Crust Glass Formation Reveals the Neutron Star Birth Properties in IGR J17480-2446

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IGR J17480-2446 is a low-mass x-ray binary, harboring an exceptional accreting pulsar (a neutron star) with an unusual spin frequency of 11 Hz and a very slow postoutburst crust cooling. The former may imply that it is observed at an early stage of recycling, while the latter was shown to indicate the presence in the outer crust of a low thermal conductivity layer, possibly made of glass. Here we argue that the glass layer formation is a natural result of accretion induced failure of pristine cold crystalline crust. This allows us to determine the mass of the accreted material as $\Delta M \approx 2.4 \times 10^{-6} M_\odot$, confirming a very early accretion stage for this neutron star. An analysis of the spin and thermal state reveals a peculiar set of neutron star birth properties which is commonly associated with "recycled" neutron stars, i.e., those that have been experiencing prolonged periods of accretion from a companion. We speculate that such birth properties may represent the outcome of neutron star formation in an electron-capture supernova.

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Introduction—Neutron stars (NSs) are fascinating objects with a typical mass of $\sim 1.4 M_\odot$ and a radius of just ~ 12 km $\approx 1.7 \times 10^{-5} R_\odot$. With central densities well in excess of the standard nuclear saturation density and main properties of matter (composition, thermodynamics, kinetics) governed by strong interaction, NSs are excellent laboratories for nuclear physics [1–6].

The superdense core of an NS is enshrouded by a relatively thin (~ 1 km thick) crust whose density decreases outward, from nuclear to typical terrestrial densities, covering a wide range of physical conditions not accessible in a laboratory. These include matter composed of exotic (neutron-rich, possibly extremely aspherical) atomic nuclei, unbound neutrons, and electrons; strongly coupled fully ionized Coulomb plasma; and various partially ionized states with strong magnetic field effects. Understanding crust properties is crucial for correct interpretation of the information related to the NS core and for verification of theoretical plasma models [7–10].

NSs are also of fundamental importance for astrophysics as they are thought to be responsible for numerous striking observational phenomena such as radiopulsars, millisecond pulsars (MSPs), magnetars, central compact objects, etc. [11–16]. Several scenarios of NS formation have been proposed. The most standard is a collapse of the iron core of a massive star (hereafter FeCC) promoted by iron photodisintegration and giving rise to spectacular supernovae events [17,18]. An alternative route to an NS is a collapse of a degenerate ONeMg core [19,20]. The latter is promoted by a loss of degenerate electron pressure due to electron captures on ^{24}Mg and ^{20}Ne, if the core exceeds a critical mass of $\sim 1.37 M_\odot$. Presently, it is poorly understood whether particular NS formation scenarios produce certain observational classes of NSs or NSs with specific birth properties. See the End Matter for a brief review of the available information.

NSs can be found in binary systems with or without mass transfer onto an NS [21]. Among mass-transferring systems, of special interest are low-mass x-ray binaries (LMXBs). These systems are thought to be progenitors of MSPs via the recycling scenario [22,23]. There is however an LMXB/MSP birthrate problem [24–26], and an alternative scenario of direct MSP birth via accretion-induced collapse (AIC) has been proposed [27,28].

Transiently accreting LMXBs constitute an important subclass [29–33]. For ~ 10 of these sources, real time cooling was observed [34], which provided an opportunity for a detailed analysis. This cooling was interpreted as thermal relaxation of the crust, which was heated out of thermal equilibrium with the core by accretion; modeling this process allows one to deduce dense matter properties such as heat capacity and thermal conductivity [34–37].

IGR J17480-2446 (hereafter J1748; $\mathrm{NS}_{\mathrm{J1748}}$ its NS), the main object of the present Letter, belongs to this last group [38–42]. An interpretation of unique kinetic properties of the crust of $\mathrm{NS}_{\mathrm{J1748}}$ [43,44], allows us to address NS birth properties in a degenerate collapse as well as to contribute to understanding of magnetorotational evolution of NSs in binaries.

J1748—Only one episode of active accretion can be certainly attributed to J1748 [43]. It took place in 2010 [45]. Subsequent x-ray observations detected thermal emission with decreasing temperature interpreted as crust cooling [43,46]. The observed cooling was unusually slow.

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This was shown to be inconsistent with theoretical models of thermal evolution, unless a region with low thermal conductivity was phenomenologically introduced in the outer crust [43,44].

