1E 161348–5055 IN THE supernova remnant RCW 103: A magnetar in a young low-mass binary system?

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Received 2007 July 30; accepted 2008 March 9

ABSTRACT

We suggest that the unique X-ray source 1E 161348–5055 at the center of the supernova remnant RCW 103 consists of a neutron star in close orbit with a low-mass main-sequence star. The time signature of 6.67 hr is interpreted as the neutron star’s spin period. This requires the neutron star to be endowed with a high surface magnetic field of ~10^{15} G. Magnetic or/and material (propeller) torques are able to rapidly spin the young neutron star down to an asymptotic, equilibrium spin period in close synchronism with the orbital period, similar to what happens in the Polar cataclysmic variables. 1E 161348–5055 could be the first case of a magnetar born in a young low-mass binary system.

Subject headings: binaries: general — stars: magnetic fields — stars: neutron — novae, cataclysmic variables — supernovae: individual (RCW 103) — X-rays: binaries

1. INTRODUCTION

The soft X-ray source 1E 161348–5055 (hereafter 1E) was discovered by Tuohy & Garmire (1980) with the Einstein observatory, close to the geometrical center of the young supernova remnant (SNR) RCW 103. Owing to its thermal-like spectrum and to the lack of a counterpart at radio or optical wavelengths, 1E was classified as the first example of a radio-quiet, isolated neutron star. The host SNR is very young (<100 yr; Carter et al. 1997) and it is located at a distance of ~3.3 kpc (Caswell et al. 1975; Reynoso et al. 2004). The association of 1E and RCW 103 is very robust. The point source lies within 15′ of the SNR center; moreover, 1E and RCW 103 have consistent distance measurements (Reynoso et al. 2004).

The original interpretation of 1E as an isolated neutron star was later questioned by the observation of a large variability on few years’ timescale (Gottshelf et al. 1999). This has been confirmed in recent years, when a factor ~100 brightening (from ~9 × 10^{-13} to ~7 × 10^{-10} erg cm^{-2} s^{-1}) was discovered in comparing the first and second Chandra observations of the source, performed in 1999 September and in 2000 March, respectively (Garmire et al. 2000). Even more puzzling, a possible long periodicity at ~6 hr was hinted by the first Chandra observation (Garmire et al. 2000a), but it was not recognized in the second Chandra data set. Subsequent observations with Chandra (Sanwal et al. 2002) and XMM-Newton (Becker & Aschenbach 2002) failed to ultimately settle the periodicity issue.

A breakthrough came with the deep (90 ks) observation of 1E performed by XMM-Newton in 2005. De Luca et al. (2006) caught the source in a low state (1.7 × 10^{-12} erg cm^{-2} s^{-1}) and reported conclusive evidence for a strong (~50%), nearly sinusoidal modulation of 1E at $P = 6.67 \pm 0.03$ hr. The source spectrum, thermal-like and well described by the sum of two black bodies, varies along the 6.67 hr cycle and is harder at the pulse maximum (De Luca et al. 2006). The 6.67 hr periodicity was also recognized in the new XMM-Newton data set (collected in 2001, when the source was a factor ~5 brighter), although with a much smaller pulsed fraction (~12%) and with a remarkably different, complex light curve, featuring two narrow minima (or “dips”) per period, separated by 0.5 in phase. No faster periodicities were found, with a 3σ upper limit of 10% on the pulsed fraction down to $P = 12$ ms. De Luca et al. (2006) also reported a complete picture of the peculiar long-term variability of 1E. The source has been continuously fading for more than 5 yr, since the 1999–2000 re-brightening. In 2005 August 1E was seen with a flux similar to the pre-outburst level. XMM-Newton data clearly show that the spectrum of the source is also evolving. While a double-blackbody model yields in any case the best description of the spectrum, the emission is harder and more absorbed when the source is brighter.

On the optical side, deep observations of the field with the ESO VLT and with the Hubble Space Telescope (HST) showed nothing but two or three very faint IR sources (H ~ 22), consistent with the accurate Chandra position (De Luca et al. 2006 and references therein). No firm conclusions could be drawn about the possible association of any of such sources with 1E.

1.1. Isolated vs. Double Star Scenario

The combination of long-term variability, 6.67 hr periodicity, young age, and underluminous optical/IR counterpart makes 1E a unique source among compact objects in SNRs. As discussed by De Luca et al. (2006) 1E could be either a very young low-mass X-ray binary (LMXB) system, featuring a 2000 yr old compact object and a low-mass companion star (M4, or later) in an eccentric orbit, or a very peculiar isolated neutron star. In the latter case, 1E could only be a magnetar, dramatically slowed down, possibly by a propeller interaction with a debris disk. Both scenarios require highly nonstandard assumptions.

The binary scenario is appealing, since it provides an easy explanation of the long 6.67 hr periodicity, as due to the orbital period of the system. However, the observed properties of 1E are dramatically different from those of any known LMXB. The system is orders of magnitude dimmer than persistent LMXBs. Moreover, a very peculiar “double accretion” mechanism is required to explain the unusual large flux and spectral variations along the 6.67 hr cycle as an orbital modulation, as well as the dramatic long-term variability. Indeed, wind accretion along an eccentric orbit could naturally explain the quasi-sinusoidal modulation observed in the low state. However, the light curve of the source in its “active state” is very complex. In order to explain its features within such a scenario (as due to varying occultation of the central source), very peculiar structures at the rim of a (possibly transient) accretion disk are required.

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The isolated magnetar scenario could easily explain the phenomenology of the periodic modulation, with the star rotation bringing into view or hiding different emission regions. Indeed, the light-curve profile is very similar to the one observed in most magnetars. The long-term variability (in both luminosity and pulse profile) is also reminiscent of that seen in a few “transient” anomalous X-ray pulsars (AXPs). However, all known magnetars spin thousands of time faster than 1E, with periods well clustered in the 5–12 s range. The spin history of 1E should have been dramatically different. De Luca et al. (2006) and Li (2007) showed that propeller interaction with a supernova debris disk could provide a very efficient slowing-down mechanism. Such a mechanism could brake the neutron star to 6.67 hr in 2000 yr, provided that the star was born with a mass exceeding 1.4 M⊙ and a low main-sequence star (a red dwarf, RD) of mass M2 = 0.4 M⊙ (consistent with the current optical/infrared constraints). This system survived the supernova explosion which led to the birth of RCW 103. The orbit is expected to be born quite eccentric, with e most likely in the range 0.2–0.5 (De Luca et al. 2006); this eccentricity deserves some comments. In a narrow system like that we envisage for 1E, the tidal torques are expected to be quite effective in synchronization the companion star’s rotation with the orbital period and in circularizing the orbit. Unfortunately, there is no general consensus about the efficiency of these two processes (see, e.g., Zahn 1977; Tassoul 1988; Rieutord & Zahn 1997; Tassoul & Tassoul 1997; see also Meibom & Mathieu 2005 and references therein). The most conservative estimates of the synchronization and circularization timescales are those derived by Zahn (1977; his eqs. [4.12] and [4.13])

\[
t_{\text{sync}} = \frac{1}{6kq^2} \left( \frac{M_2 R_2}{L_2} \right)^{1/3} \frac{I_2}{M_2 R_2^2} \left( \frac{a}{R_2} \right)^6
\]

\[
t_{\text{circ}} = \frac{1}{63 q (1 + q/k^2)} \left( \frac{M_2 R_2^2}{L_2} \right)^{1/3} \left( \frac{a}{R_2} \right)^8,
\]

where \( L_2 \approx 2 \times 10^{32} \text{ergs}^{-1} \) is the main-sequence star’s luminosity, \( q = M_1/M_2 \) is the mass ratio, and \( I_2 \) and \( k_2 \) are the main-sequence star’s momentum of inertia and the apsidal motion constant, respectively (see Zahn 1977 for more details). The quantities \( k_2 \) and \( I_2/M_2 R_2^2 \) are both of the same order (see, e.g., Motz 1952), and therefore should approximately cancel each other, so a stellar system with \( P_{\text{orb}} = 6.67 \text{ hr} \) and \( M_1 = 1.4 M_\odot \), and \( M_2 = 0.4 M_\odot \) would take \( t_{\text{sync}} \approx 10^2 \text{ yr} \) to synchronize and \( t_{\text{circ}} \approx 2 \times 10^3 \text{ yr} \) to circularize. The companion star is thus expected to be synchronous, but the orbit may still be eccentric.

Although the eccentricity may contribute to modulate the X-ray pulsations of 1E, its introduction at the present level of discussion would complicate the model without any real benefit. As it will be discussed in § 5, our results are not expected to be too sensitive to the use of a circular orbit approximation. For this reason, we suppose that the orbit is circular, with the two stars separated by the fixed distance \( a \). If the NS’s period \( P_1 = 6.67 \text{ hr} \) is equal (or close) to the orbital period \( P_{\text{orb}} \), a is given by Kepler’s third law,

\[
a = 1.5 \times 10^{11} \left( \frac{P_{\text{orb}}}{6.67 \text{ hr}} \right)^{2/3} \left( \frac{M}{1.8 M_\odot} \right)^{1/3} \text{ cm},
\]

where \( M = M_1 + M_2 \) is the binary’s total mass.

