A Deeply Bound Dibaryon is Incompatible with Neutron Stars and Supernovae

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(Dated: September 19, 2018)

We study the effect of a dibaryon, S, in the mass range 1860 MeV < mS < 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 \( \sim O(10) \) s. A deeply bound dibaryon is therefore incompatible with the observed supernova explosion, unless the cross section for S production is severely suppressed.

I. INTRODUCTION

The possibility that six light quarks form the QCD bound state \( uuddss \), known as the \( H \) dibaryon with binding energy \( B_H \equiv 2m_\Lambda - m_H \geq 0 \), has been considered for several decades [1]. The suggestion that a much more deeply bound state \( \tilde{S} \) called the \( S \) sexaquark [3, 4], with \( B_S \equiv 2m_\Lambda - m_S \geq m_\Lambda - (m_p + m_e) = 176.9 \text{ MeV} \) and which nontrivially avoids observational bounds [5], deserves further scrutiny. Although lattice studies support the existence of a weakly bound dibaryon with \( B_H \sim O(10) \) MeV [6–9], we distinguish this from the more tightly bound and thus stable or cosmologically metastable sexaquark with \( B_S \sim O(\text{few} \times 100) \) MeV, which cannot be ruled out at the current level of understanding of lattice systematics [10] or dedicated searches [11].

In this work, we assume that the \( S \) is light enough to be metastable and massive enough that it cannot be produced as a fusion product of two nucleons. This gives the constrained mass range 1860 MeV < \( m_S \) < \( m_\Lambda + m_p + m_e \approx 2054 \) MeV, which in turn implies

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

Among other implications, this ineluctably leads to the conclusion that the production of dibaryons from \( \Lambda \) baryons is on-shell and exothermic. We study the implications of the production of such a deeply bound QCD state in hot proto-neutron stars, and 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 [12, 13]. Further, since temperatures of the order of tens of MeVs are sustained for a period of about 10 seconds – a time scale set by neutrino diffusion from the proto-neutron star [14] – 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 \( N = n,p, \Lambda \) states; charge conservation is implicit throughout. \( \Lambda \)'s can be produced either by the leptonic process \( e^- + p \rightarrow \Lambda + \nu_e \), or by the non-leptonic process \( NN \rightarrow N\Lambda \) and \( n\pi \rightarrow \Lambda \). 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

\[
\frac{\dot{n}_a}{n_a} = (\text{rate of } a \text{ production per unit volume}) - (\text{rate of } a \text{ disappearance per unit volume}).
\]

Because baryon number \( B \) is conserved, we expect that the rate of \( N \) decay (production) is proportional to \( n_N (n_\Lambda) \), and vice versa. With this in mind, we write:

\[
\dot{n}_N = -n_N^2 \langle \sigma_{NN\rightarrow\Lambda N}\nu \rangle - n_N n_\pi \langle \sigma_{N\pi\rightarrow\Lambda}\nu \rangle + n_\Lambda n_N \langle \sigma_{N\Lambda\rightarrow NN}\nu \rangle \tag{2a}
\]

\[
\dot{n}_\Lambda = +n_N^2 \langle \sigma_{NN\rightarrow\Lambda N}\nu \rangle + n_N n_\pi \langle \sigma_{N\pi\rightarrow\Lambda}\nu \rangle - n_\Lambda^2 \langle \sigma_{\Lambda\Lambda\rightarrow NN}\nu \rangle - 2n_\Lambda \langle \sigma_{\Lambda\Lambda\rightarrow SX}\nu \rangle + 2n_\Lambda n_X \langle \sigma_{SX\rightarrow\Lambda\Lambda}\nu \rangle \tag{2b}
\]

\[
\dot{n}_S = +n_N^2 \langle \sigma_{NN\rightarrow\Lambda N}\nu \rangle - n_\Lambda n_X \langle \sigma_{SX\rightarrow\Lambda\Lambda}\nu \rangle, \tag{2c}
\]

