A new leading contribution to neutrinoless double-beta decay

Vincenzo Cirigliano,1 Wouter Dekens,1 Jordy de Vries,2 Michael L. Graesser,1
Emanuele Mereghetti,1 Saori Pastore,1 and Ubirajara van Kolck3,4

1Theoretical Division, Los Alamos National Laboratory, Los Alamos, NM 87545, USA
2Nikhef, Theory Group, Science Park 105, 1098 XG, Amsterdam, The Netherlands
3Institut de Physique Nucléaire, CNRS/IN2P3, Université Paris-Sud, Université Paris-Saclay, 91406 Orsay, France
4Department of Physics, University of Arizona, Tucson, Arizona 85721, USA

Within the framework of chiral effective field theory we discuss the leading contributions to the neutrinoless double-beta decay transition operator induced by light Majorana neutrinos. Based on renormalization arguments in both dimensional regularization with minimal subtraction and a coordinate-space cutoff scheme, we show the need to introduce a leading-order short-range operator, missing in all current calculations. We discuss strategies to determine the finite part of the short-range coupling by matching to lattice QCD or by relating it via chiral symmetry to isospin-breaking observables in the two-nucleon sector. Finally, we speculate on the impact of this new contribution on nuclear matrix elements of relevance to experiment.

Introduction: Neutrinoless double-beta decay (0νββ) is the most sensitive laboratory probe of lepton number violation (LNV). In 0νββ L is violated by two units when two neutrons in a nucleus turn into two protons, with the emission of two electrons and no neutrinos. The observation of 0νββ would demonstrate that neutrinos are Majorana fermions [1], shed light on the mechanism of neutrino mass generation [2–4], and give insight into leptogenesis scenarios for the generation of the matter-antimatter asymmetry in the universe [5].

0νββ is actively being searched for in a number of even-even nuclei for which single-β decay is energetically forbidden. Current experimental limits [6–15] on the half-lives are at the level of $T_{1/2} > 5.3 \times 10^{25}$ y for $^{76}\text{Ge}$ [12] and $T_{1/2} > 1.07 \times 10^{26}$ y for $^{136}\text{Xe}$ [10], with next-generation ton-scale experiments aiming at improvements in sensitivity by two orders of magnitude.

0νββ can be generated by a variety of dynamical LNV mechanisms, which, in an effective field theory (EFT) approach to new physics are parametrized by $\Delta L = 2$ operators of odd dimension greater than four [16–22]. If the mass scale associated with LNV is much higher than the electroweak scale, the only low-energy manifestation of this new physics is a Majorana mass for light neutrinos, encoded in a single gauge-invariant dimension-five operator [16], which induces 0νββ through light Majorana-neutrino exchange [23, 24]. To interpret positive or null 0νββ results in this minimal scenario it is crucial to have good control over the relevant hadronic and nuclear matrix elements. Current knowledge of these is not satisfactory [25], as various many-body approaches lead to estimates that differ by a factor of two to three and most calculations are not based on a modern EFT analysis. In Ref. [26] a first step was presented towards the analysis of 0νββ induced by a light Majorana neutrino in the chiral EFT framework [27–29], which provides a systematic expansion of hadronic amplitudes in $p/\Lambda$, where $p \sim m_\pi \sim k_F \sim \mathcal{O}(100 \text{ MeV})$ and $\Lambda \sim 4\pi F_\pi \sim m_N \sim \mathcal{O}(1 \text{ GeV})$. The 0νββ transition operators were derived up to next-to-next-to-leading order ($\mathcal{N}^4\text{LO}$) in Weinberg’s power-counting scheme [30, 31].