We recall that matter in the outer crust consists of bare atomic nuclei and an almost uniform charge-neutralizing degenerate electron gas [1,7]. In the standard model, the nuclei are arranged in a crystal lattice, and electron conductivities are very high. An attempt was made [43] to link the poor conduction with multicomponent nuclear composition of the crust. The latter was described by the impurity parameter $Q \equiv \langle Z^2 \rangle - \langle Z \rangle^2$, where Z is the charge number and $\langle \ldots \rangle$ here denotes averaging over nuclear species. However, a study within the Markov chain Monte Carlo framework had shown that the observational data required $Q \sim 10^3$ [43], which implied a very peculiar composition and exceeded Q predicted in detailed models based on nuclear reaction networks [47,48] by more than 1 order of magnitude.

Accordingly, it was proposed [43,44] that the thermal conductivity was low not due to a strong impurity contamination, but because of a lack of a long-range (i.e., crystalline) order in ion positions so that the matter in the layer was in an amorphous (glass) state. To fit the cooling data, the authors of Ref. [44] placed the poorly conducting layer at densities $\rho_1 \lesssim \rho \lesssim \rho_2$, where $\rho_1 = 3 \times 10^{10}$ g/cc, $\rho_2 = 3 \times 10^{11}$ g/cc were set as fiducial values. Let us stress that Ref. [44] proposed the glass layer on pure phenomenological grounds, providing neither model of glass formation nor model of its density range.

In this Letter, we present a natural model of glass formation and analyze its astrophysical consequences.

Model—The idea of a glassy state of matter in an NS crust has been put forward previously [49,50], but the ultimate source of disorder in these works was a strong nuclear charge heterogeneity. Since we here consider consequences of the lack of a positional long-range order, we shall assume that at any given density there is only one sort of nuclei. A realistic model of matter is then a one-component Coulomb plasma or Yukawa plasma in the weak-screening regime (pointlike nuclei on an almost uniform neutralizing background of electrons).

It is well known that a glassy state of terrestrial materials can be obtained by a rapid quench of a liquid [51]. The same was demonstrated for Coulomb or Yukawa plasmas [52,53]. In these simulations, a cooled disordered plasma was not able to "find" a crystal state which clearly had lower free energy, due to the fact that too low kinetic energy of ions did not allow them to fully explore the configuration space.

We are not aware of any physical process which could produce a rapid quench of a liquid NS crust. Simulations of crust heating and relaxation [44] indicate that the layer in question is below crystallization temperature T_m both during and after the outburst. Thus, a plausible path to a glassy state must start from a cold crystalline state. We suggest that this path is associated with crystal breaking under an excessive stress. Indeed, if the crystal failure results in a disordered state and if the temperature is sufficiently low, it would take exponentially long for matter to nucleate a new crystal. The stress in question can be due to the weight of the accreted material. A similar process of stress-induced amorphization of crystals is known in terrestrial conditions [54–56] and was briefly discussed for Yukawa plasma [57].

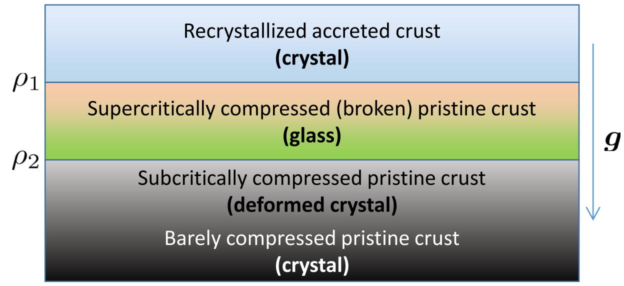

FIG. 1. A cartoon of a glass layer formation in NS crust.

The initial accretion stage can be viewed as loading of the crystallized pristine crust by the accreted material (Fig. 1). The crystal strength decreases with depth decrease. At a given time, the loading produces critical stress at a density, which can be identified with ρ_2. At $\rho < \rho_2$, the critical stress has been exceeded and the crystal is broken; at $\rho > \rho_2$, the original (albeit deformed) crystal survives (gray region in Fig. 1; darker color illustrates higher densities and weaker deformation).

To be consistent with observations, a broken crystal at $\rho < \rho_1$ should have been able to form a new crystal state. To understand the origin of a specific ρ_1 value, let us introduce the dimensionless coupling parameter

$$\Gamma = \left(\frac{4\pi e^6}{3 m_\mathrm{u}} \right)^{1/3} \frac{(Z^5 Y \rho)^{1/3}}{T}, \tag{1}$$

where e is the elementary charge, m_u is the atomic mass unit, $Y \equiv Z/A$, A is the mass number, and T is the temperature. Crystallization occurs at $\Gamma = \Gamma_\mathrm{m} \approx 175.7$ [58].

