Baryon-number violation by two units and the deuteron lifetime

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Abstract

We calculate the lifetime of the deuteron with dimension-nine quark operators that violate baryon number by two units. We construct an effective field theory for $|\Delta B| = 2$ interactions that give rise to neutron-antineutron ($n-\bar{n}$) oscillations and dinucleon decay within a consistent power counting. We calculate the ratio of the deuteron lifetime to the square of the $n-\bar{n}$ oscillation time up to next-to-leading order. Our result, which is analytical and has a quantified uncertainty, is smaller by a factor $\simeq 2.5$ than earlier estimates based on nuclear models, which impacts the indirect bound on the $n-\bar{n}$ oscillation time and future experiments. We discuss how combined measurements of $n-\bar{n}$ oscillations and deuteron decay can help to identify the sources of baryon-number violation.
At the classical level the Standard Model (SM) has two accidental global $U(1)$ symmetries associated with baryon-number ($B$) and lepton-number ($L$) conservation [1-3]. At the quantum level only $B - L$ is conserved, while $B + L$ is anomalous. Since it can be expected that all global symmetries are only approximate, it is plausible that beyond-the-SM (BSM) physics violates $B$, $L$, and $B - L$ separately. For instance, extending the SM with the only gauge-invariant dimension-five operator leads to violation of $L$ by two units [1,3]. Additional $B$- and $L$-violating operators appear at the dimension-six level, while the first gauge-invariant operators that violate $B$ by two units ($|\Delta B| = 2$) appear at dimension nine [4].

The best limits on $B$-violating interactions come from the observed stability of the proton. The limit on its lifetime translates into a scale $\Lambda_{|\Delta B|=1} \gtrsim 10^{13}$ TeV for grand unified theories [5]. Such energies are out of reach of colliders. However, models exist wherein $B$ is only violated by two units and the proton is stable [6-9]. These interactions lead to the oscillation of neutral baryons into antibaryons, in analogy to strangeness-changing SM interactions that lead to kaon-antikaon oscillations. In particular, a neutron in a beam can oscillate into an antinucleon [9] that annihilates with a nucleon in a target, producing several pions with a few hundred MeV of energy [10]. An ILL experiment sets a lower limit on the neutron-antineutron ($n-\bar{n}$) oscillation time of $\tau_{n\bar{n}} > 0.86 \cdot 10^8$ s $\simeq 2.7$ yr (90% C.L.) [11], which converts to a BSM scale $\Lambda_{|\Delta B|=2} \gtrsim 10^2$ TeV, within reach of future colliders. An experiment at the European Spallation Source can improve $\tau_{n\bar{n}}$ by two orders of magnitude [12], probing regions of parameter space relevant for the observed baryon asymmetry of the universe [13].

Apart from “in-vacuum” $n-\bar{n}$ oscillations, $|\Delta B| = 2$ interactions also induce the decay of otherwise stable nuclei. A bound neutron can oscillate inside the nucleus into an antineutron, which then annihilates with another nucleon. Since a neutron and an antineutron have very different potential energies, the typical nuclear lifetime is far greater than $\tau_{n\bar{n}}$ [4]. Alternatively, two nucleons can annihilate directly. If $n-\bar{n}$ oscillations are the dominant mechanism, one can calculate how the nuclear lifetime and $\tau_{n\bar{n}}$ are related. This relation was previously obtained from phenomenological models of the nuclear wavefunction and the nucleon-antineutron potential [10], a procedure that suffers from unknown uncertainties.

In this letter, we improve the theory of $|\Delta B| = 2$ interactions in the simplest nucleus, the deuteron, with effective field theory (EFT). EFT allows for all interactions and processes consistent with the symmetries. Figure 1 shows the two classes of diagrams that represent deuteron decay. The left diagram shows a process that converts a neutron into an antineutron in a beam, which then annihilates with another nucleon. Since a neutron and an antineutron have very different potential energies, the typical nuclear lifetime is far greater than $\tau_{n\bar{n}}$ [4]. Alternatively, two nucleons can annihilate directly. If $n-\bar{n}$ oscillations are the dominant mechanism, one can calculate how the nuclear lifetime and $\tau_{n\bar{n}}$ are related. This relation was previously obtained from phenomenological models of the nuclear wavefunction and the nucleon-antineutron potential [10], a procedure that suffers from unknown uncertainties.