In this letter we demonstrate that Weinberg’s scheme for 0νββ assumed in Ref. [26] breaks down and any consistent power counting requires a leading-order (LO) short-range $\Delta L = 2$ operator, whose effect is missing in all current calculations. Our argument is based on renormalization. Using two different schemes (dimensional regularization with minimal subtraction and a coordinate-space cutoff) we show that once the strong nucleon-nucleon scattering amplitude is made finite and regulator-independent, an additional $\Delta L = 2$ contact operator with coupling $g_{\nu\nu}^{\text{NN}}$ has to be introduced to make the $nn 	o ppee$ amplitude finite and regulator-independent. The finite part of $g_{\nu\nu}^{\text{NN}}$, which encodes hard-neutrino exchange, can be determined by (i) matching the chiral EFT $nn \to ppee$ amplitude to future lattice QCD calculations; (ii) relating it via chiral symmetry to electromagnetic low-energy constants (LECs) that control isospin-breaking in the two-nucleon sector. A combination of couplings involving $g_{\nu\nu}^{\text{NN}}$ can be fit to nucleon-nucleon charge-independence-breaking (CIB) observables, confirming the LO scaling of this coupling. Based on this, we speculate on the impact of $g_{\nu\nu}^{\text{NN}}$ on nuclear matrix elements of relevance to experiments.

The need for an LO short-range $\Delta L = 2$ interaction: We consider a scenario in which LNV at low energy is dominated by the electron-neutrino Majorana mass

$$L_{\Delta L=2} = -\frac{m_{\nu\nu}}{2} \bar{v}_L^T C v_L,$$

where $C = i\gamma_2\gamma_0$ denotes the charge conjugation matrix. The nuclear effective Hamiltonian can be written as

$$H_{\text{eff}} = H_{\text{strong}} + 2G_F^2 V_{ud}^2 m_{\beta\beta} \bar{e}_L C e_L^T v_\nu,$$
in terms of the Fermi constant $G_F$ and the $V_{ud}$ element of the CKM matrix [32, 33]. The neutrino potential $V_\nu$ can be obtained from two-nucleon irreducible diagrams mediating $nn \rightarrow ppee$ to a given order in $p/\Lambda$. Within Weinberg’s power counting the only LO contribution [26] comes from the exchange of potential neutrinos, with $q^0 \ll |q|$, 

$$V_{\nu,0}(q) = \tau^{(1) + \tau^{(2)} + \frac{1}{q^2}} \left( 1 - g_A^2 \sigma^{(1)} \cdot \sigma^{(2)} \right) + g_A^2 \sigma^{(1)} \cdot q \sigma^{(2)} \cdot \frac{2m_n^2 + q^2}{(q^2 + m_n^2)^2} , \quad (3)$$

where $g_A \simeq 1.27$ is the nucleon axial coupling, $m_\pi$ the pion mass, and $q$ the momentum transfer. $\zeta^2$LO terms arise from corrections to the single nucleon weak currents, irreducible one-loop diagrams, and contact interactions mediating $\pi \pi \rightarrow ee$, $n \rightarrow \nu e^+ e^-$, and $nn \rightarrow ppee$. In particular, the short-range potential includes a two-nucleon term [26]

$$V_{\nu,CT} = -2g_{\nu NN}^{\pi} \tau^{(1) + \tau^{(2)} +} , \quad (4)$$

where the LEC $g_{\nu NN}^{\pi}$ is $O((4\pi F_\pi)^{-2})$ in Weinberg’s counting and $F_\pi = 92.2$ MeV is the pion decay constant. However, it is known that Weinberg’s power counting leads to inconsistent results in nucleon-nucleon scattering [34–37] and nuclear processes mediated by external vectors, with $q \ll |q|$. In particular, the short-range potential includes a two-nucleon term [26]

$$V_{\nu,0}(q) = \tilde{C} + V_\pi(q), \quad V_\pi(q) = -\frac{g_A^2 m_\pi^2}{4 F_\pi^2 q^2 + m_\pi^2} , \quad (5)$$

with $\tilde{C} \sim O(F_\pi^{-2}, m_\pi^2 F_\pi^{-4})$ [31, 34, 35]. We have checked that transitions involving higher partial waves such as $3P_{0,1} \rightarrow 3P_{0,1}$ are correctly renormalized and do not require enhanced $\Delta L = 2$ counterterms.