Glass transition in a one-component plasma is characterized by a value Γ_g above which a crystal does not arise from a disordered state. Theoretical predictions for Γ_g have considerable scatter and can be roughly divided into higher and lower values, $\Gamma_\mathrm{g} \gtrsim 400$–$600$ [52,53,59] and $\Gamma_\mathrm{g} \lesssim 300$–$400$ [59,60]. However, the explicit dependence of Γ on density [Eq. (1)] is weak, and the temperature is a smooth function of ρ due to high thermal conductivity. A natural physical mechanism which would set a particular value of ρ_1 could be a jump of Z, and indeed, such a jump is predicted for the boundary between the pristine and accreted nuclear compositions of the crust.

According to Kepler or Superburst outer crust models of Ref. [48], accreted matter at ρ_1 is expected to have an average charge number $\langle Z \rangle \approx 21$–$22$ (average mass number $\langle A \rangle \approx 52$–$55$, $Y \approx 0.4$). The pristine crust in the immediate vicinity of ρ_1 has large Z [61]: starting from $Z = 36$ ($A = 86$, $Y \approx 0.42$), it gradually drops to $Z = 28$, then jumps to $Z = 44$, and eventually, at ρ_2, settles at $Z = 38$ ($A = 120$, $Y \approx 0.32$). Thus, at $T \sim (2$–$3) \times 10^8$ K and $\rho = \rho_1$, accreted matter has $\Gamma \sim 290$–430, i.e., $\lesssim \Gamma_g$, whereas for pristine matter, $\Gamma \sim 690$–$1030 > \Gamma_g$, locking supercritically compressed (i.e., broken) pristine crust matter in a glassy state.

Summarizing, we predict the glass layer to be bounded by the following conditions: (a) the layer above it is accreted and thus has low Z and (b) the region below it is made of subcritically compressed pristine matter (crystal is not broken); the glass region is shown by the green color in Fig. 1.

Other transiently accreting NSs—It is appropriate to confront our glass formation model with available constraints on the crustal transport properties in other transiently accreting LMXBs with observed NS cooling.

It has been shown [62] that crustal thermal conductivity of another transient source, KS 1731-260, is high, in particular, that an "amorphous" NS crust is not consistent with observations of its postoutburst relaxation [62–64]. Similarly, fitting thermal relaxation of the transient source MXB 1659-29 [44,65] yields impurity parameter Q in full agreement with theoretical models of multicomponent accreted crust composition [48,66] and does not require low-conductivity regions in the outer crust.

This may seem to contradict our model, but, actually, it is not so, as the model implies that the glassy layer is formed from pristine crust, and it sinks to higher densities as more matter is accreted. In particular, the glassy layer should be swept out from the outer crust, as soon as pristine material is replaced by an accreted one. The latter takes place for accreted mass $\sim 5 \times 10^{-5} M_\odot$, which is much less than a typical mass of accreted matter during the LMXB stage. Thus, it is not surprising that other transiently accreting NSs do not have glassy layers in the outer crust. In fact, KS 1731-260 and MXB 1659-29 have spin frequencies of 524 Hz and 567 Hz, respectively, plausibly implying that they have by now accreted much more matter and angular momentum.

In principle, for a fully accreted crust, a glassy layer can form in its innermost layer. In the course of pristine crust replacement by continued accretion, the glass inner boundary, ρ_2, increases until it reaches subnuclear densities where nuclei disintegrate [67,68]. The outer boundary, ρ_1, also increases, until interion spacing becomes sufficiently small to compensate [cf. Eq. (1)] for the smallness of $Z^5 Y$ of the accreted matter in the inner crust; numerically, it is expected to occur at $\sim 10^{13}$ g/cc. Note that continuing relaxation of KS 1731-260 seems to favor a poorly conducting layer at high densities presently explained by a hypothetical disordered nuclear pasta phase [64,69]. Alternatively, this low-conductivity layer could be yet another manifestation of the glass state and amorphization.

Accreted mass—The above logic allows one to determine the accreted mass ΔM in J1748 by applying the hydrostatic equilibrium condition at ρ_1:

$$4\pi R^2 P(\rho_1) = g \Delta M. \tag{2}$$

Henceforth, we neglect general relativity corrections; $g = GM/R^2$ is the gravitational acceleration, P is the pressure, M and R are the NS mass and radius, and G is the gravitational constant. This yields

$$\Delta M \approx 2.4 \times 10^{-6} M_\odot \left(\frac{1.4 M_\odot}{M}\right)\left(\frac{R}{12\ \mathrm{km}}\right)^4, \tag{3}$$

where we took $P(\rho_1) \approx 3.4 \times 10^{28}$ dyn/cm^2, corresponding to $Y = 0.4$ at the bottom of the accreted crust.

Let us note that the possibility of a partially accreted (also called hybrid) crust in $\mathrm{NS}_{\mathrm{J1748}}$ has been discussed previously [41,70] on the basis of different arguments and with a much higher anticipated boundary density (e.g., $\rho \sim 10^{12}$ g/cc was suggested [70] as a boundary between accreted and ground-state layers). Our estimate of ΔM based on the glass formation condition turns out to be 2–4 orders of magnitude lower than that predicted in these works. As we demonstrate below, it carves very specific constraints on the $\mathrm{NS}_{\mathrm{J1748}}$ formation scenario.