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Central to our analysis are the gauge-invariant $|\Delta B| = 2$ operators. As each quark field has $B = 1/3$, operators with at least six quarks are required. Since we expect $\Lambda_{|\Delta B|=2}$ to lie well above the electroweak scale, we focus on the four operators [4,16,18] that are invariant under the full SM gauge group $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$. At the QCD scale $\mu \sim 1$ GeV, we write

$$\mathcal{L}_{|\Delta B|=2} = \mathcal{C}_1 Q_1 + \mathcal{C}_2 Q_2 + \mathcal{C}_3 Q_3 + \mathcal{C}_4 Q_4 + \text{H.c.},$$  \hspace{1cm} (1)
FIG. 1. The two classes of diagrams for deuteron decay. The crossed circle denotes the deuteron. Single lines with arrows to the right (left) denote (anti)nucleon propagators. In the left diagram, the blob depicts a one- or two-nucleon process that converts two nucleons into a nucleon and an antinucleon, and the circle their annihilation into a mesonic final state (dashed lines). In the right diagram, the blob depicts the propagation of two nucleons, and the square their direct annihilation into the same final states.

\[
\begin{array}{c|c|c}
\text{Operator} & \text{Notation of Ref. [16]} & \text{Chiral irrep} \\
\hline
Q_1 & -D_R D_R D_R^T T^{AAS}/4 & O_{RRR}^3 \\
Q_2 & -D_L D_R D_R^+ T^{AAS}/4 & O_{LRR}^3 \\
Q_3 & -D_L D_L D_R^+ T^{AAS}/4 & O_{LLR}^3 \\
Q_4 & -D_R^{33+} T^{SSS}/4 & (O_{RRR}^1 + 4O_{RRR}^2)/5 \\
\end{array}
\]

TABLE I. The independent $|\Delta B| = 2$, SM gauge-invariant, dimension-nine operators with $u$ and $d$ quarks, and the irreducible chiral representations they belong to [19].

where the six-quark operators $Q_i$ are multiplied by complex Wilson coefficients $C_i$ expected to scale as $O(c_i \Lambda_{\chi}^{-5})$, with $c_i$ dimensionless constants. The operators can be expressed via diquark fields [19] as

\[
D_{L,R} \equiv q^iTCP_{L,R}i\tau^2q^i, \quad D_{L,R}^a \equiv q^iTCP_{L,R}i\tau^2\tau^a q^i,
\]

\[
D_{L,R}^{bc} = D_{L,R}^{\{a} D_{L,R}^{b} D_{L,R}^{c\}} - \frac{1}{5} \left( \delta^{ab} D_{L,R}^{\{d} D_{L,R}^{d} D_{L,R}^{c\}} + \delta^{ac} D_{L,R}^{\{d} D_{L,R}^{d} D_{L,R}^{b\}} + \delta^{bc} D_{L,R}^{\{d} D_{L,R}^{d} D_{L,R}^{a\}} \right),
\]

where $q^i = (u^i d^i)^T$ is the quark doublet with color index $i$, $C$ is the charge-conjugation matrix, $P_{L,R}$ are left- and right-handed projectors, $\tau^a$ ($a = 1, 2, 3$) are the Pauli isospin matrices, and $\{\}$ denotes symmetrization. Color singlets are formed by contracting the color indices suppressed on the left-hand side of Eq. (2) with the tensors

\[
T^{SSS} \equiv \varepsilon_{ikm}\varepsilon_{jln} + \varepsilon_{ikn}\varepsilon_{jlm} + \varepsilon_{jkm}\varepsilon_{iln} + \varepsilon_{jkn}\varepsilon_{ilm},
\]

\[
T^{AAS} \equiv \varepsilon_{ikm}\varepsilon_{jln} + \varepsilon_{ikn}\varepsilon_{jlm}.
\]

The resulting gauge-invariant operators are given in Table I where $\tau^\pm = (\tau^1 \pm i\tau^2)/2$.

