Jets, winds and bursts from coalescing neutron stars

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ABSTRACT
Recent high-resolution calculations of neutron star coalescences are used to investigate whether $\nu\bar{\nu}$ annihilation can provide sufficient energy to power gamma-ray bursts, especially those belonging to the short duration category. Late time slices of the simulations, where the neutrino emission has reached a maximum, stationary level are analyzed to address this question. We find that $\nu\bar{\nu}$ annihilation can provide the necessary driving stresses that lead to relativistic jet expansion. Maximum Lorentz factors of the order of 15 are found along the binary rotation axis, but larger values are expected to arise from higher numerical resolution. Yet the accompanying neutrino-driven wind must be absent from the axis when the burst occurs or prohibitive baryon loading may occur. We argue that even under the most favorable conditions, $\nu\bar{\nu}$ annihilation is unlikely to power a burst in excess of $\sim 10^{48}$ erg. Unless the emission is beamed into less than one percent of the solid angle, which we argue is improbable if it is collimated by gas pressure, this may fail to satisfy the apparent isotropic energies inferred at cosmological distances. This mechanism may nonetheless drive a precursor fireball, thereby evacuating a cavity into which a later magnetic driven jet could expand. A large range of time delays between the merger and the black hole formation are to be expected. If the magnetic driven jet occurs after the black hole has formed, a time span as long as weeks could pass between the neutrino powered precursor and the magnetic driven GRB.

Key words: dense matter; hydrodynamics; neutrinos; gamma rays: bursts; stars: neutron; methods: numerical

1 INTRODUCTION
Double neutron star (NS) binaries, such as the famous PSR1913+16, will eventually coalesce due to angular momentum and energy losses to gravitational radiation, provided that their orbital separation is small enough. When a NS binary coalesces, the rapidly-spinning merged system could be too massive to form a single NS; on the other hand, the total angular momentum is probably too large to be swallowed immediately by a black hole (BH). The expected outcome would then be a spinning BH, orbited by a torus of NS debris, whose accretion can release sufficient gravitational energy, $\approx 10^{52}$ erg, to power a gamma-ray burst (GRB; Lattimer & Schramm 1976; Paczyński 1986, 1991; Eichler at al. 1989; Narayan, Paczyński & Piran 1992; Mochkovitch et al. 1993; Kluźniak & Lee 1998; Rees 1999; Rosswog et al. 1999; Ruffert & Janka 1999; Salmonson, Wilson & Mathews 2001; Rosswog & Davies 2002). The energy released in the accretion process is expected to be transformed into a fireball, which may provide the driving stress necessary for relativistic expansion (see Piran 1999 and Mészáros 2002 for recent reviews). Possible forms of this outflow include the following. Relativistic particles generated by $\nu\bar{\nu}$ annihilation – the thermal energy released as neutrinos is reconverted, via collisions, into $e^\pm$ pairs and photons. Electromagnetic Poynting flux – as in pulsars, strong magnetic fields anchored in the dense matter convert the rotational energy of the system into a Poynting-dominated outflow. In either case, the duration of the burst is determined by the viscous timescale of the accreting gas, which is substantially longer than the dynamical or orbital timescale, thus providing a simple interpretation for the large difference between the durations of bursts and their fast variability.

