A Two-Step Model for Gamma-Ray Bursts Associated with Supernovae

K. S. Cheng\textsuperscript{1} and Z. G. Dai\textsuperscript{2}

\textsuperscript{1}Department of Physics, University of Hong Kong, Hong Kong, China
\textsuperscript{2}Department of Astronomy, Nanjing University, Nanjing 210093, China

ABSTRACT

We here propose a two-step model for gamma-ray bursts (GRBs) associated with supernovae. In the first step, the core collapse of a star with mass $\geq 19M_\odot$ leads to a massive neutron star and a normal supernova, and subsequently hypercritical accretion of the neutron star from the supernova ejecta may give rise to a jet through neutrino annihilation and/or Poynting flux along the stellar rotation axis. However, because of too much surrounding matter, this jet rapidly enters a nonrelativistic phase and evolves to a large bubble. In the second step, the neutron star promptly implodes to a rapidly rotating black hole surrounded by a torus when the mass of the star increases to the maximum mass and meanwhile its rotation frequency increases to the upper limit due to the accreted angular momentum. The gravitational binding energy of the torus may be dissipated by a magnetized relativistic wind, which may then be absorbed by the supernova ejecta, thus producing an energetic hypernova. The rotational energy of the black hole may be extracted by the Blandford-Znajek’s mechanism, leading to another jet. This jet is relatively free of baryons and thus may be accelerated to an ultrarelativistic phase because the first jet has pushed out of its front matter and left a baryon-free exit. We expect that the second jet generates a GRB and its afterglow. Our two-step model may alleviate the baryon-contamination problem suffered possibly from in the hypernova models. Furthermore, this model not only accounts for association of several GRBs with supernovae but also explains well the features of the afterglows of these bursts.

\textit{Subject headings:} gamma-ray: bursts – supernovae: general – stars: neutron – black holes: physics
1. INTRODUCTION

It has been widely believed that gamma-ray burst (GRB) events do indeed occur at cosmological distances (Metzger et al. 1997; Kulkarni et al. 1998a, 1999; Andersen et al. 1999), which implies that a successful model for progenitors of cosmological GRBs must satisfy two essential requirements: (1) The model should produce an extremely relativistic fireball, which should subsequently emit an amount of gamma-ray isotropic energy \( E_{\gamma} \sim 10^{51}-10^{54} \) ergs implied by the observed fluences and the cosmological distance scale. The recent multi-wavelength observations of GRB afterglows support the so-called fireball shock model (Piran 1999). According to this scenario, GRBs are produced as a result of internal shocks when fast moving shells catch up with slower shells that were ejected at earlier times. The relative kinetic energy of motion of the shells is converted into observed gamma-ray emission in relativistic shocks via synchrotron radiation or inverse-Compton scattering mechanisms. The observed afterglow emission is produced when the shell decelerates as a result of interaction with the ambient matter. Based on the fireball shock model, several authors (Panaitescu, Spada & Mészáros 1999; Kumar 1999; Lazzati, Ghisellini & Celotti 1999) recently found that the efficiency for producing gamma-rays in internal shocks is a few percent. Therefore, the isotropic energy of fireballs in some bursts (e.g., GRB 990123) must be up to a few times \( 10^{55} \) ergs (Kulkarni et al. 1999). If anisotropic emission with a beaming factor of \( \Delta \Omega / 4 \pi \sim 0.01 \) is assumed, this energy can be reduced to \( E_{\text{jet}} \sim \) a few times \( 10^{53} \) ergs. (2) The rapid variability of GRBs and their nonthermal spectra (Woods & Loeb 1995) and the low radiative efficiency in internal shocks (Lazzati et al. 1999) requires that the Lorentz factor of the fireball be 100-1000. This implies that the fraction of contaminating baryons must be less than 1%. If the emission is anisotropic, the mass of loading matter \( \Delta M \leq 0.01 E_{\text{jet}}/c^2 \sim 10^{-3} M_\odot \).

Two currently popular models of GRB progenitors are the merger of two compact objects (neutron stars and black holes) and the collapse of massive stars. The former model is the plausibly baryon-clean one, but compact objects would be expected to have such a significant space velocity that their merger would take place outside their birthplaces (Paczynski 1998; Bloom, Sigurdsson & Pols 1999). The observational evidence for the association of several GRBs with star forming regions then provided weak evidence against the compact object merger as GRB progenitors and favored massive star progenitors. The
population synthesis of Fryer, Woosley & Hartmann (1999) supported this conclusion.