where \( \langle \sigma,\nu \rangle \) indicates the thermally averaged cross section times velocity for the process \( i \); we discuss the values of the various \( \langle \sigma,\nu \rangle \) in the ensuing sections. The particle \( X \) in the process \( \Lambda\Lambda \rightarrow SX \) is chosen to conserve strong isospin [15]. 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)
through (2c) since \( \tilde{B} \propto n_N + n_\Lambda + 2n_S = 0 \). We use initial conditions \( n_N(t = 0) = n_0 \), \( n_\tau(t = 0) = T^3 \exp(-m_\tau/T) \), and \( n_S(t = 0) = n_S(t = 0) = 0 \). We assume that the core has a constant temperature \( T = 30 \) MeV and is at the nuclear saturation density \( n_0 = 0.16 \) fm\(^{-3} \). The \( N \to \Lambda \) and \( \Lambda \to N \) transition rates in Eqs. (2a) and (2b) each contain two contributions. Because the \( \pi \) 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 \( \Lambda \) lifetime in the medium, denoted \( \tilde{\tau}_\Lambda \). This is true in vacuum, where \( \tilde{\tau}_\Lambda \approx 2.6 \times 10^{-10} \) s, but in a dense medium we expect that direct \( \Lambda \) decay is affected by Pauli blocking; we find that the decay width is reduced, \( \tilde{\tau}_\Lambda \approx 4\tilde{\tau}_\Lambda \). Because \( N\Lambda \) collisions are so frequent, \( \Lambda \) disappearance can be dominated by a process analogous to collisional de-excitation, e.g., \( N\Lambda \to NN \) may be more rapid than spontaneous decay. For the nucleon densities we consider, \( n_N(\sigma_{NN\to NN}) \gtrsim \tilde{\tau}_\Lambda^{-1} \) if \( (\sigma_{NN\to NN}) \gtrsim 10^{-29} \) cm\(^2\) / 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 \( n_\gamma(E_\gamma > B_S) \sim T^3 \exp(-B_S/T) \). 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 \( n_S \approx n_N \). 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) through (2c) approximately yields to analytic solution. First, consider the limiting scenario (\( \sigma_{NN\to NN} \to 0 \)). It is clear that \( n_N, n_\Lambda \) reach an equilibrium where \( \tilde{n}_\Lambda = \tilde{n}_N = 0 \) when the \( \Lambda \) abundance has increased to

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

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

\[
n_\Lambda(t) = \tilde{n}_\Lambda \frac{2\tan(\gamma t/2)}{\tan(\gamma t/2) + \sqrt{1 + r}}.
\]

with \( \gamma \equiv (\tilde{\tau}_\Lambda^{-1} + n_N(\sigma_{NN\to NN}))\sqrt{1 + r} \) and \( r \equiv 8\tilde{n}_\Lambda(\sigma_{NN\to S\gamma}) / \tilde{\tau}_\Lambda^{-1} + n_N(\sigma_{NA\to N\Lambda}) \).

The asymptotic \( \Lambda \) abundance is \( n_\Lambda^\infty \equiv \tilde{n}_\Lambda(t \gg \gamma^{-1}) = 2\tilde{n}_\Lambda/1 + \sqrt{1 + r} \), 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 \( n_\Lambda^\infty \), Eq. (2a) dictates that the \( S \) abundance will rise linearly as long as fission is unimportant, \( n_S(t) \propto n_\Lambda(t) \). 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 \( \tilde{n}_S(t_S) = n_0 \), to an accuracy of 10%, or

\[
t_S = \frac{n_0}{(n_\Lambda^\infty)^2 \langle \sigma_{NN\to N\Lambda} \rangle}. \tag{5}
\]

Plugging \( n_\Lambda^\infty \) into Eq. (5) and assuming a hierarchy of rates: \( \tilde{n}_\Lambda(\sigma_{NN\to S\gamma}) \ll n_0(\sigma_{NA\to N\Lambda}) \sim \tilde{\tau}_\Lambda^{-1} \), we find that \( S \) production equilibrates at a time \( t_S \approx s \left[ 4 \times 10^{-31} \mathrm{cm}^3 / s \right] \left[ 1 + 2 \times 10^{-32} \mathrm{cm}^3 / s \right]^{-2} \). 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, \( \gamma S \to \Lambda \) will become important only deep in the back-reacted 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. \( \Lambda \) 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 \( t_\nu \sim \mathcal{O}(10) \) 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 \( (\sigma_{NN\to S\gamma}) \gtrsim 10^{-34} \) cm\(^2\) / s, \( S \) production is fatal for the proto-neutron star. Here, we calculate \( \langle \sigma_{NN\to N\Lambda} \rangle \) and \( \langle \sigma_{NN\to S\gamma} \rangle \).

