Deeply bound dibaryon is incompatible with neutron stars and supernovae

Samuel D. McDermott,[1] Sanjay Reddy,[2] and Srimoyee Sen[2]
[1]*Fermi National Accelerator Laboratory, Theoretical Astrophysics Group, Batavia, 60510 Illinois, USA*
[2]*Institute for Nuclear Theory, University of Washington, 98105 Seattle, Washington, USA*

(Received 5 November 2018; published 12 February 2019)

We study the effect of a dibaryon S in the mass range $1860 < m_S < 2054$ MeV, which is heavy enough not to disturb the stability of nuclei and light enough to possibly be cosmologically metastable. Such a deeply bound state can act as a baryon sink in regions of high baryon density and temperature. We find that the ambient conditions encountered inside a newly born neutron star are likely to sustain a sufficient population of hyperons to ensure that a population of S dibaryons can equilibrate in less than a few seconds. This would be catastrophic for the stability of neutron stars and the observation of neutrino emission from the proto-neutron star of Supernova 1987A over $\sim\mathcal{O}(10)$ s. A deeply bound dibaryon is therefore incompatible with the observed supernova explosion, unless the cross section for S production is severely suppressed.

DOI: 10.1103/PhysRevD.99.035013

I. INTRODUCTION

The possibility that six light quarks form the QCD bound state $uuddss$, known as the H dibaryon with binding energy $B_H \equiv 2m_\Lambda - m_H \gtrsim 0$, has been considered for several decades [1]. Direct searches from accelerator-based experiments have ruled out the possibility that such a state has weak decays that are easily detected [2–6] or that such a state is more massive than approximately 2 GeV [7,8]. The suggestion that a much more deeply bound state [9] called the S sexaquark [10,11], with $B_S \equiv 2m_\Lambda - m_S \geq m_\Lambda - (m_p + m_e) = 176.9$ MeV and which nontrivially avoids these observational bounds [12,13], deserves further scrutiny. Lattice studies will eventually be able to test the full spectrum of six-quark states and conclusively decide if such a state exists. Present studies support the existence of a weakly bound dibaryon with $B_H \sim \mathcal{O}(10)$ MeV [14–17], but the more tightly bound and thus stable or cosmologically metastable sexaquark with $B_S \sim \mathcal{O}(\text{few} \times 100)$ MeV, cannot be ruled out at the current level of understanding of lattice systematics [18].

In this work, we consider an S that is light enough to be metastable but massive enough that it is not exothermically produced as a fusion product of two nucleons. This gives the constrained mass range $1860 < m_S < m_\Lambda + m_p + m_e \simeq 2054$ MeV, which in turn implies

$$176.9 < B_S < 361 \text{ MeV}. \tag{1}$$

Due to its electric neutrality and its (meta)stability, such a particle would be a candidate for the dark matter of the Universe [11]. Such a state would avoid detection in underground direct detection experiments due to the overburden of Earth, and may inefficiently deposit energy in the only relevant high-altitude direct detection search [19]. The sexaquark would further have a small enough elastic scattering cross section to avoid present-day cosmological constraints from the power spectrum of the cosmic microwave background radiation [20] or from astrophysical gamma ray searches [21].

The range of binding energies in Eq. (1) ineluctably leads to the conclusion that the production of dibaryons from Λ baryons is on shell and exothermic, however. We study the implications of the production of such a deeply bound QCD state in hot proto-neutron stars. We conclude that observations are in grave tension with the hypothesis of a deeply bound S unless the S production cross section is highly suppressed.

II. BARYONS AND DIBARYONS IN A PROTO-NEUTRON STAR

Production and decay of the S dibaryon is suppressed under ordinary conditions, because creating two units of strangeness requires a doubly weak process. However, the temperature and densities encountered in a proto-neutron star formed during a core-collapse supernova are large enough to produce a thermal population of hyperons through weak reactions [22,23]. Further, since temperatures of the order of tens of MeVs are sustained for a period of about 10 s—a

Published by the American Physical Society under the terms of the Creative Commons Attribution 4.0 International license. Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

timescale set by neutrino diffusion from the proto-neutron star [24]—we will demonstrate that reactions involving hyperons equilibrate the number density of the S dibaryon except under the most extreme possible assumptions.

We begin by writing the coupled differential equations for the number density of different species of baryons. We include only the $N = n, p$, and Λ states; charge conservation is implicit throughout. Λ's can be produced either by the leptonic process $e^- + p \to \Lambda + \nu_e$, or by the nonleptonic process $NN \to N\Lambda$ and $n\pi \to \Lambda$. Due to the high baryon density expected in the neutron star, we shall ignore leptons for simplicity. The time evolution of the number density of each species a is of the schematic form $\dot{n}_a =$ (rate of a production per unit volume) – (rate of a disappearance per unit volume). Because baryon number B is conserved, we expect that the rate of N decay (production) is proportional to n_N (n_Λ), and vice versa. With this in mind, we write

$$\dot{n}_N = -n_N^2\langle\sigma_{NN\to\Lambda N}v\rangle - n_N n_\pi\langle\sigma_{N\pi\to\Lambda}v\rangle + \frac{n_\Lambda}{\tilde{\tau}_\Lambda} + n_\Lambda n_N\langle\sigma_{N\Lambda\to NN}v\rangle \tag{2a}$$

$$\dot{n}_\Lambda = +n_N^2\langle\sigma_{NN\to\Lambda N}v\rangle + n_N n_\pi\langle\sigma_{N\pi\to\Lambda}v\rangle - \frac{n_\Lambda}{\tilde{\tau}_\Lambda} - n_\Lambda n_N\langle\sigma_{N\Lambda\to NN}v\rangle - 2n_\Lambda^2\langle\sigma_{\Lambda\Lambda\to SX}v\rangle + 2n_S n_X\langle\sigma_{SX\to\Lambda\Lambda}v\rangle \tag{2b}$$

$$\dot{n}_S = +n_\Lambda^2\langle\sigma_{\Lambda\Lambda\to SX}v\rangle - n_S n_X\langle\sigma_{SX\to\Lambda\Lambda}v\rangle, \tag{2c}$$

where $\langle\sigma_i v\rangle$ indicates the thermally averaged cross section times velocity for the process i; we discuss the values of the various $\langle\sigma_i v\rangle$ in the ensuing sections. The particle X in the process $\Lambda\Lambda \to SX$ is chosen to conserve strong isospin [25]. We assume $X = \gamma$ in what follows and discuss the rate in detail Sec. III B.