The massive star progenitor model has become more favorable since the discovery of the Type Ic supernova SN 1998bw in the error box of GRB 980425 (Galama et al. 1998). The high energy inferred for the optical supernova, \((2 - 3) \times 10^{52}\) ergs (Iwamoto et al. 1998; Woosley, Eastman & Schmidt 1999), and the high expansion velocity inferred for the radio supernova (Kulkarni et al. 1998b) strengthen the GRB-SN connection. Recently, this connection was further confirmed by Bloom et al. (1999), Reichart (1999) and Galama et al. (1999), who found the dramatical brightening and extreme reddening of the optical afterglows of GRB 980326 and GRB 970228 at late times, respectively. However, it is quite difficult to understand such a connection. From the above second essential requirement, too many baryons possibly exist in the vicinity of the collapsing core in the hypernova models (Woosley 1993; Paczyński 1998) so that an ultrarelativistic jet forming during the collapse of the core rapidly become nonrelativistic. This conclusion is consistent with the numerical studies (MacFadyen & Woosley 1999): an ultrarelativistic jet converts to a nonrelativistic large bubble. A simple reason for this may be that a large amount of radiative energy is released impulsively but the mass of contaminating baryons is in fact of the order of \(M_\ast(\Delta \Omega/4 \pi) \sim 0.1M_\odot(M_\ast/10M_\odot)(10^2\Delta \Omega/4 \pi)\), where \(M_\ast\) is the mass of the matter around the core.

In this Letter, we propose a scenario, in which a supernova explosion first may produce a massive neutron star, and about two hours later the star will start to accrete the fall-back supernova ejecta at a hypercritical rate, while an accretion disk will form near the stellar surface due to its large angular momentum. In particular neutrino annihilation and/or Poynting flux along the rotation axis of the star may lead to a jet which will pushes out of its front matter. The mass of the accreting neutron star will eventually reach the maximum mass about several hours after the supernova explosion and thus will promptly collapse to a rapidly rotating black hole surrounded by a torus. The gravitational binding energy of the torus, which is of the order of several \(10^{52}\) ergs, may be dissipated into the supernova ejecta, which may in turn give rise to an energetic hypernova. Another jet with energy of a few times \(10^{53}\) ergs and with low-mass baryon contamination will occur along the rotation axis of the black hole by extracting its rotational energy via the Blandford-Znajek’s (1977) mechanism. We expect that the second jet can produce a GRB and its
afterglow.

Delayed formation of black holes in supernovae has been widely discussed (Woosley 1988; Chevalier 1989; Brown & Weingartner 1994; Brown & Bethe 1994; Woosley & Weaver 1995; Fryer 1999). Recently, it has been shown numerically that the accretion of such a black hole may lead to a relativistic jet required by a typical GRB (Woosley, MacFadyen & Heger 1999; MacFadyen, Woosley & Heger 1999). The present model is an analytical one, in which we suggest that hypercritical accretion of a newborn neutron star and delayed formation of a black hole could produce two discrete jets. The first jet will push out of its front matter and leave an exit for the second jet, which could thus be relatively free of baryons.

2. HYPERCRITICAL ACCRETION OF A NEWBORN NEUTRON STAR AND FORMATION OF A BUBBLE

It is well known that the core collapse of massive stars with 10-25$M_\odot$ produces neutron stars accompanying Type II supernovae. Timmes et al. (1996) numerically studied the initial mass function of newborn neutron stars and found that their initial mass distribution is bimodal with peaks at 1.27 and 1.76$M_\odot$. The principal reason for this bimodal distribution is the difference in the presupernova structure of stars above and below 19$M_\odot$, the mass separating stars that burn carbon convectively from those that produce less carbon and burn radiatively. Here we consider neutron stars with 1.76$M_\odot$ as the starting point of our work. When such a massive neutron star first forms in a supernova explosion, it is surrounded by a dense gas (supernova ejecta), some of which falls onto the neutron star and cools by neutrino emission (Colgate 1971; Zeldovich, Ivanova & Nadëzhin 1971). It is the neutrino emission that allows accretion of the star at a high rate. From simple analytical arguments, Chevalier (1989) and Brown & Weingartner (1994) estimated a lower limit to steady neutron star accretion with neutrino losses assuming spherical symmetry. However, the accreted matter may have a large angular momentum which leads to an accretion disk. In this case, the lower limit with neutrino losses is estimated as $\dot{M}_{\text{tr}} \sim 1.1 \times 10^{-3}M_\odot\text{yr}^{-1}$ (Chevalier 1996).