#### A. \( \Lambda \) Production Cross Section

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

\[
\tau_\Lambda^{-1} \approx \Gamma_{\Lambda\to N\pi} \approx \frac{g_{N\Lambda}^2}{8\pi m_\Lambda} \left[ (m_\Lambda - m_N)^2 - m_\pi^2 \right] \tag{6}
\]

giving \( g_{N\Lambda}^2 \approx 7 \times 10^{-11} \). Because strangeness-changing processes are weak, this small dimensionless number can be interpreted as coming from \( (Gm_\Lambda)^2 \approx 10^{-10} \). Assuming a constant matrix element, appropriate in the limit of small \( m_\pi \), and assuming that the momentum released to the nucleons is large compared to the Fermi momentum, we may write \( \langle \sigma_{NN\to N\Lambda} \rangle \equiv a g_{N\Lambda}^2 \alpha_{N\pi} \sqrt{T/\pi m_\Lambda} m_\Lambda \approx a \times 10^{-22} \) cm\(^3\) / s, where \( \alpha_{N\pi} \approx 15 \) and \( a \) is a function of temperature and density that parameterizes 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
App. \ref{app: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 $\langle \sigma_{NN \rightarrow AN} v \rangle \simeq 10^{-20}$ cm$^3$/s. To be conservative, we will use $\langle \sigma_{NN \rightarrow AN} v \rangle = 3 \times 10^{-30}$ cm$^3$/s as a default value for the rest of this note, 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 $m_s$-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 $\langle \sigma_{NN \rightarrow AN} v \rangle$ 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 $\Xi$ population experiences a Boltzmann suppression such that $n_\Xi / n_N \lesssim (n_\Lambda / n_N)^2$. 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 $\Lambda \Lambda \rightarrow S \gamma$. 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

$$\mathcal{L} \supset d_\Lambda \bar{\Lambda} \sigma^{\mu \nu} A_F \mu \nu + g_{\Lambda S} \bar{\Lambda} \Lambda S \mathbb{1} + \text{h.c.},$$

(7)

where the dipole moment $d_\Lambda = -0.613 \pm 0.001 \mu_N \simeq 10^4$ MeV$^{-1}$, $\Lambda^c$ is the $\Lambda$ 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

$$\langle \sigma_{\Lambda \Lambda \rightarrow S \gamma} v \rangle \simeq 3 \times 10^{-23} \frac{g_{\Lambda S}^2 B_S}{176.9 \text{MeV}} \frac{T}{30 \text{MeV}} \text{cm}^3 / \text{s},$$

(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 wavefunction overlap [13, 19]. We will follow these works and use the Isgur-Karl [20] and Brueckner-Bethe-Goldstone [21] models to calculate the overlap of the $\Lambda$ and the $S$. This is, of course, only one model of the complicated nuclear quantum mechanics involved.

As discussed in more detail in App. \ref{app:B}, the wavefunction overlap has a striking dependence on the dibaryon radius, $r_S$, and the $\Lambda$ radius, $r_\Lambda$. 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 [11]. This gives

$$r_S \geq \frac{1}{m_S} + \frac{x}{m_\pi} = 0.1 \text{fm} \frac{2054 \text{MeV}}{m_S} + 0.34 x \text{fm}.$$  \hspace{1cm} (9)

We will show results for $x = 0, 0.1$ in our final plots. If instead we required that the non-relativistic 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 $m_\pi \leq m_S$ would translate to a bound $r_S \geq 0.22 \text{fm} \sqrt{2054 \text{MeV}/m_S}$. Taking a constituent quark mass $m_\pi \simeq m_S/6$, we would have $r_S \gtrsim 0.53 \text{fm}$. This latter value roughly matches the constituent quark Compton wavelength, $6/m_\pi \gtrsim 0.58 \text{fm}$. For this reason, restricting to the range $0.1 \text{fm} \leq r_S \leq 1.0 \text{fm}$ is very conservative, and the choice $r_S \approx 0.1 \text{fm}$ would be an extremely novel feature for a QCD bound state.

