Are Gravitational Waves Spinning Down PSR J1023 + 0038?

B. Haskell
Nicolaus Copernicus Astronomical Center, Polish Academy of Sciences, ulica Bartycka 18, 00-716 Warsaw, Poland

A. Patruno
Leiden Observatory, Leiden University, Neils Bohrweg 2, 2333 CA Leiden, The Netherlands
and ASTRON, The Netherlands Institute for Radio Astronomy, Postbus 2, 7900 AA Dwingeloo, The Netherlands

(Received 24 March 2017; published 17 October 2017)

The pulsar J1023 + 0038 rotates with a frequency $\nu \approx 592$ Hz and has been observed to transition between a radio state, during which it is visible as a millisecond radio pulsar, and a low-mass x-ray binary (LMXB) state, during which accretion powered x-ray pulsations are visible. Timing during the two phases reveals that during the LMXB phase the neutron star is spinning down at a rate of $\dot{\nu} \approx -3 \times 10^{-15}$ Hz/s, which is approximately 27% faster than the rate measured during the radio phase, $\dot{\nu} \approx -2.4 \times 10^{-15}$ Hz/s, and is at odds with the predictions of accretion models. We suggest that the increase in spin-down rate is compatible with gravitational wave emission, particularly with the creation of a "mountain" during the accretion phase. We show that asymmetries in pycnonuclear reaction rates in the crust can lead to a large enough mass quadrupole to explain the observed spin-down rate, which thus far has no other self-consistent explanation. We also suggest two observational tests of this scenario, involving radio timing at the onset of the next millisecond radio pulsar phase, when the mountain should dissipate, and accurate timing during the next LMXB phase to track the increase in torque as the mountain builds up. Another possibility is that an unstable r mode with an amplitude $\alpha \approx 5 \times 10^{-8}$ may be present in the system.

DOI: 10.1103/PhysRevLett.119.161103

The system PSR J1023 + 0038 (referred to hereafter as J1023) is a peculiar binary that has been observed to transition back and forth between a radio millisecond pulsar (RMSP) state and a low-mass x-ray binary (LMXB) state [1]. The neutron star spins at a rate of ≈592 Hz, and the companion star is a main sequence star of $\approx 0.2\ M_\odot$. Timing of the radio pulsations has led to a precise measurement of the spin-down of the pulsar $\dot{\nu} = -2.3985 \times 10^{-15}$ Hz/s [2]. After the last transition, which occurred in June 2013 [3,4], x-ray observations during the LMXB state of J1023 allowed us to measure the accretion powered pulsations [5] and the spin-down of the accreting pulsar $\dot{\nu} = -3.0413(90) \times 10^{15}$ Hz/s [6], which is approximately 27% faster than the rate measured from timing during the radio state.

The interpretation of the enhanced spin-down due to the interaction between the accretion disk and the neutron star magnetosphere is somewhat problematic because several inconsistencies remain in each model considered (see Ref. [6] for an extended discussion). For example, a propeller model (see, e.g., Ref. [7]) with a standard α disk where the inner portions are truncated and ejected from the system, or an enhanced pulsar wind model [8], would require a careful fine-tuning of the model parameters to explain the close match between the observed radio and LMXB spin-down rates. The most promising alternative scenario is the trapped-disk model (see Refs. [9,10]), which was instead proposed to explain the presence of outflows in the system and the peculiar low luminosity of J1023 [6]. However, in this case the spin-down needs to be enhanced during both the RMSP and the LMXB stage, meaning that no difference in $\dot{\nu}$ should be observed.

Here, we propose an alternative scenario that would solve this dilemma. We suggest that the additional spin-down is due to gravitational wave (GW) emission, triggered during the LMXB state. Evidence for the presence of GWs in accreting neutron stars has recently mounted due to the lack of submillisecond pulsars (see, e.g., Refs. [11,12]). This may be due, specifically, to the formation of "mountains," i.e., asymmetries in the mass distribution, supported either by crustal or magnetic strains, or unstable modes of oscillation [13]. Note that these mechanisms have been considered before in LMXBs [14–17], but always in the context of spin balance, and it was generally found that it is not easy to build a large enough mountain to balance the spin-up torque due to accretion in these systems [18] (with the notable exception of some persistently accreting systems where the mountain could potentially be large enough for the neutron star to be spinning down during an outburst), and that detection of these signals would require next generation gravitational wave detectors, such as the Einstein Telescope [19,20].

