. It is possible that such cases have been witnessed as “hypernovae,” but strong asymmetries can affect estimates of energetics (Höflich, Wheeler, & Wang 1999). We will return to this topic in the conclusions. Another alternative is that very extreme conditions lead to especially fast, narrow jets that actually put less energy into the star (Khokhlov & Höflich 2001). We will refer to these as “fast toroidal jets.” In this case, the MHD explosion mechanism may fail in the first instance and lead to the formation of a black hole. The for-
mation of a rapidly accreting black hole may, in turn, lead to the formation of an even faster, highly relativistic jet that interacts with the first, slower jet to produce a γ-ray burst. We will explore this possibility in § 4.3.

Wheeler et al. (2000) discussed the possible physical processes, associated timescales, and energetics that could lead to the production of pulsars, jets, and asymmetric supernovae and to a 10^{15} G magnetar in more extreme circumstances in the collapse of a massive stellar core. Wheeler et al. (2000) noted that when the neutron star contracts and speeds up two significant things may happen. One is that the rotational energy increases. The energy becomes significantly larger than that required to produce a supernova and, especially for recent estimates (Frail et al. 2001; Panaitescu & Kumar 2001), to drive a cosmic γ-ray burst. In addition, the light cylinder may contract from a radius large compared to the Alfvén radius to a radius comparable to that of the neutron star. This could disrupt the structure of an organized dipole field and promote the generation of intense MHD waves. The frequency of the MHD waves, \( \sim \Omega_{\text{NS}} \sim 10^4 \) Hz, would always be less than the plasma frequency, so these waves would be reflected internally and perhaps also channeled up the rotation axis. Collimation of the flow of energy may thus be expected.

A jet emerging from deep in the collapsing iron core of a star will emerge from the helium core in 5–10 s. Wheeler et al. (2000; see also Nakamura 1998) argued that a jet launched after the deleptonization contraction of the neutron star could arrive at the surface of the helium star at about the same time. This argument is incorrect. The 5–10 s associated with the loss of neutrinos in the deleptonization phase represents the characteristic timescale for the luminosity, but the luminosity is enhanced at later phases by the contraction, which results in a higher temperature at the neutrinosphere. The timescale for the contraction of the radius and hence the spin-up of the neutron star is actually considerably shorter, of order 1 s (Wilson & Mayle 1993). This means that any stronger jet launched during the contraction will catch up to and dominate any magnetic “bubble” or jet launched promptly during the initial collapse before the first jet can emerge from the helium core. If this results in a supernova explosion, then any \( r \)-process material ejected, however slowly, from the neutron star in the first phase will be at least partially ejected in the explosion. Furthermore, if the collapse results in the formation of a black hole with the generation of a third, relativistic, jet, this jet will, in turn, catch up with any jet matter ejected earlier. The time when this happens will depend on the delay between the first iron-core collapse and the second proto-neutron-star collapse to form the black hole (§ 4.3).

3. MAGNETOCENTRIFUGAL JETS

3.1. The Blandford-Payne Mechanism: “Spring” and “Fling”

The steady state theory of magnetorotationally driven outflows was discussed by Blandford & Payne (1982) in the context of magnetized disks. Numerical treatments of such collimating winds have been performed by Ustyugova et al. (1995); Ouyed et al. (1997); Krasnopolsky, Li, & Blandford (1999); and others. In the steady state model, a magnetic field anchored at the base in a rotating object transfers energy and momentum to the plasma frozen onto the field lines. Differential rotation between the inner and outer regions of the outflowing plasma causes the field lines to be swept backward in an open helix; and the overall rotation of this structure pushes plasma upward parallel to the rotation axis and slowly compresses it toward that axis. Eventually this creates a collimated bipolar outflow at large distances from the central object. This process is sometimes referred to as the “fling” model, since it employs centrifugal action to initially accelerate plasma along field lines. The forces that eventually complete the acceleration and collimation are magnetic pressure upward parallel to the axis and hoop stress.

A similar, but time-dependent, process was investigated by Uchida & Shibata (1985) and more recently by Kudoh, Matsumoto, & Shibata (1998); Uchida et al. (1999); and Koide et al. (2000). In this case an initially poloidal magnetic field (e.g., uniform or dipole) threads a differentially rotating medium, such as a thick disk or torus. As the evolution proceeds, the differential rotation coils the field in a barber-pole fashion, lifting and pinching material along the rotation axis. Depending on the strength of the field and the differential rotation, the early evolution of this system can be quite dynamic. The field will not have had time to expand to a broad open structure and approach the Blandford-Payne steady state. Instead, it may wind quickly to very large toroidal strengths and then expand vertically along the rotation axis, dynamically lifting and pinching plasma along the way. For this reason, this more dynamical process is sometimes referred to as the “spring” mechanism for jet outflow. Both spring and fling, however, use the same magnetic pressure and hoop stresses to perform the final acceleration and collimation. While both processes were first discussed in the context of accretion disks, they will have close analogies in collapsing supernova cores, especially if the core rotation results in significant core flattening and an accretion-torus—like structure surrounding the core.

For nonrelativistic outflow, the power output of an MHD flow is given by (Blandford & Payne 1982)

\[
L_{\text{MHD}} = B^2 R^3 \Omega / 2 ,
\]  

where \( B, \Omega, \) and \( R \) are, respectively, the magnetic field at the base of the jet, the jet angular velocity, and the size scale of the region where outflowing matter is injected along the field lines.