The contributions to $A_{\Delta L = 0}$ from the exchange of a light neutrino ($A_{\Delta L = 0}^{(\nu)}$) are shown in Fig. 1. The blue ellipse denotes the iteration of the Yukawa potential $V_\pi(q)$. The diagrams in the second and third rows involve an infinite number of bubbles, dressed with iterations of $V_\pi$. Without loss of generality for our arguments, we use the kinematics $n(p) n(-p) \rightarrow p(p') p(-p') e(p_{e1} = 0, p_{e2} = 0)$, with $|p| = 1$ MeV and correspondingly $|p'| = 38$ MeV.

$A_{\Delta L = 0}^{(\nu)}$ can be expressed in terms of the Yukawa “in” and “out” wavefunctions $\chi_\nu^\pm(r)$ and the propagators

\[
G_E^\nu(r, r') = \langle r' | (E - T - V_\pi \pm i0^+)^{-1} | r \rangle \] [34, 37]. Observing that the bubble diagrams in Fig. 1 are related to $G_E^\nu(0, 0)$, while the triangles dressed by Yukawas are related to $\chi_\nu^+(0)$ and $\chi_\nu^-(0) = \chi_\nu^+(0)$ [34], the LO amplitude reads

$$A_{\Delta L = 2}^{(\nu)} = A_A + K_E A_B + \tilde{A}_B K_E + K_E A_C K_E , \quad \tilde{C} = \frac{\chi_\nu^+(0) \tilde{C}}{1 - GG_E^\nu(0, 0) , \quad (6)}$$

where $A_A, A_B,$ and $A_C$ denote the first diagram in the first, second, and third rows of Fig. 1, respectively (without the wavefunctions at 0, in the case of $A_B$ and $A_C$). $A_B$ is similar to $A_B$ and not shown in Fig. 1.

To study the renormalization of the $\Delta L = 2$ amplitude, we now discuss the divergence structure of $A_{\Delta L = 0}^{(\nu)}$. $\chi_\nu^+(0)$ is finite and the divergence in $G_E^\nu(0, 0)$ is absorbed by $\tilde{C}^{-1}$, so that $K_E$ is finite and scheme-independent [34]. We note that:

(i) All diagrams in $A_A$ are finite. The tree level is finite and each $V_\pi$ insertion improves the convergence by bringing in a factor of $d^2k/(k^2)^2$, where one $k^2$ comes from the pion propagator and the other from the two-nucleon propagator.

(ii) All the diagrams in $A_B$ and $A_B$ are finite. The first loop goes as $d^2k/(k^2)^2$, while $V_\pi$ insertions further improve the convergence.

(iii) The first two-loop diagram in $A_C$ has a logarithmic divergence, which stems from an insertion of the most
singular component of the neutrino potential, namely

\[ \tilde{V}_\nu(q) = \varphi^{(1)+\varphi^{(2)}+\frac{1}{2q^2} \left( 1 - \frac{2}{3}g^2_{\Delta} \sigma^{(1)} \cdot \sigma^{(2)} \right). \]  

The two-loop diagram with insertion of \( V_{\nu,0} - \tilde{V}_\nu \) and higher-loop diagrams are convergent.

We focus on \( A_C \) and write \( A_C = A_C^{(\text{div})} + \delta A_C \). In dimensional regularization,

\[ A_C^{(\text{div})} = - \left( \frac{m_N}{4\pi} \right)^2 (1 + 2g^2_{\overline{\Delta}}) \left[ \Delta + L_{p,p'}(\mu) \right], \]

\[ L_{p,p'}(\mu) = \frac{1}{2} \left( \log \frac{-m^2 + p^2 + \mu^2}{2} + i0^+ + 1 \right), \]

where \( \Delta \equiv (1/(4-d) - \gamma + \log 4\pi)/2 \). The divergence for \( d \to 4 \) can be removed by introducing \( g_{\nu,NN}^{(\text{div})} \) at LO. The counterterm amplitude, shown in the fourth line of Fig. 1, reads