The situation here is, however, radically different. The accretion rate is much lower (on the order of $10^{-13}\ M_\odot\ \mathrm{yr}^{-1}$ [6,7]) and while, on the one hand, this reduces the amount of accreted mass that can build the mountain, on the other it ensures that the spin-up torque is weak enough to not contaminate the spin-down measurement. Furthermore, the precise spin-down rate obtained from radio timing allows for a detailed comparison of the rates during the radio and accretion phase, which is not possible for other LMXBs.

In conclusion, the spin-down rate we attribute to GW emission is the difference between the enhanced rate during the LMXB state and the previous rate during the radio state, i.e.,

$$\dot{\nu}_{\rm diff} = -6.428 \times 10^{-16}\ {\rm Hz/s}. \quad (1)$$

The spin-down rate due to GW emission is

$$\dot{\nu}_{\rm GW} \approx -1.4 \times 10^{-13} \nu_{500}^5 I_{45}^{-1} \left(\frac{Q_{22}}{10^{37}\ {\rm g\,cm^2}}\right)^2\ {\rm Hz/s}, \quad (2)$$

where Q_{22} is the mass quadrupole moment, I_{45} the moment of inertia of the neutron star in units of 10^{45} g cm^2, and ν_{500} the spin frequency in units of 500 Hz.

We can see that, to explain the additional spin-down in Eq. (1) for the spin-frequency of J1023, one requires a quadrupole of

$$Q_{22} = 4.4 \times 10^{35} I_{45}\ {\rm g\,cm^2}, \quad (3)$$

which corresponds to an ellipticity $\varepsilon \approx 5 \times 10^{-10}$ (corresponding to a strain $h_0 \approx 6 \times 10^{-28}$ for a distance to the source of 1.4 kpc [21]), well below the maximum that can be sustained without breaking the crust, $\varepsilon_{\rm max} \approx 10^{-5}$ [22]. Note that this is a conservative estimate of the GW contribution, as we have neglected the spin-up torque due to accretion which, although weak, may contribute to the spin-up at a level of $\dot{\nu}_m \approx 10^{-16}\ {\rm Hz\,s^{-1}}$ for the maximum accretion rate of $\dot{M} = 6 \times 10^{-13}\ M_\odot\,{\rm yr}^{-1}$. Let us thus consider some of the most likely models to establish whether they may lead to such a quadrupole in J1023.

First of all, we will consider the scenario in which asymmetries in the local accretion rate and crustal composition can lead to asymmetric heat release due to pycnonuclear reactions in the crust, i.e., "deep crustal heating" [23], that will source a mass quadrupole [14,24]:

$$Q_{22} = 3 \times 10^{35} R_{12}^4 \left(\frac{\delta T_q}{10^5\ {\rm K}}\right) \left(\frac{E_{\rm th}}{30\ {\rm MeV}}\right)^3\ {\rm g\,cm^2}, \quad (4)$$

where R_{12} is the radius in units of 12 km, $E_{\rm th}$ is the threshold energy for the pycnonuclear reactions responsible for deep crustal heating (corresponding to the electron chemical potential in equilibrium), and T_q is the quadrupolar temperature increase due to the reactions (which will be only a fraction of the total heating δT). Rearranging, we see that we require a quadrupolar temperature increase of

$$\delta T_q \approx 1.5 \times 10^5 R_{12}^{-4} I_{45} \left(\frac{E_{\rm th}}{30\ {\rm MeV}}\right)^{-3}\ {\rm K}. \quad (5)$$

Is such a quadrupolar temperature increase possible in J1023? The total local increase in temperature due to pycnonuclear reactions is [25]

$$\delta T \approx 10^2 C_k^{-1} p_{30}^{-1} Q_M \Delta M_{21}\ {\rm K}, \quad (6)$$

where C_k is the heat capacity in units of the Boltzmann constant per baryon, p_{30} is the pressure in units of 10^{30} erg/cm^3, Q_M is the heat released locally by the reactions per accreted baryon in units of MeV, and ΔM_{21} is the accreted mass in units of 10^{21} g. To obtain an estimate from the above expression, we will take an accretion rate of $5 \times 10^{-14}\ M_\odot/{\rm yr} \lesssim \dot{M} \lesssim 6 \times 10^{-13}\ M_\odot/{\rm yr}$ (estimated by Ref. [7]) and thus consider that, in a year of accretion, the system can accrete $\Delta M \approx 10^{21}$ g.