Low-energy hadronic and nuclear observables such as $\tau_{\bar{n}n}$ and $\Gamma^{-1}$ are difficult to calculate due to the breakdown of the perturbative expansion in the strong coupling constant. We use Chiral EFT ($\chi$EFT) [20, 21], the low-energy EFT of QCD with nucleons and pions as effective degrees of freedom. Pions play an important role as pseudo-Goldstone bosons of the spontaneously-broken, approximate $SU(2)_L \otimes SU(2)_R$ symmetry of QCD. In $\chi$EFT one can calculate observables at momenta $Q \lesssim m_\pi \approx 140$ MeV, the pion mass, in an expansion in powers of $Q/\Lambda_\chi$, where $\Lambda_\chi \sim 2\pi F_\pi \sim m_N$ is the chiral-symmetry-breaking scale, with $F_\pi \approx 185$ MeV the pion decay constant and $m_N \approx 940$ MeV the nucleon mass.
The first step towards the calculation of $\Delta B = 2$ observables is to construct the chiral Lagrangian of QCD supplemented by Eq. (1). The EFT includes all chiral interactions that transform as the terms in this extended Lagrangian. Each term comes with a low-energy constant (LEC) that subsumes the nonperturbative QCD dynamics. These LECs have to be calculated with nonperturbative methods, preferably lattice QCD, or estimated, e.g. by naive dimensional analysis (NDA) [22].

In the single-baryon sector, we write the $B$-conserving Lagrangian for nonrelativistic (anti)nucleons $N = (p, n)^T$, $(N^c = (p^c, n^c)^T)$ interacting with pions $\pi^a$ as

$$
L^{(2)}_{\Delta B=0} = N^\dagger \left( i\partial_0 + \frac{\nabla^2}{2m_N} \right) N + N^{c\dagger} \left( i\partial_0 + \frac{\nabla^2}{2m_N} \right) N^c + \frac{g_A}{F_\pi} (N^{c\dagger} \sigma_k \tau^a N + N^{c\dagger} \sigma_k \tau^a T N^c) \nabla_k \pi^a - \frac{1}{2} \pi^a (\partial^2 + m_\pi^2) \pi^a + \ldots, 
$$

(4)

where $\sigma_k$ ($k = 1, 2, 3$) are the Pauli spin matrices and $g_A \approx 1.27$. Here and below the dots stand for terms that only contribute at higher orders in our calculation.

The chiral Lagrangian relevant for $n\bar{n}$ oscillations has recently been constructed in Refs. [23, 25], viz.

$$
L^{(2)}_{\Delta B=2} = -\delta m n^{c\dagger} n + \text{H.c.} + \ldots, 
$$

(5)

where $\delta m$ is a LEC that can be made real by a $U(1)$ transformation on the nucleon and antinucleon fields. Due to the chiral properties of the operators in Table I only $Q_i$ with $i = 1, 2, 3$ contribute at lowest orders [18], $\delta m$ scaling as $O(c_i \Lambda_\chi^2 F_\pi^2 / \Lambda_\chi^5 |_{\Delta B=2})$. The case $Q_4$ is discussed below. The $n\bar{n}$ oscillation time reads [23, 25]

$$
\tau_{n\bar{n}} = (\delta m)^{-1} \left[ 1 + O(m_\pi^2 / \Lambda_\chi^2) \right]. 
$$

(6)

$\delta m$ has recently been calculated in lattice QCD [26, 27].

To calculate deuteron decay we exploit the fact that the deuteron binding energy is only $B_d \approx 2.225$ MeV. The fine-tuning represented by the small binding momentum $\kappa \equiv \sqrt{m_N B_d} \approx 45$ MeV can be incorporated by assigning to $NN$ LECs an enhanced scaling with respect to NDA [23]. The deuteron arises as a bound state when the leading $NN$ interaction is iterated to all orders. Pion exchange between nucleons can be treated as subleading interactions in a perturbative expansion in $Q / \Lambda_{NN}$ [29], where $\Lambda_{NN} \equiv 4\pi F_\pi^2 / g_A m_N \sim F_\pi$ and $Q \sim m_\pi \sim \kappa$. This scheme has been applied successfully to the electromagnetic form factors of the deuteron [30, 33]. Since the nucleon-antinucleon $(NN)$ isospin-triplet $^3S_1$ scattering length $a_{\bar{n}n}$ has a natural value, similar enhancements of $NN$ interactions are not necessary.