As required, baryon number is conserved in Eqs. (2a)–(2c) since $\dot{\mathsf{B}} \propto \dot{n}_N + \dot{n}_\Lambda + 2\dot{n}_S = 0$. We use initial conditions $n_N(t=0) = n_0$, $n_\pi(t=0) = T^3\exp(-m_\pi/T)$ and $n_\Lambda(t=0) = n_S(t=0) = 0$. We assume that the core has a constant temperature $T = 30$ MeV and is at the nuclear saturation density $n_0 = 0.16\ \mathrm{fm}^{-3}$. The $N \to \Lambda$ and $\Lambda \to N$ transition rates in Eqs. (2a) and (2b) each contain two contributions. Because the π population is Boltzmann suppressed, however, $N\pi \to \Lambda$ is unlikely to be important in this environment. Similarly, one may assume that the $\Lambda \to N$ transition rate $\Gamma_{\Lambda\to N}$ is dominated by the Λ lifetime in the medium, denoted $\tilde{\tau}_\Lambda$. This is true in vacuum, where $\tau_\Lambda \simeq 2.6 \times 10^{-10}$ s, but in a dense medium we expect that direct Λ decay is affected by Pauli blocking; we find that the decay width is reduced, $\tilde{\tau}_\Lambda \simeq 4\tau_\Lambda$. Because $N\Lambda$ collisions are so frequent, Λ disappearance can be dominated by a process analogous to collisional deexcitation, e.g., $N\Lambda \to NN$ may be more rapid than spontaneous decay. For the nucleon densities we consider, $n_N\langle\sigma_{N\Lambda\to NN}v\rangle \gtrsim \tilde{\tau}_\Lambda^{-1}$ if $\langle\sigma_{N\Lambda\to NN}v\rangle \gtrsim 10^{-29}\ \mathrm{cm}^3/\mathrm{s}$.

One important feature of Eq. (2c) is that S disappearance has only one channel, which is suppressed by the large binding energy of the S, since $n_\gamma(E_\gamma > B_S) \sim T^3\exp(-B_S/T)$ is small. Thus, the same features that guarantee the S is cosmologically metastable ensure that it cannot be efficiently destroyed in the proto-neutron star environment: S decay is doubly weak, and S fission is suppressed by its large binding energy, $B_S \gg T$. For this reason, S acts as a sink for baryon number until $n_S \simeq n_N$. If S formation is efficient, all baryon number in the hot proto-neutron star core will be processed into S particles.

The S abundance from Eqs. (2a)–(2c) approximately yields to analytic solution. First, consider the limiting scenario $\langle\sigma_{\Lambda\Lambda\to S\gamma}v\rangle \to 0$. It is clear that n_N, n_Λ reach an equilibrium where $\dot{n}_\Lambda = \dot{n}_N = 0$ when the Λ abundance has increased to

$$\bar{n}_\Lambda = n_N\frac{\langle\sigma_{NN\to\Lambda N}v\rangle}{\langle\sigma_{N\Lambda\to NN}v\rangle + \tilde{\tau}_\Lambda^{-1}/n_N}. \tag{3}$$

The $N - \Lambda$ cross sections are related by detailed balance, such that $\bar{n}_\Lambda/n_N \leq \langle\sigma_{NN\to\Lambda N}v\rangle/\langle\sigma_{N\Lambda\to NN}v\rangle = (m_\Lambda/m_N)^{3/2}\times \exp[-(m_\Lambda - m_N)/T]$. Next, we note that, for constant n_N, Eq. (2b) has an analytic solution even with $\langle\sigma_{\Lambda\Lambda\to S\gamma}v\rangle \neq 0$,

$$n_\Lambda(t) = \bar{n}_\Lambda\frac{2\tanh(\gamma t/2)}{\tanh(\gamma t/2) + \sqrt{1+r}},$$
$$\text{with}\quad \gamma \equiv (\tilde{\tau}_\Lambda^{-1} + n_N\langle\sigma_{N\Lambda\to NN}v\rangle)\sqrt{1+r}$$
$$\text{and}\quad r \equiv \frac{8\bar{n}_\Lambda\langle\sigma_{\Lambda\Lambda\to S\gamma}v\rangle}{\tilde{\tau}_\Lambda^{-1} + n_N\langle\sigma_{N\Lambda\to NN}v\rangle}. \tag{4}$$

The asymptotic Λ abundance is $n_\Lambda^\infty \equiv n_\Lambda(t \gg \gamma^{-1}) = 2\bar{n}_\Lambda/(1+\sqrt{1+r})$, where the time constant satisfies $\gamma^{-1} \leq \tilde{\tau}_\Lambda$. Crucially for our purposes, this happens promptly on the timescales of relevance for a supernova explosion.

