Baryon-Number Nonconservation and
the Stability of Strange Matter

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Abstract

If baryon-number nonconservation exists in the form of an effective six-quark operator which also changes strangeness by four units, it could have a tremendous impact on the absolute stability of strange quark matter. We show that such an operator is negligible in the supersymmetric standard model with $\lambda_{ijk} u^c_i d^c_j d^c_k$ terms in the superpotential, but may be of importance in models with exotic particle content. From the experimental lower limit of $10^{25}$ years on the stability of nuclei, we find a model-independent lower limit of the order $10^5$ years on the stability of strange matter against such decays.
1 Introduction

Matter containing a large number of strange quarks may have a lower energy per baryon than ordinary nuclei and be absolutely stable\[1, 2\]. This intriguing possibility has generated a great deal of interest across many subfields of physics. A crucial implicit assumption for the stability of this new kind of matter, called strange matter, is that the baryon number $B$ is exactly conserved. Of course, we know nuclei are stable against $\Delta B \neq 0$ decays. The best mode-independent experimental lower limit to date is $1.6 \times 10^{25}$ years\[3\]. However, there may be effective $\Delta B \neq 0$ interactions which are highly suppressed for nuclei but not for strange matter. We examine such a hypothesis here and show that whereas such exotic decays are possible, the relevant lifetime should be longer than $10^5$ years.

In Sec. 2 we review briefly the present experimental constraints on $\Delta B \neq 0$ interactions. These come mainly from the nonobservation of proton decay and of neutron-antineutron oscillation. We then point out the possible consequences of an effective $\Delta B = 2$, $\Delta N_s = 4$ operator (where $N_s$ denotes the number of $s$ quarks) on the stability of strange matter. We proceed to discuss some possible theoretical origins of such an operator in Sec. 3 and Sec. 4.

We deal first with the supersymmetric standard model with $R$-parity nonconserving terms of the form $\lambda_{ijk} u^c_i d^c_j d^c_k$ and show that their effects are negligible. We then discuss two possible extended models where the effects may be large. In Sec. 5 we consider this operator in conjunction with the usual weak interaction and show that the suppression coming from the stability of ordinary nuclei translates to a phenomenological lower limit of $10^5$ years on the lifetime of strange matter. Finally in Sec. 6, there are some concluding remarks.
2 Baryon-Number Nonconserving Interactions

Given the standard $SU(3) \times SU(2) \times U(1)$ gauge symmetry and the usual quarks and leptons, it is well-known that the resulting renormalizable Lagrangian conserves both the baryon number $B$ and the lepton number $L$ automatically. However, new physics at higher energy scales may induce effective interactions (of dimension 5 or above) which do not respect these conservation laws. The foremost example is the possibility of proton decay. For this to happen, there has to be at least one fermion with a mass below that of the proton. Since only leptons are known to have this property, the selection rule $\Delta B = 1, \Delta L = 1$ is applicable. The most studied decay mode experimentally is $p \to \pi^0 e^+$, which requires the effective interaction

$$\mathcal{H}_{\text{int}} \sim \frac{1}{M^2} (uud) (uud).$$

(1)

It is now known that $\tau_{\text{exp}} (p \to \pi^0 e^+) > 9 \times 10^{32}$ years, hence $M > 10^{16}$ GeV. [Here $M$ should be considered an effective mass, because in some scenarios the denominator of the right-hand side of Eq. (1) is really the product of two different masses.] This means that proton decay probes physics at the grand-unification energy scale.

The next simplest class of effective interactions has the selection rule $\Delta B = 2, \Delta L = 0$. The most well-known example is of course

$$\mathcal{H}_{\text{int}} \sim \frac{1}{M^5} (udd)^2,$$

(2)

which induces neutron-antineutron oscillation and allows a nucleus to decay by the annihilation of two of its nucleons. Note that since 6 fermions are now involved, the effective operator is of dimension 9, hence $M$ appears to the power $-5$. This means that for the same level of nuclear stability, the lower bound on $M$ will be only of order $10^5$ GeV. The present best experimental lower limit on the $n - \bar{n}$ oscillation lifetime is $8.6 \times 10^7$ s. Direct search for $NN \to \pi\pi$ decay in iron yields a limit of $6.8 \times 10^{30}$ y. Although the above two
numbers differ by 30 orders of magnitude, it is well established by general arguments as well as detailed nuclear model calculations that

\[ \tau(n\bar{n}) = 8.6 \times 10^7 \text{s} \Rightarrow \tau(NN \rightarrow \pi\pi) \sim 2 \times 10^{31} \text{y}. \]  