3.2. The Magnetic Switch

If, in the above cases, the magnetic field is sufficiently strong or the rotation sufficiently fast, then the power output can exceed a critical dynamical value, given by the ratio of the escape energy to the free-fall time. For the case of a constant density corona of radius \( R \), this can be expressed as

\[
L_{\text{crit}} = \frac{E_{\text{esc}}}{\tau_{\text{ff}}} = m_{\text{ej}} \left( \frac{GM}{R} \right)^{3/2},
\]

where the escape energy is taken to be \( GMm_{\text{ej}}/R \). Simulations of magnetized accretion disk coronae with a height comparable to the radial extent (\( m_{\text{ej}} \sim 4 \pi R^2 \rho \); see Meier et al. 1997; Meier 1999) have shown that under these conditions much of the MHD power may be converted quickly into kinetic energy, creating a fast, highly collimated jet close to the injection region, and that this condition can per-
sist well into the steady state. In the current situation, the mass to be ejected that determine the escape energy is that in a cone of fractional solid angle $f_{\Delta t} = \Delta \Omega / 4\pi$, where $\Delta \Omega$ is the solid angle of the jets emerging from the poles of the neutron star. Assuming that the mass in this cone is concentrated near the lower portions where the density is $\rho$ at radius $R$, we take $m_{\Delta t} \sim 4\pi R^2 f_{\Delta t} \rho$ and hence

$$L_{\text{crit}} = \frac{E_{\text{esc}}}{\tau_{\text{ff}}} = \frac{4\pi}{3} \rho R^2 f_{\Delta t} \left( \frac{GM}{R} \right)^{3/2}. \tag{4}$$

For $L_{\text{MHD}} < L_{\text{crit}}$, the outflow will be a classical Blandford-Payne–type outflow, collimating slowly with the jet thrust spread over a broad region. (See, e.g., the steady state simulations of Krasnopolsky et al. 1999.) This criterion for a slow jet can be expressed in terms of a dimensionless parameter, $\nu$, as

$$\nu \equiv \left( \frac{L_{\text{MHD}}}{L_{\text{crit}}} \right)^{1/2} = \frac{v_A}{v_{\text{esc}}} f_{\Delta t}^{-1/2} \left( \frac{3\Omega}{K} \right)^{1/2} < 1. \tag{5}$$

For $\nu > 1$, the jet may have a much higher velocity (with $v_{\text{jet}} \sim v_A > v_{\text{esc}}$) early on, with the jet thrust concentrated in a very narrow axial spindle (Meier et al. 1997).

4. MAGNETOCENTRIFUGAL JETS IN SUPERNOVAE

We will now explore the application of the “spring” and “fling” and fast and slow jets to the supernova collapse problem and consider in more detail the three stages of the collapse that may lead to MHD jets as outlined in § 2.

4.1. The LeBlanc-Wilson Jet

The first type of MHD outflow we consider is that discovered by LeBlanc & Wilson (1970, hereafter LW), also treated later by Symbalisty (1984). LW studied the magnetorotational collapse of a 7 $M_\odot$ iron core by numerically solving the two-dimensional MHD equations coupled to the equation for neutrino transport. Their simulations showed that a “circulation vortex” formed along the rotation axis, at about the time of core bounce, amplified the magnetic field there to of order $10^{15}$ G, and created two oppositely directed, magnetically driven, high-density, supersonic jets of material emanating from inside the proto–neutron star. In the calculation of LW, the jet formed at about 50 km, in that of Symbalisty, at about 30 km. LW estimated that their jet carried away $\sim 10^{52}$ g with $\sim (1-2) \times 10^{51}$ ergs in $\sim 1$ s. Khokhlov et al. (1999) assumed characteristics of the LeBlanc-Wilson jet for their proof of principle calculations.

The calculation of LW was extended by Meier et al. (1976) to treat stellar cores with mass, central density, and rotational characteristics more commensurate with the results of modern advanced stellar evolutionary calculations. One of their main constraints on the LW mechanism was that it not overproduce $r$-process material in the Galaxy, a limit of roughly $2 \times 10^{-4} M_\odot$ of highly neutronized ejecta per supernova. Meier et al. pointed out that this constraint was not compatible with the notion that the LW mechanism provides the explosion energy for most supernovae: if it were to do so, then the $r$-process elements would be overproduced by a factor of $\sim 1000$. Therefore, either the LW jet provides the explosion energy (and most of the Galaxy’s $r$-process material) in only a very few rare supernovae, or LW jets or bubbles are quite common in most core-collapse supernovae, but eject only $\sim 10^{-4} M_\odot$ per event, and are not important in the supernova’s overall energetics because the progenitor did not have especially large initial rotation and magnetic field.

In the latter, more plausible case, Meier et al. (1976) showed that the LW events would proceed as “buoyancy instabilities,” with the magnetized material rising subsonically up the rotation axis of the proto–neutron star and bursting into the mantle above. In the modern language of MHD winds and jets, these LW events will result in Blandford-Payne–type MHD wind flows emanating from the upper interior of the proto–neutron star. These MHD “fountains” will be initially subsonic, relative to the core sound speed, and rather confined by the high pressure, but when they burst into the mantle, they may be supersonic with respect to the sound speed there. In a completely bare core collapse, with little or no mantle to disrupt these weak outflows, these events eventually could accelerate to supersonic speeds and form jets of $r$-process material, but in most supernovae the LW ejecta are likely to quickly mix with the polar regions of the infalling mantle material.