To obtain the heat capacity, we first need to estimate the temperature of the neutron star which is currently unconstrained from x-ray observations (which are dominated by the thermal emission of the hot polar caps during the RMSP state and by the accretion induced x-ray radiation during the LMXB state [26,27]). To do this, let us consider heating due to deep crustal heating at a rate [28]

$$W_{\rm CH} = 6 \times 10^{30} \left(\frac{\dot{M}}{10^{-13}\ M_\odot/{\rm yr}}\right)\ {\rm erg/s}, \quad (7)$$

which will be balanced by photon cooling at the surface,

$$L_{\rm ph} = 1.7 \times 10^{33} R_{12}^2 \left(\frac{T_s}{10^6\ {\rm K}}\right)^4\ {\rm erg/s}, \quad (8)$$

with T_s being the surface temperature, which, for an iron envelope, can be related to the core temperature T by the relation [29]

$$\left(\frac{T_s}{10^6{\rm K}}\right)^4 = 2.42 g_{14} \left(18.1 \frac{T}{10^9{\rm K}}\right)^{2.42}, \quad (9)$$

with g_{14} being the surface gravity in units of 10^{14} cm/s^2, or by Urca reactions if the star is massive enough, at a rate

$$L_{\rm Urca} = 10^{33} \left(\frac{T}{2 \times 10^7\ {\rm K}}\right)^6 \left(\frac{R_c}{3\ {\rm km}}\right)^3\ {\rm erg/s}, \quad (10)$$

with R_c being the radius of the core region in which Urca reactions can proceed. For both cooling mechanisms, and taking the maximum estimated accretion rate during outburst, we obtain $T \lesssim 10^7$ K for the star. At these temperatures, the heat capacity per baryon in units of the Boltzmann constant is [30] $C_k \approx 10^{-6}$ at $\rho \approx 10^{12}$ g/cm^3, which is approximately the density close to the neutron drip point, where most of the heating occurs (with $E_{\rm th} = 30$ MeV, $Q_M = 0.5$ MeV, and $p_{30} = 1$).

From Eq. (6), we obtain a total heating rate of $\delta T \approx 5 \times 10^6$ K for an accreted mass of $\Delta M = 10^{20}$ g, which is what J1023 is expected to have accreted on the order of a month during the LMXB state.

In order to build a large enough quadrupole, we see from Eq. (5) that we would need (although note that deeper layers will also contribute to the quadrupole, thus reducing the required heating in a single layer at neutron drip)

$$\frac{\delta T_q}{\delta T} \gtrsim 3 \times 10^{-2}. \tag{11}$$

There is no firm constraint on this quantity, with the only limits coming from the nondetection in x rays of quadrupolar flux perturbations in quiescence in transiently accreting LMXBs [20,24], which sets $\delta T_q/\delta T \lesssim 0.1$.

To obtain a first estimate, we may assume that asymmetries in the accreted mass at the surface are confined on a Rossby adjustment radius [31], $R_a = \sqrt{(p/\rho)}/4\pi\nu \approx 3 \times 10^5$ cm, for J1023 with $p = 10^{30}$ erg/cm^3 and $\rho = 10^{12}$ g/cm^3. The rapid rotation rate of the source may thus allow for asymmetries in composition imprinted by accretion at the surface to also persist deep in the crust. It is important to note, however, that the flow in the outer crust is unlikely to be geostrophic, and that friction may be acting on a faster time scale than accretion.

Nevertheless, compositional asymmetries may be frozen in [24] and would allow us to "build" the mountain over successive accretion phases, as would stable magnetic field burial [32]. If this were the case, we would predict the increase in spin-down rate to remain even after the LMXB state, and for the measured value in radio during the next quiescent state to be the same as the current rate in x rays (although we note that this scenario would be more fine-tuned, as one would expect previous outbursts to be similar to the current one, and for the mountain size to already be saturated. Another possibility is that the mountain dissipates on time scales of years, thus persisting in quiescence, but dissipating before the next outburst).

In conclusion, it is likely that a large enough quadrupole can be built on J1023 to explain the additional spin-down. After the accretion phase is over, the mountain will be washed away on a thermal time scale for the crust [28] $\tau_{\rm th} \approx 0.2 p_{30}^{3/4}$ yr, although we note that deeper layers, at higher pressures than the ones we have considered, may also contribute to the quadrupole and thus may dissipate on longer time scales.