These binary encounters are thought to transpire outside star forming regions (Fryer, Woosley & Hartmann 1999) and are not currently thought to be appropriate for long GRBs ($\geq 20$ s), whose afterglows have been predominantly localised within the optical image of the host galaxy (Bloom, Kulkarni & Djorgovski 2001). However, these calculations are uncertain because they are sensitive to a number of parameters that are not well understood (e.g. distribution of initial separations; see Belczynski et al. 2002). Such models may also be in difficulties producing adequate duration, collimation, and energy (Katz & Canel 1996; Ruffert et al. 1997; Popham, Woosley & Fryer 1999; Ruffert & Janka 1999;
Lee & Ramirez-Ruiz (2002), but each of these problems is model-dependent and might be cured by different assumptions regarding the poorly known input physics such as disk viscosity or the jet production mechanism. In this Letter, we study the neutrino emission produced in the shocked and dissipatively heated merger remnant of a NS binary encounter. Late time slices of three dimensional (3D), high resolution calculations of the last stages of the coalescence are used (Rosswog & Davies 2002; hereafter RD) to investigate the potential of $\nu\bar{\nu}$ annihilation as a viable source for GRB production, at least for bursts belonging to the short category. Neutrinos could give rise to a relativistic $e^\pm$ pair dominated outflow if they annihilate in a region of low enough baryon density. To that effect, high resolution simulations seem to be required, since even a very small amount of baryons polluting the outflow could severely limit the attainable Lorentz factor. Information regarding how much $\nu\bar{\nu}$ energy and momentum has been injected, its confinement and collimation are at the forefront of our attention. We address all three issues here, along with the type of predictions that would help us to discriminate between this mode of energy extraction and that of strong magnetic fields anchored in the dense matter.

2 NUMERICAL METHODS

The 3D hydrodynamics simulations of the mergers of binary NS were performed by RD using a smoothed particle hydrodynamics method with up to $\sim 10^6$ particles. A realistic equation of state for hot, dense nuclear matter has been used, which is based on the tables provided by (Shen et al. 1998a,b) and smoothly extended to the low density regime with a gas consisting of neutrons, alpha particles, photons and $e^\pm$ pairs (see RD for details). This allows one to closely follow the thermodynamic evolution of the neutron-rich debris after the merger. Under the conditions encountered in the debris – temperatures of several MeV and densities around $10^{12}$ g cm$^{-3}$ – neutrinos are emitted copiously. Their effect on the cooling and the changes in the composition of the material are taken into account via a detailed multi-flavour neutrino treatment which is described in detail in Rosswog et al. (2002). The results presented here are based on the close analysis of late time segments, typically a time $t_{\text{sim}} \approx 15$ ms after the start of the merger simulations, where the neutrino luminosities have reached their maximum, stationary level (Rosswog et al. 2002). Since we are interested in the maximum possible effect, we assume the system to maintain this level of emission for as long as 1 s (see §3). We describe results from three representative runs. First, an initially corotating system with twice 1.4 $M_\odot$ (run D from RD, referred to in the following as c1.4), which yields the lowest temperatures of all runs; it can therefore be considered as a lower limit on the neutrino luminosities. Second, a system with twice 1.4 $M_\odot$ and no initial NS spins (“irrotational”; run C from RD, hereafter i1.4), which we regard as the generic case and finally, an extreme system of twice 2.0 $M_\odot$ and no initial spins (run E from RD, hereafter i2.0), which we consider to be an upper limit for the neutrino luminosity.

3 $e^\pm$ FIREBALLS PRODUCED BY $\nu\bar{\nu}$ ANNIHILATION

The $\nu\bar{\nu} \rightarrow e^+e^-$ process (which scales with the square of the neutrino luminosity) can tap the thermal energy of the hot debris disk. For this mechanism to be efficient, the neutrinos must both escape before being advected and not being produced too gradually. Typical estimates for coalescing NS suggest a limit of $\lesssim 10^{51}$ erg for the neutrino energy dumped into $e^\pm$-pairs (Ruffert et al. 1997; Ruffert & Janka 1999; Popham et al. 1999). If the $e^\pm$-dominated plasma were collimated into a solid angle $\Omega_{\nu\bar{\nu}}$, then of course the apparent “isotropized” energy would be larger by a factor of $4\pi/\Omega_{\nu\bar{\nu}}$.