Taking the collapsing proto–neutron star to be about 1.5 $M_\odot$ in mass, with initial and final radii of $\sim 10^8$ cm and $\sim 5 \times 10^6$ cm, respectively, we estimate from Meier et al. (1976) that the magnetic field still will be rather high (few times $10^{15}$ G) in the circulation-vortex region near the axis, but that the differential rotation rate there will be rather low, less than 100 rad s$^{-1}$. From the $r$-process constraints and equation (5) we find that the magnetic switch parameter $\nu$ will be of order 0.03 or less (for $f_{\Delta t} \sim 1$), nearly independent of the exact values of the collapse parameter and the efficiency of conversion of rotational energy to magnetic energy. This result that the MHD outflow is well below the magnetic switch point is consistent with the LW fountain being an initially subsonic, Blandford-Payne–like, polar outflow into the mantle.

4.2. The Slow Toroidal Jet

4.2.1. Generation of the Toroidal Field

Rotation alone can inhibit the contraction of the newly formed neutron star (Symbalisty 1984; Mönchmeyer et al. 1991; Fryer & Heger 2000). However, the addition of magnetic fields provides a source of dissipation that will allow further contraction, thus enhancing the rotational energy, the differential rotation, and ultimately the field energy. The presence of even a small field, therefore, can significantly increase the possibility of creating strong axial jets. As discussed in the introduction, Wheeler et al. (2000) considered the linear wrapping phase of field amplification due to differential rotation within the neutron star (Meier et al. 1976; Kluzniak & Ruderman 1998; Ruderman, Tao, & Kluzniak 2000), but they did not emphasize the resulting toroidal field, only the dipole component of that field (taken to be $\sim 1$% of the toroidal field) and its possible effects in concert with the Alfvén surface and speed of light circle. Here we argue that the principle focus should, rather, be on the strong toroidal field itself. By analogy to the work on the creation of jets by MHD processes in accretion disks (§ 3), the toroidal field can act as a “spring” to launch jets. We expect this sort of jet to overtake and supplant the original LeBlanc-Wilson jet or bubble, as illustrated schematically in Figure 1.
The original rotational energy of the proto–neutron star (PNS) is approximately

\[ E_{\text{rot},\text{PNS}} \approx \frac{1}{2} I_{\text{PNS}} \Omega_{\text{PNS}}^2 \approx 9 \times 10^{50} \text{ ergs} \]

and the rotational energy after the deleptonization and contraction is about

\[ E_{\text{rot},\text{NS}} \approx \frac{1}{2} I_{\text{NS}} \Omega_{\text{NS}}^2 \approx 2 \times 10^{52} \text{ ergs} \]

It is not clear how much of the latter energy can be dissipated in the magnetic field, but there is substantial energy for a supernova for any final rotation period of less than about 7 ms. Strictly speaking only the energy associated with the differential rotation can be tapped, but the neutron star is always in strong differential rotation compared to the progenitor star. There may be constant torques on the neutron star until the explosion succeeds or fails, so this issue is complex and involves the magnetic fields we will now explore.

If the precollapse progenitor core has a field strength comparable to that of a magnetized white dwarf, \( \sim 10^8 \text{ G} \), then a field of \( \sim 10^{12} \text{ G} \) could arise during the collapse simply from flux-freezing. This field can be amplified further by differential rotation in the neutron star (Meier et al. 1976; Kluzniak & Ruderman 1998) to produce a strong toroidal field. For a collapsing white dwarf, the proto–neutron star is likely to be differentially rotating with angular velocity \( \Omega \) increasing outward since the proto–neutron star will be “stiffer” than the highly degenerate white dwarf progenitor (Ruderman et al. 2000). For a neutron star formed by homologous iron core collapse, the angular velocity of the homologous core will increase outward within the homologous core at core bounce, but after bounce the angular velocity is nearly constant in the inner core and decreases with radius beyond the boundary of the initial homologous core (S. Akiyama et al., in preparation). The process of field generation and expulsion will depend on the details of the differential rotation.

For this “wrapping” mechanism, the field will grow linearly with time and be proportional to the seed field. After \( n_0 \) revolutions of the neutron star, the initial seed (poloidal) field will produce a toroidal field,

\[ B_\phi \approx 2\pi n_0 B_p, \]

where \( B_p \) is the initial seed poloidal field. The field will be limited by buoyancy, which will operate to expel the field.
from the site where it is generated. The field at this buoyancy limit, $B_b$, is given by

$$\frac{B_b^2}{8\pi} \geq f_b \rho c_s^2,$$  \hspace{1cm} (9)

where $f_b$ is the fractional difference in density between the rising flux tube elements and the stellar material.

There are two regions where there is likely to be large shear and hence field growth. One is within the neutron star. The angular velocity can be large there, but, as remarked above, the gradient is rather flat after core bounce. A complication is that the dynamics of the LW phase may already have amplified the field within the neutron star, so the seed field for the subsequent wrapping phase may be larger. The other region where strong field amplification is expected is at the boundary of the proto–neutron star, where the angular velocity is less than in the neutron star, but the difference in the angular velocity with respect to the post-shock, settling matter results in a rather large gradient.

