MERGING NEUTRON STARS I:
INITIAL RESULTS FOR COALESCENCE OF NON-COROTATING SYSTEMS

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ABSTRACT

We present 3D Newtonian simulations of the coalescence of two neutron stars, using a Smoothed Particle Hydrodynamics (SPH) code. We begin the simulations with the two stars in a hard, circular binary, and have them spiral together as angular momentum is lost through gravitational radiation at the rate predicted by modeling the system as two point masses. We model the neutron stars as hard polytropes ($\gamma = 2.4$) of equal mass, and investigate the effect of the initial spin of the two stars on the coalescence. The process of coalescence, from initial contact to the formation of an axially symmetric object, takes only a few orbital periods. Some of the material from the two neutron stars is shed, forming a thick disk around the central, coalesced object. The mass of this disk is dependent on the initial neutron star spins; higher spin rates resulting in greater mass loss, and thus more massive disks. For spin rates that are most likely to be applicable to real systems, the central coalesced object has a mass of $2.4M_\odot$, which is tantalizingly close to the maximum mass allowed by any neutron star equation of state for an object that is supported in part by rotation. Using a realistic nuclear equation of state, we estimate the temperature of the material after the coalescence. We find that the central object is at a temperature of $\sim 10\text{MeV}$, whilst the disk is heated by shocks to a temperature of 2-4MeV.

1. INTRODUCTION

It has been suggested that $\gamma$-ray bursts arise at cosmological distances in the merger of binaries consisting of either two neutron stars, or a black hole and a neutron star (Eichler et. al. 1989, Piran, 1990, Paczyński, 1991, Piran et. al. 1992, Narayan, Paczyński & Piran 1992). Two different processes may provide the electromagnetic energy for the bursts: neutrino-antineutrino annihilation into electron-positron pairs during the merger, and magnetic flares generated by Parker instabilities in a postmerger differentially-rotating disk. Merging binary neutron stars are also a source of gravitational radiation, whose waveform and amplitude can be predicted with reasonable confidence (see e.g. Lincoln & Will, 1990; Cutler, et. al., 1993) and are prime candidates for detection by LIGO (Schutz, 1986, Thorne, 1987; Abramovici et. al. 1992). It has been estimated that $\sim 3$ such mergers will occur in a year within 200 Mpc (Narayan, Piran & Shemi 1991, Phinney 1991). These events should be observable by the Laser Interferometer Gravitational Wave Observatory (Abramovici et. al. 1992). A merger of two neutron stars may lead to the ejection of neutron-rich material that may be a good site for the production of r-process elements. It is possible that such mergers account for all heavy r-process material in the Galaxy (Eichler et. al. 1989).

The production of binaries containing two neutron stars has been discussed in the literature (see for examples, Narayan, Piran & Shemi 1991; van den Heuvel 1993). For completeness, the scenario is outlined here. Begin with two massive stars in a binary. The primary – the more massive of the two stars – will evolve first, forming a neutron star in a supernova. If the binary is tight enough, and the secondary sufficiently massive, the binary may not be disrupted by this violent event. Mass transfer between the secondary and neutron star will ensue as the former evolves beyond the main sequence, the star either filling its Roche lobe, or losing mass via a wind. At this stage in its evolution, the system will manifest itself as a massive X-ray binary. If the secondary were to evolve to a supernova without significant mass loss and if the explosion is symmetric, the binary
would be assured of disruption, as the mass lost from the exploding star would be much
greater than that remaining in the two neutron stars. However, we may speculate upon the
following panacea: because the bloated secondary is more massive than the neutron star,
mass transfer from the former will be unstable, if the mass transfer rate is large enough,
bringing the two stars closer together until they form a common envelope system. Such a
system represents an efficient way to remove the nefarious excess matter in the secondary’s
envelope. By finely tuning the separation of the neutron star and core of the secondary
at the beginning of the common envelope phase, we are able to speculate that one may
remove a large fraction of the gaseous envelope, leaving the core of the secondary and the
neutron star in a tight binary. When the core finally explodes, the relative mass loss from
the system may thus have been reduced sufficiently to avoid disruption. Alternatively,
the supernova explosion might be assymetric leaving the neutron stars in an eccentric
orbit. It is unclear how likely the binary is to remain intact. Estimates based on the
current known neutron stars binaries suggest that this happens in about one percent of
the cases (Narayan, Piran & Shemi, 1991). The fragility of such systems is evidenced by
the eccentricities of the four known neutron star binaries, suggesting these systems were
close to disruption when they were formed. It should be noted that systems formed in this
manner will have circularized before the two neutron stars come into contact (C. Cutler,
private communication).

Two neutron stars in a binary will spiral together as angular momentum is lost via
the emission of gravitational radiation. The timescale for merger of two neutron stars,
with $M_1 = M_2 = 1.4M_{\odot}$, is less than $10^{10}$ years if the initial separation satisfies $d \lesssim 5R_{\odot}$
(Iben and Tutukov 1984). Recent work (Kochanek 1992; Bildsten & Cutler, 1992; Lai,
Rasio & Shapiro, 1993) has shown that the viscosity of the neutron-star fluid is unlikely to
be sufficient to spin-up the stars to achieve tidal-locking as the two stars spiral together.
One therefore has to ask how the merger process will change as a function of neutron star
spin.

Hydrodynamic simulations of neutron-star mergers are a formidable task. Beyond
the 3D hydrodynamic calculations it should include general relativitistic effects, employ a
realistic equation of state, contain neutrino transport and neutrino cooling, and possible
nuclear reactions. So far there have been several attempts to address various aspects of
the problem. Most of these attempts focused on the gravitational radiation emission
(Oohara & Nakamura, 1989, 1990; Nakamura & Oohara 1989, 1991; Shibata, Nakamura
& Oohara, 1992). These were Newtonian calculations using a finite difference code, a
polytropic equation of state and including a gravitational radiation backreaction formula.
Rasio & Shapiro (1992) focused on the hydrodynamic evolution of the rotating core and
began with the two neutron stars corotating at contact. Here we present a study that
focuses on the thermodynamics and nuclear physics of the coalescence. As was mentioned
above, the low viscosity of the neutron-star material leads us to consider the merger of non-
corotating neutron stars. In order to achieve a realistic estimate of the radial velocity at
the moment that the neutron star come in to contact we begin with the two neutron stars
being far apart and follow the spiral-in phase. We follow the evolution until the neutron
stars merge and we examine both the matter distribution and the thermal properties of the
system after the two neutron stars merge. Our goal is to estimate the thermal conditions
in the coalesced core and the disk that forms around it after the two neutron stars merge,
and to study the possible subsequent nuclear processes and neutrino radiation processes. As a first stage we do not attempt to include relativistic effects, or even the complete equation of state in these calculations. At this stage, we perform Newtonian calculations with a polytropic equation of state, with a gravitational radiation backreaction formula. We employ, however, a realistic equation of state to estimate the thermodynamic conditions during the computations. We compare the conditions prevailing in our simulations to those required to produce a gamma-ray burst by neutrino-antineutrino annihilation (Eichler et al. 1989) or by accretion of the disk onto the compact core (Narayan, Paczyński & Piran 1992).