Likewise, the $\Lambda$ radius carries some uncertainty. It is reasonable to assume that increasing strangeness leads to a more compact baryon, $r_\Lambda \lesssim r_\pi$. The strong interaction radius extracted from experimental data $\sqrt{r_\Lambda^2} = 0.76 \pm 0.01 \text{fm}$ [22] is somewhat larger than the naive value in the constituent quark model, $r_\Lambda \simeq |2A_{QCD} m_\Lambda/3|^{-1/2} \approx 0.51 \text{fm}$. Being cautious once again, we decide to show the relatively wide range $0.5 \leq r_\Lambda \leq 0.8 \text{fm}$, where the lower limit is chosen to account for the possibility that the $\Lambda$ charge radius is smaller than the strong interaction radius.

**IV. FATE OF THE PROTO-NEUTRON STAR**

We show our final results in Fig. \ref{fig:1}, fixing $\langle \sigma_{NN \rightarrow AN} v \rangle = 3 \times 10^{-30}$ cm$^3$/s. The left panel of Fig. \ref{fig:1} depicts the lifetime as a function of $r_S$, fixing $r_\Lambda = 0.1 \text{fm}$. The strong interaction radius extracted from experimental data $\sqrt{r_\Lambda^2} = 0.76 \pm 0.01 \text{fm}$ [22] is somewhat larger than the naive value in the constituent quark model, $r_\Lambda \simeq |2A_{QCD} m_\Lambda/3|^{-1/2} \approx 0.51 \text{fm}$. Being cautious once again, we decide to show the relatively wide range $0.5 \leq r_\Lambda \leq 0.8 \text{fm}$, where the lower limit is chosen to account for the possibility that the $\Lambda$ charge radius is smaller than the strong interaction radius.
In the right panel of Fig. 1 we depict $t_S$ for $r_\Lambda = 0.76 \text{ fm}$ 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 $^{16}\text{O}$ nuclei unstable \cite{19}. In the purple shaded region, and at larger masses, the dibaryon cannot possibly be cosmologically metastable, since it has a singly weak decay \cite{2}. In the dark (light) gray region, $r_S$ violates Eq. \eqref{eq:9} for $x = 0 (0.1)$. In all of the heretofore phenomenologically viable parameter space, we find that $t_S \ll 10 \text{ s}$, unless $r_S$ is very close the minimum value allowed by Eq. \eqref{eq:9}.

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 wavefunction overlap calculation, we find that this is possible only if the $S$ radius is very close to its Compton wavelength $\approx 0.1 \text{ fm}$. 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 $\nu \approx 10 \text{ s}$, unless $r_S$ is very close the minimum value allowed by Eq. \eqref{eq:9}.

For instance, the fact that $S$ dibaryons do not couple to a weak current is certainly in conflict with the fact that a neutrino cooling phase of $\approx O(10 \text{ s})$ was observed from the hot interior of Supernova 1987A \cite{25–27}.

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

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig1}
\caption{\textbf{Left:} Contours of $t_S$ as defined in Eq. \eqref{eq:5} for $m_S = 1900 \text{ MeV}$ as a function of the $\Lambda$ and $S$ sizes. The gray region violates Eq. \eqref{eq:9} for $x = 0.0.1$. \textbf{Right:} Contours of $t_S$ for $r_\Lambda = 0.76 \text{ fm}$. In both panels, we have assumed $\langle \sigma_{\Lambda\Lambda\rightarrow NNv} \rangle = 3 \times 10^{-30} \text{ cm}^3 / \text{s}$. The gray region violates Eq. \eqref{eq:9} for $x = 0.0.1$. In the blue region, $^{16}\text{O}$ nuclei are destabilized. In the purple region, the dibaryon has a singly weak decay. All of the parameter space depicted in each panel has $t_S \ll 10 \text{ s}$, and is thus ruled out by the observation that SN1987A continued to emit neutrinos for $t_\nu \approx 10 \text{ s}$, unless $r_S$ is very close the minimum value allowed by Eq. \eqref{eq:9}.}
\end{figure}
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. SDM was supported by Fermi Research Alliance, LLC under Contract No. DE-AC02-07CH11359 with the U.S. Department of Energy, Office of Science, Office of High Energy Physics. The United States Government retains and the publisher, by accepting the article for publication, acknowledges that the United States Government retains a non-exclusive, paid-up, irrevocable, world-wide license to publish or reproduce the published form of this manuscript, or allow others to do so, for United States Government purposes. SR and SS are supported by Department of Energy grant DE-FG02-00ER41132.

Note Added: As our paper was being finalized, we received a draft of [28], which critically addresses the possibility that the $S$ is present in the QCD spectrum, which makes it complementary to the present work.