After bounce, the central density of the neutron star is about $2 \times 10^{14}$ g cm$^{-3}$. The boundary of the homologous core at $R \sim 10$ km is initially at about $10^{13}$ g cm$^{-3}$ at bounce, but settles to somewhat higher densities, $\sim 3 \times 10^{13}$ g cm$^{-3}$, after some relaxation (Burrows, Hayes, & Fryxell 1995; Janka & Müller 1996; Mezzacappa et al. 1998). For sound speed $c_s \sim 10^{10}$ cm s$^{-1}$, density $\rho \sim 10^{14}$ g cm$^{-3}$, and $f_b \sim 0.01$ the critical value of the field at the onset of buoyancy is

$$B_b \simeq 6 \times 10^{16} G f_{b-2}^{1/2} \rho_{14}^{1/4},$$  \hspace{1cm} (10)

where $f_{b-2} = f_b/0.01$ and $\rho_{14} = \rho/10^{14}$ g cm$^{-3}$. The number of revolutions to reach the buoyancy limit $B_b$ is

$$n_f \simeq \frac{B_b}{B_0} \frac{1}{2\pi} \simeq 10,000 B_{12}^{-1} f_{b-2}^{1/2} \rho_{14}^{1/4}.$$  \hspace{1cm} (11)

For $B_0 \sim 10^{12}$ G and $\rho \sim 10^{14}$ g cm$^{-3}$, the linear amplification timescale before the field is expelled by buoyancy is thus

$$t_f \simeq n_f P \sim 10 \ s \ P_{ms} B_{12}^{-1} f_{b-2}^{1/2} \rho_{14}^{1/4},$$  \hspace{1cm} (12)

where $P_{ms}$ is the rotational period in ms.

The second region, around the boundary of the proto–neutron star at $R \sim 50$ km, has a typical density $\sim 10^{11}$ g cm$^{-3}$ after the shock stalls at about 200 km. For sound speed $c_s \sim 3 \times 10^9$ cm s$^{-1}$, density $\rho \sim 10^{11}$ g cm$^{-3}$, and $f_b \sim 0.01$, the critical value of the field at the onset of buoyancy at the proto–neutron-star boundary is

$$B_b \simeq 5 \times 10^{14} G f_{b-2}^{1/2} \rho_{11}^{1/4},$$  \hspace{1cm} (13)

the number of revolutions to reach the buoyancy limit is,

$$n_f \simeq 80 B_{12}^{-1} f_{b-2}^{1/2} \rho_{11}^{1/4},$$  \hspace{1cm} (14)

and the linear amplification timescale before the field is expelled by buoyancy is

$$t_f \sim 0.08 \ s \ P_{ms} B_{12}^{-1} f_{b-2}^{1/2} \rho_{11}^{1/4}. $$  \hspace{1cm} (15)

Linear wrapping, therefore, is slow at first, but is likely to accelerate with the contraction and spin-up of the neutron star. Because there will be a gradient in density and sound speed from within the neutron star to the post-
can sustain a breakout buoyant field of $10^{16}$ G from deep within, or a field of $10^{13}$ G near the proto–neutron-star boundary, for $\gtrsim 0.1$ s as the differential rotational energy is tapped, a strong supernova could result.

For the conditions of interest here, the critical parameter of the magnetic switch mechanism (eq. [5]) can be written as

$$\nu \simeq 0.05 B_{15} (\rho_{15} f_{\text{M}})^{-1/2} R_{6}^{5/4} M_{\text{NS}}^{-3/4} \ .$$

where for convenience here and below, the mass of the neutron star has been normalized to $1.5 M_{\odot}$. This equation says that there is a locus of critical values at which the switch would be activated and an especially high-speed jet created, defined by $\nu = 1$, namely,

$$B_{\text{crit}} \simeq 2 \times 10^{16} G (\rho_{15} f_{\text{M}})^{1/2} R_{6}^{-5/4} M_{\text{NS}}^{1/4} \ .$$

The value $\nu = 1$ corresponds to a (critical) MHD power of (eq. [4])

$$L_{\text{crit}} \simeq 1.2 \times 10^{54} \text{ ergs s}^{-1} \rho_{15} f_{\text{M}} R_{6}^{1/2} M_{\text{NS}}^{3/2} \ .$$

The conditions to make this switch in jet character could be found in collapse conditions for especially rapid rotation and hence large field strength or if a strong field becomes buoyant and “floats” into lower density conditions, or as the inflow density declines. Note the potentially important role of the collimation of the jet as reflected in the parameter $f_{\text{M}}$. This parameter would be of order unity for a slow Blandford-Payne jet with a broad base and size of order the neutron star has been normalized to $1.5 M_{\odot}$. This tendency may also help to broaden the fan of the wind, but at larger radii where the opening angle of the slow toroidal jet, there will be relatively less neutronization will not be as extreme as it is deep within the proto-neutron star where the LW jet forms and where large toroidal fields will form in the subsequent wrapping, contraction phase. Nevertheless, the infalling matter will still be neutron rich (the electron fraction $Y_e$ plummets from $\lesssim 0.5$ at the standing shock to $\gtrsim 0.2$ for $R \gtrsim 50$ km; e.g., Mezzacappa et al. 1998), and excessive ejection would again violate limits on neutron-rich species. The broad fan of the Blandford-Payne MHD wind that is conducive to a robust explosion thus may also be counterproductive from a nucleosynthesis point of view. There are several factors that may mitigate this production of neutron-rich ejecta.