Another hydrostatic equilibrium condition can be written at ρ_2, where the same accreted mass causes an additional stress which is equilibrated by a critical compression of the crystal. This allows one to infer a crustal breaking parameter at ρ_2 which turns out to be in a good agreement with theory (see the End Matter).

Thermal and spin state of $\mathrm{NS}_{\mathrm{J1748}}$—According to binary evolution models, accretion starts a long time after NS formation, meaning that the NS had enough time to cool down to very low temperatures. After the beginning of accretion, the long-term thermal state of the NS is determined by a balance between heating processes in the crust (e.g., deep crustal heating) and cooling due to photon emission from the surface and neutrino emission from the interior [29,71,72]. Our model predicts an extremely low amount of accreted matter, prompting a question of whether it was possible to heat up the NS by such ΔM to the observed preoutburst state.

To estimate the thermal content of an NS, one chiefly needs to specify its mass and equation of state, in particular, adopt a model of nucleon superfluidity which suppresses nucleon heat capacity (see the End Matter). Under the assumptions of a relatively low NS mass ($M \lesssim 1.4 M_\odot$), strongly suppressed nucleon heat capacity, and partial core heating by the shallow heating mechanism, there is no contradiction between the observed thermal state of NS_{J1748} and the ΔM value Eq. (3).

It is widely accepted that newly born NSs spin down until accretion from a binary companion accelerates their rotation [73,74]. Based on the time-average accretion rate estimate of $\langle \dot{M} \rangle \sim 2 \times 10^{-11} M_{\odot}/\mathrm{yr}$ deduced from 30-year observations [32,43], accretion onto the NS has started very recently, $\Delta M/\langle \dot{M} \rangle \sim 120$ kyr ago. During the 2010 outburst, the spin-up rate of $\dot{\nu} \approx 1.4 \times 10^{-12}$ Hz/s was measured [38,75]. Assuming that it was proportional to the accretion rate during the outburst $\dot{M}_{\mathrm{b}} \sim 3 \times 10^{-9} M_{\odot}/\mathrm{yr}$ and integrating over the accreted mass, we obtain a spin frequency increase during the entire accretion phase as $\Delta\nu = \dot{\nu}(\Delta M/\dot{M}_{\mathrm{b}}) \sim 0.035$ Hz. This is much lower than the observed NS_{J1748} spin frequency $\nu_{\mathrm{obs}} \approx 11$ Hz (see also the End Matter).

We therefore have an NS that had spun down essentially to its present frequency during the previous evolutionary phase and is now at the very beginning of a spin-up. In view of the low $\Delta\nu/\nu_{\mathrm{obs}}$ ratio, this conclusion cannot be altered by uncertainties of ΔM and $\Delta\nu$, stemming from the lack of specific information on the NS mass, radius, and moment of inertia. Owing to weak B dependence of the torque, $N \propto B^{2/7}$, the conclusion also cannot be altered by any reasonable supposition of a much greater B field before the accretion onset. It crucially constrains the spin-down phase and thus the birth parameters of NS_{J1748}.

Indeed, adopting a model of spin-down (while allowing for a possible magnetic field evolution) and taking into account the source age and its present day spin frequency, one can constrain (see the End Matter) the magnetorotational parameters of the NS at birth, B_{ini} and ν_{ini}. In the standard scenario, thanks to rapid evolution of massive stars, NS_{J1748} must have been born not later than 2–4 Gyr ago, the time of the last star formation episode in Terzan 5 [76]. A lower NS age would imply a nonstandard delayed collapse scenario (see the End Matter), whereas the maximum allowed NS age is comparable to the Hubble time.

In Fig. 2, we plot possible NS_{J1748} birth parameters (B_{ini} vs ν_{ini}) by thin, intermediate, and thick curves, for three fiducial NS ages, $t = 2 \times 10^8$, 2×10^9, and 1.3×10^{10} yrs, respectively, and two spin-down models characterized by braking indices $n = 3$ (solid) and $n = 5$ (dashes). Also shown are present day parameters [77] of isolated (asterisks) and binary (circles) pulsars from the ATNF catalog [80]. The cluster of sources in the top left corner shows normal rotation-powered pulsars whereas MSPs accumulate on the lower right.