2. SETTING UP THE MODEL

We assume that 1E is a binary system made up of a neutron star (NS) of mass \( M_1 = 1.4 M_\odot \) and a low main-sequence star (a red dwarf, RD) of mass \( M_2 = 0.4 M_\odot \) (consistent with the current optical/infrared constraints). This system survived the supernova explosion which led to the birth of RCW 103. The orbit is expected to be born quite eccentric, with e most likely in the range 0.2–0.5 (De Luca et al. 2006); this eccentricity deserves some comments. In a narrow system like that we envisage for 1E, the tidal torques are expected to be quite effective in synchronizing the companion star’s rotation with the orbital period and in circularizing the orbit. Unfortunately, there is no general consensus about the efficiency of these two processes (see, e.g., Zahn 1977; Tassoul 1988; Rieutord & Zahn 1997; Tassoul & Tassoul 1997; see also Meibom & Mathieu 2005 and references therein). The most conservative estimates of the synchronization and circularization timescales are those derived by Zahn (1977; his eqs. [4.12] and [4.13])
the well-known formula $L_X = GM_1 \dot{M}_1/R_1$, where $R_1 = 10^6$ cm is the NS radius. It is easy to check that the requirement $R_A \gtrsim \alpha \approx 10^{11}$ cm implies a NS magnetic moment $\mu_1 \approx 10^{33} - 10^{34}$ G cm$^3$, or a magnetar-like surface magnetic field $B_1 \approx \mu_1/R_1^3 \approx 10^{15} - 10^{16}$ G.

The Roche lobe of the NS is entirely filled by the star’s own magnetosphere, and under this circumstance the formation of an accretion disk is impossible: any mass flow from the RD through L1 is immediately channeled along the NS’s field lines and accreted. Independently of the accretion pattern, the balance between the material accretion spin up torque and the magnetic spin down propeller torque fixes the equilibrium condition $R_{eq} \sim R_A$ between the Alfven radius and the corotation radius,

$$R_{eq} = \left(\frac{GM_1}{\omega_1^2}\right)^{1/3},$$

where $\omega_1 = 2\pi/P_1$ is the NS rotation angular frequency. The rotation equilibrium condition $R_{eq} \sim R_X$ sets the NS final spin period

$$P_1 \rightarrow P_{eq} \approx 6 \left(\frac{\mu_1}{10^{33} \text{ G cm}^3}\right)^{6/7} \left(\frac{|M_2|}{3 \times 10^{13} \text{ g s}^{-1}}\right)^{-3/7} \left(\frac{M_1}{1.4 M_\odot}\right)^{-5/7} \text{ hr.} \quad (6)$$

In a Polar system (and also in 1E) the RD is endowed with a large magnetic moment. This may be either induced by the compact star or intrinsic. This last alternative is quite possible, as it is known that M dwarf stars may harbor sizeable magnetic fields, even of a few kG (see, e.g., Donati et al. 2006), and magnetic moments above $10^{34}$ G cm$^3$. In this paper we assume for the RD the magnetic moment $\mu_2 = 10^{34}$ G cm$^3$.

The X-ray emission from Polars is powered by mass accretion from the secondary. In our case, this may still be true, but some important contributions may come from other sources as well. Currently known magnetar candidates (anomalous X–ray pulsars [AXPs] and soft gamma-ray repeaters [SGRs]; see Woods & Thompson 2006 for a review) are observed to emit X–ray radiation with luminosities ranging from $\sim 10^{33}$ to $\sim 10^{36}$ erg s$^{-1}$. This emission is believed to be powered by the decay of their ultrastrong magnetic field; this mechanism might be also responsible for the $\sim 10^{35} - 10^{36}$ erg s$^{-1}$ X-ray luminosity of 1E. In addition, the spectrum and the long-term variability of 1E are remarkably similar to the characteristics of some AXPs and SGRs (for example, the AXP XTE J1810–197; Gotthelf & Halpern 2005). Therefore, the X-ray emission of 1E might be either induced or partially explained by typical magnetar emission rather than by accretion, as the main role of the companion star invoked in this paper is to explain the discrepancy between the 6.67 hr pulsation period of 1E and the much shorter periods $5 - 12$ s observed in AXPs and SGRs. If this is the case, the X-ray luminosity $L_X$ only provides an upper limit to the mass accretion rate $M_1$ on the NS: in the low state of 1E, for instance, $M_1 \lesssim 10^{13}$ g s$^{-1}$. Of course, the companion star mass transfer to the compact object only provided that it is the same size as the Roche lobe. According to the similarity relation valid for low ZAMS stars (e.g., Kippenhahn & Weigert 1994), the radius of the companion star is

$$R_2 = 3.3 \times 10^{10} \left(\frac{M_2}{0.4 M_\odot}\right)^{0.8} \text{ cm}, \quad (7)$$

and is comparable to the star’s Roche lobe radius $R_L \sim 4 \times 10^{10}$ cm, as computed from equation (2) of Eggleton (1983). Since $R_2$ is close to $R_L$, in view of the uncertainties, we need to explore both the possibility that the system is detached ($R_2 < R_L$) and the possibility that there is a Roche lobe overflow from the RD to the NS ($R_2 \gtrsim R_L$). The geometry of the system (i.e., detached or not) plays a fundamental role in the NS spin evolution. Indeed, Polars may reach exact synchronism provided that the magnetic torque is stronger than the accretion spin-up torque $\propto M_1$ (e.g., Campbell 1986). If most of the X-ray luminosity of 1E is powered by magnetar emission and not by accretion, $M_1$ will be small (the binary may be detached), so the magnetic torques are dominant and may synchronize the system. Note that in order to attain synchronism the magnetic torque must not (only) be stronger the accretion torque $\text{now}$, but it must also have been so over a significant fraction if the system’s life. Incidentally, we note that if deeper observations show the mass of the companion to be significantly smaller than its current upper limit ($0.4 M_\odot$), then the companion is most likely to lie well within its own Roche lobe. In this case the system is detached, and the X-ray luminosity is clearly due to the magnetar emission. On the other hand, if the accretion torque is strong enough, it may keep the NS away from exact synchronism. In this case 1E should be more similar to an intermediate Polar (IP). The spin evolution of these systems is dominated by the material torques, and their action leads to rotation equilibrium periods distributed over a fairly wide range of values, all shorter than the orbital period (Norton et al. 2004).

In the next section we set up all the formal machinery to deal with both these cases, aiming to follow the spin evolution of 1E.

### 3. DYNAMICS OF THE SYSTEM

In this section we focus on the evolution of the NS spin and the orbital period after the supernova explosion. In order to keep our analysis as simple as possible, we assume that all the torques are directed perpendicular to the orbital plane. Any misalignment would lead to a rotation of the orbital angular momentum; this just complicates the equations, without any relevant benefit for the physical insight into the problem.

We evaluate the torques acting on the NS, as well as those acting on the orbit and the RD, as they may redistribute angular momentum among the different components.

#### 3.1. Torques Acting on the NS

The torques acting on the NS are the following.

1. **The spin-down torque due to magneto-dipole losses** (Lipunov 1992),

$$N_{md} = -\frac{2}{7} \frac{\mu^2 \omega^3}{c^5} \sin^2 \chi \hat{k}_1, \quad (8)$$

where $\hat{k}_1$ is the direction of the NS spin, $\chi$ is the angle between the magnetic and the spin axes, and $\mu_1 = B_1 R_1^3$ is the NS magnetic moment’s modulus. In the following we assume $\hat{k}_1$ coincident with the unit vector $\hat{k}$ perpendicular to the orbital plane, and $\chi = \pi/2$, i.e., the NS’s magnetic moment $\mu_1$ lies in the orbital plane.

We may assign a characteristic timescale to the torque,

$$t_{md} = \frac{\omega_1 I_1}{N_{md}} \approx 3 \times 10^{-3} \left(\frac{I_1}{10^{45} \text{ g cm}^2}\right) \times \left(\frac{\mu_1}{10^{33} \text{ G cm}^3}\right)^{-2} \left(\frac{P_1}{10 \text{ ms}}\right)^2 \text{ yr}, \quad (9)$$

where $P_1$ and $I_1$ are NS spin period and moment of inertia, respectively. On account of the strong dependence $t_{md} \propto P_1^2$, this
torque is relevant in the early NS life, when $P_1$ is small. It is fundamental in braking the initially fast NS rotation, but it soon becomes negligible and does not affect the final equilibrium rotation period. In the present analysis we assume a constant magnetic field, since a possible field decay does not affect our results, at least qualitatively.