The numerical methods applied to this problem and the initial conditions used are described in section 2. The results of our simulations are presented in section 3 and in section 4 we discuss the various physical processes that could take place in this system. The implications of these results to the question of formation of gamma-ray bursts and to r-process nucleosynthesis are summarized in section 5.

2. NUMERICAL METHODS AND INITIAL CONDITIONS

We performed a series of simulations of the merger of two neutron stars of equal mass, using a 3D Smoothed Particle Hydrodynamics (SPH) code, based on one used previously to study collisions between stars (see for example, Benz & Hills 1992). Being a Lagrangian particle code, SPH is well suited to this problem. We have no need for a computational box (gravitational forces being computed using a tree), hence we do not waste computational resources simulating the evolution of the voids between the neutron stars as they spiral together.

Since we are mostly interested in the thermodynamic conditions of the material we are using a Newtonian code, rather than a general relativistic one. The only general relativistic effect which we include is the addition of the gravitational radiation backreaction force, which we discuss below. We modelled the cold neutron stars as hard polytropes. We employed a realistic nuclear equation of state developed by Lattimer & Swesty (1991), to compute pressure as a function of density for cold material at nuclear densities with a compressibility $K = 180$ MeV, and obtained $\gamma = 2.4$ for a polytropic approximation of the equation of state. We checked a posteriori that the Lattimer & Swesty equation of state is well approximated by a polytrope for the density range that appears in our calculations (except for a small amount of mass in the very low density region). In view of the large uncertainty in the equation of state of nuclear matter we believe that this is a reasonable first step. We later make use of the Lattimer & Swesty equation of state to estimate the temperature from the internal energy and density at the final stages of the computation. To construct the initial neutron-star model, we solve the Lane-Emden equation for $\gamma = 2.4$, and we construct a polytrope with the required density profile using the method described in Davies, Benz & Hills (1992).

We began each simulation with the two stars 5$R_{ns}$ apart and had them spiral in at the rate predicted by treating the system as two point masses and considering the energy lost through the emission of gravitational radiation. The change in energy, $E$, and angular
momentum, $J$, is given by (e.g. Shapiro & Teukolsky 1983)

\[
\frac{dE}{dt} = \alpha E, \quad \frac{dJ}{dt} = -\frac{\alpha}{2} J
\]

where $\alpha = \frac{64 G^3 M^2 \mu}{5 c^5 a^4}$

where $a$ is the binary separation, $M$ the total mass of the system, and $\mu = M_1 M_2 / (M_1 + M_2)$ is the reduced mass, where $M_1 = M_2 = (1/2)M$ is the mass of the two neutron stars. Using the above equations, we derive the following formulae for the acceleration due to the emission of gravitational radiation.

\[
\begin{align*}
\dot{x} &= \frac{-GM_1}{2r^3} x + \frac{\alpha}{2M_1 (\vec{v} \cdot \vec{r})} \left( Ex - \frac{Jy}{2} \right) \\
\dot{y} &= \frac{-GM_1}{2r^3} y + \frac{\alpha}{2M_1 (\vec{v} \cdot \vec{r})} \left( Ey + \frac{Jx}{2} \right)
\end{align*}
\]

As we apply the same acceleration to each SPH particle within each neutron star, the circulation of the fluid will be conserved, as the accelerating field has a zero curl. This in-spiral force was turned off once the two stars came into contact. At this stage it is clear that this approximation breaks down. While this force is a far cry from a sophisticated gravitational radiation emission backreaction force (Blanchet, Damour & Schafer 1990) it provides us with a simple formula that can be easily added to our Newtonian calculations which will provide a realistic radial velocity component at the moment when the two stars collide.

As the two stars spiral together, a lag angle will develop between the stars’ long-axes and the line joining their centers-of-mass. The size of this angle is partially a function of the viscosity of the neutron star fluid. As the fluid inside the neutron stars is extremely inviscid, we expect the figures of the two neutron stars to remain essentially aligned, until the very last stages of the in-spiral. However, one might be concerned that the much larger effective viscosity present in our fluid would lead to an artificially-large lag angle, with the two stars then receiving torques. Indeed, in early, lower-resolution runs, the stars were spun up close to the orbital spin rates. To circumvent this problem, we ran a second set of mergers, aligning the two neutron stars every time-step, so that their long-axes remained parallel. The shift in total angular momentum due to this procedure was measured and found to be tolerably small ($\sim 1$ in $10^5$). In fact, the two higher resolution runs presented here (runs C and D in Table I) that used this realignment mechanism show very similar results to the equivalent runs performed without the realignment (runs A and B). In other words, the problem which seemed to require the use of realignment largely disappeared when the resolution was increased in the final runs we performed. Both sets of runs are given here for completeness.

Each simulation was run for 3.5ms ($\sim 50$ dynamical times) after the stars came into contact. This was sufficient time for all mergers to produce a central coalesced object
surrounded by a disk. As we show later, our simplifying assumptions do not break down at this stage. However, the evolution of this system on longer time scales is determined by processes, which we did not include, but address in detail in section 4.

3. SPH POLYTROPIC RESULTS

We performed four high-resolution mergers (runs A – D) and two lower resolution simulations (runs E and F). In runs A – D, we used 4271 particles per neutron star, in runs E and F, we used 1256 particles per neutron star. In these six runs we keep all initial parameters the same except the initial spins of the two neutron stars. The spin angular velocities of the neutron stars when they came into contact is given in Table 1. For run A the lag angle \( \sim 4^\circ \) by the time the two stars come into contact. For run B, the value is somewhat larger at \( \sim 10^\circ \). Taking \( M_{\text{ns}} = 1.4M_\odot \) and \( R_{\text{ns}} = 10\text{km}, \Omega_{\text{spin}} = 1 \) (the synchronous spin with the two neutron stars 2\( R_{\text{ns}} \) apart) corresponds to a spin period of 0.46ms. Hence in runs A and C, the two neutron stars have spin periods \( \sim 2.5\text{ms} \). We consider this to be a reasonable upper limit to the neutron star spins for any merging system, as the stars are unlikely to be spun up to their orbital periods just before contact, owing to the low viscosity of the neutron star fluid. In runs B and D, we consider the case where the two neutron stars have negligible spin. We expect essentially all systems to lie in between these two cases. In runs E and F, we further investigate the effect of the neutron star spins, by considering cases where the stars are spinning in the opposite sense to the orbital spin.