The $B$-conserving Lagrangian for $NN$ and $N\bar{N}$ scattering we write as

$$
L^{(4)}_{\Delta B=0} = -\left( C_0 + D_2 m_N^2 \right) (N^T P_i N)^\dagger (N^T P_i N) + \frac{C_2}{8} \left[ (N^T P_i N)^\dagger (N^T P_i (\nabla - \nabla)^2 N) + \text{H.c.} \right] - H_0 (N^{cT} \tau^a Y_i^a N)^\dagger (N^{cT} \tau^2 Y_i^a N) + \ldots, 
$$

(7)

where $P_i \equiv \sigma_2 \sigma_i \tau^2 / \sqrt{8}$ ($Y_i^a \equiv \sigma_2 \sigma_i \tau^a / 2$) projects an $NN$ ($N\bar{N}$) pair onto the isospin-singlet (triplet) $^3S_1$ state. The term with $C_0$ is the leading $NN$ interaction, the real part of which scales as $\text{Re} C_0 = O(4\pi/m_N \kappa)$. One-pion exchange and one insertion of the subleading
LECs \( ReC_2 \sim ReD_2 = \mathcal{O}(4\pi/m_N\kappa^2\Lambda_{NN}) \) appear at relative \( \mathcal{O}(\kappa/\Lambda_{NN}) \). Neglecting small imaginary parts discussed below, these LECs are determined from \( NN \) observables \([29,34]\), e.g.

\[
C_0 = \frac{4\pi}{m_N(\kappa - \mu)} + \ldots,
\]

\[
C_2 = \frac{4\pi}{m_N(\kappa - \mu)^2\Lambda_{NN}} \left[ \frac{r_{np}\Lambda_{NN}}{2} - 1 + \frac{8\xi}{3} - 2\xi^2 \right] + \ldots,
\]

where \( \mu \) is the renormalization scale, \( r_{np} \approx 1.75 \) fm \([35]\) is the \( ^3S_1 \) \( np \) effective range, and \( \xi \equiv \kappa/m_\pi \approx 0.32 \). The LEC \( H_0 \) is the leading interaction in the \( ^3S_1 \bar{n}p \) channel, which is complex due to annihilation \([36,37]\). Calculating the \( ^3S_1 \bar{n}p \) scattering amplitude and matching to the effective-range expansion, we obtain up to NLO

\[
H_0 = \frac{4\pi a_{np}}{m_N} \left[ 1 + \mu a_{\bar{np}} + \frac{2(\mu - m_\pi)}{3\Lambda_{NN}} \right] + \frac{4\pi}{m_N\Lambda_{NN}^2} \left( \frac{4\mu}{9} - \frac{3m_\pi}{2} \right) + \ldots
\]

We use the value \( a_{\bar{np}} = (0.44 - i 0.96) \) fm obtained with a chiral potential \([38,39]\) fitted to state-of-the-art \( NN \) partial-wave amplitudes \([40,41]\). The natural size of \( a_{np} \) justifies the assignment \( H_0 = \mathcal{O}(4\pi/m_N\Lambda_{NN}) \) and the use of perturbation theory.

The presence of \( |\Delta B| = 2 \) interactions has two consequences. First, the \( B \)-conserving LECs get imaginary parts because two nucleons can now annihilate. The imaginary part is strongly suppressed, since it requires two \( |\Delta B| = 2 \) insertions; for example, \( \text{Im} C_0 = \mathcal{O}(\delta m^2 \Lambda_\chi^2/\kappa^2 \Lambda_{NN}^4) \), where we account for a \( (\Lambda_{NN}/\kappa)^2 \) enhancement due to renormalization by LO \( NN \) scattering on both sides of the vertex \([28,30]\). Second, we need to consider the \( NN \leftrightarrow N\bar{N} \) interactions

\[
\mathcal{L}_{|\Delta B|=2}^{(4)} = i\bar{B}_0 \left[ \left( N^T P_i N \right) \dagger \left( N^T \pi^2 Y_i - N \right) - \text{H.c.} \right] + \ldots,
\]