2. The torque due to the interaction between the magnetic moments of the NS and of the secondary star.—The torque acting on the NS due to this interaction is (see, e.g., Jackson 1975)

$$N_{dd} = \mu_1 \times B_2(1),$$  \hspace{1cm} (10)

where $B_2(1)$ is the RD magnetic field in correspondence of the NS; this can be written in terms of the dwarf’s magnetic moment $\mu_2$ as (Jackson 1975)

$$B_2(1) = \frac{3(\mu_2 \cdot \hat{n}) \hat{n} - \mu_2}{a^3},$$  \hspace{1cm} (11)

where $\hat{n}$ is the unit vector pointing from the RD to the NS. We fix a nonrotating frame of reference centered on the NS, whose $xy$ plane coincides with the orbital plane. The torque (eq. [10]) also works on the orbital angular momentum to make $\mu_1$, $\mu_2$, and $\hat{n}$ coplanar (King et al. 1990). We therefore assume that these vectors all lie in the orbital plane $xy$, so

$$\hat{n} = \begin{bmatrix} \cos \vartheta \\ \sin \vartheta \\ 0 \end{bmatrix},$$

$$\frac{\mu_1}{\mu_1} = \begin{bmatrix} \cos \phi_1 \\ \sin \phi_1 \\ 0 \end{bmatrix}, \quad \frac{\mu_2}{\mu_2} = \begin{bmatrix} \cos \phi_2 \\ \sin \phi_2 \\ 0 \end{bmatrix},$$  \hspace{1cm} (12)

where $\vartheta$ is the orbital true anomaly, and $\phi_1$ and $\phi_2$ are the inclinations of the NS’s and RD’s magnetic moments with respect to the $x$-axis. After some algebra, the torque equation (10) reads

$$N_{dd} = -\frac{\mu_1 \mu_2}{a^3} \left[ 2 \cos (\phi_2 - \vartheta) \sin (\phi_1 - \vartheta) + \sin (\phi_2 - \vartheta) \cos (\phi_1 - \vartheta) \right] \hat{k}.$$  \hspace{1cm} (13)

As a final remark concerning $N_{dd}$, we observe that it does not contribute to lock the period of the NS to the orbital value, but rather its phase angle $\phi_1$ (see also the discussion in Appendix A). Phase locking is a stronger condition than simple period locking. In the first case, the spin period and the orbital period coincide, but $\phi_1$ may have any value. In the second case the orientation of $\phi_1$ is fixed by the dynamics.

Omitting the geometrical factor, we define the characteristic timescale of the dipole-dipole torque given by equation (13) as

$$t_{dd} = \frac{\omega_1 I_1}{|N_{dd}|} \simeq 7 \times 10^6 \left( \frac{I_1}{10^{45} \text{ g cm}^2} \right) \left( \frac{P_1}{10 \text{ ms}} \right)^{-1} \times \left( \frac{\mu_1 \mu_2}{10^{67} \text{ G}^2 \text{ cm}^6} \right)^{-1} \left( \frac{P_{\text{orb}}}{6.67 \text{ hr}} \right)^2 \text{ yr},$$  \hspace{1cm} (14)

This torque is weak when the NS rotates fast, but when its period is in the range of hours, the timescale (eq. [14]) drops to few years, and is therefore important for the final equilibrium.

3. A dissipative torque acting to synchronize the NS spin with the orbital period.—The nature and strength of this torque have been investigated by a host of authors, but there is still no general consensus about it. In their pioneering paper, Joss et al. (1979) considered that the asynchronous orbital evolution of the secondary in the primary’s magnetosphere would induce strong electrical currents in the secondary’s atmosphere, which are dissipated by Ohmic decay, leading the compact star to rotate synchronously with the orbit. In a series of papers Campbell (1983, 1984, 1999, 2005) considered the role of turbulence in the secondary star, and concluded that the turbulent dissipation would be much more effective than Ohmic decay in dissipating the currents (but see Lamb 1985). Yet this torque becomes quite inefficient if $\omega_1 \sim \Omega_{\text{orb}}$ (where $\Omega_{\text{orb}} = 2\pi / P_{\text{orb}}$ is the orbital frequency), and therefore the compact star’s initial spin period must not be too different from the orbital period to reach synchronism. Other authors suggest different mechanisms for the dissipative torque, such as unipolar induction (Channugam & Dulk 1983; Wu et al. 2002) or MHD torques (Lamb et al. 1983; Kaburaki 1986). Despite their great differences, most models write the dissipative torque as

$$N_{\text{diss}} = f \frac{\mu_1 \mu_2}{a^3},$$  \hspace{1cm} (15)

with a typical timescale

$$t_{\text{diss}} = (\omega_1 - \Omega_{\text{orb}}) I_1 / N_{\text{diss}}$$  \hspace{1cm} (16)

necessary to synchronize the primary; here $f \simeq 1$ is a model dependent dimensionless factor. The estimates of $t_{\text{diss}}$ given by the different models disagree by orders of magnitude; typical values range from about $10^3$ s (Kaburaki 1986; Channugam & Dulk 1983; Wu et al. 2002) to more than $10^6$ yr (Campbell 1983). In view of the large uncertainty about $N_{\text{diss}}$, we have chosen to parameterize it. We would like to single out the dependence of this torque on the asynchronism between the NS’s rotation and the orbital period, so we write

$$N_{\text{diss}} = -I_1 \frac{\omega_1 - \Omega_{\text{orb}}}{\tau_{\text{syn}}} \hat{k},$$  \hspace{1cm} (17)

where the synchronization timescale $\tau_{\text{syn}}$ is a free parameter of our model. We will calculate the evolution of the NS’s period for several values of $\tau_{\text{syn}}$. An important caveat is in order here. The time $t_{\text{diss}}$ strongly depends on the orbital separation $a$ (see eq. [15]), so the dissipation torque is inefficient in wide binaries. Essentially all the models for Polar CVs call for the condition $R_A \gtrsim a$ for synchronism. As the strong dependency on $a$ is not explicit in the parameter $\tau_{\text{syn}}$, we must check for consistency that in our synchronous solutions the condition $R_A > a$ is fulfilled.

If the dissipation torque is inefficient (i.e., $\tau_{\text{syn}}$ exceeds the age of RCW 103), the magnetic torques are unable to synchronize the NS, contrary to what happens in a Polar CV. In this case, 1E would be more similar to an IP system, governed by material torques, rather than a Polar-like system ruled by purely magnetic interactions.

These times are meant as scaled from a white dwarf to a magnetar.
4. The accretion and propeller torques.—In its youth the NS spins too fast to accrete matter; this is the so-called “propeller” regime (see e.g., Romanova et al. 2003). The mass coming from the companion star flies around the NS’s magnetosphere, extracts angular momentum, and spins the NS down. The exact form of the torque in the propeller stage is not well understood; we adopt the formula suggested by Lipunov (1992, p. 159),

\[ N_p = -k_t \frac{\mu_1^2}{R_{\text{co}}^2} \dot{\kappa}, \tag{18} \]

where \( k_t \) is a constant to be determined later.

This torque is very efficient, as shown by its short characteristic timescale

\[ t_p = \frac{\omega I_1}{|N_p|} \approx 3 \times 10^2 k_t^{-1} \left( \frac{I_1}{10^{45} \text{ g cm}^2} \right) \left( \frac{\mu_1}{10^{33} \text{ g cm}^3} \right)^{-2} \left( \frac{M_1}{1.4 M_\odot} \right) \left( \frac{P_1}{10 \text{ ms}} \right) \text{s}. \tag{19} \]

The accretion torque depends on the mass accretion rate \( \dot{M}_1 \) on the primary, as well as on the arm’s length of the material stress. In the case of a Polar, the flow from the secondary is directly accreted as soon as it crosses the inner Lagrangian point L1: if the orbit is only moderately eccentric, we can adopt the formula (Plavec & Kratochvíl 1964)

\[ b_1 = a(0.500 - 0.227 \log_{10} q) \approx 9.5 \times 10^{10} \text{ cm} \tag{20} \]

to approximate the distance \( b_1 \) between the inner Lagrangian point L1 and the NS center. The torque’s arm \( b_1 \) yields

\[ N_a = \dot{M}_1 (GMb_1)^{1/2} \hat{k}, \tag{21} \]

whose associated characteristic timescale is (up to factors of order unity)

\[ t_a \approx 1.7 \times 10^2 \left( \frac{I_1}{10^{45} \text{ g cm}^2} \right) \left( \frac{\dot{M}_1}{10^{13} \text{ g s}^{-1}} \right)^{-1} \left( \frac{P_{\text{orb}}}{6.67 \text{ hr}} \right)^{1/3} \left( \frac{P_1}{6.67 \text{ hr}} \right)^{-1} \text{ yr}. \tag{22} \]

We describe both the accretion and the propeller regimes with a unique accretion/propeller torque,

\[ N_{ap} = N_p + N_a. \tag{23} \]

We write \( \dot{M}_1 \) as a function of the mass transfer rate \( \dot{M}_2 \) from the secondary. We postulate the relation

\[ \dot{M}_1 = \alpha |\dot{M}_2|, \tag{24} \]

where \( \alpha \approx 0 \) if \( N_a \ll N_p \) and \( \alpha \approx 1 \) if \( N_a \gg N_p \). Since the accretion torque is dominant if \( b_1 \ll R_{\text{co}} \) and the opposite inequality is true if the propeller is dominant, we choose

\[ \alpha = \frac{R_{\text{co}}^3}{b_1^3 + R_{\text{co}}^3}, \tag{25} \]

which can be rewritten

\[ \alpha = \frac{1}{1 + \hat{\lambda}(\omega_1/\Omega_{\text{orb}})^2}, \tag{26} \]

where

\[ \hat{\lambda} = \lambda(q) \equiv (b_1/a)^3(1 + q) \tag{27} \]

is a function of the mass ratio \( q \) only (cf. eq. [20]).