Given n_Λ^∞, Eq. (2c) dictates that the S abundance will rise linearly as long as fission is unimportant, $n_S(t) \simeq \bar{n}_S(t) \equiv (n_\Lambda^\infty)^2\langle\sigma_{\Lambda\Lambda\to S\gamma}v\rangle t$. This is true until an $\mathcal{O}(1)$ fraction of baryons are in S dibaryons, which happens at a time t_S defined by $2n_S(t_S) = n_N(t_S)$. We find that t_S defined in this way is equivalent to solving for $\bar{n}_S(t_S) = n_0$, to an accuracy of 10%, or

$$t_S = \frac{n_0}{(n_\Lambda^\infty)^2\langle\sigma_{\Lambda\Lambda\to S\gamma}v\rangle}. \tag{5}$$

Plugging n_Λ^∞ into Eq. (5) and assuming a hierarchy of rates $\bar{n}_\Lambda\langle\sigma_{\Lambda\Lambda\to S\gamma}v\rangle \ll n_0\langle\sigma_{N\Lambda\to NN}v\rangle \sim \tilde{\tau}_\Lambda^{-1}$, we find that S production equilibrates at a time $t_S \simeq \mathrm{s}\frac{4\times10^{-34}\ \mathrm{cm}^3/\mathrm{s}}{\langle\sigma_{\Lambda\Lambda\to S\gamma}v\rangle}\times [1 + \frac{2\times10^{-32}\ \mathrm{cm}^3/\mathrm{s}}{\langle\sigma_{NN\to\Lambda N}v\rangle}]^2$. After t_S has elapsed, backreaction will become non-negligible due to the heat dumped by the exothermic S fusion process. Due to the large binding energy, $\gamma S \to \Lambda\Lambda$ will become important only deep in the backreacted regime. By this time, however, the assumption of thermal equilibrium will have long since broken down, and the proto-neutron star will either combust or decay entirely to S particles.

III. Λ AND S PRODUCTION

If t_S given in Eq. (5) is short compared to the neutrino burst from SN1987A, which was observed to last for $t_\nu \sim \mathcal{O}(10\ \mathrm{s})$, S production equilibrates quickly on the timescales of relevance to the proto-neutron star. As we discuss in the next section, a proto-neutron star composed entirely of S dibaryons is incompatible with observations. Our analysis indicates that, for $\langle\sigma_{\Lambda\Lambda\to S\gamma}v\rangle \gtrsim 10^{-34}\ \mathrm{cm}^3/\mathrm{s}$, S production is fatal for the proto-neutron star. Here, we calculate $\langle\sigma_{NN\to\Lambda N}v\rangle$ and $\langle\sigma_{\Lambda\Lambda\to S\gamma}v\rangle$.

A. Λ production cross section

To obtain $\langle\sigma_{NN\to\Lambda N}v\rangle$, we first observe that all rates $N\ldots \leftrightarrow \Lambda\cdots$ share a strangeness-changing coupling $g_{\Lambda N\pi}$. We obtain this coupling from the in-vacuum Λ lifetime,

$$\tau_\Lambda^{-1} \simeq \Gamma_{\Lambda\to N\pi} \simeq \frac{g_{\Lambda N\pi}^2}{8\pi m_\Lambda}\frac{|\vec{p}_N|}{m_\Lambda}[(m_\Lambda - m_N)^2 - m_\pi^2], \quad (6)$$

giving $g_{\Lambda N\pi}^2 \simeq 7\times 10^{-11}$. Because strangeness-changing processes are weak processes, this small dimensionless number can be interpreted as coming from $(G_F m_N^2)^2 \sim 10^{-10}$. Assuming a constant matrix element, appropriate in the limit of small m_π [27–29], and assuming that the momentum released to the nucleons is large compared to the Fermi momentum, we may write $\langle\sigma_{NN\to\Lambda N}v\rangle \equiv a g_{\Lambda N\pi}^2 \alpha_{N\pi}\sqrt{T/\pi m_N^3 m_\Lambda^2} \simeq a\times 10^{-27}\ \mathrm{cm}^3/\mathrm{s}$, where $\alpha_{N\pi}\simeq 15$ and a is a function of temperature and density that parametrizes our ignorance of complicated higher-order physics that may become important in the proto-neutron star environment. A more complete calculation including the effects of nucleon degeneracy, described in Appendix A, gives $a \simeq 0.3$–0.5 for the temperatures and densities of interest if single-pion exchange is a good description of the scattering.

It is well known that pion exchange is nonperturbative, so it is possible that higher-order diagrams have a non-negligible interference with the tree-level scattering. If there is a cancellation to 10% in the matrix element, then $a \simeq 10^{-2}$, and the cross section is $\langle\sigma_{NN\to\Lambda N}v\rangle \simeq 10^{-29}\ \mathrm{cm}^3/\mathrm{s}$. To be conservative, we will use $\langle\sigma_{NN\to\Lambda N}v\rangle = 3\times 10^{-30}\ \mathrm{cm}^3/\mathrm{s}$ as a default value for the rest of this paper, corresponding to a 10% cancellation in the matrix element for this process that is sustained for the entirety of the proto-neutron star explosion, on top of the $\sim\mathcal{O}(50\%)$ suppression from m_π effects and nucleon degeneracy. We emphasize that, although such cancellations are known to exist at the $\sim\mathcal{O}(50\%)$ level in the context of $N-N$ scattering, a cancellation of $\sim\mathcal{O}(90\%)$ would be extremely unusual. But a larger value of $\langle\sigma_{NN\to\Lambda N}v\rangle$ will hasten the rate at which baryon number is processed into S particles, so we choose this value to ensure that our results are indeed conservative.

We also mention here that we have neglected additional baryon species. This is reasonable because baryons of increasing strangeness are increasingly massive. For instance, the equilibrium Ξ population experiences a Boltzmann suppression such that $n_\Xi n_N \lesssim (n_\Lambda^\infty)^2$. Including such additional baryons would marginally increase the S production rate, but more importantly would make the cancellation we implicitly absorb even more unlikely. Thus, our analysis is conservative, but this contributes subdominantly to the calculation of t_S.