\[ A_{\Delta L=2}^{(\text{NN})} = K_{E} 2g_{\nu,NN}^{(\text{NN})} C_{\nu}, \]  

and we can renormalize \( A_{\Delta L=2} \) by replacing \( A_C \to A_C + 2g_{\nu}^{(\text{NN})}/C^2 \) in Eq. (6). In the \( \overline{\text{MS}} \) scheme,

\[ A_C \to \delta A_C + \left( \frac{m_N}{4\pi} \right)^2 \left[ 2g_{\Delta}^{(\nu)}(\mu) - (1 + 2g^2_{\overline{\Delta}}) L_{p,p'}(\mu) \right] \]

after defining the dimensionless coupling

\[ \tilde{g}_{\nu}^{(\text{NN})} = \left( \frac{4\pi}{m_N C} \right)^2 g_{\nu,NN}^{(\text{NN})}. \]  

This coupling obeys the renormalization-group equation (RGE)

\[ \frac{d \tilde{g}_{\nu}^{(\text{NN})}}{d \mu} = \frac{1}{2} \left( 1 + 2g^2_{\overline{\Delta}} \right), \]  

confirming that \( \tilde{g}_{\nu}^{(\text{NN})} \sim O(1) \). Since \( \tilde{C}(\mu = m_\tau) \approx -0.9/F^2_\pi \), we find that \( g_{\nu}^{(\nu,NN)} \sim O(F^{-2}_\pi) \) instead of \( O((4\pi F_\pi)^{-2}) \). A similar enhancement also occurs in

\[ \text{four-nucleon couplings induced by higher-dimensional LNV operators. Treating } V_{\nu} \text{ as a subleading correction the } \overline{\text{MS}} \text{ scheme, is equivalent to working to LO in pionless EFT, and does not affect our conclusions about the importance of } g_{\nu}^{(\nu,NN)} \text{ [26]. Details on how to obtain } \delta A_C \text{ will be provided in future work [40].} \]

\[ A_{\Delta L=2} \text{ in a cutoff scheme: The need for an LO counterterm can be demonstrated also in a coordinate-space scheme that makes no direct reference to Feynman diagrams. In this approach we regulate the short-range part of } V_0 \text{ with a smeared } \delta \text{-function,} \]

\[ \tilde{C} \delta^{(3)}(r) \to \left( \frac{\tilde{C}(R_S)}{\sqrt{\pi R_S}} \right)^3 \exp \left( -\frac{r^2}{2R_S^2} \right) \equiv \tilde{C}(R_S) \delta^{(3)}(r), \]  

and obtain \( \psi^-_p(r) \) and \( \psi^+_p(r) \) by solving the Schrödinger equation. We determine \( \tilde{C}(R_S) \) by requiring that the \( S_0 \) scattering length be reproduced as a function of \( R_S \). The solid line corresponds to a fit that includes \( O(R_S, \log R_S) \) power corrections. The dash-dotted line shows \( A_{\Delta L=2}^{(\nu)} \) as a function of \( 1/\mu \). The horizontal bands represent the total amplitude \( A_{\Delta L=2} \) with \( g_{\nu}^{(\nu)} = (C_1 + C_2)/2 \), as discussed in the main text.

\[ \frac{d^3 r}{4\pi} \psi_{-p}(r)^* V_{\nu,0}(r) \psi_{p}(r), \]

and \( V_{\nu,0}(r) \) is obtained by Fourier-transforming the \( S_0 \) projection of Eq. (3). In Fig. 2 we plot \( A_{\Delta L=2}^{(\nu)} \) as a function of \( R_S \). The plot displays a logarithmic dependence on \( R_S \) (analogous to the \( \log \mu \) dependence in Eq. (10)) as well as milder power-like behavior. Therefore, to obtain a physical, regulator-independent amplitude one needs to include an LO counterterm, given in \( r \)-space by \( V_{\nu,CT}(r) = -2g_{\nu}^{(\nu,NN)}(R_S)\delta^{(3)}(r). \) The corresponding amplitude,