We determine the constant \( k_t \) so that \( N_{ap} = 0 \) when \( R_{\text{co}} = b_1 \), i.e., when the centrifugal barrier opens. After some algebra, we retrieve the formula

\[ N_{ap} = |\dot{M}_2| (GMb_1)^{1/2} \left[ \frac{1}{1 + \hat{\lambda}(\omega_1/\Omega_{\text{orb}})^2} \right] \frac{\omega_1}{2 \Omega_{\text{orb}}} \hat{k}. \tag{28} \]

In the accretion limit \( \omega_1 \ll \hat{\lambda}^{-1/2} \Omega_{\text{orb}} \), the propeller torque is small, and this expression reduces to that of pure accretion (eq. [21]). On the other hand, in the propeller regime \( \omega_1 \gg \hat{\lambda}^{-1/2} \Omega_{\text{orb}} \) the first term in square brackets is negligible, and the torque reduces to pure propeller \( N_{ap} \propto -\omega_1^2 \), as in equation (18). Incidentally, we note that the NS spin equilibrium period \( P_{\text{eq}} = \lambda^{-1/2} P_{\text{orb}} \), defined by \( N_{ap} = 0 \) is the Keplerian period at L1 of a mass orbiting the NS, and coincides with that predicted by the diamagnetic blobs accretion model put forward by King (1993) and Wynn & King (1995) for highly magnetized CVs.

The torques analyzed in this section are the building blocks for the overall torque acting on the NS: this will be made explicit in § 4.

3.2. Torques Acting on the Secondary Star

The reason we also consider the torques on the secondary star is that they may couple to the orbit to alter \( P_{\text{orb}} \), thus affecting synchronism (see also King et al. 1990 for a fuller discussion). In general, the torques on the secondary star are:

1. The magnetic braking torque.—If the RD is magnetic and blows a wind, this is coupled to the magnetic field and may extract angular momentum from the secondary. This torque only depends on the RD period and not on the orbital one, so magnetic braking does not directly affect the orbital period. However, it may couple with the tidal torque (discussed immediately below) to transfer angular momentum from the orbit to the RD.

2. The tidal torque.—This torque reduces the relative rotation of the RD and the orbit, and of course vanishes if the secondary rotates synchronously. It may also redistribute angular momentum between the orbit and the RD. Stars with convective envelopes (which is most likely to be the case of the our companion star) in binary systems with orbital periods of a few hours have tidal synchronization times \( \approx 10^7 \text{ yr} \) (eq. [1]), so we shall assume that the RD is always locked to the orbital period. The magnetic braking extracts angular momentum from the secondary, which is then taken away from synchronism. The tidal torque, on the other hand, is very effective in locking the secondary back to orbital synchronism. As a consequence, the synergetic action of these torques pumps angular momentum from the orbit to the secondary, keeping it in synchronous rotation. (Verbunt & Zwaan 1981).

If the secondary is not synchronous, there are two effects. First, the dipole-dipole torque (eq. [13]) keeps oscillating from positive to negative, so it cannot lock the NS’s phase. Second, the secondary can exchange some of its intrinsic angular momentum with the orbit. In view of our estimates (§ 3.3), we do expect that the amount of transferred angular momentum is too small to affect
orbit. Summarizing, the secondary’s possible lack of synchronism does not directly alter the primary’s period locking, but may invalidate its phase locking.

3. The dipole-dipole torque.—This is exactly the same kind of torque acting on the primary; its expression is readily derived by swapping \( \phi_1 \) and \( \phi_2 \) in equation (13):

\[
N'_{dd} = -\frac{\mu_1 \mu_2}{a^3} \left[ 2 \cos(\phi_1 - \theta) \sin(\phi_2 - \theta) + \sin(\phi_1 - \theta) \cos(\phi_2 - \theta) \right] \frac{\dot{J}}{M}.
\]

This torque depends on the true anomaly as well as on the stars’ magnetic moments, and is able to exchange angular momentum between the secondary and the orbit. The synchronization timescale for this torque is

\[
t'_{dd} = \frac{\omega \dot{J}_{dd}^2}{|N'_{dd}|} \approx 3 \times 10^3 \left( \frac{I_2}{10^{54} \text{ g cm}^2} \right) \left( \frac{P_2}{6.67 \text{ hr}} \right)^{-1} \times \left( \frac{\mu_1 \mu_2}{10^{67} \text{ G}^2 \text{ cm}^6} \right)^{-1} \left( \frac{P_{\text{orb}}}{6.67 \text{ hr}} \right)^2 \text{ yr},
\]

where \( P_2 \) is the period of secondary. Since \( t'_{dd} \) is much longer than the age of RCW 103, the RD phase is not likely to be locked, although its period is most likely synchronous with the orbit.

In summary, we assume that the secondary’s rotation is synchronous with the orbit. The instantaneous phase \( \phi_2 \) (i.e., the position of the RD magnetic moment \( \mu_2 \)) is

\[
\phi_2 = \Omega_{\text{orb}} t + \phi'_2,
\]

where the phase \( \phi'_2 \) is arbitrary due to the lack of phase locking.

3.3. Orbital Torques

In this section we derive and discuss the torques acting on the orbit of the binary system. Since some torques and the mass transfer alter the orbital period, they must be taken into account to study the synchronization. As usual, we suppose that the orbit is circular, so the orbital angular momentum reads

\[
J_{\text{orb}} = \left( \frac{G a}{M} \right)^{1/2} M_1 M_2 \frac{\dot{J}}{M}.
\]

The rate of change of \( J_{\text{orb}} \) is due to the action of all the stellar torques containing the true anomaly \( \theta \) (or the orbital frequency \( \Omega_{\text{orb}} \)), plus the tidal/magnetic braking torque \( N_{\text{mag}} \) due to the angular momentum exchange with the secondary discussed in § 3.2:

\[
\dot{J}_{\text{orb}} = -N_{dd} - N'_{dd} - N_{sp} - N_{\text{diss}} - N_{\text{mag}}.
\]

Gravitational radiation has not been included; it is generally important for systems with shorter periods (\( P_{\text{orb}} \approx 2 \) hr), and it is immaterial here.

The torque \( N_{\text{mag}} \) is (Verbunt & Zwaan 1981)

\[
N_{\text{mag}} \approx 0.5 \times 10^{-28} I_2 R_2^2 \Omega_{\text{orb}}^3
\]

(in cgs units). This torque drags orbital angular momentum on the timescale

\[
t_{\text{mag}} = \frac{J_{\text{orb}}}{N_{\text{mag}}} \approx 10^8 \left( \frac{m}{0.3 \, M_\odot} \right) \left( \frac{P_{\text{orb}}}{6.67 \text{ hr}} \right)^{10/3} \times \left( \frac{I_2}{10^{54} \text{ g cm}^2} \right)^{-1} \left( \frac{R_2}{3.3 \times 10^{10} \text{ cm}} \right)^{-2} \text{ yr},
\]

where \( m \) is the reduced mass of the binary system. Since this timescale is very long, the orbital angular momentum yields to the secondary may be neglected.