where \( \bar{B}_0 \) is a complex LEC. The \( NN \) interaction enhances it over NDA, \( \bar{B}_0 = \mathcal{O}(4\pi \delta m/\kappa \Lambda_{NN}^2) \), and implies that \( \text{Im} \bar{B}_0/C_0 \propto (\kappa - \mu) \text{Im} \bar{B}_0 \) is \( \mu \)-independent, which becomes important below. Some interactions are further enhanced thanks to the one-body character of \( \delta m \). For example, we find \( \text{Re} \bar{B}_0 = \mathcal{O}(4\pi \delta m/\kappa \Lambda_{NN}) \): requiring the \( NN \) and \( N\bar{N} \) scattering amplitudes to be independent of the renormalization scale leads to a renormalization-group equation whose solution is

\[
\text{Re} \bar{B}_0 = -m_N \delta m \text{Re} C_2/\sqrt{2} + \ldots
\]

The mesonic final states in Fig. 1 contain hard pions with energies outside of the regime of \( \chi \)EFT. Instead of directly calculating deuteron-decay diagrams we determine the imaginary part of the pole of the deuteron propagator. The hard pions then only appear as intermediate states that can be integrated out. Following Ref. \([30]\), we write the propagator for a deuteron with four-momentum \( p^\mu = (2m_N + \vec{p}^2/4m_N + \vec{E} + \ldots, \vec{p}) \) in terms of the irreducible two-point function \( \Sigma(\vec{E}) \), which contains all diagrams that do not fall apart when cutting any \( \text{Re} C_0 \) vertex, and expand around \( \bar{B}_d \), i.e.

\[
G(\vec{E}) = \frac{\Sigma(\vec{E})}{1 + i \text{Re}(C_0) \Sigma(\vec{E})} = \frac{iZ_d}{\vec{E} + B_d + i\Gamma_d/2} + \ldots,
\]
where the wavefunction renormalization $Z_d$ is real and

$$\Gamma_d = \frac{2 \text{Im}(i\Sigma(\bar{E}))}{\text{Re}(d\Sigma(\bar{E})/dE)} \bigg|_{E=-B_d} + \ldots$$

is the deuteron decay rate. Up to NLO \[30\],

$$\text{Re} \left( \frac{d\Sigma(\bar{E})}{d\bar{E}} \right) \bigg|_{E=-B_d} = \frac{m_N^2}{8\pi\kappa} \left\{ 1 + \frac{m_N}{2\pi} (\kappa - \mu) \left[ C_2\kappa(\mu - 2\kappa) + D_2m^2_\pi \right] 
+ \frac{2}{\Lambda_{NN}} \left( \kappa - \mu + \frac{m_\pi}{1 + 2\xi} \right) \right\}. \quad (14)$$

Figure 2 shows diagrams that give nonvanishing contributions to $\text{Im}(i\Sigma(\bar{E}))$ up to NLO. We power-count diagrams with the following rules \[28, 30\]: $Q^5/(4\pi m_N)$ for each loop integral, $m_N/Q^2$ for each nucleon propagator, $1/Q^2$ for each pion propagator, and the product of the LECs appearing in each diagram. The dominant contribution to deuteron decay is due to diagram 2(a), which is $O(\delta m^2m^2_N/\kappa^2)$. A diagram with two $\delta m$ insertions but no NN vertex is real and does not contribute to $\text{Im}(i\Sigma)$. Diagrams 2(b,c,d,e,f) are $O(\kappa/\Lambda_{NN})$ relative to 2(a) and give NLO corrections. Diagrams 2(b,c,d) are similar to 2(a) but involve an additional insertion of a subleading NN vertex ($C_2$ or $D_2$), a $B$-conserving pion exchange, or the leading NN vertex $H_0$. Diagrams 2(a,b,c,d) come from the left diagram in Fig. 1; they depend solely on $\delta m$ and are directly related to the free $n-\bar{n}$ transition. Diagrams 2(e,f) are a mixture of both diagrams in Fig. 1. They are proportional to, respectively, $\text{Im} \tilde{B}_0$ and $\text{Re} \tilde{B}_0$, the latter being related to $\delta m$ by Eq. (11). From these diagrams we obtain