The supernova explosion scenarios involve an outgoing shock wave. When this shock enters the hydrogen envelope, the deceleration of matter occurs. This deceleration sharp-
ens into a reverse shock. When the reverse shock reaches the neutron star surface, the star starts to accrete the fall-back supernova ejecta at a hypercritical rate. The time scale for the reverse shock to pass through the core is about the same as the time that the outgoing shock front takes to reach the stellar surface, which is $t_0 \sim 2\,\text{hr}(R_*/3 \times 10^{12}\,\text{cm})(M_{ej}/10M_\odot)^{1/2}(E_{sn}/10^{51}\text{ergs})^{-1/2}$, where $R_*$ is the presupernova stellar radius, $M_{ej}$ is the mass of the ejected matter and $E_{sn}$ is the supernova explosive energy (Shigeyama, Nomoto & Hashimoto 1988). For SN1987A, $t_0$ is about 2 hr (Shigeyama et al. 1988). In order to derive the rate of mass accretion onto the neutron star after $t_0$, we follow Brown & Weingartner (1994). By assuming the neutron star to be at rest with respect to its ambient matter and using the Bondi’s (1955) spherical accretion theory, we obtain the accretion rate

$$
\dot{M} = 5.63 \times 10^{-4} \left( \frac{M}{1.8M_\odot} \right)^2 \left( \frac{v_f}{10^8\,\text{cm}\,\text{s}^{-1}} \right)^{-15/8} \left( \frac{t}{1\,\text{yr}} \right)^{-15/8} \dot{M}_\odot\,\text{yr}^{-1},
$$

where $M$ is the neutron star mass ($\sim 1.8M_\odot$) and $v_f$ is the final velocity (after being slowed down by the reverse shock) of the carbon-oxygen core. For SN1987A, $v_f \sim 6 \times 10^7\,\text{cm}\,\text{s}^{-1}$ (Woosley 1988). In deriving equation (1), we have assumed that the ambient matter of the neutron star is radiation-dominated due to the effect of the outgoing shock. Owing to this assumption, our accretion rate is slightly larger than that of Brown & Weingartner (1994). The time at which radiation significantly affects accretion can be estimated based on $\dot{M}(t_{cr}) = \dot{M}_{cr}$. Hence, we obtain this timescale $t_{cr} = 0.7v_{f,8}^{-1}\,\text{yr}$, where $v_{f,8} = v_f/10^8\,\text{cm}\,\text{s}^{-1}$. Below, we take $v_{f,8} = 1$. When $t < t_{cr}$, the accretion is hypercritical and the total accreted baryon mass is given by

$$
\Delta M_{acc} = \int_{t_0}^{t} \dot{M} \, dt = 1.81M_\odot v_{f,8}^{-15/8}[(t_0/1\text{hr})^{-7/8} - (t/1\text{hr})^{-7/8}] .
$$

For the modern realistic equation of state for neutron matter chosen in the next section, an accreting neutron star with initial mass of $1.76M_\odot$ will collapse to a rapidly rotating black hole when $\Delta M_{acc} = 0.55M_\odot$.

Before the collapse, the accreted matter forms a disk near the neutron star because the accreted angular momentum may be up to $3^{1/2}R_s c$ per gram, where $R_s$ is the Schwarzchild radius (Woosley & Chevalier 1989). The temperature of the accretion disk near the
neutron star can be estimated by the following equation: \( \eta \dot{M} c^2 = \dot{\varepsilon}_\nu \Delta \Omega_d R_{ns}^3 / 4 \), where \( \eta \) is the efficiency for the conversion of the gravitational energy to heat (\( \sim 0.1 \)), \( R_{ns} \) is the neutron star radius (\( \sim 10^6 \) cm), \( \Delta \Omega_d \) is the solid angle of the disk, and \( \dot{\varepsilon}_\nu = 1.0 \times 10^{25} (T/\text{MeV})^9 \text{erg s}^{-1} \text{cm}^{-3} \) is the neutrino pair energy production rate per unit volume (Dicus 1972). For typical values of these parameters (e.g., \( \Delta \Omega_d \sim 3 \)), we obtain \( T \sim 6 \) MeV. The total energy for neutrino losses is approximated by \( E_\nu = G M \Delta M_{\text{acc}} / R_{ns} \sim 3 \times 10^{53} \) ergs. Since anisotropic neutrino emission takes place due to the effect of the disk (Kluzniak 1998), neutrino annihilation along the rotation axis of the neutron star leads to a jet. By using the efficiency of neutrino annihilation, \( \chi \sim 0.3\% \) (Goodman, Dar & Nussinov 1987; Kluzniak 1998), we obtain the energy of the jet