Ignoring the exchanges of orbital angular momentum with the RD, we focus on the exchanges with the rapidly spinning NS. Adopting the parameterization given by equation (24) for the mass accretion rate, the logarithmic derivative of the modulus of equation (32) gives

\[
\frac{J_{\text{orb}}}{\dot{J}_{\text{orb}}} = 1 - \frac{\dot{M}_2}{M_2} - \frac{2(1 - \alpha q^2) + q(1 - \alpha)}{2(1 + q)}.
\]

Plugging the expressions of the relevant torques into equation (33), and using the Kepler relation \( a \propto \Omega_{\text{orb}}^{-2/3} \), we retrieve our equation for the evolution of the orbital frequency,

\[
\frac{\Omega_{\text{orb}}}{\dot{\Omega}_{\text{orb}}} = \frac{3 I_1}{m a^2 \Omega_{\text{orb}}} \left\{ 3 \frac{\mu_1 \mu_2}{I_1 a^2} \sin(\phi_1 + \phi_2 - 2\theta) \right. \left. + \frac{\omega - \Omega_{\text{orb}}}{\tau_{\text{syn}}} - \frac{\dot{M}_2 (G M b_1)^{1/2}}{I_1} \right\}
\]

\[
\times \left[ \frac{1}{1 + \lambda (\omega / \Omega_{\text{orb}})^2} - \frac{1}{2} \left( \frac{\omega}{\Omega_{\text{orb}}} \right)^2 \right] \left\{ 3 \frac{\dot{M}_2}{M_2} \frac{2(1 - \alpha q^2) + q(1 - \alpha)}{2(1 + q)} \right. \left. + 3 \frac{\dot{M}_2}{M_2} \frac{2(1 - \alpha q^2) + q(1 - \alpha)}{2(1 + q)} \right\}.
\]

We briefly discuss the magnitude of the terms on the right-hand side of equation (37). The terms in curly braces are the same dipole-dipole, dissipative, and accretion-propeller torques (with allowance of signs and geometrical factors) acting on the NS; all these terms are multiplied by the small factor \( \zeta = 3 I_1 / m a^2 \Omega_{\text{orb}} \approx 10^{-6} \), showing that the timescales for the orbital evolution are longer than the timescales for the NS’ spin evolution by a factor \( \zeta^{-1} \approx 10^6 \).

Mass transfer unbalances the relative masses of the two components, and only becomes significant over the long timescale

\[
t_2 \approx \frac{M_2}{M_2} = 2.5 \times 10^{12} \left( \frac{M_2}{0.4 \, M_\odot} \right) \left( \frac{\dot{M}_2}{10^{-13} \text{ g s}^{-1}} \right)^{-1} \text{ yr}.
\]

In summary, the orbital torques are characterized by timescales much longer than the current age of RCW 103; the orbit may be then be regarded as fixed, and in the following we only consider the torques on the NS. The results presented in the next section are based on this hypothesis; however, we had an additional check of its validity by running a numerical calculation where the orbital parameters were allowed to evolve according to equation (37).

4. EVOLUTION OF THE NS ROTATION

We adopt \( \phi_1 \) (the phase of the NS magnetic moment) as the NS Euler free rotation angle (e.g., Landau & Lifshitz 1969).
Putting all the torques of equations (8), (13), (17) and (28) together, the equation of motion for the NS’s phase angle \( \phi_1 \) reads

\[
\begin{align*}
\dot{\phi}_1 &= \omega_1, \\
I_1 \dot{\omega}_1 &= -\frac{\mu_1 \mu_2}{a^3} \left[ 2 \cos (\phi_2 - \vartheta) \sin (\phi_1 - \vartheta) + \sin (\phi_2 - \vartheta) \cos (\phi_1 - \vartheta) \right] - \frac{2}{3} \frac{\mu_1^3}{c^3} \omega_1^3 \\
&- I_1 \frac{\omega_1 - \Omega_{\text{orb}}}{\tau_{\text{syn}}} + \left[ \frac{1}{1 + \lambda (\omega_1 / \Omega_{\text{orb}})^2} \frac{2}{3} \left( \frac{\omega_1}{\Omega_{\text{orb}}} \right)^2 \right].
\end{align*}
\]

where the true anomaly \( \vartheta = \Omega_{\text{orb}} t \) is assumed to be zero at \( t = 0 \).

As noted above, we keep the orbital parameters constant, as their variation due to the angular momentum exchanged with the stars is negligible.

As shown in Appendix A, equation (39) admits an asymptotic constant solution, i.e., a stable NS equilibrium spin period. In the absence of material torques \( (M_2 = 0) \), this period is attained thanks to the second and the third terms in equation (39), leading to full synchronism with the orbit. The magneto-dipole torque brakes the initial fast NS rotation, and later on the dissipative torque brings the NS spin to orbital synchronism. In the presence of material torques, it is the interplay between propeller and accretion torques that drives the young magnetar into equilibrium (as implied by eq. [6]). In this case, the orbital period does not coincide with the NS’s spin equilibrium \( P_{\text{eq}} \), but it is quite close to it \( (P_{\text{eq}} \simeq 0.5-0.7 \, P_{\text{orb}}) \). Even in this case the NS magnetosphere keeps filling the binary Roche lobe, warranting the consistency of our results.

We integrated equations (39) with a step-adaptive fourth-order Rosenbrock algorithm (Press et al. 2002). This choice is motivated by the possible stiffness of the problem due the diverse timescales of the torques. The NS and RD magnetic moments are the free parameters of the model, together with \( M_2 \) and \( \tau_{\text{syn}} \). The magnetic moments of the NS and the RD are \( \mu_1 = 10^{33} \, \text{G cm}^3 \) and \( \mu_2 = 10^{34} \, \text{G cm}^3 \), respectively. The NS’s moment of inertia has the standard value of \( I_1 = 10^{45} \, \text{g cm}^2 \). We let the mass-transfer rate and the synchronization time vary in each run, in order to explore different evolution paths. The initial phases \( \phi_1(t = 0) \) and \( \phi_2(t = 0) \) have been randomly chosen, with an initial NS rotation period of 3 ms. The RD has been always considered synchronous. In all our runs the NS final equilibrium spin period has been set equal to 6.67 hr (see Appendix A), and we have derived the orbital period in a self-consistent manner.

For consistency, we have also checked that the condition that \( R_A > a \) is always met in our runs. The NS Roche lobe is dominated by the star’s magnetic field, and the mass from the secondary is effectively channeled from L1 to the star’s polar caps without forming an accretion disk, consistent with the torque given by equation (28) used in our calculations.

It is convenient to first present our result in the case where the mass-transfer rate \( M_2 \) vanishes. The material (i.e., accretion and propeller) torques are always zero, and the evolution is controlled uniquely by the magnetic torques. Figure 1 plots the evolution of the NS period for different values of the synchronization time, \( \tau_{\text{syn}} = 10, 100, \) and 1000 yr.

The dissipative magnetic torque in this case plays a key role in driving the NS spin into orbital synchronization, attained within \(~2000 \, \text{yr}\) if \( \tau_{\text{syn}} \) is not longer than few hundreds years. A magnetar-like field and a close orbit are necessary, and this is consistent with estimates given in the literature, possibly with the exception of Campbell (1983, 2005; see also the discussion in § 3.1).

The situation changes if the mass-transfer rate \( \dot{M}_2 \) is nonnegligible. Figure 2 plots the evolution of \( P_1 \) and of the synodic normalized frequency

\[
\xi = (\omega_1 - \Omega_{\text{orb}})/\Omega_{\text{orb}},
\]

for \( |M_2| = 10^{13} \, \text{g s}^{-1} \). Here the material torques are relevant, and lead to the equilibrium spin period within few hundreds years, fairly independent of the exact value of \( \tau_{\text{syn}} \). The orbital period is longer than the NS’s equilibrium spin period; we find \( P_{\text{orb}} = 6.9 \, \text{hr} \) if \( \tau_{\text{syn}} = 10 \, \text{yr} \), \( P_{\text{orb}} = 8.7 \, \text{hr} \) if \( \tau_{\text{syn}} = 100 \, \text{yr} \), and \( P_{\text{orb}} = 11.5 \, \text{hr} \) for \( \tau_{\text{syn}} = 1000 \, \text{yr} \). An additional run with \( \tau_{\text{syn}} = 10^7 \, \text{yr} \) shows that the magnetar attains its spin equilibrium within \(~10^3 \, \text{yr}\), with a somewhat longer orbital period of \(~12 \, \text{hr}\) so that in all cases we find \( P_{\text{spin}}/P_{\text{orb}} \approx 0.5 \). In magnetic CVs period ratios larger than 0.5 correspond to stream-fed accretion onto the white dwarf (Norton et al. 2004). Consistently, we also find here that the condition \( R_A > a \) is always met; the NS’s Roche lobe is filled by the magnetosphere, and no accretion disk can form. In addition, for \( \tau_{\text{syn}} < 1000 \, \text{yr} \), the stronger condition \( R_A > a \) is also met; this warrants that this (parametric) synchronization time is consistent with the requirement of a strong magnetic interaction between the stars (§ 3.1).