The presence of a region of very low density along the rotation axis, away from the equatorial plane of the debris must be present when the burst occurs or otherwise prohibitive baryon loading will occur – for instance an energy of $10^{51}$ erg deposited into $e^\pm$ pairs could not accelerate an outflow to $\Gamma > 100$ if it had to drag more than $E/(\Gamma c^2) \approx 5 \cdot 10^{-4} M_\odot$ baryons with it. In Figure 1 we plot the ratio of $e^\pm$ energy deposition to baryon rest mass energy, $\eta = Q_{\nu\bar{\nu}}/(\rho c^2)$, in the region above the poles of the merged remnant. $\rho$ denotes the matter density and, for simplicity, an energy injection time, $\tau_{\text{ini}} = 1$ s, well within the distribution of the short-duration bursts has been assumed. $\eta$ can be understood as an indicative of the terminal bulk Lorentz factor, $\Gamma$. The above $\eta$ estimates assume, also for simplicity, that the energy deposition rate by $\nu\bar{\nu}$ annihilation into $e^\pm$ pairs at $t_{\text{sim}}$ is both representative of the subsequent phases and steady for at least one second (this is clearly an upper limit to the total injected energy over the assumed period $\tau_{\text{ini}}$). Due to the finite resolution of the simulations the densities along the rotation axis are determined by particles located in the inner parts of the disk. We therefore expect to overestimate the densities along the rotation axis. Given the above assumptions, higher numerical resolution would yield higher values for $\eta$. This is supported by calculating these densities for two equivalent runs of different numerical resolution (run A and B of RD). The densities in well-resolved case are lower by approximately one order of magnitude.

The resulting constraints on the injected $e^\pm$ energy as a function of $\eta$ are illustrated in Figure 2. It shows the fast
A drop of the contained energy as a function of the maximally reachable Lorentz factor. The copiously emitted electron neutrinos and anti-neutrinos dominate the annihilation process by \( \sim 95\% \), their mean energies \( \langle E \rangle \) are \( \sim 10\) MeV for \( \nu_e \) and \( \sim 15\) MeV \( \bar{\nu}_e \). The highest annihilation rates occur in the uninteresting (since extremely baryon-rich) hot inner disk regions. Even in the most optimistic case \((i2.0)\) only \( \sim 10^{48} \left( \frac{4\pi}{\Omega_{\text{ch}}} \right) \) erg are found able to blow out material near the axis with a Lorentz factor larger than 10 (indeed no SPH-particles are enclosed in this region, but we are limited in our density calculation by the finite resolution). A broad spread in the Lorentz factor will therefore be present in the outflow – close to the rotation axis \( \Gamma \) may be high (i.e. \( \geq 10 \)); at large angles away from the axis an increasing degree of entrainment will correspond to a decrease in \( \Gamma \) (see Fig. 1). Even if the outflow is not narrowly collimated, some beaming is expected because the disk geometry (see Fig. 15 RD) channels the energy preferentially along the rotation axis. But unless \( \Omega_{\text{ch}} \) is \( \leq 10^{-1} \) – \( 10^{-2} \) this may fail to satisfy the apparent isotropized energies of \( \sim 10^{54} \) erg implied for “short-hard” bursts at a redshift \( z = 1 \) (Panaitescu, Kumar & Narayan 2001; Lazzati, Ramirez-Ruiz & Ghisellini 2001).

3.1 Beaming and Confinement

When a large amount of energy is suddenly released into a compact region, an opaque “fireball” is created due to the prolific creation of \( e^{\pm} \) pairs (Cavallo & Rees 1978; Goodman 1986; Paczyński 1986; Shemi & Piran 1990). A diagram similar to Figure 2 can be drawn to illustrate the volume that contains material above a given \( \eta \) (see Fig. 3). Consider a homogeneous fireball of energy \( E_{\text{ch}} \), total mass \( M_0 \) initially confined to a compact region, whose volume is given by \( V_{\text{ch}} \).

Clearly, since the optical depth \( \tau > 1 \), the initial fireball will be an opaque sphere in thermal equilibrium, characterised by a single temperature:

\[
T_{\text{ch}} \approx 10^{9/4} \frac{E_{\text{ch}}}{1.4 \times 10^{48} / \eta} \frac{V_{\text{ch}}}{1.8^{-1/4}} \text{ MeV},
\]

where we adopt the convention \( Q = 10^7 \) \( Q_{b} \), using cgs units. This radiative sphere expands and cools rapidly until the energy of the photons degrades below the \( e^{\pm} \) production threshold. When some amount of baryonic matter is mixed with the fireball and the radiation energy dominates the evolution (i.e. \( \eta > 1 \)), the fluid expands under its own pressure such that its Lorentz factor grows initially linearly with radius, \( \Gamma \propto r \), and reaches a final value of \( \Gamma \approx \eta \) (Piran 1999).