\[ A_{\Delta L=2}^{(\text{NN})} = - \int d^3r \psi_{-p}(r)^* V_{\nu,CT}(r) \psi_{p}(r), \]

is also regulator-dependent. As expected from Eq. (9), we find its leading divergent behavior to be well reproduced by \( 1/C(R_S)^2 \).
performed at the same kinematic point, as it is done in
the strong-interacting sector [41]. First lattice results relat-
ed to double-beta decay are starting to appear [42, 43].

We now discuss a complementary estimate based on
the fact that the short-range operators and associated
LECs arising in 0νββ and electromagnetic processes are
closely related [26]. In the electromagnetic case, the
short-range hadronic operators arise from amplitudes in
the underlying theory involving two insertions of the elec-
 tromagnetic current with exchange of hard virtual pho-
tons [44, 45]. In the ΔL = 2 case, up to a proportionality
factor, the same operators are generated by the inser-
tion of two weak currents with exchange of hard neu-
trinos. This comes about because the neutrino propa-
gator and weak vertices combine to give a massless
gauge-boson propagator in Feynman gauge, multiplied
by $SO_1^2V_2^2m_β\bar{e}_L\bar{e}_L$. The LECs needed for 0νββ
are therefore related to the LECs associated with the isospin
I = 2 component of the product of two electromagnetic
currents, which belongs to the 5x1 irreducible repre-
sentation of chiral SU(2)$_L \times$ SU(2)$_R$.

Only two independent four-nucleon operators that
transform as I = 2 objects exist:

$$O_1 = \bar{N}Q_LN\bar{N}Q_RN - \frac{Tr[Q_L]}{6}\bar{N}\tau N + \{L \leftrightarrow R\},$$

$$O_2 = 2\left(\bar{N}Q_LN\bar{N}Q_RN - \frac{Tr[Q_LQ_R]}{6}\bar{N}\tau N + \bar{N}\tau N\right)\ (16)$$

where $Q_L = u^\dagger Q_Lu$, $Q_R = uQ_Ru^\dagger$, $u = \exp(i\tau \cdot
\pi/(2F_\pi))$, and $Q_L$, $Q_R$ are “spurious” transforming un-
der the chiral group as $Q_L \rightarrow LQ_LL^\dagger$, $Q_R \rightarrow RQ_RR^\dagger$. In
the electromagnetic case $Q_L = Q_R = \tau^2/2$, while in 0νββ
$Q_L = \tau^+_2$, $Q_R = 0$. In our conventions $O_1$ enters the
ΔL = 2 Lagrangian with coefficient $2G_F^2V_2^2m_β\bar{g}_{\nu\nu}^N$. The
LECs needed for 0νββ are starting to appear [42, 43].

We have performed a similar analysis for $A = 6, 12$
uclei, using Variational Monte Carlo nuclear wavefunc-
tions [51] based on the AV18 two-nucleon [50] and IL7
three-nucleon [52] interactions. The mismatch between
the short-range behaviors of existing strong-interaction
potentials and our 0νββ interaction introduces additional
model dependence, which we mitigate by: (i) Considering
an alternative extraction of $(C_1 + C_2)/2$ from the phase-
shift analysis of Refs. [47, 48] 2, which employs the same
regulator (13) with $R_S \simeq 0.6 - 0.8$ fm, approximately the
range of AV18’s short-range part. (ii) Simply replacing our
$V_{\nu\beta\beta}(r)$ with AV18’s short-range CIB potential.