\[
E_{\text{jet,1}} \sim E_\nu \chi \sim 10^{54} \text{ergs}. \tag{3}
\]

It should be pointed out that this efficiency is much larger than that of Popham, Woosley & Fryer (1999), in which an accretion disk surrounding a black hole is advection-dominated. In the present case, however, matter is accreted onto the neutron star surface where the conversion of the gravitational energy to heat (\( \eta \)) is much more efficient. Since this thermal energy is released via neutrino emission, the efficiency of neutrino annihilation should increase substantially. In addition to the neutrino annihilation mechanism, there is another possible mechanism to produce a jet proposed by Katz (1997), in which the magnetic field amplified by the differential rotation of the disk may result in a strong Poynting flux. The jet produced by these mechanisms will push its front baryonic matter whose velocity is given by \( v_{\text{jet}} = (2E_{\text{jet,1}} / \Delta M')^{1/2} \sim 3 \times 10^9 (M_*/10 M_\odot)^{-1/2} (10^2 \Delta \Omega / 4 \pi)^{-1/2} \text{cm s}^{-1} \), where \( \Delta M' = M_*(\Delta \Omega / 4 \pi) \) is the mass of baryons loading with the jet. It should be noted that this velocity is much larger than that of the outgoing shock. As numerically studied by MacFadyen & Woosley (1999), this nonrelativistic jet will expand laterally, producing a bubble. When \( t \sim 5 \) hr, at which the collapse of the neutron star may take place as argued in the next section, the radius of the bubble is estimated as \( R_b \sim v_{\text{jet}}(t-t_0) \sim 3 \times 10^{13} (M_*/10 M_\odot)^{-1/2} (10^2 \Delta \Omega / 4 \pi)^{-1/2} \text{cm} \), which is much larger than the presupernova stellar radius \( R_* \).

3. COLLAPSE OF THE NEUTRON STAR AND GENERATION OF A GRB
It has been argued by Brown & Bethe (1994) and Bethe & Brown (1995) that hypercritical accretion of neutron stars in supernovae (e.g., SN1987A) may lead to collapse of the stars to low-mass black holes. These authors considered a soft equation of state like kaon condensation for neutron matter as the starting point of their discussions. Here we adopt a more realistic modern equation of state named UV14+TNI (Wiringa, Fiks & Fabrocini 1988). This equation of state gives the property of a nonrotating neutron star at the maximum mass: maximum gravitational mass $M_{\text{max}} = 1.84M_\odot$, corresponding baryon mass $M_b = 2.17M_\odot$, radius $R_{\text{ns}} = 9.51$ km and moment of initia $I = 1.5 \times 10^{45}$ g cm$^2$ (Wiringa et al. 1988). According to this property, the maximum frequency of rotation for this neutron star equation of state is given by $\Omega_{\text{max}} = 7.84 \times 10^3 (M_{\text{max}}/M_\odot)^{1/2} (R_{\text{ns}}/10 \text{ km})^{-3/2} \text{s}^{-1} = 1.17 \times 10^4 \text{s}^{-1}$ (Cook, Shapiro & Teukolsky 1994). The property of a rotating neutron star at this maximum frequency and at the allowable maximum mass is as follows (Cook et al. 1994): maximum gravitational mass $M_{\text{max}} = 2.19M_\odot$, corresponding baryon mass $M_b = 2.55M_\odot$, radius $R_{\text{ns}} = 12.7$ km and angular momentum $J = 2.85 \times 10^{49}$ erg s or the Kerr rotation parameter $a = cJ/GM_{\text{max}}^2 = 0.67$.