Formation scenario—The range of NS_{J1748} birth properties shown in Fig. 2 is at variance with theoretical expectations for a newly born NS, originating via the standard path of NS formation, i.e., via a FeCC. These NSs, as exemplified by a population of normal radiopulsars [81–83], are not expected to have at birth weak B fields $\ll 2 \times 10^{11}$ G and/or millisecond periods. The pulsars with

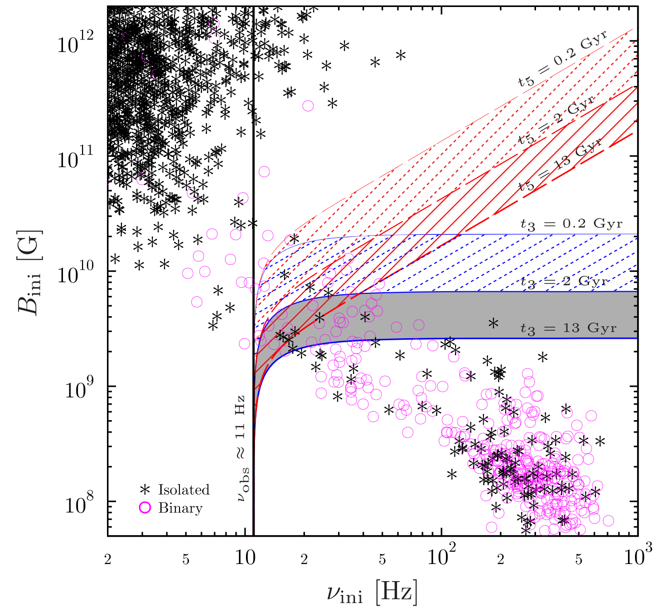

FIG. 2. Initial magnetic field B_{ini} vs initial spin frequency ν_{ini} for three fiducial ages, $t = 2 \times 10^8$, 2×10^9, and 1.3×10^{10} yrs (from thinner to thicker lines) assuming braking index $n = 3$ (solid) and $n = 5$ (dashes). The allowed NS_{J1748} birth age ($2 \times 10^9 < t < 1.3 \times 10^{10}$ yrs) is filled by gray color (hatched by solid lines) for $n = 3$ ($n = 5$) models. The region $2 \times 10^8 < t < 2 \times 10^9$ yrs is hatched by dashes, indicating a significant NS birth delay. Asterisks (circles) are present day parameters of isolated (binary) pulsars from the ATNF catalog [78,80].

similar parameters are observed (cf. Fig. 2), but they are believed to be recycled (i.e., spun up by accretion).

This can be taken as a sign of a nonstandard formation path for NS_{J1748} which would be in line with the long held view [23,27,84–86] that the majority of NSs retained in globular clusters (GCs) originate in the collapse of a degenerate ONeMg core (see the End Matter), predominantly via AIC of a white dwarf (WD) or electron-capture supernova (ECSN). In these cases, the expected NS mass is $\sim 1.25 M_{\odot}$ [87]; thus our tentative conclusion regarding low NS_{J1748} mass is compatible with this idea.

The typically discussed mechanism of NS magnetic field origin in AIC relies on the conservation of a WD magnetic flux [28,88,89]. In this way, NS magnetic fields of $\sim 10^{11}$–10^{13} G typical for regular pulsars can be obtained from the upper range of observed WD fields of $\sim 10^6$–10^8 G (moreover, simulations of AIC of ultramagnetized WDs predict ultramagnetized NSs [90,91]). But lower resulting NS fields of 10^8–10^{10} G should plausibly be more common [27]. We are not aware of studies of magnetic field generation in ECSN. They also cannot be constrained from observations of Galactic field pulsars [83]. It seems reasonable to assume that, in the degenerate precollapse core, the field is not greater than that of an ordinary ONe WD. Consequently, after the collapse, the field is likely lower than those expected for AIC NSs.

NS_{J1748} birth via AIC has been suggested previously [40]; however no restrictions on the NS birth properties were deduced. The AIC scenario was further elaborated [89] by an argument that the presently observed tight binary was not the one in which AIC occurred. The latter had to be disrupted, and the NS had to acquire a new companion via a dynamical interaction.

In our model, delayed collapse is neither required nor excluded. However, AIC is typically followed by a brief detachment phase, after which an intense accretion (already onto NS) resumes [89]. This has to produce an NS with a substantially replaced crust (with low Z) which is incompatible with the glass layer in the outer crust and thus with our interpretation of the cooling data. To avoid that, the post-AIC binary must have been disrupted by a dynamical interaction shortly after the collapse or by a natal NS kick (strong enough to unbind the binary but weak enough to keep the NS in the GC) with NS picking up the present companion later on. Both possibilities seem to require some fine-tuning, making AIC less likely.

If NS_{J1748} were formed via a merger-induced collapse (MIC), it also had to capture the present companion. While the probability of such an event in Terzan 5 may be acceptable, NS formation via MIC is a less common process in GCs [86]. It plausibly results in a higher NS mass and B field [92–94], rendering this path also unlikely.