(3) 

Hence the two limits are comparable. To estimate the magnitude of \( M \), we use

\[ \tau(n\bar{n}) = M^5 |\psi(0)|^{-4}, \]  

(4) 

where the effective wavefunction at the origin is roughly given by

\[ |\psi(0)|^{-2} \sim \pi R^3, \quad R \sim 1 \text{ fm}. \]  

(5) 

Hence we obtain \( M > 2.4 \times 10^5 \) GeV. Bearing in mind that \( \mathcal{H}_{\text{int}} \) is likely to be suppressed also by products of couplings less than unity, this means that the stability of nuclei is sensitive to new physics at an energy scale not too far above the electroweak scale of \( 10^2 \) GeV. It may thus have the hope of future direct experimental exploration.

Consider next the effective interaction

\[ \mathcal{H}_{\text{int}} \sim 1 \frac{1}{M^5} (uds)^2. \]  

(6) 

This has the selection rule \( \Delta B = 2, \Delta N_s = 2 \) (\( \Delta S = -2 \)). In this case, a nucleus may decay by the process \( NN \rightarrow KK \). Although such decay modes have never been observed, the mode-independent stability lifetime of \( 1.6 \times 10^{25} \) y mentioned already is enough to guarantee that it is very small. However, consider now

\[ \mathcal{H}_{\text{int}} \sim 1 \frac{1}{M^5} (uss)^2, \]  

(7) 

which has the selection rule

\[ \Delta B = 2, \quad \Delta N_s = 4. \]  

(8)
To get rid of four units of strangeness, the nucleus must now convert two nucleons into four kaons, but that is kinematically impossible. Hence the severe constraint from the stability of nuclei does not seem to apply here, and the following interesting possibility may occur.

Strangelets (i.e. stable and metastable configurations of quarks with large strangeness content) with atomic number $A$ (which is of course the same as baryon number $B$) and number of strange quarks $N_s$ may now decay into other strangelets with two less units of $A$ and one to three less $s$ quarks. For example,

\begin{align}
(A, N_s) & \to (A - 2, N_s - 2) + KK, \\
(A, N_s) & \to (A - 2, N_s - 1) + KKK, \text{ etc.}
\end{align}

For the states of lowest energy, model calculations show\cite{11} that $N_s \sim 0.8A$, hence the above decay modes are efficient ways of reducing all would-be stable strangelets to those of the smallest $A$.

Unlike nuclei which are most stable for $A$ near that of iron, the energy per baryon number of strangelets decreases with increasing $A$. This has led to the intriguing speculation that there are stable macroscopic lumps of strange matter in the Universe. On the other hand, it is not clear how such matter would form, because there are no stable building blocks such as hydrogen and helium which are essential for the formation of heavy nuclei. In any case, the above exotic interaction would allow strange matter to dissipate into smaller and smaller units, until $A$ becomes too small for the strangelet itself to be stable. A sample calculation by Madsen\cite{11} shows that

\begin{equation}
m_0(A) - m_0(A - 2) \simeq (1704 + 111A^{-1/3} + 161A^{-2/3}) \text{ MeV},
\end{equation}

assuming $m_s = 100$ MeV, and the bag factor $B^{1/4} = 145$ MeV. Hence the $KK$ and $KKK$ decays would continue until the ground-state mass $m_0$ exceeds the condition that the energy per baryon number is less than 930 MeV, which happens at around $A = 13$. 