3.1 Phenomenology of Mergers

In all cases we observe the appearance of a single rotating coalesced object surrounded by a disk. The size of the disk varies greatly from one run to another. In some cases additional long tidal tails form as well. Density contour plots for runs C and D are given in Figures 1 and 2. Logarithmic contours are plotted at intervals of 0.25 dex, beginning at a density of 0.001 in code units (1 code unit \( \equiv 2.786 \times 10^{15} \text{g/cm}^3 \)). The sequences shown begin just after the onset of coalescence, times being given in code units (1 time unit \( \equiv 73\mu\text{s} \)) with \( t = 0 \) when the two stars were 5\( R_{\text{ns}} \) apart. In both cases, a single coalesced object forms on the timescale of one orbit (i.e. \( \sim 1\text{ms} \)) after the objects collide. The system quickly becomes axisymmetric after the two stars come into contact, hence only a small amount of energy will be emitted as gravitational radiation at this stage, as has been calculated in earlier simulations by Shibata, Nakamura & Oohara (1992) and by Rasio & Shapiro (1992). Hence our approximation of switching off the gravitational radiation when the two stars come into contact has been vindicated.

In run C (where the neutron stars are rapidly rotating), far larger tidal tails develop behind the neutron stars. These tails are clearly visible by \( t = 306 \), and are less visible at the same time for run D. By a time, \( t = 314 \), some of the material removed from the neutron stars in the tidal tails has begun to form a disk around the central object. As we will see shortly, this material is shock-heated. By the end of the simulation, a large fraction of the mass removed from the two stars remains in the disk: only some of the material thrown off initially remains, unshocked, in the two tidal tails. We thus observe a difference between runs C and D: in the latter case, the lower neutron-star spin results in a delay in the formation of tidal tails, the material has thus been ejected less far when the two objects coalesce, and any material that has been removed from the two neutron stars
is located in the disk by the end of our simulation.

Figure 3 illustrates the distribution of SPH particles, in the plane of the original orbits, at the end of Run C. This plot illustrates the three components of the final system produced in our simulation: the central coalesced object (out to a radius \( \sim 2R_{\text{ns}} \)), a disk (out to a radius \( \sim 6R_{\text{ns}} \)), and two elongated tidal tails, which have expanded to \( \sim 20R_{\text{ns}} \) by the end of the simulation. One has to address different questions regarding each of these regions. We wish to know whether the energetics of the core or the disk are sufficient to fuel cosmological gamma-ray bursts. For this we need to examine on one hand the thermal energies and the temperatures of the central coalesced object and the disk as well as the mass of the disk, and on the other hand the opacity for neutrino emission in various regions. We are concerned with the stability and subsequent evolution of the central object and for this we are interested in the rotational and gravitational energies of the core. Finally we must also estimate how much material is ejected from the system. This question has some implication to the emission of gamma-ray bursts by these events and to the possibility that r-process nucleosynthesis is taking place in these sites.

In Figure 4, we plot density contours of the final configuration for run C in a plane perpendicular to that of the neutron star trajectories. Logarithmic contours are plotted, at intervals of 0.25 dex, starting at a density of 0.0001 in code units (recalling 1 code unit \( \equiv 2.786 \times 10^{15} \text{ g/cm}^3 \)). The maximal density decreases by about 10% as the two stars are tidally distorted and then subsequently increases by about 5% relative to the initial density as a central coalesced object forms. The final central density is \( \approx 5 \) times the nuclear density. This is only slightly higher than the central density of the original neutron stars. At these densities the polytropic index derived from the Lattimer & Swesty (1991) equation of state is practically unchanged. Figure 4 shows the form of the central, coalesced object and the disk of material surrounding it. It is clear from this plot that the disk is thick, almost toroidal; the material having expanded on heating through shocks. This disk surrounds a central object that is somewhat flattened due to its rapid rotation. It is apparent from this figure that an almost empty centrifugal funnel forms around the rotating axis and there is almost no material above the polar caps. The funnel is also apparent in Figure 5, where we produce a contour plot of the integrated column density, \( \int \rho dz \). This funnel might be important for the formation of gamma-ray bursts as it provides a region in which a baryon free radiation-electron-positron plasma could form. Such plasma is an essential step in any cosmological gamma-ray burst model.

3.2 Angular Velocity and Mass Distribution

For each run, we computed the ratio of system kinetic energy to the gravitational potential energy, \( T/|W| \), as a function of time. All runs show somewhat similar behavior, with the initially more spun-up neutron stars having slightly greater kinetic energy. In all runs, \( T/|W| \) increased from \( \sim 0.1 \) to \( \sim 0.17 \) between \( t = 290 \) and \( t = 300 \). As mass was shed from the neutron stars when they produced tails, \( T/|W| \) decreased, reaching a value of \( \sim 0.13 \) at the end of the simulations. For rigidly-rotating polytropes of \( \gamma \gtrsim 2.2 \), a secular instability occurs for \( T/|W| \approx 0.14 \). Under such an instability, a rotationally-symmetric ellipsoid (Maclaurin spheroid) will transform into a triaxial ellipsoid (Jacobi ellipsoid). However, such a transformation will occur on a relatively slow timescale (Press and Teukolsky 1973), \( i.e. \) many dynamical timescales. Hence, even though the polytrope considered in our simulations is sufficiently hard, and the system reaches a large
enough value of $T/|W|$, the swirling mass has resolved itself into a central coalesced object surrounded by a thick disk before the system has time to transform into a Jacobi ellipsoid. A discussion concerning the subsequent stability of the central coalesced object is given in section 4.

In Figure 6 we plot the spin angular velocities of the SPH particles as a function of cylindrical radius. For both runs C and D, again, we see evidence for the three regions produced in run C. The material in the central object ($r \lesssim 1.5R_{ns}$) rotates rigidly; an effect of the viscosity of our SPH fluid (which is much larger than that of the real neutron star fluid). In run D, the rotation curve is somewhat broadened in the same region. This is due to the lower spin rates of the material within the neutron stars before they coalesced.