B. S production cross section

We now calculate the cross section for $\Lambda\Lambda \to S\gamma$. Given the range of dibaryon masses considered, this process is exothermic and involves no change of strangeness. The effective Lagrangian that allows this process is

$$\mathcal{L} \supset d_\Lambda \bar{\Lambda}\sigma^{\mu\nu}\Lambda F_{\mu\nu} + g_{\Lambda S}\overline{\Lambda^c}\Lambda S^\dagger + \mathrm{H.c.}, \quad (7)$$

where the dipole moment $d_\Lambda = -0.613 \pm 0.001\mu_N \simeq (10^4\ \mathrm{MeV})^{-1}$, Λ^c is the Λ charge conjugate, and $g_{\Lambda S}$ is a function of inherent dibaryon properties discussed in more detail below. From direct calculation, we find that for the temperatures and binding energies of interest the cross section due to the Lagrangian in Eq. (7) is

$$\langle\sigma_{\Lambda\Lambda\to S\gamma}v\rangle \simeq 3\times 10^{-23}\frac{g_{\Lambda S}^2 B_S}{176.9\ \mathrm{MeV}}\frac{T}{30\ \mathrm{MeV}}\frac{\mathrm{cm}^3}{\mathrm{s}}, \quad (8)$$

where we have assumed that the fraction of final states with the quantum numbers of the S is $1/1440$. The magnitude of $g_{\Lambda S}$ introduces the largest uncertainty into our calculations.

The coupling $g_{\Lambda S}$ is in principle a low-energy output of QCD. Since strongly coupled QCD is not currently amenable to analytic calculation, and since lattice studies are difficult for a large number of light quarks, we must choose a model to calculate $g_{\Lambda S}$. In prior work, $g_{\Lambda S}$ has been determined by a geometric factor given by the integrated wave function overlap [9,30]. We will follow these works and use the Isgur-Karl [31] and Brueckner-Bethe-Goldstone [32] models to calculate the overlap of the Λs and the S. This is, of course, only one model of the complicated nuclear quantum mechanics involved.

As discussed in more detail in Appendix B, the wave function overlap has a striking dependence on the dibaryon radius r_S and the Λ radius r_Λ. The S radius is entirely unknown, so to be maximally conservative we simply require that r_S exceed the Compton wavelength of the dibaryon plus some fraction x of the Compton wavelength of the lightest meson to which it couples, as advocated in [13]. This gives

$$r_S \geq \frac{1}{m_S} + \frac{x}{m_{f^0}} = 0.1\ \mathrm{fm}\frac{2054\ \mathrm{MeV}}{m_S} + 0.34x\ \mathrm{fm}. \quad (9)$$

We will show results for $x = 0, 0.1$ in our final plots. Since the dibaryon is a boson, it has no inherent exclusion principle to provide pressure against collapse, so a large coupling to a vector mediator satisfying $g_\omega/m_\omega \geq g_\sigma/m_\sigma$ is necessary [33]. We return to this point below. If instead we required that the nonrelativistic zero-point kinetic energy $r^{-2}/2m$ of quarks localized within the dibaryon of radius r_S should not exceed the energy scale of QCD confinement, we would find a sharper bound. Asserting only that $m_q \leq m_S$ would translate to a bound $r_S \geq 0.22\ \mathrm{fm}\sqrt{2054\ \mathrm{MeV}/m_S}$. Taking a constituent quark mass $m_q \simeq m_S/6$, we would have $r_S \gtrsim 0.53$ fm. This latter value roughly matches the constituent quark Compton wavelength, $6/m_S \gtrsim 0.58$ fm. For this reason, restricting to the range $0.1 \leq r_S \leq 1.0$ fm is very conservative, and the choice $r_S \simeq 0.1$ fm would be an extremely novel feature for a QCD bound state.

Likewise, the Λ radius carries some uncertainty. It is reasonable to assume that increasing strangeness leads to a more compact baryon, $r_\Lambda \lesssim r_N$. The strong interaction radius extracted from experimental data $\sqrt{\langle r_\Lambda^2 \rangle_{\mathrm{st}}} = 0.76 \pm 0.01$ fm [34] is somewhat larger than the naïve value in the constituent quark model, $r_\Lambda \simeq [2\Lambda_{\mathrm{QCD}} m_\Lambda/3]^{-1/2} \simeq 0.51$ fm. Being cautious once again, we decide to show the relatively wide range $0.5 \leq r_\Lambda \leq 0.8$ fm, where the lower limit is chosen to account for the possibility that the Λ charge radius is smaller than the strong interaction radius.

Finally, we note that if the binding energy is near the extreme of the range in Eq. (1), then the process $\Lambda\Lambda \to S\pi\pi$ is on shell and exothermic as well. Emission of two pions is likely dominated by quark rearrangement processes, which occur at long distances due to the small pion mass. Because the light quarks in the initial state can escape to distances of order the pion Compton wavelength, the cross section should be $\langle \sigma_{\Lambda\Lambda\to S\pi\pi} v \rangle \sim \mathcal{O}(m_\pi^{-2})$, which does not suffer from an exponential wave function overlap suppression factor. There will be $\sim\mathcal{O}(0.1)$ hadronization and mass-dependent phase-space suppression factors that we cannot calculate, however. Regardless, for masses $m_S \lesssim 1950$ MeV, we expect that the timescale $t_S \ll$ ns is unsuppressed and independent of r_S. This strengthens the argument considerably in the mass range 1850–1900 MeV, which is of particular interest in recent studies [13].