5
3 Supersymmetric Standard Model

If the standard model of quarks and leptons is extended to include supersymmetry, the $\Delta B \neq 0$ terms $\lambda_{ijk} u_i^c d_j^c d_k^c$ are allowed in the superpotential. In the above notation, all chiral superfields are assumed to be left-handed, hence $q^c$ denotes the left-handed charge-conjugated quark, or equivalently the right-handed quark, and the subscripts refer to families, i.e. $u_i$ for $(u, c, t)$ and $d_j$ for $(d, s, b)$. Since all quarks are color triplets and the interaction must be a singlet which is antisymmetric in color, the two $d$ quark superfields must belong to different families. In the minimal supersymmetric standard model, these terms are forbidden by the imposition of $R$-parity. However, this assumption is not mandatory and there is a vast body of recent literature exploring the consequences of $R$-parity nonconservation[12].

Starting with Yukawa terms of the form $\lambda_{ijk} u_i^c d_j^c d_k^c$, where two of the fields are quarks and the third is a scalar quark, we can obtain the effective interaction of Eq. (7) in several ways[13]. In Ref. [10] the $dd \rightarrow \tilde{b}\tilde{b}$ box diagram (where $\tilde{q}$ denotes the supersymmetric scalar partner of $q$) is considered to obtain the effective interaction of Eq. (2) for $n-\bar{n}$ oscillation using $\lambda_{udb}$. Here we take

$$ss \rightarrow \tilde{b}\tilde{b}$$

and use $\lambda_{usb}$ instead. The box diagram involves the exchange of the $(u, c, t)$ quarks, the $W$ boson, and their supersymmetric partners. Since only gauge couplings appear, the Glashow-Iliopoulos-Maiani (GIM) mechanism[14] is operative and this diagram vanishes if the $(\tilde{u}, \tilde{c}, \tilde{t})$ scalar quarks have the same mass and that the $(u, c, t)$ quark masses can be neglected. Thus in Fig. 3 of Ref. [10], there is a peak at $M_{\tilde{t}} = 200$ GeV. However, there are actually two scalar quarks $(\tilde{q}_L, \tilde{q}_R)$ for each quark $q$. Consider just $t$ and $(\tilde{t}_L, \tilde{t}_R)$. The $\tilde{t}$ mass matrix is given by

$$M^2_{\tilde{t}} = \begin{bmatrix} m^2_{\tilde{t}_L} + m^2_t & Am_{\tilde{t}} \\ Am_{\tilde{t}} & m^2_{\tilde{t}_R} + m^2_t \end{bmatrix}.$$  

(13)
Since only $\tilde{t}_L$ couples to $b_L$ through the gauge fermion $\tilde{w}$ and $\tilde{t}_L$ is not a mass eigenstate, its approximate effective contribution is given by

$$M_{\tilde{t}_L} = \frac{\bar{m}_R^2 + m_t^2}{(\bar{m}_L^2 + m_t^2)(\bar{m}_R^2 + m_t^2) - A^2 m_t^2}.$$  \hfill (14)$$

Unless this somehow cancels the $\tilde{u}$ and $\tilde{c}$ contributions accidentally, the box diagram will not vanish. Furthermore, it is often assumed that the soft supersymmetry-breaking terms \(\tilde{m}_L^2\) and \(\tilde{m}_R^2\) are universal, in which case this amplitude would be zero if the $(u, c, t)$ quark masses were neglected. Rewriting Eq. (7) more specifically as

$$H_{int} = T_1 \epsilon_{\alpha\beta\gamma} \epsilon_{\alpha'\beta'\gamma'} u_{R\alpha} s_{R\beta} L\gamma u_{R\alpha'} s_{R\beta'} L\gamma',$$  \hfill (15)$$

where the color indices and chiralities of the quarks are noted, the contribution of Eq. (12) is then given by [13]

$$T_1 \approx \frac{g^4 \lambda_{usb} A^2 m_R^2 m_R}{32\pi^2 \bar{m}^4} \rho^2 F(m_R^2, M_W^2, \bar{m}^2, m_t, A),$$  \hfill (16)$$

where we have assumed universal soft supersymmetry-breaking scalar masses, $V_{ts}$ is the quark-mixing entry for $t$ to $s$ through the $W$ boson, and