If the flow is assumed to be stationary and isentropic, then the fluid velocity will increase as the external pressure of the surrounding medium decreases. When the pressure has halved the flow becomes transonic and the cross-sectional area is minimised. In this way a directed de Laval nozzle can be established. Unlike the static situation envisaged by Blandford & Rees (1974), the \( e^{\pm} \) injected in this plasma do already have a relativistic outward motion and the flow is then unlikely to be recollimated in a fluid nozzle. Although it will become free and supersonic (indeed relativistic, if it was so initially) since the pressure drop at \( r_{\text{ch}} \) is very steep. In the simulations, the external pressure is found to vary with distance along the rotation axis as \( p_{\text{ext}} \propto r^{3/4} \). If the pressure drops faster than \( r^{-2} \) onwards, sound waves will eventually be unable to cross the jet, which will become overpressured with respect to the external medium, and therefore expand out over an angle \( \Gamma^{-1} \). If the free expansion starts just outside \( r_{\text{ch}} \), where \( \Gamma \sim 2 \), then it will spread over a wide angle (i.e. \( \theta_{e,\Gamma} \sim 30^\circ \)) and develop into a roughly semi-spherical blast wave. The fireball will then accelerate until the entire energy is converted into kinetic energy at \( r_{\text{ch}} \). If \( \eta > \eta_b \approx 150E_{\text{ch}}^{3/2} / \theta_{e,\Gamma}^{3/2} \), the fireball continues to be radiation dominated when it becomes optically thin and most of the energy escapes as photons. An observer will detect them with a characteristic thermal peak frequency of \( T_{\text{ch}} \approx T_{\text{ch}} \) (Goodman 1986), which is likely to be outside the BATSE [20-600] keV spectral window (see equation 1). Alternatively, if \( \eta < \eta_b \) the fireball becomes matter dominated, and most of the initial energy is converted into bulk kinetic energy with \( \Gamma \approx \eta < \eta_b \). The inertia of the swept-up external matter will then decelerate these ejecta significantly by the time it reaches \( \Gamma_1 \approx 10^{16} E_{\text{ch}}^{2/3} \eta^{-1/3} \theta_{e,\Gamma}^{-2/3} \eta_b^{-1/3} \text{ cm} \) (Mészáros & Rees 1997). For an approximately smooth distribution of external matter, the bulk Lorentz factor of the fireball decreases as an inverse power of the time (\( \propto t^{-3/8} \)).
As a consequence, the spectrum softens in time as the synchrotron peak corresponding to the minimum Lorentz factor and magnetic field decreases, leading to the possibility of multi-wavelength afterglow radiation.

3.2 Accompanying winds

The neutrinos that are emitted from the inner shock-heated regions of the torus will deposit part of their energy in the outer parts of the thick disk that has formed around the central object (see Fig. 15 of RD). The possibility of such a wind from a NS coalescence has already been realized by Ruffert et al. (1997). Using the typical numbers from our simulations, we find that the neutrinos will drive a mass outflow at a rate (Qian & Woosley 1996):

\[ \dot{M} \approx 2 \cdot 10^{-2} M_\odot / s L_{\nu,83}^{5/3} \left( \frac{\nu_e}{15 \text{MeV}} \right)^{\frac{4}{5}} \left( R_{\text{em}} \right)^{\frac{2}{5}} \left( \frac{2.5 M_\odot}{M_{\text{em}}} \right)^{\frac{2}{5}} \]

where \( L_{\nu,83} \) is the luminosity in \( \nu \)-neutrinos, \( \nu_e = \left( \frac{\nu}{15 \text{MeV}} \right) \). \( R_{\text{em}} \) and \( M_{\text{em}} \) are the radius of and the mass inside the neutrino emitting region. This estimate is consistent with those of other authors (Duncan, Shapiro & Wasserman 1986, Woosley 1993). This wind accelerates the blown-off material to its asymptotic velocities of \( \approx 10^9 \) cm s\(^{-1}\). A wind of this strength represents a possible danger for the emergence of a highly relativistic jet. The jet can only accelerate to high Lorentz factors if the wind is kept out of the funnel region along the rotation axis via centrifugal forces. However, from the bulged geometry (see Fig. 15 RD) and the assumption that the neutrinos will be emitted preferentially in the direction \( \dot{n} = -\nabla \rho / | -\nabla | \rho | \), we conclude that most of the wind will be still directed away from the jet.