IV. FATE OF THE PROTO-NEUTRON STAR

We show our final results in Fig. 1, fixing $\langle \sigma_{N\Lambda\to NN} v \rangle = 3 \times 10^{-30}\ \mathrm{cm}^3/\mathrm{s}$. The left panel of Fig. 1 depicts the lifetime as a function of r_Λ and r_S for $m_S = 1900$ MeV. In the dark (light) gray region, r_S violates Eq. (9) for $x = 0(0.1)$. S production equilibrates in the proto-neutron star much faster than 10 s for most of the range of r_Λ and r_S

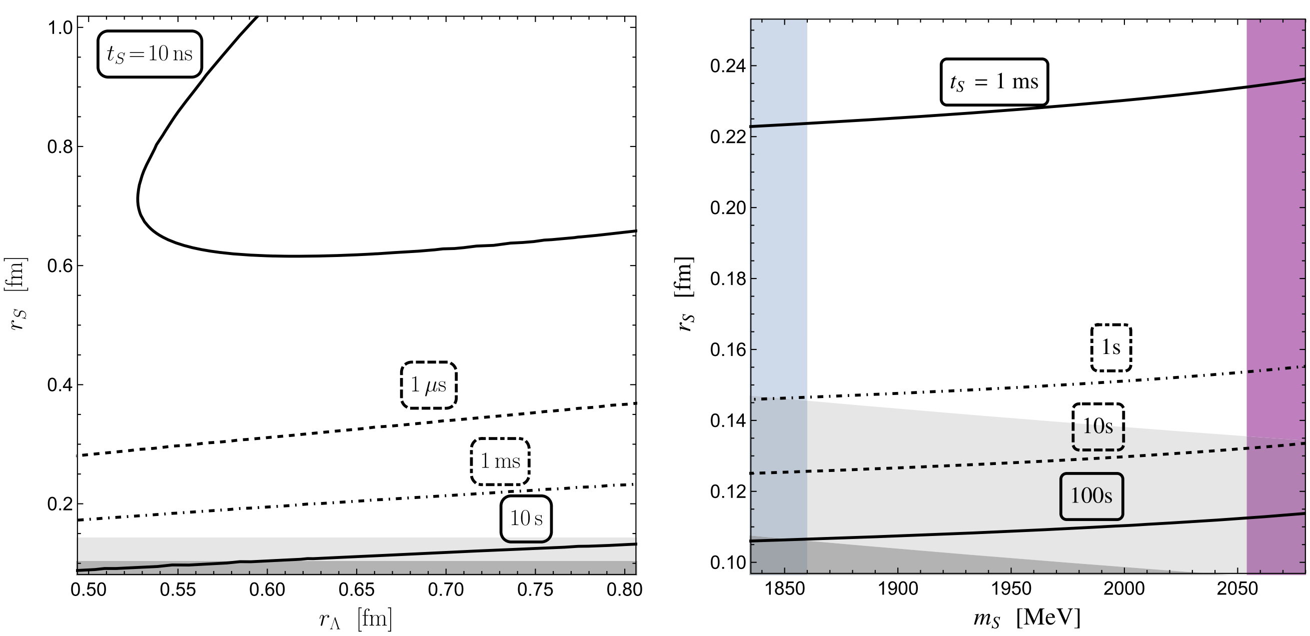

FIG. 1. (Left) Contours of t_S as defined in Eq. (5) for $m_S = 1900$ MeV as a function of the Λ and S sizes. The gray region violates Eq. (9) for $x = 0$, 0.1. (Right) Contours of t_S for $r_\Lambda = 0.76$ fm. In both panels, we have assumed $\langle \sigma_{N\Lambda\to NN} v \rangle = 3 \times 10^{-30}\ \mathrm{cm}^3/\mathrm{s}$. The gray region violates Eq. (9) for $x = 0$, 0.1. In the blue region, ^{16}O nuclei are destabilized. In the purple region, the dibaryon has a singly weak decay. All of the parameter space depicted in each panel has $t_S \ll 10$ s and is thus ruled out by the observation that SN1987A continued to emit neutrinos for $t_\nu \simeq 10$ s, unless r_S is very close the minimum value allowed by Eq. (9).

that we consider, unless r_S is very close to 0.1 fm. For such a small radius, the coupling can be as small as $g^2_{\Lambda S} \sim 10^{-11}$–$10^{-14}$ by the wave function overlap calculation discussed in Appendix B.

In the right panel of Fig. 1, we depict t_S for $r_\Lambda = 0.76$ fm as a function of dibaryon mass m_S and radius r_S. In the blue shaded region, and at smaller masses, the existence of an S dibaryon renders ^{16}O nuclei unstable [30]. In the purple shaded region, and at larger masses, the dibaryon cannot possibly be cosmologically metastable, since it has a singly weak decay [9]. In the dark (light) gray region, r_S violates Eq. (9) for $x = 0(0.1)$.

In all of the heretofore phenomenologically viable parameter space, we find that $t_S \ll 10$ s, unless r_S is very close to 0.1 fm. Such a fast equilibration of the S number density implies that all baryons in the proto-neutron star interior rapidly find themselves inside S dibaryons. This would have catastrophic consequences. Since the S dibaryon is a compact boson, its equation of state would be characterized by a pressure that is much smaller than the pressure of the neutron-rich matter it replaces. Fermi degeneracy and strong interactions between neutrons produce enough pressure to support neutron stars up to a maximum mass $>2\ M_\odot$, compatible with observations of massive neutron stars [35,36]. In contrast, matter composed of the S dibaryon, where pressure is solely due to short-range repulsion, would be too compressible to support such a large maximum mass. We have estimated the strength of repulsive interactions needed to support a maximum mass of $2\ M_\odot$ and found that, in a simple model where dibaryons interact by exchanging vector mesons with mass $m \simeq m_\omega 800$ MeV, the coupling strength needed to produce adequate repulsion to support observed neutron star masses is unnaturally large. Treating the dimensionless dibaryon-vector meson coupling strength g_S as free parameter, we calculated the equation of state of the interacting dibaryon system in mean field theory and found that to support a maximum mass $>2\ M_\odot$ we require unnaturally large values of $g_S > 10$. As discussed above, a coupling large enough to ensure stability would also increase the characteristic size of the dibaryon and would preclude $r_S \simeq 1/m_S$. Interestingly, in this simple model with large repulsive couplings we also find that the radius of typical neutron stars (with masses in the range 1.2–1.5 $M_\odot$) would be greater than 15 km. This is in conflict with the constraints from GW1701817 [37–39]. Taken together, this suggests that interactions between dibaryons is unlikely to change our conclusion that the star composed mostly of tightly bound dibaryons is incompatible with observations.