$$F = \frac{1}{2} J(m_R^2, M_W^2, \bar{m}^2 + m_t^2 + A m_t, m_t^2) + \frac{1}{2} J(m_R^2, M_W^2, m_t^2 + m_t^2 - A m_t, m_t^2)$$

$$- J(m_R^2, M_W^2, \bar{m}^2, m_t^2) - \{m_t^2 \rightarrow 0\ \text{in the last entry of each term}\},$$  \hfill (17)$$

with the function $J$ given by [10, 13]

$$J(a_1, a_2, a_3, a_4) = \sum_{i=1}^4 \frac{a_i^2 \ln(a_i)}{\prod_{k \neq i} (a_i - a_k)}.$$  \hfill (18)$$

Using the correspondence of $n - \bar{n}$ oscillation to $NN$ annihilation inside a nucleus, we estimate the lifetime of strangelets from the above effective interaction assuming $A = 200$ GeV, $\lambda_{usb} < 1$, and $\bar{m} > 200$ GeV, to be

$$\tau > 10^{22} \text{y}.$$  \hfill (19)$$
This tells us that such contributions are negligible from the supersymmetric standard model.

Recently, another contribution to Eq. (2) has been identified involving $\lambda_{tds}$ and $\lambda_{tdb}$ without the GIM suppression of Eq. (16). The form of its contribution to Eq. (7) is

$$\mathcal{H}_{\text{int}} = T_2 \epsilon_{\alpha \beta \gamma} \epsilon_{\alpha' \beta' \gamma'} u_L s L_{\alpha} s R_{\gamma} u_L s L_{\beta} s R_{\gamma},$$

which is not the same as Eq. (15). The important coupling is now $\lambda_{tsb}$, but $T_2$ is also suppressed relative to $T_1$ by $V_{ub}^2$, hence this contribution to Eq. (7) is even more negligible.

4 Extended Models Involving Strangeness

Although the supersymmetric standard model cannot have a sizable contribution to the $\Delta B = 2, \Delta N_s = 4$ operator of Eq. (7), an extended model including additional particles belonging to the fundamental $27$ representation of $E_6$, inspired by superstring theory, may do better. Consider a slight variation of the model of exotic baryon-number nonconservation proposed some years ago. In addition to the usual quark and lepton superfields

$$Q \sim (3, 2, 1/6), \quad u^c \sim (\bar{3}, 1, -2/3), \quad d^c \sim (\bar{3}, 1, 1/3);$$

$$L \sim (1, 2, -1/2), \quad e^c \sim (1, 1, 1), \quad N^c \sim (1, 1, 0);$$

we have

$$h \sim (3, 1, -1/3), \quad h^c \sim (\bar{3}, 1, 1/3);$$

$$E \sim (1, 2, -1/2), \quad \bar{E} \sim (1, 2, 1/2), \quad S \sim (1, 1, 0);$$

each transforming under the standard $SU(3) \times SU(2) \times U(1)$ gauge group as indicated. Imposing the discrete symmetry $Z_2 \times Z_2$ under which

$$Q, u^c, d^c, N^c \sim (+, -),$$

$$L, e^c \sim (-, +),$$

$$h, h^c, E, \bar{E}, S \sim (+, +),$$

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we get the allowed terms
\[ QQh, \ u'c'd'h, \ d'hNc, \ Nc'Nc, \] (28)
in the superpotential, but not \(d'h\). We also assume that \(\tilde{N}c\) does not develop a vacuum expectation value, so that \(B\) is broken explicitly by the above \(Nc'Nc\) term only. [Note that without this last term, we can assign \(B = -2/3, 2/3, 1\) to \(h, h^c, Nc\) respectively and \(B\) would be conserved.] This model allows us to obtain an effective \((uss)^2\) operator without going through a loop, as shown in Fig. 1.