For mountains sustained by magnetic stresses, one has [33]

$$Q_{22} \approx 5 \times 10^{32} \Delta M_{21} \mathcal{A} \left(1 + \frac{\Delta M}{M_c}\right)^{-1} \text{ g cm}^2, \tag{12}$$

where $\mathcal{A}$ is a constant of order unity that depends on the equation of state [34] and $M_c \approx 10^{-7}(B/10^{12} G)^{4/3}\ M_\odot$ is the critical mass at which the amplitude of the quadrupole saturates. Note that, close to the critical mass, the simple estimate in Eq. (12) is no longer accurate and numerical simulations are necessary [35]. In general, close to the critical mass one finds that the external dipolar magnetic field is reduced by approximately an order of magnitude by field burial, although numerical simulations seem to indicate that, while the quadrupole saturates, magnetic burial does not, and it may reduce the field even further [32,36]. Despite the uncertainties, the estimate in Eq. (12) suggests that a large enough magnetic mountain cannot be built on J1023 during an accretion phase, as the required amount of mass would take much longer to be accreted. We will thus not consider this mechanism further.

Another possibility is that modes of oscillation of the star may grow unstable during the accretion phase and may provide the additional gravitational wave spin-down torque. The main candidate for this mechanism is the r mode [15], as the f-mode instability will be stabilized by superfluid mutual friction for temperatures below $\approx 10^9$ K [37]. For an internal temperature of $T \approx 10^7$ K and $\nu \approx 592$ Hz, standard models of hadronic neutron stars would predict J0123 to be r-mode unstable (although see Refs. [38–40] for a discussion of why additional physics is probably required in these models). The spin-down rate due to an unstable r mode of dimensionless amplitude α is, if we assume an $n = 1$ polytrope for the equation of state [41],

$$\dot{\nu} \approx -6.7 \times 10^{-16} \left(\frac{\alpha}{10^{-7}}\right)^2 M_{1.4} R_{12}^4 \nu_{500}^7 \text{ Hz/s}, \tag{13}$$

where $M_{1.4}$ is the neutron star mass in units of 1.4 $M_\odot$. For our source (taking $M = 1.71\ M_\odot$, as reported in Ref. [42]), we thus require

$$\alpha \approx 5 \times 10^{-8} R_{12}^{-2}, \tag{14}$$

which is well below theoretical estimates of saturation amplitudes [43] and is consistent with observational upper limits on r-mode amplitudes in LMXBs [39,44,45]. It is also well below current upper limits set by LIGO [46] (although we note that these limits have been set from observations of nonaccreting stars, and the situation may be different in older stars with an accreted crust).

We can also estimate the heating that the r mode would produce

$$W_r \approx 4.5 \times 10^{33} \left(\frac{\alpha}{10^{-7}}\right)^2 M_{1.4}^2 R_{12}^6 \nu_{500}^8 \text{ erg/s}, \tag{15}$$

which, balanced by direct Urca reaction, gives $T \approx 2 \times 10^7$ K, thus potentially contributing to reheating the system more than deep crustal heating.

It is thus possible that the system lies close to the instability curve and is pushed into the unstable region by heating due to deep crustal reactions. The r mode can then grow unstable and contribute to the observed spin-down increase, heating the system further. This scenario is, however, problematic. If the damping time scale increases with temperature, the r mode can indeed enter the instability window, but it is expected to undergo a thermal runaway [47], which will only be halted if, at higher temperatures, the damping time scale decreases with temperature. Such a scenario, involving resonances with other inertial modes, was suggested by Gusakov *et al.* [48] and may lead to oscillations of the mode amplitude around the equilibrium value [49,50].

Another possibility is that the saturation amplitude of the mode is actually small, on the order of $\alpha \approx 10^{-8}$–10^{-7}, which challenges most theoretical models [51] and requires additional physics to be included in the picture, such as, for example, the existence of a phase transition to quark matter in the core [52]. Given our current understanding of r-mode physics, this scenario would, therefore, appear to be somewhat fine-tuned, although it cannot be excluded.