Finally, the large energy released by the exothermic reactions, $B_S \sim 100$ MeV per baryon, is comparable to the gravitational binding energy. S production likely unbinds the stellar remnant, but even if the proto-neutron star remains intact, this heat dump disrupts the standard evolution of the proto-neutron star.

V. CONCLUSIONS

In this work, we have shown that the hot interior of a proto-neutron star provides a valuable laboratory for probing the nature of the proposed deeply bound S dibaryon. The S can be produced on shell in $\Lambda\Lambda$ collisions, and this exothermic reaction equilibrates quickly on the timescales of relevance to the neutron star explosion unless the dibaryon production cross section is suppressed by 11 orders of magnitude. In the context of a wave function overlap calculation, we find that this is possible only if the S radius is very close to its Compton wavelength $\simeq$0.1 fm. Absent this suppression, rapid equilibration of S density implies that all baryon number inside of the proto-neutron star is processed into S number much more quickly than the observed neutrino burst of Supernova 1987A. Indeed, the energy released in the hard gamma rays that accompany the formation of the new population of S dibaryons is large and could unbind the proto-neutron star entirely. Finally, if such an object were to survive, an entire star composed entirely of S particles would have a much softer equation of state than a neutron star. Thus, the existence of proto-neutron stars and old neutron stars with properties roughly similar to those predicted from standard nuclear astrophysics seems to be in grave tension with the presence of a dibaryon in the QCD spectrum.

ACKNOWLEDGMENTS

We thank Nikita Blinov, Glennys Farrar, Rocky Kolb, and Michael Turner for discussions. S. D. M. was supported by Fermi Research Alliance, LLC under Award No. DE-AC02-07CH11359 with the U.S. Department of Energy, Office of Science, Office of High Energy Physics. The U.S. Government retains and the publisher, by accepting the article for publication, acknowledges that the U.S. Government retains a nonexclusive, paid-up, irrevocable, worldwide license to publish or reproduce the published form of this manuscript, or allow others to do so, for U.S. Government purposes. S. R. and S. S. are supported by Department of Energy Award No. DE-FG02-00ER41132.

Note Added.—Recently, we received a draft of [26], which critically addresses the possibility that the S can attain an interesting cosmological abundance. The underlying assumption of [26] is that S is present in the QCD spectrum, which makes it complementary to the present work. Recently, we also became aware of [40], which finds no candidate events from a search for the S in Υ decays.

APPENDIX A: $NN \to \Lambda N$ CALCULATION

The cross section for $\langle \sigma_{NN\to\Lambda N} v \rangle$ determines the equilibrium Λ abundance, which in turn determines t_S. Assuming a trivial matrix element for single-pion exchange and integrating over nondegenerate phase space, in agreement with calculations of nucleon-nucleon scattering in the single-pion-exchange limit [27–29], gives $\langle\sigma_{N\Lambda\to NN}v\rangle \equiv g^2_{\Lambda N\pi}\alpha_{N\pi}\sqrt{T/\pi m_N^3 m_\Lambda^2} \simeq \times 10^{-27}\ \mathrm{cm}^3/\mathrm{s}$, where $g_{\Lambda N\pi}$ is obtained from Eq. (6) and $\alpha_{N\pi} \simeq 15$. Effects of degeneracy are expected to be mild in this environment [27], but should have effects at the $\sim\mathcal{O}(1)$ level [28]. Here we confirm this expectation with explicit calculation.

The rate per unit volume for production of Λ baryons in NN collisions is

$$\frac{\Gamma}{\mathrm{Vol}} = \int \prod_{i=1}^{4} \frac{d^3\vec{p}_i}{(2\pi)^3 2E_i}(2\pi)^4\delta^{(4)}(p_1+p_2-p_3-p_4) \times f(N_1)f(N_2)[1-f(N_3)][1-f(\Lambda_4)]|\mathcal{M}_{NN\to\Lambda N}|^2, \tag{A1}$$

where $f(B_i) = \{\exp[(E_i-\mu_i)/T]+1\}^{-1}$ is the Fermi-Dirac distribution function for the baryon i. The matrix element $\mathcal{M}_{NN\to\Lambda N}$ follows from the Lagrangian $\mathcal{L} \supset g_{NN\pi}\bar{N}\gamma_5 N\pi + g_{\Lambda N\pi}\bar{\Lambda}\gamma_5 N\pi + (\mathrm{H.c.})$, where $g_{NN\pi}$ is given by the Goldberger-Treiman relation. The chemical potential and temperature are related by the requirement that $n_0 = \int \frac{2\times d^3\vec{p}}{(2\pi)^3} f(N_i)$. The chemical potentials satisfy $\mu_\Lambda = \mu_N$ by detailed balance. We find that $\mu_N \gtrsim m_N$, and thus the N are mildly degenerate, for $T \lesssim 50$ MeV.