It is unlikely that the real neutron star fluid will be as orderly within the coalesced object, as it is so inviscid. For $1.5R_{ns} \lesssim r \lesssim 6R_{ns}$ (i.e. the disk), $\Omega_{\text{spin}}$ decreases almost as $r^{-3/2}$; the slight flattening of the power law caused by the mass contained in the disk. For run C, the material located at greater radii is contained in the two tails. Note that, as discussed earlier, no tails remain in run D, and a mere few, stray particles, are found at such large distances from the center-of-mass. The material in the tails produced in run C closely follows a power law, $\Omega \propto r^{-\beta}$, where $\beta \approx 1.8$, i.e. the rotational velocity falls off faster than Keplerian. This may, initially, seem a rather enigmatic result, for which we furnish the following explanation. The essential point to realize is that the particles contained in the tails are travelling on eccentric paths. Imagine first the simple case where all these particles were ejected with the same angular velocities and at the same distance from the center-of-mass, but at different times. We would then have an ensemble of particles located at different positions on the same ellipse – or rather a set of identical ellipses with different orientations of the semi-major axes in the orbital plane. Kepler’s Second law tells us that $\Omega r^2 = \text{constant}$ for a given ellipse, hence if our simple model were correct, we would expect to see $\Omega \propto r^{-2}$. However, imagine now that the particles ejected first (and now located at the largest radii) were ejected with more angular momentum. Their “constant” would be therefore larger than that for particles ejected later, and thus now closer to the central object. This will have the effect of softening the power law, as is observed.

In Figure 7 we plot the enclosed mass (as a fraction of total mass) as a function of cylindrical radius for all runs. Again, we see the three components of the systems produced in Runs A and C: the central object contains $\sim 85\%$ of the system mass, the disk some $13\%$, with the remaining $2\%$ of the material being found in the elongated tidal tails at $r_{cyl} \gtrsim 6R_{ns}$. Runs B and D show somewhat similar profiles to those of A and C, except they leave no material in tidal tails, and the disk is slightly less massive, containing $\sim 10\%$ of the system mass. For the two cases where the neutron stars were spinning initially in the direction opposite to the orbital motion, $M_{\text{disk}} \lesssim 0.05M_{\text{system}}$. The dependence of this function on the initial neutron star spins is clear: higher spin rates (in the same sense as orbital spin) result in greater mass-loss, and more massive disks.

3.3 Temperatures and Thermal Energy

In Figure 8 we plot the energy per unit mass of each SPH particle against the density as the particle site for run C (both are in code units: units of energy/mass being $1.858 \times 10^{20}$ ergs g$^{-1}$). We computed isotherms using Lattimer and Swesty’s nuclear equation of state, assuming nuclear matter in beta equilibrium, and were thus able to estimate the temperatures at the SPH-particle locations. As stated earlier, this is an
approximate method, which is particularly sensitive for the high density regions in which the isotherms are very near each other, that is a small change in \( u \) produces a large change in \( T \). The isotherms derived from this method are shown in Figure 9. A lot of the material around the Z-axis in the final configuration has come from outer regions of the neutron stars, has thus been compressed considerably, and thus heated, increasing its \( T \). Typical temperatures in the central, coalesced object are around \( T \sim 10 - 15 \text{MeV} \). It should be noted that even if we are overestimating the temperature of the central coalesced object by a factor of two here, the conclusions drawn regarding the timescale for neutrino escape in section 4.1 are not radically altered. When material is tidally shredded from the two neutron stars, it initially decompresses adiabatically, then subsequently shock heats as the material from the two tidal tails come into contact, forming a disk. This shocked material now has densities in the range \( 3 \times 10^{11} \text{g cm}^{-3} \lesssim \rho \lesssim 3 \times 10^{13} \text{g cm}^{-3} \). Again the same polytropic index for the EOS applies as for higher densities, with the exception of the very outermost low density layers. In the simulation, the disk is heated by shocks to a temperature of 2-4 MeV, and contains \( \sim 0.25 \times 10^{53} \text{ergs of kinetic energy, and } 2 \times 10^{51} \text{ergs in thermal energy.} \)

For reasonable uncertainties in the equation of state, Lattimer and Yahil (1989) obtained a relation between gravitational mass and binding energy of non-rotating neutron stars, which reads

\[
E_{\text{bin}} = (1.5 \pm 0.15) \left( \frac{M_g}{M_\odot} \right)^2 \times 10^{53} \text{erg.} \tag{3}
\]

If we apply it to two 1.4 M\(_\odot\) neutron stars and a 2.4 M\(_\odot\) central object (assuming it is stable), we find a difference of roughly \( 2.76 \times 10^{53} \text{erg} \) in binding energy which would be expected to go into internal energy (heat), as long as the system did not loose energy by radiation yet. However, the system is not in its ground state, but in an excited state due to its rotation. The binding energy is decreased by the rotation energy. In the Newtonian limit we have \( E_{\text{rot}} = 0.5I\Omega^2 \), with \( I \) denoting the moment of inertia. In an order of magnitude estimate we make use of the moment of inertia for a homogeneous object \( (0.4MR^2 \text{ for a sphere and } 0.5MR^2 \text{ for a disk}) \). The central coalesced object is slightly elongated and we take a medium factor 0.45, a radius of 17km (cutting at \( \rho = 5 \times 10^{13} \text{g cm}^{-3} \), see Figure 4), and \( \Omega = 8.6 \times 10^3 \text{s}^{-1} \). This leads to \( E_{\text{rot}} = 2.32 \times 10^{53} \text{erg} \) or a reduced binding energy increase between two single neutron stars and the rotating coalesced object of \( 4.4 \times 10^{52} \text{erg} \) or \( 9.2 \times 10^{18} \text{erg g}^{-1} = 9.55 \text{MeV per nucleon, which should show up as additional thermal energy in the formed compact object and is only due to stronger gravitational binding. This agrees well with the temperature (internal energy) of the central object of 10-15 MeV (see Figure 9).} \)

This excercise also illustrates that the object is a relatively fast rotator (although not at break-up and rotationally supported), which leads to the fact that its central density is only increased by roughly 5% over that of the initial neutron stars (to about \( 1.7 < 10^{15} \text{g cm}^{-3} \)). In that regime the same polytropic index applies to the Lattimer and Swesty (1991) EOS as at lower densities and the object is still stable.
4. ADDITIONAL PHYSICAL PROCESSES

The previous sections discussed the hydrodynamical behavior during the first few milliseconds of the merger of two neutron stars, resulting in one central object with densities up to \(10^{15}\) g cm\(^{-3}\) and a disk, containing about 0.2-0.4M\(_\odot\) of matter with densities of \(3 \times 10^{11} - 1 \times 10^{13}\) g cm\(^{-3}\). This part of the calculation included only (adiabatic) hydrodynamics and a partial loss of gravitational radiation. There are two questions which have to be asked at this point: (i) do energy loss or energy production mechanisms exist which work on shorter time scales, thus making the present result invalid, and (ii) which physical processes will shape the future behavior of that system, its stability, energy balance between nuclear energy generation and neutrino losses, and its composition. We address these questions below.