Appendix A: $NN \rightarrow \Lambda N$ Calculation

The cross section for $\langle \sigma_{NN \rightarrow \Lambda N} \rangle$ determines the equilibrium $\Lambda$ abundance, which in turn determines $t_s$. Assuming a trivial matrix element for single pion exchange and integrating over non-degenerate phase space, in agreement with calculations of nucleon-nucleon scattering in the single-pion-exchange limit [10–13], gives $\langle \sigma_{NN \rightarrow \Lambda N} \rangle \equiv g_{\Lambda N}^3 \sigma_{NN} \sqrt{T/\pi m_{\pi} m_{\Lambda}^2} \approx \times 10^{-27} \text{cm}^3/\text{s}$, where $g_{\Lambda N}$ is obtained from Eq. (6) and $\alpha_{N\pi} \approx 15$. Effects of degeneracy are expected to be mild in this environment [18], but should have effects at the $\sim O(1)$ level [17]. Here we confirm this expectation with explicit calculation.

The rate per unit volume for production of $\Lambda$ baryons in $NN$ collisions is

$$\frac{\Gamma}{\text{Vol}} = \int \prod_{i=1}^{4} \frac{d^3p_i}{(2\pi)^3 2E_i} \delta(4)(p_1 + p_2 - p_3 - p_4) \times f(N_1) f(N_2) [1 - f(N_3)] [1 - f(A_4)] |M_{NN \rightarrow \Lambda N}|^2,$$

where $f(B_i) = \exp \left[ (E_i - \mu_\ell)/T \right] + 1 \right]^{-1}$ is the Fermi-Dirac distribution function for the baryon $i$. The matrix element $M_{NN \rightarrow \Lambda N}$ follows from the Lagrangian $\mathcal{L} \supset g_{NN\pi} N^\pi \bar{N} \pi + g_{\Lambda N\pi} \Lambda^\pi \bar{N} \pi + (\text{h.c.})$, where $g_{NN\pi}$ is given by the Goldberger-Treiman relation. The chemical potential and temperature are related by the requirement that $n_\Lambda = \int \frac{2\pi^2 p^3}{(2\pi)^3} f(N_\Lambda)$. The chemical potentials satisfy $\mu_\Lambda = \mu_N$ by detailed balance. We find that $\mu_\Lambda \gtrsim m_N$, 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

$$\langle \sigma_{NN \rightarrow \Lambda N} \rangle \equiv \frac{\Gamma}{\text{Vol}} n_\Lambda^2.$$

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

Appendix B: Wavefunction Overlap Calculation

Following [2], we integrate the Isgur-Karl wavefunctions of two initial-state baryons against a relative wavefunction that incorporates the $\Lambda - \Lambda$ potential. In agreement with [2] [19], we have

$$g_{\Lambda S}^{\text{(ovp.)}} = 32 \left( \frac{3}{2\pi} \right)^{3/4} \frac{(r_S/r_\Lambda)^{9/2}}{1 + (r_S/r_\Lambda)^2} r_\Lambda^{-3/2} \times \int d^3\alpha \psi_{rel}^2 \psi_{\Lambda} \exp^{-3\alpha^2/4r_\Lambda^2},$$

where $\psi_{\Lambda}$ has mass dimension $-3/2$. We assume that the $\gamma$ 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 $\Lambda \Lambda \rightarrow S\pi\pi$ or $N\Xi \rightarrow S\pi$, the presence of the $\pi$ 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 vacuum, which is at a distance much larger than $r_S$, meaning that the wavefunctions need not coincide as exactly as in our model calculation, and the cross section may be as large as $m_\pi^{-2}$. However, such effects are difficult to quantify in the absence of a calculable model of hadronization, so we restrict to $\Lambda \Lambda \rightarrow S\pi$, 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 $\psi_{rel}$, we use the relative wavefunctions depicted in Fig. 5 of [29]. These wavefunctions are generated from potentials calibrated on the Nagara 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 [30] [31]. Needless to say, an attractive potential would lead to a relative wavefunction 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 $\Lambda\Lambda \rightarrow S\gamma$ event may change strangeness while they are within range of each other’s potential. Thus, the correct relative wavefunction may be a linear combination of relative $\Lambda - N$ and $\Lambda - \Lambda$ wavefunctions. For this reason, the slightly repulsive potentials of [29] provide a conservative model of this process.

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

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