SUPRANOVA EVENTS FROM SPUN-UP NEUTRON STARS: AN EXPLOSION IN SEARCH OF AN OBSERVATION

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ABSTRACT

We consider a formation scenario for supramassive neutron stars (SMNSs) that takes place through mass and angular momentum transfer from a close companion during a low-mass X-ray binary phase, with the ensuing suppression of the magnetic field. After the end of the mass transfer phase, SMNSs will lose, through magnetic dipole radiation, most of their angular momentum, triggering the star’s collapse to a black hole. We discuss the rate of occurrence of these collapses and propose that these stars, because of the baryon-clear environment in which the implosion/explosion takes place, are the originators of gamma-ray bursts.

Subject headings: black hole physics — gamma rays: bursts — instabilities — relativity — stars: neutron

1. INTRODUCTION

One of the key requirements of gamma-ray burst (GRB) models is that they make contact with the fireball model (Rees & Mészáros 1992), which has proven so successful in predicting and interpreting the observed properties of GRB afterglows. In particular, this entails that a large explosion is to take place in a region with small baryon contamination: for $E = 10^{50}$ ergs, the baryon contamination must be at most $E/\gamma c^2 \approx 10^{-4} M_\odot$ for $\gamma = 300$, the bulk Lorentz factor of the explosion. Vietri & Stella (1998, hereafter Paper I) presented a model that could accomplish this, involving a supramassive neutron star (SMNS), i.e., a neutron star that has a larger baryon number than any normal neutron star because it derives part of its support against self-gravity from the centrifugal force; these supramassive stars cannot be slowed down to zero spin rate because they are so massive that, as they lose angular momentum, they become unstable to black hole formation before reaching zero spin rate (Cook, Shapiro, & Teukolsky 1994a, 1994b; Salgado et al. 1994). The model consists of the implosion/explosion of a supramassive neutron star which has lost through magnetic dipole radiation so much angular momentum that it must then collapse to a black hole; the rotational energy of the small amount of equatorial mass left behind because it is already in near centrifugal equilibrium provides the energy source that powers the burst.

In Paper I we proposed that SMNSs are formed in the SN explosion of a core with too much mass and angular momentum to end up in a normal neutron star. Although nothing yet stands against this possibility and we are not retracting it, we have now realized that a different channel exists: mass and angular momentum accretion from a companion in a low-mass X-ray binary (LMXB). The following discussion is relevant to the formation of millisecond pulsars (MSPs) and will mimic arguments used in discussing the evolution of LMXBs into MSPs. We first discuss how the accretion of large amounts of mass and angular momentum may be realized in nature and then apply this scenario to GRBs.

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2. MASS AND ANGULAR MOMENTUM ACCRETION

The main obstacle to accretion onto a normal neutron star of large amounts of angular momentum from a close companion via an accretion disk is the neutron star’s magnetic field: the neutron star can rotate only so fast as to make the corotation and Alfvén radii coincide, lest a propeller phase set in, which would actually entail angular momentum loss (Ghosh & Lamb 1978; Illarionov & Sunyaev 1975). The coincidence of these two radii leads to an equilibrium period $P_{eq} = 1.3(B/10^{12} \, \text{G})^{6/7}(M/M_\odot)^{-3/7} \, \text{s}$ (Ghosh & Lamb 1992), which clearly shows that $B$ must decrease before significant spin-up can occur. Although no unique model has emerged yet, the current consensus is that the neutron star magnetic field decays by at least 3–4 decades as a direct result either of mass accretion (Phinney & Kulkarni 1994) or of the ensuing spin-up (Ruderman, Zhu, & Chen 1998).

The strongest constraints on field decay in neutron stars come from LMXBs and MSPs. In only one LMXB, SAX J1808.4–3658, a coherent 2.5 ms signal has been detected in the persistent X-ray emission, providing direct evidence for the presence of a small magnetosphere; the inferred magnetic field is in the $B \approx 10^8–10^{10}$ G range (Psaltis & Chakrabarty 1999). All other LMXBs have undetectably small coherent pulsations in their persistent emission, if at all. Yet spin periods in the 2–4 ms range have been deduced for about 10 LMXBs from the X-ray flux oscillations that are present during type I bursts emitted by these sources (see van der Klis 1999). Spin-up through accretion can have occurred in these neutron stars only if their magnetic field is lower than $\sim 10^9$ G. Further evidence that the field might have decayed to dynamical insignificance derives from the modeling of the kilohertz quasi-periodic oscillations (QPOs), a common phenomenon observed in LMXBs. In the sonic point model (Miller, Lamb, & Psaltis 1998), the Alfvén radius is located at a radius corresponding to a Keplerian frequency of $\approx 350$ Hz, corresponding to a magnetic field $B \approx 8 \times 10^8$ G: this already bears witness to a thousand-fold reduction of the magnetic field below that of a typical newborn pulsar. A better model explains QPOs in terms of the fundamental frequencies of test particle motions in the general-relativistic potential well in the vicinity of the neutron star (Stella & Vietri 1999). The model is capable of explaining the observed relation between peak QPO frequency and its lower frequency counterpart over 3 orders of magnitude in peak QPO frequency and several distinct classes of sources, including candidate black holes and LMXBs (Stella, Vietri, & Morsink...
implying that the magnetic field has been reduced already to $\leq 2 \times 10^8$ G during the LMXB phase.