4.1 Neutrino Escape

Let us first assume that all constituents are in thermal equilibrium and that 2 MeV neutrinos populate the disk and 10 MeV neutrinos the central object. The neutrino nucleon scattering cross section (see e.g. Tubbs and Schramm 1975, Shapiro and Teukolsky 1983) is given by

\[
\sigma_{\nu,n} = \frac{1}{4} \sigma_0 \left( \frac{E_\nu}{m_e c^2} \right)^2 \approx \sigma_0 E_\nu^2 (\text{MeV}) \quad \sigma_0 = 1.76 \times 10^{-44}\text{cm}^2.
\]

The mean free path between scattering events \(\lambda_\nu = 1/(n\sigma_{\nu,n})\) depends on the scattering cross section \(\sigma_{\nu,n}\) and the nucleon number density \(n = \rho N_A\). This can be expressed in the following form

\[
\lambda_\nu = 9.9 \times 10^5 \left( \frac{10\text{MeV}}{E_\nu} \right)^2 \frac{10^{12}\text{gcm}^{-3}}{\rho} \text{cm}.
\]

Thus, for a typical density of \(10^{12}\) g cm\(^{-3}\) and 2 MeV neutrinos in the disk we find a mean free path of 246 km and for an average density of the central object of about \(8 \times 10^{14}\) g cm\(^{-3}\) and 10 MeV neutrinos we find 12.4 m. In case the disk settled already to a nuclear statistical equilibrium for that temperature and density, it will consist partially of alpha particles (see Lamb et al. 1978), which leads however to the same result (we have to multiply \(\lambda_\nu\) by \(\approx N^2/A = 4/A\), making use of the coherent scattering cross section for nuclei and approximating the Weinberg angle by \(\sin^2 \theta_W = 0.25\)).

When scattering events follow a random walk in three dimensions, the timescale for
travelling an absolute distance $d$ is

$$t(d) = \frac{3d^2}{\lambda_\nu c}.$$  \hspace{1cm} (6)

It has been pointed out by Burrows and Lattimer (1986) and Burrows (1988) that neutrinos do not follow a scattering random walk, but that their behavior is more complex and absorptions can occur as well. Nevertheless, equation (6) gives a good order of magnitude estimate, sufficient for our discussion. When employing a typical height of the disk of 15 km and 10 km for the central object, this leads to typical diffusion time scales of $9.1 \times 10^{-6}$s in the disk and $8.1 \times 10^{-2}$s in the central object. The first number indicates that neutrino escape would already be substantial in the disk at a time of 1ms, while this is not the case for the central object that cools on a longer time scale. The difference with respect to a supernova core collapse, where typical neutrino diffusion time scales of about 1s occur, is due to the somewhat higher temperatures obtained in the proto neutron star after stellar collapse and bounce. With thermal energies being higher by about a factor of 2.5-3.0, leading to a factor of 6-9 in mean free path and diffusion time scale, where the energies enter quadratically, we have $t_{diff} = 0.5 - 0.7$s.

4.2 Neutrino Emission and Cooling

Now we have to consider which processes produce these neutrinos and whether their production time scales are as fast as the loss time scales. The very high densities and high temperatures of the central object produce neutrinos of all families essentially instantaneously in comparison to the diffusion time scales. Neutrinos exist in thermal equilibrium and all neutrino families carry similar energies. The total internal energy gain of $\approx 4.4 \times 10^{52}$erg will be released with a diffusion time scale of roughly 0.08s (as derived above). This differs from a supernova core collapse, where the proto neutron star releases about $2.5 \times 10^{53}$erg on time scales of 1s.

Accretion onto the proto neutron star in a supernova core collapse with a rate $\dot{M}$, before the formation of the delayed shock via neutrino heating, will increase the neutrino luminosity by (see e.g. Burrows 1988)

$$L_\nu = \dot{M} \frac{GM}{R} \left( \frac{2}{1 + \sqrt{1 - 2GM/c^2R}} \right) \approx 3 \times 10^{52} \left( \frac{\dot{M}}{0.1M_\odot s^{-1}} \right) \left( \frac{M}{M_\odot} \right) \left( \frac{10\text{km}}{R} \right) \text{erg s}^{-1}. \hspace{1cm} (7)$$

A similar accretion phase, fed by matter in the disk, can occur onto the central object in systems of merged neutron star binaries.

While in the central object energy loss due to neutrinos is limited by the diffusion time scale controlling the possible escape, we noticed that the neutrino diffusion timescale
in the disk is very short (of the order of $10^{-5}s$). All neutrino production processes occur on longer time scales, and thus a free escape for neutrinos from these neutrino cooling processes is guaranteed. Positron capture on neutrons and electron capture on protons which emit anti-neutrinos and neutrinos dominate over all other neutrinos losses (pair, photo, plasma, and bremsstrahlung neutrinos – Schinder et al. 1987, Itoh et al. 1989, 1990). The rates for the disk conditions are of the order of $10^{20} - 10^{22} \text{erg g}^{-1} \text{s}^{-1}$ (Fuller et al. 1980, 1982). The next largest energy loss rate due to pair and plasma neutrinos amounts only to about $10^{15} - 10^{18} \text{erg g}^{-1} \text{s}^{-1}$ (Itoh et al. 1989, 1990). The time scale for electron and positron captures (and accompanying neutrino production) is of the order $10^{-3} - 10^{-4}s$, which underlines our previous conclusion of free escape (with diffusion time scales of the order of $10^{-5}s$), i.e. the full energy loss is encountered and no trapping occurs.

4.3 Thermonuclear Energy Generation in the Disk

In addition to the cooling processes, we have to consider nuclear energy generation in the disk. Simple nuclear network calculations show that any initial protons combine with neutrons to form alpha particles on a time scale of less than $10^{-17}s$ for the density and temperature conditions in the disk. This releases an energy of 7.07 MeV per nucleon or $6.8 \times 10^{18} \text{erg g}^{-1}$ or $1.36 \times 10^{51} \text{erg}$ per $0.1M_{\odot}$ of disk matter burned to helium. Disk material in beta equilibrium would consist of less than 1% protons. Thus, a disk in weak equilibrium would burn 2% of its mass (equal amounts of neutrons and protons) instantaneously, releasing about $10^{20} \text{erg}$ (on a $10^{-17}s$ time scale). This would heat and somewhat expand the disk because of a higher energy generation than neutrino cooling rate.