Using the formalism of Ref. [10] for the process
\[ (A, N_s) \rightarrow (A - 2, N_s - 2) + KK, \] (29)
we estimate the lifetime to be
\[ \tau \sim \frac{32\pi m_N^2}{9\rho_N} \left[ \frac{M}{\Lambda} \right]^{10} \sim 1.2 \times 10^{-28} \left[ \frac{M}{\Lambda} \right]^{10} \text{y}, \] (30)
where \(m_N \sim 1 \text{ GeV}, \rho_N \sim 0.25 \text{ fm}^{-3}\) is the nuclear density, and \(\tilde{\Lambda} \sim 0.3 \text{ GeV}\) is the effective interaction energy scale corresponding to Eq. (5). If we assume \(M = 1 \text{ TeV}\), then
\[ \tau \sim 2 \times 10^{7} \text{y}. \] (31)
In the above, we have assumed that \(\lambda_{ush}\) is unconstrained phenomenologically. However, consider the term \(Q_1Q_2h\) where \(Q_1 = (u, d')\) and \(Q_2 = (c, s')\) so that \(Q_1Q_2 = us' - cd'\). Hence from this term alone, \(\lambda_{udh} = (V_{cd}/V_{cs})\lambda_{ush}\). Since \(\lambda_{udh}\) may now combine with \(\lambda_{s'h,Nc}\) to induce \(NN \rightarrow KK\) decays, its magnitude is seriously suppressed and \(\lambda_{ush}\) would be too small for Eq. (31) to be valid. However, there is also a third generation, allowing us the \(Q_1Q_3 = ub' - td'\) term which may then be fine-tuned to eliminate the \(\lambda_{udh}\) component, thus saving Eq. (31). This is of course not a very natural solution, but cannot otherwise be ruled out.
We may also consider the following tailor-made extension. Let there be a new exotic scalar multiplet $\tilde{Q}_6 \sim (\bar{6}, 1, 2/3)$ with $B = -2/3$ and a new exotic fermion multiplet $\psi_8 \sim (8, 1, 0)$ with $B = 1$. Assume the existence of Yukawa terms $s^c s^c \tilde{Q}_6^*$, $u^c \psi_8 \tilde{Q}_6$, and the $B$-nonconserving Majorana mass term $\psi_8 \psi_8$, then an effective $(uss)^2$ term is possible, as shown in Fig. 2. Note that the $uss$ combination is now a color octet. This may in fact be a more efficient way to dissipate strange matter which is presumably not clumped into color-singlet constituents as in ordinary nuclei.

5 Stability Limit of Strange Matter

Since $\tau$ depends on $M/\Lambda$ to the power 10 in Eq. (30), it appears that a much shorter lifetime than that of Eq. (31) is theoretically possible. However, there is a crucial phenomenological constraint which we have yet to consider. Although two nucleons cannot annihilate inside a nucleus to produce four kaons, they can make three kaons plus a pion. The effective $(uss)^2$ operator must now be supplemented by a weak transition $s \rightarrow u + d + \bar{u}$. We can compare the effect of this on $NN \rightarrow KKK\pi$ versus that of the $(uds)^2$ operator on $NN \rightarrow KK$ discussed in Ref. [10].

First let us look at the phase-space difference. Consider a nucleus with atomic number $A$ decaying into one with atomic number $A - 2$ plus two kaons with energy-momentum conservation given by $p = p' + k_1 + k_2$. The decay rate is proportional to

$$F_1 = \frac{1}{(2\pi)^5} \int \frac{d^3k_1}{2E_1} \int \frac{d^3k_2}{2E_2} \int \frac{d^3p'}{2E'} \delta^4(p - p' - k_1 - k_2)\left(\frac{p'}{E'}\right)^5 \int \frac{d^3k_1}{2E_1} \int \frac{d^3k_2}{2E_2} \delta(M - M' - E_1 - E_2)\right.$$

$$\approx \frac{1}{2M'(2\pi)^5} \int \frac{d^3k_1}{2E_1} \int \frac{d^3k_2}{2E_2} \delta(M - M' - E_1 - E_2)\right.$$

$$\approx \frac{1}{2M'(2\pi)^3} \int_{E_{max}}^{E_{max}} k_1k_2dE_1, \quad (32)$$

where $k_1 = (E_1^2 - m_K^2)^{1/2}$, $k_2 = [(2m_N - E_1)^2 - m_K^2]^{1/2}$, and $E_{max} = 2m_N - m_K$. 