Because the nucleon densities are fixed to the saturation value, we may determine the cross section by

$$\langle\sigma_{NN\to\Lambda N}v\rangle = \frac{\Gamma}{\mathrm{Vol}} n_0^{-2}. \tag{A2}$$

We plot the results of Eq. (A2) and the value $g^2_{\Lambda N\pi}\alpha_{N\pi}\sqrt{T/\pi m_N^3 m_\Lambda^2} \simeq \sqrt{T/30\ \mathrm{MeV}} \times 10^{-27}\ \mathrm{cm}^3/\mathrm{s}$ for $15\ \mathrm{MeV} \le T \le 80\ \mathrm{MeV}$ in Fig. 2, left panel. The result with the assumption of a trivial phase space is a factor of ~ 3 higher at $T = 30$ MeV. The discrepancy shrinks at large T, where corrections due to $m_\pi \neq 0$ are less important.

APPENDIX B: WAVE FUNCTION OVERLAP CALCULATION

Following [9], we integrate the Isgur-Karl wave functions of two initial-state baryons against a relative wave function that incorporates the $\Lambda-\Lambda$ potential. In agreement with [9,30], we have

$$g^{(\mathrm{ovp})}_{\Lambda S} = 32\left(\frac{3}{2\pi}\right)^{3/4}\frac{(r_S/r_\Lambda)^{9/2}}{[1+(r_S/r_\Lambda)^2]^6} r_\Lambda^{-3/2} \times \int d^3a\,\psi_{\mathrm{rel}}\psi_\gamma \exp^{-3a^2/4r_S^2}, \tag{B1}$$

where ψ_{rel} has mass dimension $-3/2$. We assume that the γ is a plane wave whose presence allows conservation of energy and momentum. It is possible that, in processes where strong mesons are emitted, such as $\Lambda\Lambda \to S\pi\pi$ or $N\Xi \to S\pi$, the presence of the π has qualitative significance for the process of S formation. For instance, if quark rearrangement is important, then some of the quarks in the initial state may escape to the π, which is at a distance much larger than r_S, meaning that the wave functions need not coincide as exactly as in our model calculation, and the cross section may be as large as m_π^{-2}. However, such effects are difficult to quantify in the absence of a calculable model of hadronization, so we restrict to $\Lambda\Lambda \to S\gamma$, where such considerations are irrelevant. Nonetheless, we stress that a complete picture should include all rearrangement effects, and may lead to substantially larger cross sections.

For numerical values of ψ_{rel}, we use the relative wave functions depicted in Fig. 5 of [41]. These wave functions are generated from potentials calibrated on the Nagara

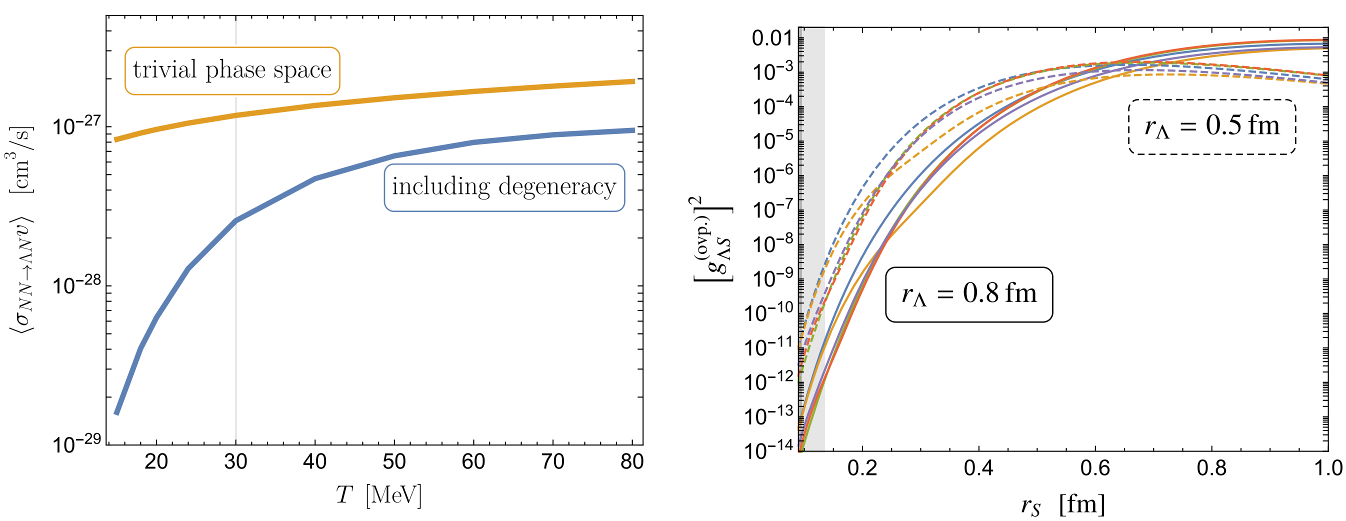

FIG. 2. (Left) The value of $\langle\sigma_{NN\to\Lambda N}v\rangle$ with and without phase-space degeneracy effects. (Right) The coupling $g^2_{\Lambda S}$ from wave function overlap, as a function of r_S for two different values of r_Λ, as in Eq. (B1). In the gray region, r_S violates Eq. (9) for $x = 0.1$.

event, which requires a slightly repulsive interaction. The inverse scattering length is small and negative, while consistency should require that the inverse scattering length for a very deeply bound dibaryon is large and positive [42,43]. Needless to say, an attractive potential would lead to a relative wave function that was larger near the origin. On the other hand, $\Lambda \leftrightarrow N$ transitions can occur more quickly than $\Lambda\Lambda$ fusion for small $g_{\Lambda S}$, meaning that the two baryons involved in a single $\Lambda\Lambda \to S\gamma$ event may change strangeness while they are within range of each other's potential. Thus, the correct relative wave function may be a linear combination of relative $\Lambda - N$ and $\Lambda - \Lambda$ wave functions. For this reason, the slightly repulsive potentials of [41] provide a conservative model of this process.