Hoop stresses and the large ambient pressure may restrict the opening angle of the initial jet. This would decrease the mass of neutron-rich matter directly ejected in the jet. We note that the jet should broaden as it passes though the standing shock at about 200 km due to the sudden decrease of the external pressure. The external pressure drops by about an order of magnitude at that point, whereas the pressure of the plasma confined within the jet remains the same. The opening angle of the jet might then expand by about the ratio of the sound speeds, or a factor of about 3. This may help to broaden the fan of the wind, but at larger radii where matter is somewhat less neutronized. On the other hand, a broad jet fan may encompass all the very neutron-rich matter originally expelled in an LW bubble leading to expulsion of all that neutron-rich matter even if the LW phase did not result in an explosion.

Another factor is that the jet is likely to go up the rotation axis, where the density will be relatively low owing to the effects of rotation. After bounce, rotating models give densities at 100 km about a factor of 5–10 lower on the rotation axis compared to the equator (e.g., Symbalisty 1984; Mönchmeyer et al. 1991). That means for a given opening angle of the slow toroidal jet, there will be relatively less matter to be intercepted by about the same factor. The attendant lower pressure will also tend to promote a broader fan, so this must be checked quantitatively. The rotation will also hinder convection in the equatorial direction, and this may promote neutrino emission up the rotation axis (Fryer & Heger 2000). This tendency may also help to decrease the density along the rotation axis, the direction of propagation of the jet.

We also note that the complementary effect is a higher density along the equator. For extreme values of rotation,
this matter will form an accretion torus with an accretion time long compared to the dynamical timescale. The large rotation and distorted geometry of this protopulsar phase will evolve through shedding of angular momentum to the jet and environment to the more familiar, more nearly spherical, pulsar phase.

Further out in the star, the rotation will cause less of an effect on the isodensity contours, but there will still be some distortion. In spherical core collapse supernovae models, iron is produced and ejected by the action of the shock on the silicon layers and the associated production of $^{56}\text{Ni}$ that subsequently decays to form $^{56}\text{Fe}$. After several seconds, this silicon will be part of the infalling matter. The jet will create bow shocks that converge on the equator and drive overlying matter out preferentially in the equatorial plane (Khokhlov et al. 1999). The bow shocks will thus tend to drive some of the silicon layer and the newly formed nickel outward along the equator. This effect would be enhanced by any rotational distortion. Some of the preexisting silicon may be directly ejected in this way and some silicon will be produced and ejected by shocking the oxygen layer. Silicon produced in the explosion may thus also be blown off in an equatorially preferential manner, but at somewhat lower velocities. These dynamics may partially account for the observation of high-velocity iron lying outside lower velocity silicon within the body of Cas A (Hughes et al. 2000; Hwang et al. 2000) and for high-velocity knots of silicon found beyond the reverse shock (Fesen & Gunderson 1996). The jets themselves should be formed of neutron-rich iron-peak material, but how much of this matter escapes the star will depend on the envelope structure and shock dynamics.

4.3. The Magnetic Switch, Failed Supernovae, and the Gamma-Ray Burst Jet

An important aspect of the dynamics of the jet propagation through a star is that the faster the jet, the less likely it is to trigger an explosion. According to the calculations of Khokhlov & Höflich (2001) and Höflich et al. (2001) (see also MacFadyen et al. 2001), faster jets have narrower opening angles, allowing less time for transverse shocks to affect the star before the jet emerges from the boundary of the core. The jet energy will tend to stay focused and directed along the jet trajectory rather than powering the bow shock and its associated transverse sheath of shocks that are critical for spreading the jet energy throughout the mantle and exploding the star. (The difference is akin to the result of aiming a baseball and a needle with the same energy at a loaf of bread.) This leads to a suggestion for a different ultimate evolution for those conditions that lead to especially fast jets, namely, a failed supernova, collapse to form a black hole, and the attendant possibility of creating a $\gamma$-ray burst. These ideas are summarized schematically in Figure 2.

We do not fully understand the conditions of either magnetic field strength or rotation of the progenitor core. It is plausible that these properties have a finite distribution such that a small fraction of progenitors have particularly large values of magnetic field and rotation. For a subset of all the core-collapse events, e.g., in a given mass range, some could have properties that satisfy the criterion that the MHD luminosity exceeds the critical luminosity of the magnetic switch, $L_{\text{MHD}} > L_{\text{crit}}$ (§ 4.2.2). In this case, we expect the generation of an especially fast, narrow, jet with $v_{\text{jet}} \sim v_{A}$, which may be well above $v_{\text{esc}}$. Such a jet may be less effective in creating a jet-induced supernova explosion. If this is the case, the mantle will continue to collapse and a black hole would form. The extreme version of this suggestion is that the narrow jet fails to form any supernova. It is plausible, however, that some matter is expelled to trigger a supernova, but an insufficient fraction of the matter is ejected to prevent the ultimate fallback of enough mass to form a black hole.