The subsequent evolution is complex and depends on many details, i.e. the resulting density and temperature of the disk and the free proton fraction $Y_p$. High densities (large electron Fermi energies) and/or temperatures favor electron captures on protons. For the disk densities electron capture time scales of about $10^{-1} - 10^{-4}s$ with a neutrino cooling rate close to $10^{20} - 10^{24} \text{erg g}^{-1} \text{s}^{-1}$ are obtained in the temperature range of 1-10MeV. The positron capture rates on neutrons are generally lower, because positrons have a negative Fermi energy. Similar values as for electron captures on protons can only be attained at high temperatures and lower densities. These rates would usually balance in order to keep beta equilibrium. But each proton produced will burn to helium on much shorter time scales, thus dominating over electron captures. The depletion in the free proton abundance $Y_p$ is replenished by further positron captures on neutrons, and gradually the free neutrons would burn to He via conversion to protons.

A total energy gain would occur if the density and proton abundance (determining the electron Fermi energy and electron capture rate on protons) become small enough that neutrino losses from electron captures are negligible, and the temperatures are low enough that the neutrino losses in each positron capture reaction (producing one proton) become smaller than 14MeV, which is $1/2$ of the energy released when two resulting protons burn with two available neutrons to He. This is the case for temperatures of $\leq 1\text{MeV}$, $Y_e \leq 10^9g \text{cm}^{-3}$, and $Y_p \leq 10^{-3}$ (tables 2 and 3 in Fuller et al. 1982 and private communication). The low temperatures are required as the weak capture cross sections scale with $E^2$ and thus the average energy of the positron and the escaping neutrino are higher than the thermal value. For the conditions discussed above the time scale of burning neutrons via protons to He is of the order 100s.
On the other hand, the low densities ensure electron Fermi energies of less than 5MeV. Thus, the formation of heavier elements via neutron captures and beta decays with typical Q-values of 13-15MeV will not be inhibited by electron captures on neutron-rich nuclei. An r-process similar to the r-process suggested for inhomogeneous big bang nucleosynthesis (Applegate, Hogan, Scherrer 1988; Thielemann et al. 1991) with light nuclei embedded in a neutron bath will occur. Then the long time scales required to convert neutrons via protons to helium can be overcome by this feature of the burning process. Once the first heavy seeds are available, neutron captures and beta decays (not temperature dependent) will drive the burning on beta decay time scales of 0.1 to 0.01s and cause a net energy deposition in the low density and temperature parts of the disk. The maximum energy of photons emitted in the nuclear reactions would result from beta decays far from stability with decay energies up to 15-20 MeV. The scenario is very similar to that of “cold decompression” (Lattimer et al. 1977, Meyer 1989), with the exception that no energy loss terms and an adiabatic expansion fueled by the energy generation was considered in the latter. The total expected energy release is of the order $f \times (5 - 6) \times 10^{51}$ erg, dependent on the distribution of nuclei produced and the fraction $f$ of the disk involved.

In the interior of the colliding neutron stars one finds 5-30% protons, depending on the equation of state (Weber and Weigel 1989). This increase is due to the fact that neutrons and protons are also degenerate at such high densities and the neutron and proton Fermi energies are of importance as well. If such matter would be ejected during the collision on shorter than weak equilibrium time scales, up to 60% of the disk matter could burn instantaneously, releasing about $3 \times 10^{51}$ erg in the disk. Our calculations show that the ejection of matter into the disk occurs on time scales comparable to weak equilibrium time scales and a behavior closer to the earlier discussion is expected.

In none of the cases would the energy release be sufficient to unbind and “evaporate” the whole disk, which has a total binding energy of $\sim 4 \times 10^{52}$ erg but possibly the outer lower density parts can escape. Further calculations which treat energy loss rates and thermonuclear reactions consistently as part of the hydro code will hopefully be able to answer the open “details”.

4.4 Stability of the Central Coalesced Object

Our hydro calculations produced a central object which was stable on the time scales considered in this calculation. We have to consider several points to understand that outcome and its generality. The main question concerns the limiting neutron star mass for a given equation of state. The object is rotating slower than, but close to the, Kepler frequency of typically $\Omega_K = 10^4$ s$^{-1}$. This allows an increase beyond the limiting mass of non-rotating neutron stars of 14-17% (Friedman, Ipser, Parker 1986; Friedman and Ipser 1987; Weber and Glendenning 1992), the largest values are related to the softest equations of state.

The core is rapidly rotating but $Jc/GM^2 \approx .6 < 1$. (similar results were obtained by Shibata Nakamura & Oohara (1992) and by Rasio & Shapiro (1992)). Thus the core is not supported completely by rotation. This can be understood intuitively from the following argument. The angular momentum just before the stars are in contact is

$$J \approx 2M \Omega R^2 = 2M (GM/4R^3)^{1/2} R^2 = (GM^3 R)^{1/2}. \quad (8)$$
The ratio \( \frac{J}{(4GM^2/c)} \) determines whether the object is rotationally supported. We find

\[
\frac{J_c}{4GM^2} = \frac{1}{4} \left( \frac{R}{r_g(M)} \right)^{1/2},
\]

where \( r_g(M) \) is the gravitational radius of one of the stars, i.e. \( r_g \approx 2(M/1.4M_\odot) \) km. With \( R \approx 10 \) km we have that \( R/r_g \approx 5 \) and hence \( Jc/GM^2 \approx .5 \). The ratio is slightly higher in our case because of the spin of the neutrons stars. However, it is quite impossible to reach \( Jc/GM^2 > 1 \). The intuitive explanation for the phenomenon is that while orbiting each star is attracted only by its companion and the centrifugal force has to balance this gravitational attraction. However, once the two objects coalesce the gravitational force towards the center results from a total mass of \( 2M \) which is too large to be stopped purely by the rotation.

The behavior of the equation of state at high densities is dominated by the compressibility \( K \) and the possible appearance of new degrees of freedom, i.e. existence of excited baryons (hyperons) in addition to neutrons and protons, kaon condensates, etc. If it is permitted to populate these states rather than neutrons and protons with their huge Fermi energies, a smaller pressure and softer EOS results, which causes a smaller limiting neutron-star mass. Our polytropic index was adjusted to the Lattimer and Swesty EOS with a compressibility \( K = 180 \) MeV. Only neutrons, protons, and leptons are considered at high densities. This corresponds to a large limiting neutron star mass of \( 2.2M_\odot \).