The remaining NS formation channel is ECSN. It was proposed decades ago [19] and has been actively studied ever since [95–105]. Quite recently, SN2018zd has been argued to be of ECSN type [106–108]. Thanks to low kick velocity it is predicted to be the main source of NSs retained in GCs [85,86]. As a result, we consider ECSN the most appropriate NS_{J1748} formation channel.

Conclusions—The observed slow NS cooling after an outburst in J1748 is explained by supercritical deformation and amorphization of a pristine crystal NS crust caused by the weight of the accreted material. This idea relies on the fact that the glass transition temperature of the ground-state matter is noticeably higher than that of the accreted matter, and is further supported by a good agreement of the crustal breaking parameter $\delta_{\rm crit}$ derived from observations of J1748 with its theoretical value. Our reasoning is fully compatible with earlier studies of other transiently accreting LMXBs, which have found high crustal thermal conductivity for their NSs. It also hints at a natural mechanism of late-time cooling, observed in some of these sources.

Moreover, the proposed picture allows us to determine total accreted mass ($\Delta M \approx 2.4 \times 10^{-6} M_{\odot}$), which strongly constrains the evolutionary scenario of NS_{J1748}. In particular, we demonstrate that the NS birth properties were such as those typically attributed to recycled pulsars. It is plausible that these birth properties represent those of ECSN NSs, indicating that the initial spin and magnetic field distributions may be different for degenerate-collapse NSs and FeCC NSs.

We expect the proposed evolutionary scenario for NS_{J1748} to stimulate renewed interest in NS cooling in LMXBs and in statistics of LMXB and MSP populations. Beyond astrophysics, it motivates further research in condensed-matter and plasma physics on the glass transition and structure of matter stressed supercritically.

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Data availability—There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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End Matter

NS formation: Observations and theory—The degenerate collapse may occur in three different configurations which are an AIC of an ONe WD, a MIC of an ONe WD upon a merger with another WD, or a collapse of the core of a super asymptotic giant branch star (also known as an ECSN) [19,20,86,109–111]. NSs produced in ECSN or AIC are expected to have a low natal kick [97,104]. By contrast, normal rotation-powered pulsars, presumably originating in FeCC, have a birth velocity distribution well described [112] by a Maxwellian with a mean three-dimensional velocity of ~400 km/s. In view of that, it has been long suspected [23,27,84–86] that the majority of NSs in GCs originated via a degenerate collapse as, for an NS to be retained in a GC after birth, it is required to receive a natal kick of no more than ~50 km/s. Besides that, there was a large body of work, suggesting that SN1054 which produced the poster-child rotation-powered pulsar Crab, was of the ECSN type based on composition and dynamics of its supernova remnant [102,113–115]. However, the pulsar kick velocity of ~160 km/s was shown to be impossible to obtain by the preferred hydrodynamic mechanism in ECSN models [104].

NSs produced in ECSN or AIC are expected to have a low mass: the gravitational mass of an NS formed by these mechanisms should be equal to the critical mass limit minus the binding energy, which yields $\sim 1.25 M_\odot$ [87]. On account of this and the low natal kick argument, ECSN appears to be required for the formation of double NS systems, where it has been proposed as the origin of the double pulsar member J0737-3039B and of a number of similar objects [87,97,116–119], as well as for certain classes of high-mass x-ray binaries [120,121].

Unlike FeCC and ECSN, AIC and MIC can produce NS with a substantial delay (up to the Hubble time [94]) with respect to a star formation episode. In view of that, AIC was invoked to explain six seemingly young, moderately spinning pulsars, observed in GCs [88,122,123]. All six have B fields of $10^{11} - 2 \times 10^{12}$ G, and if they were born during star formation episodes in GC (i.e., some Gyrs ago), they would have had enough time to spin down to much lower frequencies. Their inferred young ages thus require delayed formation, implying an AIC (or MIC [89,124]) scenario.

Both AIC and MIC can produce rapidly spinning NS by tapping just a fraction of the angular momentum reservoir available in a binary. MIC likely results in a higher NS mass as only a minor amount of mass, $\lesssim 0.1 M_\odot$, is expected to be ejected in a merger [94,125,126]. Simulations of the AIC of a super-Chandrasekhar mass WD supported by rapid rotation have been also performed [127,128].