The delayed formation of a black hole would give a new ambience for the generation of a third, very possibly highly relativistic, jet, as widely discussed in the literature (MacFadyen & Woosley 1999; Aloy et al. 2000; MacFadyen 2001; Zhang & Woosley 2001). Whether this relativistic jet can induce a supernova explosion when the earlier, slower jet failed to do so is an open question, but as we remarked above, the first jet could have generated a supernova through partial ejection of the mantle and still have allowed the formation of a black hole. In this case, the highly relativistic jet would form within an already exploding configuration, an assumption made in some of the computations (e.g., Aloy et al.).

What does seem clear is that if a new, relativistic jet can propagate out of the infalling matter, it will collide with the iron-rich matter expelled in the earlier jet. For a related scenario, see Pugliese, Falcke, & Biermann (1999). Depending on the timescale for the fallback to trigger the collapse to form the black hole (MacFadyen et al. 2001), the initial jet could have propagated well away from the star. For a typical velocity of the fast, but mildly relativistic protopulsar jet of $v \sim 10^{10}$ cm s$^{-1}$, the matter will have reached a distance of order $D \sim 10^{12} f_{\text{BH,2}}$ cm, where $f_{\text{BH,2}}$ is the delay time between the initial collapse and the formation of the black hole and relativistic jet in units of 100 s. The material of the protopulsar jet may thus act as the “target” for the collision of the second jet and this collision could create the fireball and the internal shocks that represent the onset of the $\gamma$-ray burst. The interaction between the $\gamma$-ray burst jet and the target would occur at a distance $D(c/v) - v$, which is of order $D$, but could be significantly greater as $v \rightarrow c$ and the $\gamma$-ray burst jet takes longer to catch up to the fast protopulsar jet.

In the observer frame, the duration of the $\gamma$-ray burst would be associated with the time for the relativistic $\gamma$-ray burst jet to propagate along the length of the mildly relativistic jet, and this would be determined by the longer of two timescales: the duration of the nonrelativistic, protopulsar jet (since this determines the length of the iron-rich lobe that is targeted) and the duration of the $\gamma$-ray burst jet itself. The former timescale will be approximately the spin-down time of the jet-emitting protopulsar:

$$\tau_{\text{spindown}} \approx E_{\text{rot,NS}} / L_{\text{MHD}} \approx 8 \text{ s} B_{15}^{-2} P_{\text{mer}}^{-1} M_{\text{NS}}^{-1} R_{\text{NS,6}}^{-1} .$$

In this scenario, where the tripping of the magnetic switch triggers an especially fast, inefficient jet, the field must exceed the critical value (eq. [18]), so the spin-down time will be

$$\tau_{\text{spindown}} \lesssim 0.02 \text{ s} P_{\text{mer}}^{-2} P_{11}^{-1} f_{\text{Al}}^{-1} M_{\text{NS}}^{-1/2} R_{\text{NS,6}}^{3/2} .$$

This spin-down time can be as long as a typical $\gamma$-ray burst,
for rotation at less than breakup, tight collimation, and appropriate choice of density, so this could be the controlling factor in the duration of a $\gamma$-ray burst if this scenario of colliding jets proves to have merit.

Using a strong magnetic field as the prime carrier of momentum and energy in the second, relativistic jet, with the jet material just along for the ride, might explain how high Lorentz factors can be achieved and sustained in the black-hole–collapse jet. For example, for a “tenuous” corona around the accreting black hole of $10^6$ g cm$^{-3}$ in a magnetosphere of $10^{16}$ G, equations (20) and (21) give relativistic Alfvén/jet speeds of order

$$\Gamma_A \approx \left( \frac{B^2}{4\pi \rho} \epsilon \right)^{1/2} \approx 100 B_{16} \rho^{-1/2}_{-6}. \quad (24)$$

The corresponding power from the Blandford-Znajek mechanism from a rapidly rotating hole (Blandford & Znajek 1977; Meier 2001) is

$$L_{\text{GRB}} \approx 10^{52} \text{ ergs s}^{-1} \epsilon B_{16}^2 M_{\text{BH,0.5}}^2, \quad (25)$$

where $\epsilon$ is an efficiency factor and $M_{\text{BH,0.5}}$ is the black hole mass in units of $3 M_{\odot}$. If this luminosity is pumped into the relativistic jet for a time $\tau_{\text{rel}}$ s, the total energy emitted in the jet is

$$E_{\text{GRB}} \approx L_{\text{GRB}} \tau_{\text{rel}} \approx 10^{52} \text{ ergs} \epsilon B_{16}^2 M_{\text{BH,0.5}}^2 \tau_{\text{rel}}. \quad (26)$$

Recent analyses suggest that the typical energy emitted in $\gamma$-ray bursts is around $3 \times 10^{50}$ ergs (Panaitescu & Kumar 2001; Frail et al. 2001). If this low energy is typical, then very little of the expected black hole binding energy or rotation energy can be tapped to make $\gamma$-ray bursts. For the picture we sketch here and, indeed, for any picture based on the production of jets from black holes we must have $\epsilon B_{16}^2 \tau_{\text{rel}} \approx 0.03$. In our particular magnetic switch scenario, we must satisfy equation (18) for the critical switch field and equation (24) to produce a suitably relativistic jet. This suggests that either the efficiency for the Blandford-Znajek process is relatively low (perhaps because the field is swallowed by the black hole rather than producing electromagnetic flux) or the duration of the process is short. We note that if the latter is the case, then the duration of the first jet (eq. [23]) will control the duration of the $\gamma$-ray burst.