$$\text{Im}(i\Sigma(-B_d)) = -\left( \frac{m_N^2\delta m}{8\pi\kappa} \right)^2 \text{Im} H_0 \left\{ 1 + \frac{m_N}{2\pi} (\kappa - \mu) \left[ C_2\kappa(\mu - 2\kappa) + D_2m^2_\pi \right] 
+ \frac{8}{3\Lambda_{NN}} \left( \kappa - \mu + \frac{m_\pi}{1 + 2\xi} \right) \right. 
+ \frac{m_N}{2\pi} (\kappa - \mu) \text{Re} H_0 
\left. - \frac{2\sqrt{2}\kappa(\kappa - \mu)}{m_N\delta m} \left[ \frac{\text{Im} \tilde{B}_0}{\text{Im} H_0} + \frac{m_N}{4\pi}(\kappa - \mu) \text{Re} \tilde{B}_0 \right] \right\}. \quad (15)$$
The deuteron decay rate up to NLO is then

\[ \Gamma_d = -\frac{m_N}{\kappa^2 \tau_{nn}} \text{Im} a_{\bar{n}p} \left[ 1 + \kappa \left( \tau_{np} + 2 \text{Re} a_{\bar{n}p} \right) - \frac{g_\pi^2 m_N \left( 2 - 2\xi - 5\xi^2 + 6\xi^3 \right)}{3\pi F_\pi^2} \right. \]

\[ \left. \frac{1 + 2\xi}{\kappa^2 \Lambda_{NN}} - \frac{(\kappa - \mu) \text{Im} \tilde{B}_0}{\sqrt{2} \delta m \text{Im} a_{\bar{n}p}} \right]. \]  

(16)

This result is independent of the renormalization scale, as it should be. It relates \( \Gamma_d \) to \( \tau_{nn} \) through known quantities and one unknown NLO constant \( (\kappa - \mu) \text{Im} \tilde{B}_0 \), encoded in the ratio \( R_d \equiv \Gamma_d^{-1} / \tau_{nn}^2 \). Numerically we find

\[ R_d = -\left[ \frac{m_N}{\kappa} \text{Im} a_{\bar{n}p} \left( 1 + 0.40 + 0.20 - 0.13 \pm 0.4 \right) \right]^{-1} \]

\[ = (1.1 \pm 0.3) \cdot 10^{22} \text{ s}^{-1}. \]  

(17)

The NLO corrections from known LECs affect the result by roughly 50% and are dominated by the effective-range correction. We account for the unknown value of \( \text{Im} \tilde{B}_0 \) as an uncertainty of the same size as the effective-range correction. The explicit pion contributions amount to only 13%. The limit \( m_\pi \to \infty \) recovers the result of Pionless EFT \[36\], where pions are integrated out and \( \text{Im} \tilde{B}_0 \) absorbs the surviving pion term. Use of an auxiliary dibaryon field \[36\] automatically accounts for the enhancement of Eq. (11).

Our central value for \( R_d \) is smaller by a factor \( \simeq 2.5 \) than the often-used result from Ref. \[42\] based on nuclear models for the nucleon-(anti)nucleon interactions. We have checked that when the expressions from Refs. \[42, 43\] are applied to a zero-range potential we recover our LO term in Eq. (17). The difference therefore stems from the smaller \( \text{Im} a_{\bar{n}p} \) \[44\] of the \( N\bar{N} \) potentials of Ref. \[42\], and from corrections to the zero-range limit. The diagrammatic approaches of Refs. \[43, 46\] also reduce to our LO for a zero-range potential. (Reference \[47\] disagrees from these results by a factor of 2.)

Our result is based on a systematic and improvable framework for all interactions, and we showed that NLO corrections are significant but of the expected size. In addition, we have used an up-to-date \( \bar{n}p \) scattering length \[39\]. We therefore propose to use Eq. (17) in comparisons of deuteron stability and \( n\bar{n} \) oscillation beam experiments. Taking the largest value of \( R_d \) allowed by Eq. (17), the SNO limit on \( \Gamma_d^{-1} \) \[15\] gives

\[ \tau_{n\bar{n}} = 1/\sqrt{R_d \Gamma_d} > 5.1 \text{ yr} = 1.6 \cdot 10^8 \text{ s}, \]  

(18)

about a factor two stronger than the direct ILL limit.