We note that this is not the case for crustal mountains, which provide a natural explanation without requiring any fine-tuning of the unknown model parameters. The strength of the GW torque scales as the square of the quadrupolar temperature δT_q [see Eqs. (4) and (5)]. Since our ignorance of the value of $\delta T_q/\delta T$ suggests that it can be anything between 0 and 0.1, if we use a uniform distribution for our priors, then the expectation value is $\langle \delta T_q/\delta T\rangle \approx 0.05$. This means that the observed value $\dot{\nu}_{\rm diff} \approx (0.03/0.05)^2 \langle \dot{\nu}_{\rm GW}\rangle \approx 0.4\langle \dot{\nu}_{\rm GW}\rangle$; i.e., $\dot{\nu}_{\rm diff}$ is more than one third of the expected value $\langle \dot{\nu}_{\rm GW}\rangle$. Since the value of the upper boundary we have used, $\delta T_q/\delta T$, is an upper limit, we expect $\dot{\nu}_{\rm diff} \gtrsim 0.4\langle \dot{\nu}_{\rm GW}\rangle$. Therefore, there is no fine-tuning required for crustal mountains.

There are two possible observational tests that can be performed to verify whether GWs are the main cause of the excess spin-down in J1023. The first relies on timing the pulsations during the RMSP state. In this case, if the quadrupolar asymmetry generating the GWs is dissipated on a specific time scale, then the excess spin-down should be observed to disappear on the same time scale. According to our estimates, the main contribution to the quadrupole is from layers close to neutron drip and will dissipate on a thermal time scale of a few months, with contributions from deeper layers dissipating on longer time scales of a few years. If the excess spin-down is instead the result of the interaction of the neutron star magnetosphere with the accretion disk, then the excess spin-down should disappear sharply once the transition to the RMSP state is completed. On the other hand, if the mountain is being built cumulatively over successive LMXB states, as may be the case if compositional asymmetries are frozen into the crust, then the enhanced spin-down will persist during the next RMSP state.

The second test can be performed during the LMXB state and involves prolonged timing of the x-ray pulsations observed during the accretion process. In this case, if the mountain builds up over time as additional mass is accreted, the enhanced spin-down should be observed to increase approximately linearly over time, until the mechanism saturates (assuming that the mass accretion rate remains relatively constant, which is a very plausible hypothesis for J1023 [3,27]).

It is also well known that intermittent pulsars show a change in their spin-down rate between *on* and *off* states, with variations between 50% [53] and 250% [54]. Such variations might be related to changes in their magnetospheric configuration [55] with the increased spin-down observed when the radio pulsar is on, i.e., the opposite of what is seen in PSR J1023 + 0038. Therefore, if something similar is occurring in our system, then the true spin-down rate induced by GWs might be larger than estimated in order to compensate for this effect.

Furthermore, if in the future the surface temperature of the neutron star were to be measured, and resulted in an estimate of the core temperature of $T > 10^7$ K, this would suggest that an additional heating mechanism, in addition to deep crustal heating, is active, supporting the hypothesis that the r-mode instability is active. Confirming the existence of an unstable r mode in J1023 would allow us to constrain the instability window and the saturation amplitude of the mode, thus constraining the interior physics of neutron stars [20,56].

Finally, we note that, for a distance to the source of 1.4 kpc [21], the measured gravitational wave strain would be $h_0 \approx 6 \times 10^{-28}$, which is below the detection limit for current interferometers but is potentially detectable by next generation interferometers such as the Einstein Telescope, if the signal is long-lived and can be integrated over outburst time scales on the order of a few years. If thermal and compositional asymmetries, such as those calculated here, are typical for LMXBs, however, other rapidly rotating sources with higher accretion rates are likely to be good targets for current GW searches [20].

This project has received funding from the European Union's Horizon 2020 research and innovation program under Grant Agreement No. 702713. A. P. acknowledges support from a Netherlands Organization for Scientific Research (NWO) Vidi grant.

[1] A. M. Archibald, I. H. Stairs, S. M. Ransom, V. M. Kaspi, V. I. Kondratiev, D. R. Lorimer, M. A. McLaughlin, J. Boyles, J. W. T. Hessels, R. Lynch, J. van Leeuwen, M. S. E. Roberts, F. Jenet, D. J. Champion, R. Rosen, B. N. Barlow, B. H. Dunlap, and R. A. Remillard, Science **324**, 1411 (2009).

[2] A. M. Archibald, V. M. Kaspi, J. W. T. Hessels, B. Stappers, G. Janssen, and A. Lyne, arXiv:1311.5161.