In Figure 2 we also show the evolution of \( P_1 \) and \( \xi \) for the much larger mass transfer rate \( |M_2| = 10^{15} \, \text{g s}^{-1} \), corresponding to the peak X-ray luminosity observed in 1E. This value should provide an upper limit to the influence of the mass accretion rate on the period evolution. For this value of \( |M_2| \), the condition \( R_A > b_1 \) is satisfied, but since \( R_A \) is still larger than the circularization radius \( R_{\text{circ}} \) of the accreting flow, the transferred mass is unable to circularize and form a disk (see, e.g., Frank et al. 2002). For such a strong material torque, the equilibrium period is attained in less than \(~10 \, \text{yr} \) during the early “propeller” phase, which extracts a great amount of angular momentum from the spinning young magnetar.

5. DISCUSSION

In the previous section we have shown that if 1E hosts a highly magnetic NS orbiting a RD, the NS is spun down rapidly to an
MAGNETAR IN A YOUNG LOW-MASS BINARY?

1E is more similar to a young accreting low-mass binary. The orbital period is 6–12 hr, then \( R_A > b_1 \), and the flow from the RD provides a material torque with arm of length \( b_1 \), which brings the NS to the observed equilibrium spin period \( P_{eq} = 6.67 \) hr within \( \sim 100 \) yr, independent of the magnetic torques. The orbital period in this case still depends on \( \tau_{syn} \), and is longer than \( P_{eq} \). In the limit of a very inefficient dissipative torque we find \( P_{orb} \sim 2 P_{eq} \) (i.e., \( P_{orb} \approx 12 \) hr).

It is worth spending a few words here on the upper limits of the orbital period. If the companion star tightly follows the mass-radius relation valid for low main-sequence stars, then the orbital period is unlikely to much exceed \( P_{orb} \approx 12 \) hr, since this binary separation brings the NS to its spin equilibrium period through the sole action of the material torques. These imply a Roche-lobe-filling donor star, which in our calculations we assumed to be of \( 0.4 M_\odot \), close to the upper limit set by the optical/IR observations. Were the system wider than this, the star should fall within its Roche lobe, making mass transfer impossible, as well as the action of the material torques. On the other hand, mass transfer is possible if the companion star is larger than a main-sequence star of the same size, e.g., because it absorbed energy from the early relativistic wind emitted by the NS. In this case the limits on the orbit’s size are more uncertain, but by analogy with the period distribution of bona fide Polars and IPs, we may estimate that the orbital period is unlikely to much exceed \( P_{orb} \approx 12 \) hr. The current observations, on the other hand, are much less constraining. Our X-ray data do not yield any evidence for a periodic modulation larger than 6.67 hr. In any case, it would be quite difficult to detect a modulation at a few percent level and a long periodicity (10–15 yr) superimposed on the very strong pulsation at 6.67 hr.

A remark about our use of the circular approximation. As can be seen from equation (39), which rules the secular evolution of the NS’s spin, the circular approximation affects (1) the distance between the components in the dipole-dipole interaction, (2) the orbital frequency \( \Omega_{orb} \) in the dissipation term and the accretion/propeller torque, and (3) the arm \( b_1 \) (the distance between the neutron star and the inner Lagrangian point), which is calculated in the Roche formalism. Note that (whatever the eccentricity), point (3) is immaterial if the RD fits within its Roche lobe. In this case, there is no mass transfer, and the accretion/propeller torque vanishes; if the companion is large enough to allow mass accretion, the arm \( b_1 \) only enters as a weak power in the torque, and a possible deviation from the circular geometry is unlikely to affect this torque to a significant amount. As far as points (1) and (2) are concerned, the introduction of a nonzero eccentricity would induce an orbital modulation in the dipole-dipole and in the dissipation torque. This modulation is not able to alter the qualitative results presented in this paper.

In this model, 1E would be the first young low-mass binary observed in X-rays in the first \( \sim 2000 \) yr of its evolution, and there might be a connection between the magnetar nature of the NS and this premature onset of an X-ray phase if it comes from accretion. Two questions arise:

1. Does the pulsar wind of the young magnetar induce a mass loss from the RD? Can this wind vaporize the RD? How does this effect constrain the initial secondary’s mass, or the NS initial period?

2. Looking ahead, what is the subsequent evolution of 1E?

As far as the first question is concerned, the strong wind of photons and relativistic particles of the magnetar may perturb the RD thermal equilibrium, and may in principle cause its partial

![Graph](image-url)

**Fig. 2.**—Evolution of the NS’s rotation period (top) and of the normalized synodic frequency \( \xi = (\omega_1 - \Omega_{orb})/\Omega_{orb} \) (bottom). The two groups of lines refer to the mass-transfer rates \( |M_2| = 10^{17} \) and \( 10^{13} \) g s\(^{-1}\). The lines’ coding is the same as in Fig. 1. For \( |M_2| = 10^{17} \) g s\(^{-1}\) the orbital periods are \( P_{orb} = 6.9 \) hr (for \( \tau_{syn} = 10 \) yr), \( P_{orb} = 8.7 \) hr (for \( \tau_{syn} = 100 \) yr), and \( P_{orb} = 11.5 \) hr (for \( \tau_{syn} = 1000 \) yr). For \( |M_2| = 10^{13} \) g s\(^{-1}\) the orbital periods are \( P_{orb} = 11.4 \) hr (for \( \tau_{syn} = 10 \) yr), and \( P_{orb} = 11.9 \) hr (for both \( \tau_{syn} = 100 \) and 1000 yr). Note that for the mass-transfer rate \( |M_2| = 10^{17} \) g s\(^{-1}\), the condition \( R_A > b_1 \) is violated (see text for more details).

Asymptotic equilibrium period, by either magnetic and/or material torques. The magnetar-like NS loses memory of the spin acquired at birth, and approaches a spin period close to the binary orbital period.

If the material/propeller torques have been small for an appreciable fraction of the system’s age, then magnetic torques have been dominating since birth, and orbital synchronism is attained within 2000 yr if \( \tau_{syn} \lesssim 100 \) yr. The X-ray emission in this case should be dominated by mechanisms similar to those seen in magnetars (§ 2).

On the other hand, if the material torques have been important, 1E is more similar to a young accreting low-mass binary. The close resemblance of the light curve of 1E with Polars and intermediate Polar cataclysmic variables suggests that mass transfer currently takes place in 1E via filling of the RD’s Roche lobe. (Note that this does not exclude the action of magnetic torques in the earlier life of the source.) Although the magnetic torques are unable to synchronize the system, the hypothesis of a magnetar-like field is still essential here, since the NS magnetosphere must dominate the star’s Roche lobe. If \( B_1 \approx 10^{15} \) G and the orbital period is 6–12 hr, then \( R_A > b_1 \), and the flow from the RD provides a material torque with arm of length \( b_1 \), which brings the NS to the observed equilibrium spin period \( P_{eq} = 6.67 \) hr within \( \sim 100 \) yr, independent of the magnetic torques. The orbital period in this case still depends on \( \tau_{syn} \), and is longer than \( P_{eq} \). In the limit of a very inefficient dissipative torque we find \( P_{orb} \sim 2 P_{eq} \) (i.e., \( P_{orb} \approx 12 \) hr).

It is worth spending a few words here on the upper limits of the orbital period. If the companion star tightly follows the mass-radius relation valid for low main-sequence stars, then the orbital period is unlikely to much exceed \( P_{orb} \approx 12 \) hr, since this binary separation brings the NS to its spin equilibrium period through the sole action of the material torques. These imply a Roche-lobe-filling donor star, which in our calculations we assumed to be of \( 0.4 M_\odot \), close to the upper limit set by the optical/IR observations. Were the system wider than this, the star should fall within its Roche lobe, making mass transfer impossible, as well as the action of the material torques. On the other hand, mass transfer is possible if the companion star is larger than a main-sequence star of the same size, e.g., because it absorbed energy from the early relativistic wind emitted by the NS. In this case the limits on the orbit’s size are more uncertain, but by analogy with the period distribution of bona fide Polars and IPs, we may estimate that the orbital period is unlikely to much exceed \( P_{orb} \approx 12 \) hr. The current observations, on the other hand, are much less constraining. Our X-ray data do not yield any evidence for a periodic modulation larger than 6.67 hr. In any case, it would be quite difficult to detect a modulation at a few percent level and a long periodicity (10–15 yr) superimposed on the very strong pulsation at 6.67 hr.

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In this model, 1E would be the first young low-mass binary observed in X-rays in the first \( \sim 2000 \) yr of its evolution, and there might be a connection between the magnetar nature of the NS and this premature onset of an X-ray phase if it comes from accretion. Two questions arise:

1. Does the pulsar wind of the young magnetar induce a mass loss from the RD? Can this wind vaporize the RD? How does this effect constrain the initial secondary’s mass, or the NS initial period?