At higher orders we find some of the nuclear effects discussed in the literature, two examples being shown in Fig. 3. Diagram 3(a) can be seen as an “in-medium” modification of the \( n\bar{n} \) oscillation required by chiral symmetry and contained in the dots of Eq. (5). It is nominally of relative \( \mathcal{O}(\kappa^2/\Lambda_{NN} m_N) \), but it actually vanishes. Corrections of this type should, therefore, be no larger than about 5%, to be compared with the 25-30% estimated for heavier nuclei in Ref. \[48\]. The effect of direct two-nucleon annihilation \[18\], the right diagram in Fig. 1, is proportional to the absorptive part of \( NN \) interactions, \( \text{Im} C_0 \), and appears at next-to-next-to-leading order (N^2LO), \( \mathcal{O}(\kappa^2/\Lambda_{NN}^2) \).

So far we have not discussed the operator \( Q_4 \). Since it belongs to the \( (1_L, 7_R) \) irrep it can only contribute to \( \tau_{n\bar{n}}^{-1} \) if additional sources of isospin violation are included. The lowest-order contribution to \( \tau_{n\bar{n}}^{-1} \) involves two insertions of the charge, leading to a suppression of
\( \alpha_{\text{em}}/4\pi \sim \mathcal{O}(m_N^3/\Lambda_{\text{NN}}^3) \) with respect to \( \delta m \), where \( \alpha_{\text{em}} \) is the fine-structure constant. \( \text{Im} C_0 \) induced by \( C_4 \) does not require the inclusion of extra isospin violation. For the case in which deuteron decay is dominated by \( Q_4 \), i.e. \( c_4 \gg c_{1,2,3} \), the imaginary part of diagram \( 3(b) \) is \( \mathcal{O}(m_N^3/\Lambda_{\text{NN}}^3 Q^4) \) relative to diagram \( 2(a) \) and is expected to dominate the deuteron decay rate. At LO,

\[
\Gamma_d|_{Q_4} = -\kappa(\kappa - \mu)^2 \text{Im} C_0/\pi = -4\kappa^3 \text{Im} a_{np}/m_N
\]  

in terms of the imaginary part of the \( ^3S_1 \) np scattering length \( a_{np} \). In this case \( \Gamma_d^{-1} \) and \( \tau_{n\bar{n}} \) are not dominated by the same \( |\Delta B| = 2 \) LECs, resulting in a larger value of \( R_d \). This implies that if deuteron decay and free \( n-\bar{n} \) oscillation are both observed, one could infer whether \( |\Delta B| = 2 \) violation is dominated by \( Q_{1,2,3} \) or by \( Q_4 \) (or strangeness-changing operators we have not considered \[18, 49\]).

In closing, we briefly comment on what our findings imply for heavier nuclei. Due to the low deuteron binding momentum, the expansion in \( Q/\Lambda_{\text{NN}} \sim \kappa/\Lambda_{\text{NN}} \) allows for analytical results. In denser nuclei, such as \(^{16}\text{O} \), this expansion is likely not valid and we need to resum \( Q/\Lambda_{\text{NN}} \) corrections by treating pion exchange nonperturbatively. While this complicates the calculations, it only partially affects the power-counting estimates. Contributions from \( Q \sim \Lambda_{\text{NN}} \) are transferred from LECs to explicit pion exchange. The infrared enhancement by \( \kappa^{-1} \) in the decay rate Eq. (17), which increases the overall sensitivity to \( \tau_{n\bar{n}} \), should become less pronounced, but intrinsic two-nucleon effects due to \( |\Delta B| = 2 \) pion exchange and short-range \( NN \) annihilation, and \( NN \to N\bar{N} \) interactions with unknown LECs appear in the chiral expansion only at N\(^2\)LO, \( \mathcal{O}(Q^2/\Lambda_{\text{NN}}^2) \), or beyond. Therefore, in conjunction with free \( n-\bar{n} \) transitions, stability experiments with denser nuclei also partially discriminate among \( |\Delta B| = 2 \) operators.

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