Most recent relativistic equations of state result in limiting masses for non-rotating neutron stars below \( 2.2M_\odot \), down to \( 1.4-1.6M_\odot \) (Glendenning et al. 1992). With a 14-17\% increase for objects close to their Kepler frequencies, we would start out with stable central objects for the upper end of this mass range. We would expect neutrino release from the hot object and mass accretion until accretion, cooling or loss of angular momentum lead to surpassing the limiting mass and cause the collapse to a black hole. We expect that this will take place on a time scale of seconds. A limiting non-rotating neutron star mass smaller than \( 2M_\odot \) would lead to an unstable coalesced object from the time of its formation. Thus, a black hole would form after about 1ms and a neutrino burst of only a few times \( 10^{50} \) erg would result during the collapse (Gourgoulhon and Haensel 1993). This will be enhanced by continuing accretion from the disk.

This clearly shows that the outcome of a neutron star merger, its neutrino signal, electromagnetic radiation, and probably also the disk behavior depend decisively on the limiting neutron star mass. If identified with observed events (i.e. a gamma ray burst), we could finally expect observational contraints for this theoretically quite uncertain quantity.

5. DISCUSSION

For all the neutron-star spins considered here, the merger results in a rapidly rotating, though not completely rotationally supported, central coalesced core surrounded by a disk of ejected material. A few percent of the total mass is ejected into long extended tails in some of the cases. During the merger event material from the two neutron stars is
compressed adiabatically, increasing the temperature of the gas to $\approx 10$ MeV and it becomes a strong source of neutrinos. The neutrino cooling time is around $10 \rightarrow 100$ms. The hot rapidly rotating core is on the edge of stability and its subsequent evolution depends on the true equation of state. The disk material gets heated via shocks and its temperature is $\approx 5$MeV. Since its density is much lower its neutrino cooling time is significantly shorter. Thermonuclear reactions can, however, heat the disk and its thermal evolution depends on a delicate balance between heating and cooling.

Our calculations included several simplifying approximations and it is worth while to examine their validity before turning to the interpretation of the results: (i) We employed a polytropic equation of state instead of a realistic one. However we found that by the end of the calculations the central density increased only by 5% and the polytropic approximation to the Lattimer & Swesty (1991) equation of state is valid at this density range. (ii) We did not include neutrino transport and cooling. The neutrino cooling time of the core ($10 \rightarrow 100$ms) is sufficiently long that there are no significant neutrino losses up to the moment that we stop the calculations. The time scale for neutrino cooling of the disk is, however, much shorter and our approximation breaks down there. (iii) We switch off the gravitational radiation backreaction force when the two neutron stars touch each other. The core becomes axisymmetric within a fraction of a millisecond after the two stars have come in contact, and only a small amount of energy is released via gravitational radiation during this phase (Shibata et. al., 1993, Rasio & Shapiro 1993). (iv) Our worst approximation is that apart from gravitational radiation backreaction we did not include other general relativistic effects. There is no physical justification for that, only a technical one, the required three dimensional hydrodynamics calculations are beyond the current scope of numerical relativity. General relativistic effects are the major uncertainty in the calculations that we present here. We have no way of estimating them quantitatively. Qualitatively we expect that they might enhance the dynamical instability of the central rotating core.

The mass of the central object ranges between $2.4 \rightarrow 2.6 M_\odot$. This is above most upper limits for cold neutron star masses even with the stiffest equations of state, (e.g. the limiting mass is $2.2M_\odot$ for non-rotating neutron star with the Lattimer-Swesty equation of state that we follow). It is clear that if the real equation of state is soft, the central object must be unstable and it will collapse quickly to a black hole. The situation is less clear with a hard equation of state. The maximum mass allowed increases by up to 17% when the neutron star is rotating rapidly. Furthermore, our hydrodynamics calculations show that the Newtonian system with a hard equation of state, seems stable and does not show any tendency to collapse. At the end of the calculation the central polytropic index is large which indicates that the system is far from the instability that sets in as the polytropic index decreases to $4/3$. Here it is possible that the combination of rotation and thermal pressure may prevent the direct collapse. The collapse might be induced later when additional mass from the disk accretes onto the central object or when it cools down. General relativistic effects, which we ignore here, might have a critical effect and cause the central merged object to collapse sooner. As the central core is on the edge of stability it might be the case that in some cases (depending e.g. on the spin angular momentum) the core remains extended initially while in others it collapses directly to a black hole. In either case the ultimate final configuration is a black hole.
The total neutrino emission depends on how fast the core collapses to a black hole. One can expect appreciable neutrino emission if the rotating configuration is stable, or if the combination of rotational and thermal energy make the system stable, halts the collapse until the core cools or accretes additional matter. If general-relativistic effects or a softer equation of state will make the system unstable a black hole will form after a few msec without much neutrino emission form the core.

The disk extends up to 60km or so. Its total mass is $\approx 0.4M_\odot$, and its total kinetic energy $\approx 2.5 \times 10^{53}$ erg. From our SPH runs, we see that the disk contains $\sim 10^{51}$ erg of thermal energy, with some of the material being shock-heated up to temperatures of $\sim 2-4$ MeV. The disk is optically thin to neutrinos, and neutrino cooling exceeds heating for temperatures $T \geq 1$ MeV, densities $\rho Y_e > 10^9 g/cm^3$ and free proton fraction $Y_p > 10^{-3}$. Hence we expect neutrino cooling to lead to a flattened disk in the inner high density part.

There occurs a thermonuclear flash of burning of the initial protons (with neutrons) in the disk to He. At locations in the disk where after that initial flash $T$, $\rho Y_e$, and $Y_p$ are equal or smaller than the values listed in the previous paragraph net energy deposition rather than neutrino losses will dominate and a further expansion is expected. The low densities ensure electron Fermi energies less than 5MeV. Thus, the build-up of heavy elements from the seed nuclei, produced in the early thermonuclear flash, via neutron captures and beta decays (with Q-values of typically 13-15MeV) is not inhibited by reverse electron captures on these neutron-rich nuclei. If a fraction $f$ of the whole disk is burned to heavy elements, this would result in a total release of $f \times (5-6) \times 10^{51}$ erg. We expect some of the material of the outer disk and extended arms to become unbound from the system due to this nuclear energy release and to contain r-process matter produced during the decompression phase.