A different argument involves a handful of MSPs with observed magnetic fields $\leq 2 \times 10^8$ G, of which there are currently about a dozen, including the lowest fields ever measured, $7 \times 10^7$ G in J2229+2643 and J2317+1439 (Camilo, Nice, & Taylor 1996). For these small fields, the Alfvén radius for disk accretion (Ghosh & Lamb 1978, 1992) is smaller than currently about a dozen, including the lowest fields ever measured.

The argument about QPOs (Stella & Vietri 1999) implies that the magnetic field in these objects is dominated by the dipole component, with negligible contribution from higher multipoles. Altogether, this means that we have already observed the result of accretion from a companion pushing the magnetic field to dynamical irrelevance (or, at least, very close to it) sometime during the mass exchange process. The argument about QPOs (Stella & Vietri 1999) implies that this may happen reasonably early in the LMXB history.

Detailed models are required to establish the exact history of a neutron star's mass, angular momentum, and magnetic field, but unfortunately these computations are currently fraught with uncertainties: Where is the B field located, in the core or in the crust? And what is an appropriate model for the field suffocation? Population-synthetic studies of this phenomenon (Possenti et al. 1999) have focused on two representative EOSs and modeled the decay of the magnetic field in two limiting cases, imposing at the crust-core boundary either complete field expulsion by the superconducting core or advection and freezing in a very highly conducting transition shell. The main result lies in the establishment of the existence of a tail in the rotation period distribution extending well beyond the shortest period observed so far ($P = 1.558$ ms), with only moderate dependence on the field suppression mechanism. For the softest EOS the period distribution is still increasing at the shortest value before the onset of mass shedding, where Possenti et al. stopped their computations, while for the stiffest one the period distribution had a wide maximum around $P = 2-4$ ms and a tail extending below this value. The fraction of objects with $P < 1.558$ ms is $\approx 1\%$ and $\approx 10\%$ for the stiff and soft EOS, respectively, all ending up with very small magnetic fields, $\leq 10^8$ G.

Although the accreted mass is larger when account is taken of the need to suppress the magnetic field than when the magnetic field is neglected, the difference is not very large (Burderi et al. 1999). So, in order to appraise whether the neutron stars thusly formed may be supramassive (or not), we simply consider Table I from Cook et al. (1994b). It shows that the total amount of mass that needs to be accreted from a companion in order to reach the supramassive stage at the initial point of mass shedding depends strongly upon the EOS. For EOS F, the neutron star collapses to a black hole even before reaching mass shedding. The soft EOSs (A, D, E, KC) have become supramassive; the intermediate EOSs (C, M, UT, FPS) are within less than 0.1 $M_\odot$ of doing so and will cross the threshold if accretion continues after the mass-shedding point is reached (see below).