Using the slower jet as the target and tying the delayed interaction between the two jets to the mantle collapse time would explain how $\gamma$-rays are produced so efficiently and routinely far from the central black hole environment. In addition, the iron-rich nature of the first jet may give appropriate conditions in which to address the iron lines observed in some $\gamma$-ray burst events (e.g., Piro et al. 2000). The strength of the iron emission is difficult to assess. It depends on numerous very uncertain factors such as the mass of the iron ejected, the density in the precursor jet, the cross sections of the first and second jets, and the ambient ionizing flux. Estimates in the literature of the amount of iron

$\sim 30$ s, for rotation at less than breakup, tight collimation, and appropriate choice of density, so this could be the controlling factor in the duration of a $\gamma$-ray burst if this scenario of colliding jets proves to have merit.
required to produce a given line flux by recombination vary by orders of magnitude (e.g., Piro et al. 2000; Rees & Mészáros 2000; Mészáros & Rees 2001).

This scenario of jet collisions does not help to resolve the issue of the low afterglow densities associated with some γ-ray bursts (Panaitescu & Kumar 2001; Scalo & Wheeler 2002) that are difficult to achieve in any model based on a massive star that expels a substantial wind. It also offers no explanation as to why the blast energy associated with the γ-ray burst is confined to a narrow and rather small value, around $3 \times 10^{50}$ ergs. The question of why such a small fraction of the available energy would be tapped in the jet is an important problem for the future.

5. SUMMARY AND DISCUSSION

We have presented the case that strong toroidal magnetic fields will be generated in stellar collapse and that these magnetic fields can generate magnetocentrifugal jets in analogy to those found in simulations of black hole accretion. The case for magnetocentrifugal jets suggests they may be frequent and robust and hence provide a good basic understanding of why all core collapse supernovae are found to be substantially asymmetric and predominantly bipolar.

There are concerns that the jets we describe arise deep in the gravitational potential where the infalling matter is very neutron rich. We argue that this constraint rules out the prompt LeBlanc-Wilson jets as the common origin of supernovae. This suggests that conditions of initial rotation and magnetic field extreme enough to produce strong jets in this phase do not arise frequently in nature. We point out that successful jets from the later, protopulsar phase must walk the line between being broad enough to provide ample energy to the overlying matter without being so broad as to eject excessive neutron-rich material. Partial evacuation of the matter on the axis by rotation may alleviate this problem, but it remains a concern. The jets will produce bow shocks that tend to expel matter, including iron and silicon, in equatorial tori. This may help to account for observations of the element distribution in Cas A.

There also is some concern that gravitational radiation may remove so much angular momentum, and so quickly, that there would not be enough rotational energy to power a substantial MHD jet. We estimate that the effects of gravitational radiation (GR) will be significant, but not severe. The most important GR mode is probably the r-mode. Detailed computations by Lindblom, Tohline, & Vallisneri (2001) on the evolution and effects of this mode estimate that only about half of the rotational kinetic energy will be lost before the gravitational radiation process switches off. The timescale for this to occur may be fairly short (a few tens of initial rotational periods, i.e. as short as, say, 30 ms), but the proto-neutron star still will be left with a significant fraction of its initial rotational energy, that can be tapped for the production of MHD jets.

The topics of core-collapse supernovae and γ-ray bursts overlapped, at least in principle, with the discovery of SN 1998bw. Although the association is still controversial, this odd supernova is likely to have been connected to GRB 980425 (Galama et al. 1998). Supernova-like excesses of light have been detected in the afterglow of two γ-ray bursts, GRB 970228 (Reichart 1999; Galama et al. 2000) and GRB 980326 (Bloom et al. 1999) about 2 weeks after the γ-ray bursts. The excess light in the afterglow of both GRB 970228 and GRB 980326 has been modeled by the addition of light from an event like SN 1998bw. While there may be other explanations for these excesses, the connection with supernovae must be pursued. We suggest here a new way to make this connection. The conditions that apply in stellar collapse suggest that a magnetic “switch” mechanism found to apply in the black hole simulations may also apply in the collapse case with subsequent affect on the speed of the jet. If the switch turns on at low density and large magnetic field, an especially fast jet, with propagation speed of order the Alfvén speed, could be produced. This jet would tend to propagate rapidly through the star and deposit relatively little momentum. The result might be a supernova explosion, but with enough infall to produce a delayed black hole. A reprise of the physics outlined here, and that explicitly invoked to produce relativistic MHD jets from stellar mass and AGN black holes, could then produce a relativistic jet that catches up to the first jet after it has emerged from the star. The interaction of these two jets could plausibly be the origin of the internal shocks thought to produce γ-ray bursts and could explain the presence of iron lines in the afterglow.