[3] A. Patruno, A. M. Archibald, J. W. T. Hessels, S. Bogdanov, B. W. Stappers, C. G. Bassa, G. H. Janssen, V. M. Kaspi,

S. Tendulkar, and A. G. Lyne, Astrophys. J. Lett. **781**, L3 (2014).

[4] B. W. Stappers, A. M. Archibald, J. W. T. Hessels, C. G. Bassa, S. Bogdanov, G. H. Janssen, V. M. Kaspi, A. G. Lyne, A. Patruno, S. Tendulkar, A. B. Hill, and T. Glanzman, Astrophys. J. **790**, 39 (2014).

[5] A. M. Archibald, S. Bogdanov, A. Patruno, J. W. T. Hessels, A. T. Deller, C. Bassa, G. H. Janssen, V. M. Kaspi, A. G. Lyne, B. W. Stappers, S. P. Tendulkar, C. R. D'Angelo, and R. Wijnands, Astrophys. J. **807**, 62 (2015).

[6] A. Jaodand, A. M. Archibald, J. W. T. Hessels, S. Bogdanov, C. R. D'Angelo, A. Patruno, C. Bassa, and A. T. Deller, Astrophys. J. **830**, 122 (2016).

[7] A. Papitto and D. F. Torres, Astrophys. J. **807**, 33 (2015).

[8] K. Parfrey, A. Spitkovsky, and A. M. Beloborodov, Astrophys. J. **822**, 33 (2016).

[9] H. C. Spruit and R. E. Taam, Astrophys. J. **402**, 593 (1993).

[10] C. R. D'Angelo and H. C. Spruit, Mon. Not. R. Astron. Soc. **406**, 1208 (2010).

[11] S. Bhattacharyya and D. Chakrabarty, Astrophys. J. **835**, 4 (2017).

[12] A. Patruno, B. Haskell, and N. Andersson, arXiv:1705.07669.

[13] P. D. Lasky, Pub. Astron. Soc. Aust. **32**, e034 (2015).

[14] L. Bildsten, Astrophys. J. Lett. **501**, L89 (1998).

[15] N. Andersson, Astrophys. J. **502**, 708 (1998).

[16] C. Cutler, Phys. Rev. D **66**, 084025 (2002).

[17] A. Melatos and D. J. B. Payne, Astrophys. J. **623**, 1044 (2005).

[18] B. Haskell and A. Patruno, Astrophys. J. Lett. **738**, L14 (2011).

[19] A. L. Watts, B. Krishnan, L. Bildsten, and B. F. Schutz, Mon. Not. R. Astron. Soc. **389**, 839 (2008).

[20] B. Haskell, M. Priymak, A. Patruno, M. Oppenoorth, A. Melatos, and P. D. Lasky, Mon. Not. R. Astron. Soc. **450**, 2393 (2015).

[21] A. T. Deller, A. M. Archibald, W. F. Brisken, S. Chatterjee, G. H. Janssen, V. M. Kaspi, D. Lorimer, A. G. Lyne, M. A. McLaughlin, S. Ransom, I. H. Stairs, and B. Stappers, Astrophys. J. Lett. **756**, L25 (2012).

[22] N. K. Johnson-McDaniel and B. J. Owen, Phys. Rev. D **88**, 044004 (2013).

[23] P. Haensel and J. L. Zdunik, Astron. Astrophys. **227**, 431 (1990).

[24] G. Ushomirsky, C. Cutler, and L. Bildsten, Mon. Not. R. Astron. Soc. **319**, 902 (2000).

[25] G. Ushomirsky and R. E. Rutledge, Mon. Not. R. Astron. Soc. **325**, 1157 (2001).

[26] S. Bogdanov, A. M. Archibald, J. W. T. Hessels, V. M. Kaspi, D. Lorimer, M. A. McLaughlin, S. M. Ransom, and I. H. Stairs, Astrophys. J. **742**, 97 (2011).

[27] S. Bogdanov, A. M. Archibald, C. Bassa, A. T. Deller, J. P. Halpern, G. Heald, J. W. T. Hessels, G. H. Janssen, A. G. Lyne, J. Moldón, Z. Paragi, A. Patruno, B. B. P. Perera, B. W. Stappers, S. P. Tendulkar, C. R. D'Angelo, and R. Wijnands, Astrophys. J. **806**, 148 (2015).