Table I from Keplerian Disk onto a 1.4 $M_\odot$ Neutron Star

| Equation of State | $M/M_\odot$ | $M_{\text{ms}}/M_\odot$ | $\Delta M/M_\odot$ | $P^{1/2}$/ms | $R^{1/2}$/km | $h^{1/2}$/km | $M_{\text{ms}}/M_\odot$
|------------------|------------|-------------------------|-------------------|--------------|--------------|--------------|-----------------|
| A                | 1.77       | 1.57                    | 0.428             | 0.604        | 9.59         | 2.82         | 1.66            |
| C                | 1.74       | 1.54                    | 0.389             | 0.894        | 12.1         | 0.27         | 1.86            |
| D                | 1.76       | 1.56                    | 0.405             | 0.730        | 10.7         | 1.70         | 1.65            |
| E                | 1.76       | 1.57                    | 0.414             | 0.656        | 10.0         | 2.38         | 1.75            |
| F                | 1.52       | 1.59                    | 0.172             | 0.715        | 9.21         | 3.20         | 1.46            |
| L                | 1.80       | 1.52                    | 0.443             | 1.250        | 15.0         | 0.00         | 2.70            |
| M                | 1.74       | 1.50                    | 0.367             | 1.490        | 16.7         | 0.00         | 1.80            |
| N*               | 1.84       | 1.53                    | 0.484             | 1.080        | 13.6         | 0.00         | 2.64            |
| KC               | 1.74       | 1.55                    | 0.385             | 0.888        | 12.1         | 0.31         | 1.49            |
| AU               | 1.79       | 1.58                    | 0.446             | 0.701        | 10.4         | 2.01         | 2.13            |
| UT               | 1.78       | 1.56                    | 0.436             | 0.784        | 11.2         | 1.26         | 2.20            |
| FPS              | 1.76       | 1.56                    | 0.416             | 0.747        | 10.9         | 1.52         | 1.84            |

Note: — Data from Cook et al. 1994b.

- $^a$ Final total mass-energy.
- $^b$ Initial rest mass.
- $^c$ Accreted rest mass.
- $^d$ Rotation period.
- $^e$ Initial circumferential height of corotating marginally stable orbit.
- $^f$ Initial circumferential radius.
- $^g$ Maximum static total mass-energy for the EOS.
in total \( \geq 1-1.1 \ M_\odot \) to become supramassive. This requirement cannot be fulfilled: thermal stability in the mass exchange process through Roche lobe overflow (Webbink, Rappaport, & Savonije 1983) requires that the companion of the neutron star have a mass below \( \frac{1}{2} \) of the neutron star’s, i.e., \( 1.17 \ M_\odot \) for an initial neutron star mass of \( 1.4 \ M_\odot \). Since the smallest He core that may be left behind is that of a star which took all Hubble time to evolve off the main sequence, \( 0.16 \ M_\odot \), this leaves a maximum transferable mass of \( 1.01 \ M_\odot \), less than required.

Recent studies of binary pulsar masses (Thorsett & Chakrabarty 1999) seem to argue against significant mass accretion, but it should be noticed that, by investigating millisecond pulsars with periods exceeding \( \approx 2 \) ms, the authors are investigating objects for which we know a priori that little mass need have been accreted, since their periods are long compared with SMNS’s. We may expect different results when pulsars are chosen otherwise: a recent redetermination of the mass of Cyg X-2 finds \( M = (1.8 \pm 0.2) \ M_\odot \) (Orosz & Kuulkers 1999), departing from the narrow range of Thorsett & Chakrabarty.

The lowest magnetic field for the formation of supramassive neutron stars may be lower than the empirical value \( (\approx 2 \times 10^{9} \ G) \) mentioned above (Possenti et al. 1999) because when mass accretion from the companion begins to taper off, or alternatively if mass accretion is intermittent, the Alfvén radius (which scales as \( M^{-2/3} \)) may expand farther than the corotation radius: the neutron star then goes through a new propeller phase which slows its rotation. The overall effect is not large so that we shall consider in the following a maximum magnetic field \( q \times 10^{9} \ G \), with \( q \leq 1 \).

3. FURTHER EVOLUTION

Supramassive neutron stars are unstable to collapse to a black hole when angular momentum losses reduce the initial angular momentum to about half of the initial value; furthermore, these stars are peculiar in that evolution at constant baryon number, but decreasing total angular momentum, makes them spin up rather than down; all of this is especially evident in Figures 7, 10, 13, and 16 of Salgado et al. (1994). Magnetic dipole losses cause a net torque which spins down the neutron star in a time (Paper I) \( t_{\beta} = (5 \times 10^9 \ yr) \left(10^6 \ G/B\right)^2 \). This timescale is not strongly dependent upon EOS, but depends strongly upon whether the model is only marginally supramassive or close to the absolute maximum mass (rotating or not) for the given EOS, so that it may be considerably shorter under many circumstances. Thus, a time \( t_{\beta} \), after becoming supramassive, the neutron star will collapse to a black hole. This time is reasonably long when compared with typical mass accretion timescales, which, as discussed above, are typically determined by sub/giant nuclear evolution timescales. Thus, mass transfer will have long since ceased, and the immediate SMNS surroundings will be reasonably baryon-free. The companion star, in the meantime, will have settled down as a low-luminosity, low-mass white dwarf, which is not expected to pollute the environment either. Furthermore, we can gauge the baryon cleanliness of the SMNS surroundings at large if we assume that MSPs are born through the same chain of events, except less extreme, for then we know the Galactic distribution of MSPs. These are often located well outside the Galactic disk, within an interstellar medium with typical densities well in defect of \( n = 1 \ cm^{-3} \), which makes the total baryon mass within, say, \( 0.1 \ pc \), less than \( 10^{-3} \ M_\odot \), more than enough to guarantee contact with the fireball model. We thus see that this version of the formation scenario also guarantees a baryon-clean environment, exactly like the scenario of Paper I.