In this paper we have discussed toroidal fields of order $10^{14}-10^{16}$ G arising routinely in the formation of a neutron star in a supernova. There is clearly an issue of the final effective dipole strength of the neutron stars left behind in the explosion. We must have a final dipole field distribution of the neutron stars consistent with those of observed pulsars, $\sim 10^{12}-10^{13}$ G. It seems plausible that the large fields generated in the collapse are dissipated in the helical jets. The very strong fields that may be generated in the depths of the neutron star will take longer to float out, but the time still could be short compared to the lifetime of the supernovae, never mind the pulsar. Stronger fields in the interior could also leave smaller surface dipole fields. Much of the field may be amplified and dissipated at the shearing boundary between the infalling matter and the neutron star. This layer and its field will almost surely substantially dissipate after the explosion. These issues should be given more careful consideration.

All of the issues raised here should be examined in the context of magnetars. However the large effective dipole fields arise in magnetars, there is a strong presumption that they are created in the formation of the neutron star. The formation of a magnetar may require special circumstances, or they may arise as the tail of the normal pulsar formation process. One possibility is that they represent that fraction of core-collapse events that, perhaps by accident of progenitor conditions, do attain an especially rapid differential rotation and amplify the field not just by field line wrapping, but by an $\alpha$-$\Omega$ dynamo (Duncan & Thompson 1992). This could lead to exponential field growth and total field strength of order $10^{17}$ G with dipole component of order $10^{15}$ G. This is approximately the regime of fast toroidal jets that we invoke to produce less ejecta and more infall to produce a black hole and perhaps a γ-ray burst. To reconcile these pictures, it may be that magnetars arise with somewhat higher than normal fields and rotation, but still in the slow toroidal jet phase with conditions where the magnetic switch is still “off.” Even more extreme conditions could lead to the flipping of the switch, producing an especially fast jet that would be less effective in driving the explosion and thus leaving a black hole. Clearly there is a great need to study the many aspects of this problem in numerical and quantitative detail.
Some models of SN 1998bw invoked especially large kinetic energies, in excess of $10^{52}$ ergs in spherically symmetric models, to account for the bright light curve and high velocities (Iwamoto et al. 1998; Woosley, Eastman, & Schmidt 1998), while others took note of the measured polarization to suggest that strongly asymmetric models could account for the observations with more "normal" energies (H"offlich et al. 1999). Large kinetic energy has also been attributed to several other supernovae (again based on spherically symmetric models), especially SN 1997cy (German et al. 2000) and SN 1997ef (Branch 2001; Iwamoto et al. 2000). Events like SN 1997cy, SN 1997ef, and SN 1998bw will help to sort out the physics of explosive events, whether such events are more closely related to "ordinary" supernovae or hypernovae, whether either of these classes leaves behind neutron stars as "ordinary" pulsars or highly magnetized magnetars or whether the remnant is a black hole, and whether any of these events are associated with classic cosmic $\gamma$-ray bursts, as suggested by the supernova-like brightening of the afterglow of GRB 970228 and GRB 980326. In terms of our magnetic switch mechanism, it may be that hypernovae are still in the switch "off" phase, so they have strong jets, but not so fast and narrow that the explosion is mitigated. In this sense, there may be a connection between the hypernovae and magnetars that, as argued above, may arise in the same regime of especially energetic protopulsar jets, but still with the switch "off". If it were not for the constraints of the supernova polarimetry, we might continue to study essentially spherically symmetric collapse models, as indeed, many people will and must. It may be that we are pessimistic and that neutrino asymmetries can account for the observed polarimetry. The fact that jets alone can account for supernova explosions and the observed asymmetries has been established in principle. We think nature is saying that jets must occur in routine supernovae. We also believe that the process of neutron star formation must inevitably involve the physics outlined here whether the resulting jets are weak or strong. Clearly, we need a more rigorous description of the process of field generation, buoyancy, emergence, and subsequent evolution.

We are grateful to Alexei Khokhlov, Peter H"offlich, Elaine Oran, and Almadena Ch"uchtelkanova for helping us to understand how jets work in stars, to L"u"un Wang for showing that core collapse supernovae are asymmetric, to Rob Duncan for discussions of magnetars, and to Ellen Zweibel for discussion of MHD and plasma physics. We would especially like to thank the Aspen Center for Physics for the hospitality during the gestation of this collaboration. This research was supported in part by NSF grant 95-28110, NASA grant NAG 5-2888, and a grant from the Texas Advanced Research Program. Some of this research was carried out at the Jet Propulsion Laboratory, California Institute of Technology, under contract to the National Aeronautics and Space Administration.

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Branch, D. 2001, in Supernovae and Gamma-Ray Bursts, as suggested by the supernova-like brightening of the afterglow of GRB 970228 and GRB 980326. In terms of our magnetic switch mechanism, it may be that hypernovae are still in the switch "off" phase, so they have strong jets, but not so fast and narrow that the explosion is mitigated. In this sense, there may be a connection between the hypernovae and magnetars that, as argued above, may arise in the same regime of especially energetic protopulsar jets, but still with the switch "off". If it were not for the constraints of the supernova polarimetry, we might continue to study essentially spherically symmetric collapse models, as indeed, many people will and must. It may be that we are pessimistic and that neutrino asymmetries can account for the observed polarimetry. The fact that jets alone can account for supernova explosions and the observed asymmetries has been established in principle. We think nature is saying that jets must occur in routine supernovae. We also believe that the process of neutron star formation must inevitably involve the physics outlined here whether the resulting jets are weak or strong. Clearly, we need a more rigorous description of the process of field generation, buoyancy, emergence, and subsequent evolution.

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