Ratnatunga & van den Bergh (1989) find core collapse supernova rate in our galaxy of $2.2 \pm 1.1$ per century, $10^{-6}$ $M_\odot$ to $10^{-4}$ $M_\odot$ of r-process material must be ejected per supernova event (Cowan, Thielemann & Truran 1991) to produce the observed abundance of r-process material. Eichler et. al.(1989) suggested that, alternatively, the r-process takes place in neutron star mergers. They estimated the production rate using the merger rate estimated by Clark van den Heuvel & Sutantyo. Recently, Narayan et. al.(1991) and Phinney (1991) estimate that there are $10^{-5.5 \pm 5}$ neutron star mergers per year per galaxy. Using the updated rate we find that to explain r-process nucleosynthesis as a by product of neutron star merger one requires the ejection of $10^{-2} M_\odot$ to $1 M_\odot$ per event. The upper limit is clearly impossible, however, the lower range is of the same order of magnitude of what we observed in the long tails in some of the mergers. Some of the r-process material is observed in very low metalicity stars and galactic evolution models would require that it is produced as early as a few $\times 10^7 - 10^8$ years after the formation of the galaxy. This is comparable with the life time of the famous binary pulsar PSR-1913+16 and one might expect that there is no problem caused by this time scale. This suggests that some r-process material could be produced via this mechanism even on such a short time scale.

We turn now to the implications to $\gamma$-ray bursts models. Cosmological $\gamma$-ray bursts require $\approx 10^{51}$ ergs with a rise time as short as tens of msec (for some of the short bursts). The hot coalesced core is one possible source of this energy. If the core does not collapse directly to a black hole it will emit its thermal energy as neutrinos. The neutrino flux is sufficiently large that $\approx 10^{-2}$ to $10^{-3}$ of it could be converted to electron-positron pairs.
via $\nu\bar{\nu} \rightarrow e^+e^-$ (Goodman et al., 1987) and those could produce a $\gamma$-ray burst (Eichler et al., 1989). The time scale for the neutrino burst is short enough to accommodate even the shortest rise times observed. Additional neutrino luminosity could arise on a longer time scale from accretion of the disk material on the central object.

An additional energy source that could power a $\gamma$-ray burst from a neutron star merger is the disk surrounding the central object. This energy source can operate regardless of the question of whether the central object collapse directly to a black hole or not. In this case the neutrino emission is not sufficient to power a burst via neutrino annihilation, however, Narayan et al. (1992) have suggested that the magnetic energy density in the disk can reach equipartition with the kinetic energy (corresponding to $10^{16}$ gauss) and recombination of such a magnetic field could produce an electromagnetic burst. It is not clear what will be the time scale for the emission from this disk. It could take several seconds or even longer, depending on the viscosity in the disk, which is unknown at present. As was stressed by Narayan et al. (1992) the large variability in observed properties of $\gamma$-ray bursts could be explained by those two alternative energy sources. An additional source of diversity, which we discover here, is the distinction between systems that collapse directly to a black hole and those that undergo a longer rotating core phase.

The initial $\gamma$ rays that are produced are not observed directly. The huge optical depth of the resulting electron-positron plasma would result in an optically thick fireball (Goodman, 1986; Shemi & Piran, 1990; Paczynski, 1990; Piran & Shemi, 1993). The observed $\gamma$-rays emerge much later when the fireball becomes optically thin (due to its rapid expansion) or when it interacts with interstellar material (Meszaros & Rees, 1993). This late stage determines both the duration and the spectrum of the bursts. Shemi & Piran (1990) have shown that to produce a $\gamma$-ray burst the fireball must be practically free from baryonic load. If more than $10^{-5}M_\odot$ baryons are injected into the fireball they will prevent the formation of a $\gamma$-ray burst. Larger amounts of mass are ejected in a neutron star merger. However all the mass is ejected into the equatorial plane. In all mergers simulated, we note the the regions above the poles of the coalesced object are relatively free of material. The accuracy of our simulation is of course not enough to estimate the amount of matter in those funnels, but it seems that those funnels may provide an avenue for neutrino escape and the production of a clean radiation fireball. This suggests that the $\gamma$-rays are beamed and appear only in certain directions.

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| Run | $\Omega_{\text{contact}}$ | Npart | Realign | $M_{\text{disk}}/M_{\text{system}}$ |
|-----|--------------------------|-------|---------|-----------------------------------|
| A   | 0.2329                   | 8542  | no      | 0.13                              |
| B   | 0.0478                   | 8542  | no      | 0.10                              |
| C   | 0.2143                   | 8542  | yes     | 0.13                              |
| D   | 0.0109                   | 8542  | yes     | 0.10                              |
| E   | -0.0439                  | 2512  | yes     | 0.05                              |
| F   | -0.1257                  | 2512  | yes     | 0.05                              |
**FIGURE CAPTIONS**

**Figure 1:** Density contour plots for run C, in the plane of the neutron star trajectories. Time and distances are given in code units (1 time unit $\equiv 73\mu s$, 1 distance unit $\equiv R_{ns} = 10\text{km}$). Logarithmic contours are plotted, at intervals of 0.25 dex, beginning at a density of 0.001 in code units (1 code unit $\equiv 2.786 \times 10^{15}\text{g/cm}^3$).

**Figure 2:** Density contour plots for run D. Units as given in Figure 1.

**Figure 3:** The distribution of SPH particles, in the plane of the original orbits, at the end of run C. Distance units as given in Figure 1.

**Figure 4:** Density contour plot of the final configuration in run C in a plane perpendicular to that of the neutron star trajectories. Logarithmic contours are plotted, at intervals of 0.25 dex, starting at a density of 0.0001 in code units (1 code unit $\equiv 2.786 \times 10^{15}\text{g/cm}^3$).

**Figure 5:** Contour plot of column densities, i.e. at a given location $(r, z)$, column density, $\rho_c = \int_0^\infty \rho dz$.

**Figure 6:** The spin angular velocities of the SPH particles as a function of cylindrical radius. For run C (figure a) and run D (figure b). Both angular velocity and distance are in code units (1 distance unit $\equiv R_{ns} = 10\text{km}$, $\Omega = 1$ is equivalent to a spin period of 0.461ms).

**Figure 7:** The enclosed mass (as a fraction of total mass) as a function of cylindrical radius (in units of the neutron-star radius).

**Figure 8:** The energy per unit mass of each SPH particle against the density as the particle site for the final configuration in run C. Both are given in code units, energy/mass units being $1.858 \times 10^{20}\text{ergs/gm}$.

**Figure 9:** Isotherms produced using Lattimer and Swesty’s nuclear equation of state, assuming beta equilibrium. Isotherms are draw for $T = 1.75, 2, 3, 5, 7, 10, 15,$ and $25\text{MeV}$.