The situation is clean even in the case in which the collapse occurs while mass transfer is still taking place. The total amount of baryons in the accretion disk is negligible: the disk crossing time is of order of \( \approx 1 \) month, which, with mass transfer rates \( 10^{-9} \) to \( 10^{-8} \ M_\odot \ yr^{-1} \), corresponds to much less than the maximum contamination value. The total amount of outlying mass from a wind is also rather small: for \( M_\odot \approx 10^{-9} \ M_\odot \ yr^{-1} \) and \( v_\infty \approx 30 \ km \ s^{-1} \), the total mass within, say, \( 0.1 \ pc \) is \( 3 \times 10^{-6} \ M_\odot \), again negligible. The highly relativistic ejecta and gamma rays from the burst will hit the companion and form a shock wave inside the star’s photosphere, so that local dissipation of the ejecta kinetic energy will lead to the companion’s inflating on the (long!) Kelvin-Helmholtz timescale and the nonthermal afterglow emission will not be contaminated by the reradiated thermal component.

The mechanism for the energy release is the same as discussed in Paper I: once the neutron star is destabilized, the innermost regions will collapse promptly to a black hole, while the equatorial matter, which is close to centrifugal equilibrium, will contract just a little bit and begin orbiting the newly formed black hole. A necessary condition that needs to be met is that this equatorial material lie outside the innermost stable orbit. This can be checked from Table I of Cook et al. (1994b), who show that neutron stars that have reached the mass-shedding regime have equatorial radii larger than the innermost stable orbit (see their column \( j \)), independent of EOS. In Paper I, we estimated the amount of matter left behind as \( \approx 0.1 \ M_\odot \); this configuration is identical to the one hypothesized in most current models (Mészáros 1999), and the debris torus is massive enough to power any burst, especially in the presence of a moderate amount of beaming.

We now discuss the rate at which spun-up SMNSs collapse to black holes. Since the timescales involved are a fair fraction of the age of the universe, and since star formation evolves strongly in the recent past (Madau et al. 1996), we have to consider cosmological evolution of the population. However, from Figure 1 of White & Ghosh (1998), it can be seen that the population of MSPs is roughly constant (within the accuracy of the present, order-of-magnitude estimates) over the redshift range \( 0 < z < 1 \) for most assumptions. There are currently an estimated \( 5 \times 10^6 \) MSPs in the disk of the Galaxy (Lorimer 1995; Phinney & Kulkarni 1994); assuming that there are as many systems in the bulge, that a fraction \( \beta \) of these are SMNSs, and that the typical timescale for collapse to black hole is given by \( t_{\beta} \), the expected rate of collapses in the Milky Way is \( r = \beta(5 \times 10^7 \ yr) \), which is to be compared with the inferred rate of GRBs, one every \( 3 \times 10^7 \) yr in an \( L_\star \) galaxy like the Milky Way. Scaling to \( \beta = 0.05 \), a value intermediate between the extremes of the simulations of Possenti et al. (1999), we find that the two rates agree for a beaming fraction \( \delta \theta/4\pi \approx 0.03(\beta/0.05) \); this is consistent with the idea that these explosions do not require extreme beaming fractions, since the explosion need not wade its way through a massive stellar envelope, but immediately breaks free into a baryon-clean environment.

This model makes an easily testable prediction because the location of bursts inside their host galaxies is the same as that of LMXBs, which are distributed at distances from the Galactic plane \( \approx 1 \ kpc \), most likely arising from kick velocities at the time of neutron star formation (van Paradijs & White 1995). A similar \( z \) distribution is observed for MSPs, \( z \approx 0.7 \ kpc \), and moderate transverse speeds. Thus we would expect GRBs to
cluster around galactic disks (contrary to the binary pulsar merger model, in which at least some 50% of all GRBs should be uncorrelated with the original birth galaxies), but should not correlate with star-forming regions (except for SMNSs that form directly during a SN event, as discussed in Paper I), contrary to all scenarios involving massive stars. Also, the redshift distribution of GRBs within this model should be flatter than the star formation distribution (again contrary to hypernovae) because the redshift distribution of the MSP population is rather flat (White & Ghosh 1998).

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