4 DISCUSSION

4.1 \( \nu \bar{\nu} \) annihilation as a viable energy source?

We now want to address the viability of NS-NS binary coalescences as central engines of GRBs. The main form of energy release from the disk we considered is neutrino emission, which produces relativistic jets but fails to produce very energetic bursts. Due to our post-processing approach the interaction between the jet and the matter evolution/neutrino emission is neglected. Not accounting for the evolution of the jet during the energy injection period leads to an overestimation of the energy density in the jet. But even under the most favourable conditions (maximum baryon free volume allowed by the SPH-particle distribution, neglecting a neutrino driven wind that might pollute the funnel region above the poles, see §3, and ignoring the jet evolution) we find typical energies of only \( \approx 10^{48} \) erg, where we have assumed a steady injection lasting for as long as 1 s. Unless the \( e^\pm \) plasma was collimated into a solid angle \( \Omega_{\nu,9} \leq 10^{-2} \) (which is unlikely if the plasma is solely confined by gas pressure; but see Levinson & Eichler 2000 for an alternative mechanism for hydrodynamic collimation), \( \nu \bar{\nu} \) annihilation could only be responsible for a rather weak energy release (which is also likely to be very short, almost impulsive; see Lee & Ramirez-Ruiz 2002) and may therefore fail to produce isotropized energies of \( \approx 10^{51} \) ergs that are estimated for “short-hard” bursts if they are at \( z \sim 1 \). Even if this is the case, the \( e^\pm \) deposited by \( \nu \bar{\nu} \) annihilation may be responsible for either precursor emission (thermal and non-thermal) and/or for creating a cavity into which a magnetised jet can subsequently expand (see below). In this latter case, the afterglow, at least over a range of directions, would not arise until the late magnetised ejecta hits the wall of the cavity. One would naturally expect a wide variety of afterglow behaviours arising from the impact of (possibly anisotropic) ejecta on an irregularly shaped cavity.

It is, of course also possible to produce more energetic bursts via \( \nu \bar{\nu} \) annihilation if one is able to increase the temperatures in the appropriate density regimes (i.e. \( \approx 10^{11} \) to \( 10^{13} \) g cm\(^{-3}\)). One such possibility would be that the “real” equation of state is much softer than the one used here, leading to higher temperatures and hence more energetic neutrino emission. The neutrino luminosities are sensitive functions of the temperature of the emitting region, the \( \beta \)-reactions are \( \propto T^6 \) and the plasmon decay is \( \propto T^9 \). Higher temperatures would translate into higher neutrino energies \( E_{\nu} \), which in turn render the debris more opaque to the neutrinos (the cross-sections scaling \( \propto E_{\nu}^2 \)). We still, however, expect higher luminosities and hence higher annihilation energies if the temperature is increased. An alternative possibility may arise from general relativistic effects, which are not included in the current 3D hydrodynamics simulation. The tendency of general relativistic effects would be to compaction the matter distribution, and so higher temperatures would be naturally expected. However, in this case the annihilation process would be additionally complicated by effects such as bending of neutrino trajectories and redshifts of the neutrino energies. These effects have been found to partly cancel out one another so that they do not alter the results substantially. For a further discussion of this point we refer to the literature (Salmonson & Wilson 2001, Asano & Fukuyama 2000, Jaroszynski 1993). A third possibility could arise from a very high disk viscosity. In our current simulation no explicit, physical viscosity has been introduced. We measured an equivalent \( \alpha \)-viscosity of the order of \( \alpha_{\text{disk}} \approx 4 \cdot 10^{-3} \) (lower values are found in the better resolved central object; see RD). The reason for this low value is a largely improved artificial viscosity scheme and the high resolution. Future investigations including explicit, physical viscosity should clarify this point. The raised possibilities should certainly be kept in mind, but at the present stage they are a matter of speculation.