2. Looking ahead, what is the subsequent evolution of 1E?

As far as the first question is concerned, the strong wind of photons and relativistic particles of the magnetar may perturb the RD thermal equilibrium, and may in principle cause its partial
evaporation. The problem we address here differs from the evaporation of a main-sequence star by a millisecond pulsar in a narrow binary system (e.g., Ruderman et al. 1989a, 1989b; Tavani 1992). In these systems, the NS magnetic field is low ($10^8$–$10^9$ G), and its weak irradiation lasts for more than $10^8$ yr. In a millisecond pulsar environment, the main-sequence star has enough time to adjust to thermal equilibrium during the evaporation. In the present case this assumption is most likely invalid, since the irradiation from the magnetar is much more intense and short-lived (as shown below).

If the magnetar in 1E has a birth period $P_0 = 2\pi/\omega_0$ of a few milliseconds, it holds a rotational energy $E_{\text{rot}} \simeq 2 \times 10^{53} (P_0/3 \text{ ms})^{-2}$ erg, which is dissipated mainly through magnetic dipole losses in the early years, when the propeller torques are still ineffective. The RD has a binding energy $E_{\text{RD}} \simeq GM_2^2/R_2 \simeq 10^{48} (M_2/0.4 M_\odot)^2$ erg, so the RD needs to absorb only $\sim 0.1\% E_{\text{rot}}$ to be evaporated.

We can estimate the mass-loss rate from the RD by comparing the NS energy loss (and impinging on the RD) with the kinetic energy necessary to sustain a wind, regardless of the physical interaction process. We note that while the wind of particles from an ordinary millisecond pulsar is composed of energetic photons (X-rays or $\gamma$-rays) and TeV $e^\pm$ pairs (Ruderman et al. 1989a), in the case of a magnetar, the occurrence of high-energy pairs is debatable, since the strong magnetic field may prevent their formation due to photon splitting (Duncan & Thompson 1992). With this caveat in mind, we estimate how deep the impinging wind penetrates the companion star. Photons are stopped in the stellar atmosphere, having a mean free path of a few cm in a gas with density $\sim 1$ g cm$^{-3}$. The electrons impinging on the companion star yield their energy chiefly for inverse-Compton losses on the radiation field, and for relativistic bremsstrahlung. The inverse-Compton and bremsstrahlung stopping lengths are $l_{\text{IC}} \simeq 10^6$ cm and $l_{\text{BS}} \simeq 10^3$, respectively (Lang 1980, p. 467); since $l_{\text{BS}}$ and $l_{\text{IC}}$ are both small in comparison to the stellar radius, essentially all the energy carried by the relativistic wind is deposited in the star’s envelope.

In more detail, the spin-down luminosity of the magnetar decays with time as

$$L_{\text{SD}}(t) = \frac{1}{2} \frac{l_1 \omega_0^2}{\tau_{\text{md}}} \left(1 + \frac{t}{\tau_{\text{md}}} \right)^{-2},$$

(41)

where $\tau_{\text{md}}$ is the magneto-dipole decay time

$$\tau_{\text{md}} = \frac{3e^2 l_1}{4\mu_1^2 \omega_0^2} \simeq 4.6 \times 10^3 \left(\frac{\mu_1}{10^{33} \text{ G cm}^3} \right)^2 \left(\frac{P_0}{3 \text{ ms}} \right)^2.$$  

(42)

Equations (41) and (42) show that (1) the magnetar is braked to a period of 1 s after $\sim 16$ yr; (2) most of the rotational energy is dissipated in the first year; and (3) the spin-down luminosity decays by a factor of $\sim 10^{10}$ over this very short time. At $P \sim 10^2$ s, the magnetar in 1E brakes under the action of the magnetic and material torques, but at this period the bulk of its rotational energy has already been dissipated, and is but a tiny fraction of the companion’s binding energy. For this reason, any later mechanism that brakes the NS and which may (in principle) transfer energy from the NS to the companion, is most likely to have little effect on the latter.

We suppose that this relativistic wind is emitted isotropically, so a fraction $(R_2^2/2a)$ hits the companion star; a fraction $\eta$ of this impinging power extracts an evaporation wind from the RD. Following van den Heuvel & van Paradijs (1988), we equate the effective power of the relativistic wind with the evaporation wind to estimated $M_2$:

$$\eta \frac{R_2^2}{2a} L_{\text{SD}} = \frac{GM_2 dM_2}{R_2} \frac{dt}{dt},$$

(43)

where we have assumed that the evaporation wind speed is close to the escape velocity from the star’s atmosphere. We obtain the differential equation

$$- \frac{GM_2 dM_2}{R_2} \frac{dt}{dt} = \frac{1}{2} \frac{l_1 \omega_0^2}{\tau_{\text{md}}} \left(1 + \frac{t}{\tau_{\text{md}}} \right)^{-2}$$

(44)

for the evolution of the star’s mass $M_2(t)$, which can be solved once we have specified the relation between $M_2$ and the star’s radius $R_2$. This relation is uncertain, so we consider two limiting cases aiming to bracket the actual behavior.

1. On the one hand, we assume that $R_2$ is given, instant by instant, by the relation (7) valid for lower main-sequence stars. This is the most optimistic limit, since it assumes that the companion star has enough time to reduce its size as a reaction to the evaporating flux. This may not be very realistic, however. The star reacts to the relativistic wind on the Kelvin-Helmholtz time for the envelope involved in the energy deposition,

$$\tau_{\text{KH}} \simeq \frac{GM_2^2 l_{\text{BS}}}{R_2 L_2^2 R_2} \simeq 10^4 \text{ yr},$$

(45)

which is much longer than the magneto-dipole timescale given in equation (42); $L_2$ is the RD’s luminosity, and $l_{\text{BS}}$ is the penetration depth of the relativistic wind, where thermal equilibrium is perturbed (see note).

2. On the other hand, we may suppose that the star’s envelope swells up, filling its own Roche lobe. In this case, we assume for $R_2$ the size of the Roche lobe at the evaporation onset. This is the most pessimistic limit, since in this case the section presented by the star to the impinging flux is the largest, resulting in a fast evaporation rate.

Figure 3 shows the final mass of the companion star versus its initial mass in these two limiting cases for several values of the parameter $\eta$. The fixed parameters are the NS’s initial period (3 ms) and the orbital separation ($a = 1.5 \times 10^{11}$ cm), although an increase of $a$ is expected on account of the mass lost by the system.

In both cases, the secondary’s final mass is $M_2 \leq 0.4 M_\odot$, provided that $\eta < 10^{-1}$, for an initial mass $M_2 \leq 0.6 M_\odot$. For efficiencies $\eta \gtrsim 10^{-1}$ and initial masses $M_2 \leq 0.6 M_\odot$, the secondary is completely evaporated in the Roche lobe filling case. If the $M_2/R_2$ relation (7) is followed, $M_2$ drops below 0.1 $M_\odot$ if $\eta \gtrsim 0.1$. We may conclude that (in either scenario) the secondary is left with a mass $M_2 \lesssim 0.4 M_\odot$ if its initial mass is $M_2 \simeq 0.6 M_\odot$; and the efficiency $\eta \simeq 10^{-2}$; this value is consistent with the estimates $\eta \simeq 10^{-2}$–10$^{-3}$ cited in the literature (e.g., Tavani 1992). The mass lost during the evaporation widened the orbit by about $\sim 10\%$. (e.g., Postnov & Yungelson 2006).

The evaporation scenario we have explored is perhaps the most unfavorable for the survival of the companion, as some assumptions may be a unnecessarily pessimistic. First of all, we have assumed that the relativistic wind is isotropic. This may not be the case; in some models for the birth of magnetars (in connection with the problem of long gamma-ray bursts), most of the rotation energy is removed by highly collimated jets (Bucciantini et al. 2007). If the collimated beam keeps away from the RD, the
The spin evolution of the NS depends on the relative importance of the material and magnetic torques. If the magnetic torque is dominant (which is the case if most of the X-ray luminosity is powered by magnetar emission), the NS’s spin is almost synchronous with the orbital period, similarly to what happens in Polar cataclysmic variables. On the other hand, if the material torques are more relevant, then the spin period is about 50%–70% of the orbital period. In this case, the system would be similar to an intermediate Polar CV, with a NS instead of a white dwarf. The magnetar-powered emission and the full orbital synchronization scenario is most likely to occur if the secondary star is quite difficult to address the structure of the companion after it has been shaken by the magnetar wind. The companion is likely to be larger than a ZAMS star with the same mass, so it may fill its Roche lobe and transfer mass to the NS, powering (part of) its X-ray luminosity.