Breaking parameter: Observations and theory—Another hydrostatic equilibrium condition can be written at ρ_2, where the same accreted mass causes an additional stress which is equilibrated by a critical compression of the crystal. The respective change of the hydrostatic pressure can be written as a sum of two terms: the first one is associated with a change of electron density in an ion crystal critically contracted in the vertical direction. The second one is due to compression and pressure anisotropy of the ion lattice [129–134]. This leads to the equation

$$P(\rho_1) = \delta_{\rm crit}\rho_2 \frac{{\rm d}P(\rho_2)}{{\rm d}\rho} + \delta_{\rm crit} C n_2 \frac{Z_2^2 e^2}{a_2}. \qquad \text{(B1)}$$

In this case, $n_2 = \rho_2/(A_2 m_u)$ is the ion number density at ρ_2, $a_2 = (4\pi n_2/3)^{-1/3}$ is the respective Wigner-Seitz radius, and Z_2 and A_2 are the charge and mass numbers at ρ_2. The constant $C \approx 0.24$ can be estimated using Voigt averaged elastic properties of the crust [134].

Assuming $P \propto n_{\rm e}^{4/3}$ [1], we obtain

$$\left(\frac{Y_1\rho_1}{Y_2\rho_2}\right)^{4/3} = \frac{4}{3}\delta_{\rm crit} + 0.04(Z_2/40)^{2/3}\delta_{\rm crit}, \qquad \text{(B2)}$$

and neglecting the last term, we arrive at

$$\delta_{\rm crit} \approx \frac{3}{4}\left(\frac{Y_1\rho_1}{Y_2\rho_2}\right)^{4/3}, \qquad \text{(B3)}$$

where $Y_i = Y(\rho_i)$, $i = 1, 2$, $\delta_{\rm crit} \equiv 1 - \xi_{\rm crit}$, and $\xi_{\rm crit}$ is the critical unidirectional contraction factor of crystal matter [133,135].

Substituting numerical values, we get $\delta_{\rm crit} \approx 0.047$. It has been established [135] that the minimum $\delta_{\rm crit}$ over crystallographic directions of contraction was 0.065. This result did not take into account various effects which reduce $\delta_{\rm crit}$: electron screening, ion vibrations, crystal imperfections. In view of this, we consider $\delta_{\rm crit} \approx 0.047$ deduced from observations to be in an adequate agreement with the theory. On top of that, there is likely a fair amount of uncertainty in the determination of ρ_1 and ρ_2 from cooling simulations. For example, if ρ_1 is set to 3.8×10^{10} g/cc, this would bring $\delta_{\rm crit}$ into a perfect agreement with the theoretical prediction [135].

Thermal state of $\mathrm{NS_{J1748}}$—The quiescent state of $\mathrm{NS_{J1748}}$ corresponds to an internal temperature $T \sim 10^8$ K [43]. Assuming a low-mass NS ($M \lesssim 1.4 M_\odot$) and strongly superfluid nucleons, the thermal energy $\mathcal{E}_{\rm th}$ of the NS can be estimated on the basis of its lepton heat capacity of $\lesssim 2\times 10^{37}(T/10^8\ \mathrm{K})$ erg/K [136]. This yields $\mathcal{E}_{\mathrm{th}} \lesssim 10^{45}$ erg. Dividing $\mathcal{E}_{\mathrm{th}}$ by ΔM [Eq. (3)], we get an estimate of the required heating energy per accreted nucleon as $E_{\mathrm{h,est}} \lesssim 0.2$ MeV/nucleon. It is a few times greater than the energy E_{h} predicted by theoretical deep crustal heating models [47,48,137–140] for the accreted layer constrained by the density $\rho < \rho_1$.

However, it is well known that short-term thermal evolution models require an additional shallow heating source E_{sh} to be consistent with observations. For $\mathrm{NS}_{\mathrm{J1748}}$, the energy of this source reaches $E_{\mathrm{sh}} \sim (1.5\text{–}2)$ MeV/nucleon [43,44]. In modeling [44], the source was placed at $\rho = 1.4\times 10^9$ g/cc, and from this location, a significant fraction of shallow heating energy goes into the NS interior [44,141]. Thus, we conclude that under the assumptions of a relatively low NS mass, strongly suppressed nucleon heat capacity, and partial core heating by the shallow heating mechanism, there is no contradiction between the observed thermal state of $\mathrm{NS}_{\mathrm{J1748}}$ and the ΔM value Eq. (3).

If M increases, ΔM decreases [cf. Eq. (3)]. Besides that, nucleons might not be superfluid in the central part of a massive NS in some models of superfluid critical temperatures [65]. All of these effects amplify $E_{\mathrm{h,est}}$ making it harder to reconcile heating requirements with available $E_{\mathrm{h}} + fE_{\mathrm{sh}}$ ($f < 1$ is a fraction of shallow heating, which goes into NS interior). In principle, quantitative simulation of accretion history in J1748 with known ΔM should help further constrain NS parameters and crustal heating models. Although more work is needed, we take this as tentative observational evidence of a relatively low mass of $\mathrm{NS}_{\mathrm{J1748}}$ and of a relatively strong nucleon superfluidity in its core.