We show the final results of integrating Eq. (B1) in Fig. 2. As is clear, $g_{\Lambda S}$ calculated in this way is largely insensitive to the details of the wave functions: all of these relative wave functions integrate to $\mathcal{O}(1)$ numbers. The more important scaling has to do with the large polynomial dependence on r_S and r_Λ and the exponential dependence on r_S, which cause the square of the overlap to vary by approximately 3 orders of magnitude.

[1] R. L. Jaffe, Phys. Rev. Lett. **38**, 195 (1977); **38**, 617(E) (1977).
[2] J. Belz *et al.* (BNL-E888 Collaboration), Phys. Rev. Lett. **76**, 3277 (1996); Phys. Rev. C **56**, 1164 (1997).
[3] A. Alavi-Harati *et al.* (KTeV Collaboration), Phys. Rev. Lett. **84**, 2593 (2000).
[4] B. H. Kim *et al.* (Belle Collaboration), Phys. Rev. Lett. **110**, 222002 (2013).
[5] J. Tscheuschner, Ph. D. thesis, Technische Universität Darmstadt, 2014.
[6] J. Adam *et al.* (ALICE Collaboration), Phys. Lett. B **752**, 267 (2016).
[7] J. Badier *et al.* (NA3 Collaboration), Z. Phys. C **31**, 21 (1986).
[8] R. H. Bernstein, T. K. Shea, B. Winstein, R. D. Cousins, J. F. Greenhalgh, M. Schwartz, G. J. Bock, D. Hedin, and G. B. Thomson, Phys. Rev. D **37**, 3103 (1988).
[9] G. R. Farrar and G. Zaharijas, Phys. Rev. D **70**, 014008 (2004).
[10] G. R. Farrar, arXiv:1708.08951.
[11] G. R. Farrar, in *Proceedings, 35th International Cosmic Ray Conference (ICRC 2017): Bexco, Busan, Korea, 2017* (Sissa Medialab srl Partita IVA: 01097780322, 2017).
[12] G. Zaharijas and G. R. Farrar, Phys. Rev. D **72**, 083502 (2005).
[13] G. R. Farrar, Phys. Rev. D **98**, 063005 (2018).
[14] S. R. Beane *et al.* (NPLQCD Collaboration), Phys. Rev. Lett. **106**, 162001 (2011).
[15] S. R. Beane *et al.*, Mod. Phys. Lett. A **26**, 2587 (2011).
[16] S. R. Beane, E. Chang, W. Detmold, H. W. Lin, T. C. Luu, K. Orginos, A. Parreno, M. J. Savage, A. Torok, and A. Walker-Loud (NPLQCD Collaboration), Phys. Rev. D **85**, 054511 (2012).
[17] S. R. Beane, E. Chang, S. D. Cohen, W. Detmold, H. W. Lin, T. C. Luu, K. Orginos, A. Parreno, M. J. Savage, and A. Walker-Loud (NPLQCD Collaboration), Phys. Rev. D **87**, 034506 (2013).
[18] M. J. Savage and S. R. Beane (private communication).
[19] M. S. Mahdawi and G. R. Farrar, J. Cosmol. Astropart. Phys. 10 (2018) 007.
[20] V. Gluscevic and K. K. Boddy, Phys. Rev. Lett. **121**, 081301 (2018).
[21] D. Hooper and S. D. McDermott, Phys. Rev. D **97**, 115006 (2018).
[22] J. A. Pons, S. Reddy, M. Prakash, J. M. Lattimer, and J. A. Miralles, Astrophys. J. **513**, 780 (1999).
[23] W. Keil and H. T. Janka, Astron. Astrophys. **296**, 145 (1995).
[24] A. Burrows and J. M. Lattimer, Astrophys. J. **307**, 178 (1986).
[25] We thank the authors of [26] for pointing out that strong isospin forbids $\Lambda\Lambda \to S\pi$.
[26] E. W. Kolb and M. Turner, arXiv:1809.06003.
[27] R. P. Brinkmann and M. S. Turner, Phys. Rev. D **38**, 2338 (1988).
[28] G. Raffelt and D. Seckel, Phys. Rev. D **52**, 1780 (1995).
[29] G. G. Raffelt, Lect. Notes Phys. **741**, 51 (2008).
[30] C. Gross, A. Polosa, A. Strumia, A. Urbano, and W. Xue, Phys. Rev. D **98**, 063005 (2018).
[31] N. Isgur and G. Karl, Phys. Rev. D **19**, 2653 (1979); **23**, 817(E) (1981).
[32] B. D. Day, Rev. Mod. Phys. **39**, 719 (1967).
[33] A. Faessler, A. J. Buchmann, and M. I. Krivoruchenko, Phys. Rev. C **56**, 1576 (1997).
[34] B. Povh and J. Hufner, Phys. Lett. B **245**, 653 (1990).
[35] P. Demorest, T. Pennucci, S. Ransom, M. Roberts, and J. Hessels, Nature (London) **467**, 1081 (2010).
[36] J. Antoniadis *et al.*, Science **340**, 1233232 (2013).
[37] A. Bauswein, O. Just, H.-T. Janka, and N. Stergioulas, Astrophys. J. **850**, L34 (2017).
[38] E. Annala, T. Gorda, A. Kurkela, and A. Vuorinen, Phys. Rev. Lett. **120**, 172703 (2018).
[39] E. R. Most, L. R. Weih, L. Rezzolla, and J. Schaffner-Bielich, Phys. Rev. Lett. **120**, 261103 (2018).
[40] J. P. Lees *et al.* (*BABAR* Collaboration), arXiv:1810.04724.
[41] K. Morita, T. Furumoto, and A. Ohnishi, Phys. Rev. C **91**, 024916 (2015).
[42] J. Haidenbauer, U.-G. Meiner, and S. Petschauer, Nucl. Phys. **A954**, 273 (2016).
[43] K.-W. Li, T. Hyodo, and L.-S. Geng, Phys. Rev. C **98**, 065203 (2018).