As far as question (2) is concerned, we expect that when the RD returns to thermal equilibrium, it shrinks, fitting inside its own Roche lobe. Mass transfer ceases and only the residual magnetar-like emission remains. This likely occurs on the Kelvin-Helmholtz time (eq. [45]). Later on, we expect that the magnetar field will also decay. Current theoretical models predict that the present field of $10^{15}$ G will decay to $10^{12}$ G (typical of the magnetic field of NSs in ordinary LMXBs) in $\lesssim 10^8$ yr (Heyl & Kulkarni 1998; Colpi et al. 2000; Geppert & Reinhard 2002) on account of nonlinear field effects like ambipolar diffusion and Hall drift. Later, accretion from the companion sets in, which can induce further field decay, driving it below $\lesssim 10^{10}$ G (Bhattacharya & van den Heuvel 1991; Geppert & Urpin 1994). As the companion star will exhaust its nuclear fuel in few $10^9$ yr, it will swell up as a giant, filling its Roche lobe and transferring mass to the NS. Since the NS’s field is weak by now, the accretion flow may settle in a disk. The ensuing spin-up torque may accelerate the NS up to few milliseconds, and at the end of this stage the system would look like a bona fide LMXB.

A speculative deduction of our model is that at least some LMXBs may have systems like 1E as progenitors. The magnetic field of a NS born as a magnetar in a low-mass binary might remain stronger than the field of a NS in a LMXB or of a binary millisecond pulsar. This kind of higher field is observed in the seven soft X-ray transient LMXBs where an X-ray pulsation is detected during a type I burst. This will be explored in a forthcoming paper.

6. SUMMARY

In this paper we have put forward a model for the unique X-ray source 1E 161348−5055 (1E) in the center of the young supernova remnant RCW 103. In our model 1E is a binary system composed of a NS and a lower main-sequence star. We interpret the signature of 6.67 hr as the NS’s spin period. This requires that the NS is endowed with a strong magnetic field of $\sim 10^{15}$ G.

The spin evolution of the NS depends on the relative importance of the material and magnetic torques. If the magnetic torque is dominant (which is the case if most of the X-ray luminosity is powered by magnetar emission), the NS’s rotation is almost synchronous with the orbital period, similarly to what happens in Polar cataclysmic variables. On the other hand, if the material torques are more relevant, then the spin period is about 50%–70% of the orbital period. In this case, the system would be similar to an intermediate Polar CV, with a NS instead of a white dwarf. The magnetar-powered emission and the full orbital synchronization scenario is most likely to occur if the secondary star is small (i.e., significantly below the current upper limits derived by De Luca et al. 2006).

Alternatively, an early relativistic wind from the NS may have stirred an evaporation wind from the companion. If a large amount of energy was injected into the RD’s envelope, the RD may be out of thermal equilibrium. It may be larger than a main-sequence star with the same mass, possibly powering by accretion the NS’s X-ray luminosity. This phase is expected to last $10^4$ yr, the Kelvin-Helmholtz timescale of the RD layers involved in the energy deposition from the early relativistic wind. As soon as the star
reovers its thermal equilibrium, the star fits again within its own Roche lobe, and mass transfer will end.
If the model we have suggested is correct, then 1E is the first case (to our knowledge) of a magnetar hosted in a low-mass binary system. This may spur some interesting developments in the studies of the evolution of narrow binary systems harboring a NS.

It is a pleasure to thank Sergey Popov for useful discussion and Ed van den Heuvel for his insightful help. This research has been supported by the Italian Ministry of University and Research under the grant PRIN 2005024090_002. We also thank an anonymous referee whose comments helped us to improve the paper.

APPENDIX A

EQUILIBRIUM SOLUTIONS

In this appendix we study in more detail the properties of the equilibrium solutions found numerically in § 4. To this end it is convenient to cast equation (39) in a dimensionless form. We measure the time in units of \( \Omega_{\text{orb}} \), i.e., we set \( \tilde{t} = t \Omega_{\text{orb}} \); we then introduce the synodic phases

\[
\phi'_1 = \phi_1 - \tilde{t}, \quad \phi'_2 = \phi_2 - \tilde{t}.
\]

(A1)

If the secondary is synchronous, then \( \phi'_2 \) is independent of time. After some algebra, equation (39) reads

\[
\frac{d\xi}{d\tilde{t}} = \xi W(\xi),
\]

(A2)

where

\[
\xi = \phi'_1, \\
E = \frac{1}{2} \xi^2 + n_{\text{dd}}(\sin \phi'_2 \sin \phi'_1 - 2 \cos \phi'_2 \cos \phi'_1), \\
W(\xi) = -n_{\text{md}}(1 + \xi)^3 - n_{\text{diss}} \xi + n_{\text{ap}} \left[ \frac{1}{1 + \lambda (1 + \xi)^2} - \frac{\lambda}{2} (1 + \xi)^2 \right],
\]

(A3-5)

and

\[
n_{\text{md}} = \frac{2 \mu_1^2 \Omega_{\text{orb}}^2}{3 I_1 c^2}, \quad n_{\text{diss}} = \frac{1}{(\Omega_{\text{orb}} \tau_{\text{syn}})}, \quad n_{\text{ap}} = \frac{|M_2|(GMb_1)^{1/2}}{I_1 \Omega_{\text{orb}}^2},
\]

(A6-8)

\[
n_{\text{dd}} = \frac{\mu_1 \mu_2}{G M_{I_1}^2}
\]

(A9)

are the dimensionless torques on the NS. Equation (A2) has the form of an energy equation in a nonconservative system, because of the forcing term \( \xi W(\xi) \) on its right-hand side. The dipole-dipole torque does not appear in the braking function \( W \), showing that this torque does not help to attain synchronism, since it does not alter the “energy” \( \xi \).

If there is no mass transfer, then \( n_{\text{ap}} = 0 \), and since \( n_{\text{md}} \ll n_{\text{diss}} \) equation (A2) reads approximately

\[
\frac{d\xi}{d\tilde{t}} \simeq -n_{\text{diss}} \xi^2.
\]

(A10)

Starting from a nonsynchronous configuration (\( \xi \neq 0 \)), the right-hand side of this equation is negative, so \( d\xi/d\tau < 0 \). The system then evolves toward the minimum of \( \xi \), given by \( \xi \equiv \phi'_1 = 0 \) and

\[
\cos \phi'_1 = \frac{2 \cos \phi'_2}{\sqrt{1 + 3 \cos^2 \phi'_2}}, \quad \sin \phi'_1 = -\frac{\sin \phi'_2}{\sqrt{1 + 3 \cos^2 \phi'_2}}.
\]

(A11)

This is the synchronous solution, where the synodic phase \( \phi'_1 \) is locked to the angle defined by the last couple of equations. It is also interesting to note that the “potential”

\[
V(\phi'_1) = n_{\text{dd}}(\sin \phi'_2 \sin \phi'_1 - 2 \cos \phi'_2 \cos \phi'_1)
\]

(A12)

in equation (A4) is proportional to the magnetic interaction potential energy between the stellar magnetic dipoles \( U = -\mu_1 \cdot \mathbf{B}_2(1) \).

The equilibrium synodic position angle of the synchronous solution therefore minimizes the interaction magnetic potential energy between the two stars.
The qualitative pattern of the evolution toward synchronism does not change if \( W \) remains negative. The only positive term in \( W \) which may change the sign of \( W \) and alter this conclusion is the first term in square brackets proportional to \( n_{ap} \). This is the accretion term, which tends to spin up the NS. We now show that for a large accretion rate, the NS tends to a stable asynchronous uniform rotation. If \( n_{ap} \) exceeds a threshold value, the equation \( W(\xi) = 0 \) has a root \( \xi_0 \). Since \( W(\xi) < 0 \) for any \( \xi \geq 0 \), this root is also unique. In addition, \( W(\xi) > 0 \) if \( \xi < \xi_0 \) and \( W(\xi) < 0 \) if \( \xi > \xi_0 \). These results show that \( d\tilde{E}/dt < 0 \) for \( \xi > \xi_0 \) and \( d\tilde{E}/dt > 0 \) for \( \xi < \xi_0 \); the system evolves toward the value of \( \tilde{E} \) characterized by \( \xi = \xi_0 \) and \( \chi = \xi_0 + \phi_0 \). Since \( \xi_0 \neq 0 \), this is an asynchronous solution; the system cannot attain synchronism when the accretion spin up torque exceeds a critical value, in agreement with what is well known in the literature for Polar systems (see, e.g., Campbell 1986).

The result is shown in Figure 4, which plots the equilibrium synodic frequency \( \xi \) against the mass-transfer rate \( |\dot{M}_z| \), for different values of the synchronization time \( \tau_{syn} \). As expected, \( \xi \) grows with \( |\dot{M}_z| \). Below the critical threshold \( |\dot{M}_z| \sim 10^{15} \text{ g s}^{-1} \) the equilibrium solution is always synchronous. As the dissipative torque decreases (i.e., for long \( \tau_{syn} \), the system is more difficult to achieve. For \( |\dot{M}_z| \sim 10^{13} \text{ g s}^{-1} \) (the mass accretion rate inferred from the source’s X-ray luminosity) and the most favorable case \( \tau_{syn} = 10 \text{ yr} \), we find that the NS has a degree of asynchronism of a few percent.

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