After having these properties, we discuss implications of hypercritical accretion of a newborn neutron star with gravitational mass of 1.76$M_\odot$. First, hypercritical accretion may produce and maintain a large bubble, as argued in the last section.

Second, hypercritically accreted matter may rapidly submerge the magnetic field of the neutron star (Muslimov & Page 1995). Once the accreted mass reaches 0.01$M_\odot$, the buried magnetic field ohmically diffuses out after $\sim 10^8$ yr (Geppert, Page & Zannias 1999). This implies that the magnetic field of the neutron star could be always weak in the accretion timescale.

Third, hypercritical accretion may spin up the neutron star (Woosley & Chevalier 1989). A 1.76$M_\odot$ neutron star has the baryon mass of 2.0$M_\odot$ and thus this star needs to accrete matter with mass of 0.55$M_\odot$ to become a maximum rotating neutron star with maximum mass. As suggested by Woosley & Chevalier (1989), the accreted angular momentum from the mixed mantle and helium core of the ejecta may be as large as $3^{1/2} R_s c$ per gram corresponding to $5.54 \times 10^{48}$ erg s per 0.1$M_\odot$ accreted mass. Thus, after accreting 0.55$M_\odot$ mass, the neutron star can obtain the angular momentum of $3 \times 10^{49}$ erg s. But,
the maximum angular momentum of a rapidly rotating neutron star at maximum mass is only \(2.85 \times 10^{49}\) erg s. The remaining angular momentum is \(\Delta J = 0.15 \times 10^{49}\) erg s. How is \(\Delta J\) dissipated? Fortunately, this remaining accreted angular momentum can be easily carried away by gravitational radiation (Wagoner 1984). Owing to hypercritical accretion, therefore, the mass of the neutron star not only reaches the maximum mass but it is spun up to the maximum rotation frequency.

Finally, once the mass of the accreted matter reaches \(0.55M_\odot\), the neutron star will promptly collapse to a rapidly rotating black hole. From equation (2), we can estimate the time at which the collapse will occur: \(t \sim 5.1\) hr. The resulting black hole has the rotational energy

\[
E_{\text{rot}} = f(a)M_{\text{BH}}c^2 \approx 3 \times 10^{53}\text{ ergs},
\]

where \(f(a) = 1 - \sqrt{(1 + \sqrt{1 - a^2})/2} \approx 0.067\) at \(a \approx 0.67\), and the mass of the black hole has been taken to be \(2M_\odot\). After the collapse, not all mass will be immediately accreted; the outmost layers with a small fraction (a few percent) of the total mass, in fact, have centrifugal accelerations beyond the local gravitational attraction (Vietri & Stella 1998), leading to a torus. In addition, a small amount of accreted ejecta still stagnate in the torus. Thus, the total mass of the torus may be \(M_t \geq 0.1M_\odot\). The presence of the torus will give rise to the following two effects: (i) The binding energy of the torus is \(E_b = GM_{\text{BH}}M_t/R_t = 3 \times 10^{52}\) ergs \((M_t/0.1M_\odot)(R_t/20\text{ km})^{-1}\), where \(R_t\) is the typical radius of the torus. Because the magnetic field of the torus may be amplified to \(10^{15}\) G due to differential rotation in the torus, this binding energy can be dissipated into a magnetized relativistic wind (Usov 1994; Mészáros & Rees 1997; Katz 1997). Because such a wind may be easily absorbed by the outgoing supernova ejecta as argued by Dai & Lu (1998a, 1998b), almost all of the binding energy may convert to the expansion energy of the ejecta, which may further give rise to a supernova with a much brighter optical luminosity. (ii) In the presence of the torus, the rotational energy of the hole can be extracted by the Blandford-Znajek (1977) mechanism (Mészáros, Rees & Wijers 1998). The power for this mechanism is \(P_{\text{BZ}} = 1.7 \times 10^{51} a^2 f(a)(M_{\text{BH}}/M_\odot)^2(B/3 \times 10^{15}\text{G})^2\text{erg s}^{-1} = 2.5 \times 10^{50}(B/3 \times 10^{15}\text{G})^2\text{erg s}^{-1}\) (Lee, Wijers & Brown 1999). The rotational energy of the black hole will be dissipated in the timescale of \(\sim 10^3\) s if the magnetic field strength is of the order of \(3 \times 10^{15}\) G.
We believe that this mechanism will produce an ultrarelativistic jet because the mass of matter loading with this jet can be estimated by

$$\Delta M = \Delta M'(\Delta \Omega/4\pi) = 10^{-3}M_\odot(M_*/10M_\odot)(10^2\Delta \Omega/4\pi)^2.$$  \hfill (5)