Accretion torque and birth parameters of $\mathrm{NS}_{\mathrm{J1748}}$—A similar $\Delta\nu$ estimate can be obtained using the standard theory of the accretion torque [73,74]. It can be written as

$$N = \dot{M}_{\mathrm{b}}(GMr_{\mathrm{A}})^{1/2}, \qquad \text{(D1)}$$

where r_{A} is the Alfven radius

$$r_{\mathrm{A}} = \mu^{4/7}(GM\dot{M}_{\mathrm{b}}^2)^{-1/7}. \qquad \text{(D2)}$$

The equatorial magnetic field at the surface of $\mathrm{NS}_{\mathrm{J1748}}$ at the present time is estimated as $B_{\mathrm{obs}} \sim 2\times 10^8\text{–}2.4\times 10^{10}$ G [38] which corresponds to magnetic moment $\mu_{\mathrm{obs}} \sim 2\times 10^{26}\text{–}2.4\times 10^{28}\ \mathrm{G\,cm^3}$ (following a tradition, we assume NS radius equal to 10 km while converting magnetic field to magnetic moment and vice versa). We integrate Eq. (D1), neglecting any variation of $\dot{M}_{\mathrm{b}}^{-1/7}$ within and between accretion episodes. Using canonical NS mass $M = 1.4M_{\odot}$ and moment of inertia $I = 10^{45}\ \mathrm{g\,cm^2}$, we obtain the total spin-up $\Delta\nu \sim 0.06$ Hz for the upper limit of the magnetic moment. Although this is almost twice the estimate based on the observed spin-up rate during the outburst, it is still almost ~200 times lower than ν_{obs}. Note that more elaborate accretion torque models predict lower magnetosphere radius and lower torque [142].

Let us denote a magnetic field (moment) at the end of the spin-down epoch as B_n (μ_n), where n is the braking index, and consider two representative spin-down models, $n = 3$ and $n = 5$. The $n = 3$ model assumes no magnetic field decay during spin-down so that the magnetic field at birth $B_{\mathrm{ini}} = B_3$ ($\mu_{\mathrm{ini}} = \mu_3$). In the $n = 5$ model [143], one assumes that the magnetic field decays proportionally to the spin frequency so that $\mu(t) = \mu_{\mathrm{ini}}\nu(t)/\nu_{\mathrm{ini}}$.

In both models, the source age is given by the formula

$$t_n = \frac{c^3 I}{4K(n-1)\pi^2\mu_n^2\nu_{\mathrm{obs}}^2}\left(1 - \frac{\nu_{\mathrm{obs}}^{n-1}}{\nu_{\mathrm{ini}}^{n-1}}\right), \qquad \text{(D3)}$$

where c is the speed of light and the parameter K describes the braking efficiency. Calculations of force-free magnetospheres [144] yield $K \approx 1 + \sin^2\alpha$, where α is the obliqueness. For estimates, we take the lower limit, $K = 1$, leading to an overestimation of μ_n (at a given age).

If $n = 3$, the magnetic field is conserved, $B_{\mathrm{ini}} = B_3$. Thus, if we take $t_3 \gtrsim 2$ Gyr and assume a rapid rotation at birth ($\nu_{\mathrm{ini}} \gg \nu_{\mathrm{obs}}$), we get the magnetic field at birth $B_{\mathrm{ini}} \lesssim 7\times 10^9$ G. For lower magnetic fields, $B_{\mathrm{ini}} \ll 7\times 10^9$ G, ν_{ini} approaches the present day value. A higher initial magnetic field, $B_{\mathrm{ini}} > 7\times 10^9$ G, would require a delay of the NS birth.

In the $n = 5$ model, assuming $\nu_{\mathrm{ini}} \gg \nu_{\mathrm{obs}}$ and $t_5 \gtrsim 2$ Gyr, one obtains $B_5 \lesssim 5\times 10^9$ G, which implies a range of birth parameters (region hatched by solid lines in Fig. 2), for instance, $\nu_{\mathrm{ini}} = 100$ Hz and $B_{\mathrm{ini}} \lesssim 4\times 10^{10}$ G. The birth of the NS during the main peak of star formation in Terzan 5, i.e., 12–13 Gyr ago [76], requires $B_5 \lesssim 2\times 10^9$ G (in this case, e.g., $\nu_{\mathrm{ini}} = 100$ Hz implies $B_{\mathrm{ini}} \lesssim 2\times 10^{10}$ G). A lower B_5 would correspond to a lower spin frequency at birth, whereas $B_5 > 5\times 10^9$ G means an NS age below 2 Gyr, i.e., a delay of NS birth with respect to the last star formation episode.

Note that both B_3 and B_5 are in full agreement with B_{obs}.