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Consider next $A \rightarrow (A - 2) + KKK\pi$ with $p = p' + k_1 + k_2 + k_3 + k_\pi$. Because of the limited phase space, we assume the kaons to be nonrelativistic. Hence

$$F_2 \approx \frac{1}{2M'(2\pi)^4} \frac{1}{m_K} \frac{1}{2m_K} \frac{1}{2m_K} \frac{1}{2E_\pi} \delta(M - M' - 3m_K - \frac{k_1^2 + k_2^2 + k_3^2}{2m_K} - E_\pi)$$

$$\approx \frac{1}{2M'(2\pi)^7} \frac{1}{m_K^3} \int k_\pi k_1^2 k_2^2 k_3^2 dk_1 dk_2 dk_3,$$  \hspace{1cm} (33)

where $k_\pi = \{[2m_N - 3m_K - (k_1^2 + k_2^2 + k_3^2)/2m_K]^2 - m_\pi^2\}^{1/2}$. The above integral can be evaluated by treating $k_{1,2,3}$ as Cartesian coordinates and then convert them to three-dimensional polar coordinates. The angular integration over the $k_{1,2,3} > 0$ octant yields a factor of $\pi/210$ and we get

$$F_2 \approx \frac{1}{2m_K^7} \frac{1}{m_K} \frac{\pi}{210} \int_{k_{\text{max}}}^{k_{\text{max}}} k_\pi k^8 dk,$$  \hspace{1cm} (34)

where $k_{\text{max}} = [2m_K(2m_N - 3m_K - m_\pi)]^{1/2}$. The effective interaction here also differs from that of Eq. (32) by the appearance of a third kaon and an extra pion which can be thought of as having been converted by the weak interaction from a fourth kaon. We thus estimate the suppression factor to be

$$\xi \sim \left[\frac{f_K^2 T_{K\pi}}{\Lambda^6}\right]^2 \frac{F_2}{F_1} \sim 10^{-20},$$  \hspace{1cm} (35)

where $f_K = 160$ MeV, and $T_{K\pi} = 0.07$ MeV$^2$ is obtained using the symmetric soft-pion reduction [3] of the experimental $K \rightarrow 2\pi$ amplitude. We have again invoked the effective interaction energy scale $\tilde{\Lambda} = 0.3$ GeV used in Eq. (30). Obviously, our estimate depends very sensitively on this parameter, but that is not untypical of many calculations in nuclear physics. Since the stability of nuclei is at least $1.6 \times 10^{25}$ years, the reduction by $\xi$ of the above yields

$$\tau > 10^5 \text{y}$$  \hspace{1cm} (36)

for the lifetime of strangelets against $\Delta B = 2, \Delta N_s = 4$ decays.
6 Concluding Remarks

We have pointed out in this paper that an effective \((uss)^2\) interaction may cause stable strange matter to decay, but the lifetime is constrained phenomenologically by the stability of nuclei against decays of the type \(A \rightarrow (A-2) + KK\pi\) and we estimate it to have a lower limit of \(10^5\) years. This result has no bearing on whether stable or metastable strangelets can be created and observed in the laboratory, but may be important for understanding whether there is stable strange bulk matter left in the Universe after the Big Bang and how it should be searched for. For example, instead of the usual radioactivity of unstable nuclei, strange matter may be long-lived kaon and pion emitters. The propagation of these kaons and pions through the matter itself and their interactions within such an environment are further areas of possible study. Since energy is released in each such decay, although the amount is rather small, it may be sufficient to cause a chain reaction and break up the bulk matter in a cosmological time scale. Furthermore, if the particles mediating this effective interaction have masses of order 1 TeV as discussed, then forthcoming future high-energy accelerators such as the Large Hadron Collider (LHC) at the European Center for Nuclear Research (CERN) will have a chance of confirming or refuting their existence.

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Figure Captions

Fig. 1. Effective \((uss)^2\) interaction from a supersymmetric model with \(E_6\) particle content.

Fig. 2. Effective \((uss)^2\) interaction from a model with exotic color-sextet scalar and color-octet fermion.
This figure "fig1-1.png" is available in "png" format from:

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This figure "fig1-2.png" is available in "png" format from:

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