In fact, since the angular momentum of the presupernova star may be rather large, the centrifugal force should reduce the density of matter along the rotation axis and thus the above estimate is an upper limit. After a short acceleration phase, the Lorentz factor of the jet should be

$$\Gamma = \frac{E_{\text{rot}}}{(\Delta Mc^2)} \geq 150(M_*/10M_\odot)^{-1}(10^2\Delta \Omega/4\pi)^{-2},$$  \hfill (6)

which is consistent with the constraints from the rapid variability of the light curves and the nonthermal spectra of GRBs (Woods & Loeb 1995). The collisions among the shells in such a jet will produce a GRB and subsequently the deceleration of the jet in its ambient medium will result in an afterglow.

4. DISCUSSIONS

We have proposed a two-step model for GRBs associated with supernovae. A hypernova, a much brighter supernova, may occur in our model. However, this model is clearly different from the current hypernova models (Woosley 1993; Paczyński 1998; MacFadyen & Woosley 1999). In the latter models, the core of a massive star directly collapses to a black hole. Recently, Cen (1998) and Wang & Wheeler (1998) proposed models for GRB-SN association, in which the matter above the neutrinosphere in a small cone around some special axis of a newborn neutron star is assumed to be preferentially first blown out of the deep gravitational potential well of the star in order to avoid too many baryons contaminating a subsequently resulting jet. Our model may provide a plausible way of how such an empty cone is produced: neutrinos from the hypercritical accretion disk annihilate to electron/positron pairs which form the first jet to push its front baryons and leave an exit for the second jet. Therefore, our model may avelliate the baryon contamination problem.
In our model, GRBs are naturally associated with supernovae because the former phenomenon takes place several hours after the latter phenomenon. This is consistent with the analysis of Iwamoto et al. (1998), in which the time of core collapse coincides with that of SN1998bw within (+0.7, −2.0) days. The unusually large explosive energy of SN1998bw, \((2 − 3) \times 10^{52}\) ergs (Iwamoto et al. 1998; Woosley et al. 1999), is very close to the gravitational binding energy between the black hole and torus, \(\sim 3 \times 10^{52}\) ergs. This consistency seems to support our model.

The features of some afterglows can also be explained in the context of our model. For example, the optical afterglow of GRB 980326 rapidly decayed as \(\propto t^{-2.0\pm0.1}\) in the first two days, subsequently brightened dramatically and reddened significantly, and finally declined (Bloom et al. 1999). These features were also seen in the light curve of the afterglow of GRB 970228 (Reichart 1999; Galama et al. 1999). The late-time afterglows were widely believed to be the contribution of unusually brighter supernovae, while the rapid decay of the early afterglows was analytically argued to be due to sideways expansion of jets (Rhoads 1999; Sari, Piran & Halpern 1999). Another interpretation for the rapidly decaying afterglows may be that a relativistic shock expanding in a dense medium has evolved to a nonrelativistic phase (Dai & Lu 1999a, 1999b). In our two-step model, this dense medium may be a presupernova steller wind.

In summary, our two-step model may alleviate the baryon-contamination problem suffered possibly from in the hypernova models. This model not only accounts for association of several GRBs with supernovae but also explains well the features of the afterglows of these bursts.

We would like to thank J. S. Bloom, B. Hansen, A. MacFadyen and S. E. Woosley for their comments and discussions. This work was supported by a RGC grant of Hong Kong government and the National Natural Science Foundation of China.