[28] E. F. Brown, L. Bildsten, and R. E. Rutledge, Astrophys. J. Lett. **504**, L95 (1998).

[29] E. H. Gudmundsson, C. J. Pethick, and R. I. Epstein, Astrophys. J. **272**, 286 (1983).

[30] A. Y. Potekhin, J. A. Pons, and D. Page, Space Sci. Rev. **191**, 239 (2015).

[31] A. Spitkovsky, Y. Levin, and G. Ushomirsky, Astrophys. J. **566**, 1018 (2002).

[32] M. Priymak, A. Melatos, and D. J. B. Payne, Mon. Not. R. Astron. Soc. **417**, 2696 (2011).

[33] N. Shibazaki, T. Murakami, J. Shaham, and K. Nomoto, Nature (London) **342**, 656 (1989).

[34] A. Melatos and D. J. B. Payne, Astrophys. J. **623**, 1044 (2005).

[35] D. J. B. Payne and A. Melatos, Mon. Not. R. Astron. Soc. **351**, 569 (2004).

[36] M. Vigelius and A. Melatos, Mon. Not. R. Astron. Soc. **395**, 1985 (2009).

[37] N. Andersson, K. Glampedakis, and B. Haskell, Phys. Rev. D **79**, 103009 (2009).

[38] W. C. G. Ho, N. Andersson, and B. Haskell, Phys. Rev. Lett. **107**, 101101 (2011).

[39] B. Haskell, N. Degenaar, and W. C. G. Ho, Mon. Not. R. Astron. Soc. **424**, 93 (2012).

[40] M. E. Gusakov, A. I. Chugunov, and E. M. Kantor, Mon. Not. R. Astron. Soc. **468**, 291 (2017).

[41] B. J. Owen, L. Lindblom, C. Cutler, B. F. Schutz, A. Vecchio, and N. Andersson, Phys. Rev. D **58**, 084020 (1998).

[42] A. T. Deller, A. M. Archibald, W. F. Brisken, S. Chatterjee, G. H. Janssen, V. M. Kaspi, D. Lorimer, A. G. Lyne, M. A. McLaughlin, S. Ransom, I. H. Stairs, and B. Stappers, Astrophys. J. Lett. **756**, L25 (2012).

[43] B. Haskell, Int. J. Mod. Phys. E **24**, 1541007 (2015).

[44] S. Mahmoodifar and T. Strohmayer, Astrophys. J. **773**, 140 (2013).

[45] K. Schwenzer, T. Boztepe, T. Güver, and E. Vurgun, Mon. Not. R. Astron. Soc. **466**, 2560 (2017).

[46] J. Aasi, B. P. Abbott, R. Abbott, T. Abbott, M. R. Abernathy, F. Acernese, K. Ackley, C. Adams, T. Adams, P. Addesso *et al.*, Astrophys. J. **813**, 39 (2015).

[47] Y. Levin, Astrophys. J. **517**, 328 (1999).

[48] M. E. Gusakov, A. I. Chugunov, and E. M. Kantor, Phys. Rev. Lett. **112**, 151101 (2014).

[49] R. V. Wagoner, J. F. Hennawi, and J. Liu, in *Proceedings of the 20th Texas Symposium on Relativistic Astrophysics, Austin, 2000*, edited by J. C. Wheeler and H. Martel (AIP, New York, 2001), p. 781.

[50] E. M. Kantor, M. E. Gusakov, and A. I. Chugunov, Mon. Not. R. Astron. Soc. **455**, 739 (2016).

[51] R. Bondarescu, S. A. Teukolsky, and I. Wasserman, Phys. Rev. D **76**, 064019 (2007).

[52] M. G. Alford, S. Han, and K. Schwenzer, Phys. Rev. C **91**, 055804 (2015).

[53] M. Kramer, A. G. Lyne, J. T. O'Brien, C. A. Jordan, and D. R. Lorimer, Science **312**, 549 (2006).

[54] F. Camilo, S. M. Ransom, S. Chatterjee, S. Johnston, and P. Demorest, Astrophys. J. **746**, 63 (2012).

[55] A. Lyne, G. Hobbs, M. Kramer, I. Stairs, and B. Stappers, Science **329**, 408 (2010).

[56] K. D. Kokkotas and K. Schwenzer, Eur. Phys. J. A **52**, 38 (2016).