4.2 Delayed BH formation and MHD jets

An alternative way to tap the debris energy is via magnetic fields threading the torus. Even before a BH forms, the merging system might lead to winding up the fields and dissipation before the merger. While the above mechanism can tap the rotational energy available in the debris torus, a BH formed from the coalescence is guaranteed to be rapidly rotating, and being more massive, could contain a larger energy reservoir than the debris itself. A magnetic configuration capable of powering bursts requires fields of a few times \( 10^{15} \) G. A weaker field would extract inadequate power; however, it only takes of order of a second for simple winding to amplify a \( 10^{12} \) G field to \( 10^{15} \) G (Kluźniak & Ruderman 1998). If magnetic fields of comparable strength than those
anchored in the torus thread the BH, its rotational energy offers an extra (an possibly even greater) source of energy than in principle can be extracted via the B-Z mechanism (Blandford & Znajek 1977). For a maximally rotating BH, this is 0.29 $M_B c^2$ erg, multiplied by some efficiency factor. Even allowing for low total efficiency ($\sim 20\%$), a jet powered by the B-Z effect would not require any beaming to produce the equivalent of an isotropic energy of $10^{34.5}$ erg. The entrained baryonic mass would then only need to be below $10^{-4}M_\odot$ (which is clearly the case along the rotation axis of the merger remnant) to allow high relativistic expansion speeds.

The masses of the central objects of the merger remnant are $\sim 2.5 M_\odot$ and therefore above the value of $\sim 2.2 M_\odot$ that is typically found for cold, non-rotating NS in beta-equilibrium (see e.g. Akmal et al. 1998). Due to the poorly known physics under these extreme conditions (strong field gravity, magnetic fields, ‘exotic’ matter in the NS cores) the time scale on which the central object is going to collapse is, however, uncertain. For the generic NS merger case, 1.4, there are no signs of a collapse – at the end of the simulation, the maximum density of the merger remnant has not even reached the central density of a cold, non-rotating single 1.4 $M_\odot$ NS. This is mainly due to the stabilising effect of differential rotation. For a discussion of further stabilising effects (some of which are not included in the simulation) we refer to RD. The merger remnant would lose angular momentum on a magnetic dipole radiation time scale

$$\tau_c \sim 10^3 s \left(\frac{M_{\text{CO}}}{2.5 M_\odot}\right) \left(\frac{10^{16} G}{B}\right)^2 \left(\frac{15 \text{ km}}{R_{\text{CO}}}ight)^4 \left(\frac{3000}{s \cdot \omega}\right)^2,$$

which is substantially shorter than the viscous timescale, $10^9 s$ (Shapiro 2000). Due to the dependence of this timescale on the input quantities, we expect a large range of possible delays between the merger and the subsequent BH formation. This will be mainly determined by the details of the coalescence, unless exotic matter appears already for very low masses, therefore leading to an immediate collapse. Note that with reasonable assumptions even a time scale of weeks is still possible. More prominent time delays could also be expected for mergers of unequal mass NS, or NS with other compact companions. Such events might also lead to repeating episodes of accretion, or to the eventual explosion of the NS which has dropped below the critical mass, all of which would provide a longer lasting energy output. A GRB produced by MHD coupling is likely to go off inside a cavity inflated by the $\gamma$ annihilation burst (which may not produce a standard GRB; see §3). Such cavities can be as large as fractions of a parsec, giving rise to a magnetic deceleration shock months after the first $\gamma$ annihilation burst.

Much progress has been made in understanding how gamma-rays can arise in fireballs produced by brief events depositing large amounts of energy in a small region, and in deriving the generic properties of the corresponding afterglows. Nonetheless, there still remain a number of mysteries, especially concerning the progenitors of short-duration bursts. As argued above, a wide diversity of behaviours may be the rule, rather than the exception. And if the class of models that we have advocated here turns out to be irrelevant, the trigger of these events has to be even more remarkable and exciting.

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