**REFERENCES**

Andersen, M. I. et al. 1999, Science, 283, 2075
Bethe, H. A., & Brown, G. E. 1995, ApJ, 445, L129
Blandford, R. D., & Znajek, R. L. 1977, MNRAS, 179, 433
Bloom, J., Sigurdsson, S., & Pols, O. 1999, MNRAS, 305, 763
Bloom, J. et al. 1999, Nature, 401, 453
Brown, G. E., & Weingartner, J. C. 1994, ApJ, 436, 843
Brown, G. E., & Bethe, H. A. 1994, ApJ, 423, 659
Bondi, H. 1952, MNRAS, 112, 195
Cen, R. 1998, ApJ, 507, L131
Chevalier, R. A. 1989, ApJ, 346, 847
Chevalier, R. A. 1996, ApJ, 459, 322
Colgate, S. A. 1971, ApJ, 163, 221
Cook, G. B., Shapiro, S. L., & Teukolsky, S. A. 1994, ApJ, 424, 823
Dai, Z. G., & Lu, T. 1998a, Phys. Rev. Lett., 81, 4301
Dai, Z. G., & Lu, T. 1998b, A&A, 333, L87
Dai, Z. G., & Lu, T. 1999a, ApJ, 519, L155
Dai, Z. G., & Lu, T. 1999b, astro-ph/9906109
Dicus, D. A. 1972, Phys. Rev. D, 6, 941
Fryer, C. 1999, ApJ, 522, 413
Fryer, C., Woosley, S. E., & Hartmann, D. H. 1999, ApJ, in press astro-ph/9904122
Galama, T. J. et al. 1998, Nature, 395, 670
Galama, T. J. et al. 1999, ApJ, submitted astro-ph/9907264
Geppert, U., Page, D., & Zannias, T. A&A, 345, 847
Goodman, J., Dar, A., & Nussinov, S. 1987, ApJ, 314, L7
Iwamoto, K. et al. 1998, Nature, 395, 672
Lazzati, D., Ghisellini, G., & Celotti, MNRAS, in press astro-ph/9907070
Lee, H. K., Wijers, R. A. M. J., & Brown, G. E. 1999, astro-ph/9905373
Katz, J. I. 1997, ApJ, 490, 633
Kluzniak, W. 1998, ApJ, 508, L29
Kulkarni, S. R. et al. 1998a, Nature, 395, 663
Kulkarni, S. R. et al. 1998b, Nature, 393, 35
Kulkarni, S. R. et al. 1999, Nature, 398, 389
Kumer, P. 1999, ApJ, in press astro-ph/9907090
MacFadyen, A., & Woosley, S. E. 1999, ApJ, 524, 262
MacFadyen, A., Woosley, S. E., & Heger, A. 1999, ApJ, submitted (astro-ph/9910034)
Mészáros, P., & Rees, M. J. 1997, ApJ, 482, L29
Mészáros, P., Rees, M. J., & Wijers, R. A. M. J. 1998, ApJ, 499, 301
Metzger, M. et al. 1997, Nature, 387, 878
Muslimov, A. & Page, D. 1995, ApJ, 440, L77
Paczyński, B. 1998, ApJ, 494, L45
Panaitescu, A., Spada, M., & Mészáros, P. 1999, astro-ph/9905026
Piran, T. 1999, Phys. Rep., 314, 575
Popham, R., Woosley, S. E., & Fryer, C. 1999, ApJ, 518, 356
Reichart, D. E. 1999, ApJ, 521, L111
Rhoads, J. 1999, ApJ, in press (astro-ph/9903399)
Sari, R., Piran, T., & Halpern, J. P. 1999, ApJ, 519, L17
Shigeyama, T. Nomoto, K., & Hashimoto, M. 1988, A&A, 196, 141
Timmes, F. X., Woosley, S. E., & Weaver, T. A. 1996, ApJ, 457, 834
Usov, V. V. 1994, MNRAS, 267, 1035
Vietri, M., & Stella, L. 1998, ApJ, 507, L45
Wagoner, R. V. 1984, ApJ, 278, 345
Wang, L., & Wheeler, J. C. 1998, ApJ, 504, L87
Wiringa, R. B., Fiks, V., & Fabrocini, A. 1988, Phys. Rev. C, 38, 1010
Woods, E., & Loeb, A. 1995, ApJ, 453, 583
Woosley, S. E. 1988, ApJ, 330, 218
Woosley, S. E. 1993, ApJ, 405, 273
Woosley, S. E., & Chevalier, R. A. 1989, Nature, 338, 321
Woosley, S. E., Eastman, R. G., & Schmidt, B. P. 1999, ApJ, 516, 788
Woosley, S. E., MacFadyen, A., & Heger, A. 1999, astro-ph/9909034
Woosley, S. E., & Weaver, T. A. 1995, ApJS, 101, 181
Zel'dovich, Ya. B., Ivanova, L. N., & Nadézhin, D. K. 1972, Soviet Astron., 16, 209