For ΔL = 0 transitions such as the $^6$He → $^6$Be shown in
Fig. 3 (middle panel), we find $A_{\Delta L=2}^{(0\nu\beta\beta)}/A_{\Delta L=2}^{(\nu\beta\beta)} \sim 10\%$, similarly
to the $nn \rightarrow p\bar{p}ee$ case. In realistic 0νββ transi-
ations, however, the total nuclear isospin changes by two
units, ΔL = 2. This implies the presence of a node in $\rho_{\nu\nu}(r)$
due to the orthogonality of the initial and final spatial
wavefunctions. The resulting partial cancellation between
the regions with $r < 2$ fm and $r > 2$ fm [51] leads to
a relative enhancement of the short-range contribu-
tion, as illustrated in Fig. 3 (bottom panel) for $^{12}$Be →
$^{12}$C. Numerically we find $A_{\Delta L=2}^{(0\nu\beta\beta)}/A_{\Delta L=2}^{(\nu\beta\beta)} \sim 25\%$ (our fit),

1 Our result is consistent with analyses based on chiral [46–49] and
phenomenological potentials such as AV18 [50], which also find
that, except at very low energies, long- and short-range compo-
nents of the CIB interaction induce effects of similar size.

2 $C_1 + C_2$ is related to the CIB coefficient $C_0^T$ of Refs. [47, 48]
by $(C_1 + C_2)/2 = -6C_0^T/e^2$. 

\[ \rho_{\nu\nu}(r) = \int dr \rho_{\nu\nu}(r), \quad A_{\Delta L=2}^{(0\nu\beta\beta)} = \int dr \rho_{\nu\nu}(r). \] (18)
FIG. 3. \(\rho_\nu(r)\) and \(\rho_{NN}(r)\) for the \(nn \to ppec\) process (top), and for nuclear transitions with \(A = 6\) (middle) and \(A = 12\) (bottom). In the middle and bottom panels the green \((\rho_{NN})\) and red \((\rho_{NN}^{(x)})\) bands correspond to \(g_{NN}^{\nu} = (C_1 + C_2)/2\) extracted from our analysis and from Refs. [47, 48], respectively.

\(\sim 55\%\) (fit from Refs. [47, 48]), and \(\sim 60\%\) (AV18 representation of the short-range CIB potential). Because the node in the density is a robust feature of \(\Delta I = 2\) transition [53, 54], we expect the effects in \(^{12}\text{Be} \to ^{12}\text{C}\) and experimentally relevant transitions to be of comparable size.

**Conclusion:** The above arguments suggest that the new short-range \(\Delta I = 2\) potential identified in this work can significantly impact 0\(\nu\)\(\beta\beta\) phenomenology and its implications for Majorana neutrino masses. We hope this will stimulate work towards a more controlled determination of \(g_{NN}^{\nu}\) from lattice QCD and an assessment of the impact of the short-range potential in nuclei of experimental interest.

**Acknowledgments:** VC, WD, MG, EM, and SP acknowledge support by the U.S. Department of Energy, Office of Science, and by the LDRD program at Los Alamos National Laboratory. VC and EM acknowledge partial support from the DOE topical collaboration on “Nuclear Theory for Double-Beta Decay and Fundamental Symmetries”. JdV acknowledges support by the Dutch Organization for Scientific Research (NWO) through a VENI grant. The work of UVK was supported in part by the U.S. Department of Energy, Office of Science, Office of Nuclear Physics, under award No. DE-FG02-04ER41338, and by the European Union Research and Innovation program Horizon 2020 under grant agreement No. 654002. We acknowledge stimulating discussions with Will Detmold and Martin Savage, which triggered this research. We thank Joe Carlson, Jon Engel, Amy Nicholson, Maria Piarulli, Petr Vogel, Andre Walker-Loud, and Bob Wiringa for discussions at various stages of this work.

[1] J. Schechter and J. W. F. Valle, Phys. Rev. D25, 2951 (1982).
[2] P. Minkowski, Phys. Lett. 67B, 421 (1977).
[3] R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980).
[4] M. Gell-Mann, P. Ramond, and R. Slansky, Supergravity Workshop Stony Brook, New York, September 27-28, 1979, Conf. Proc. C790927, 315 (1979), arXiv:1306.4669 [hep-th].
[5] S. Davidson, E. Nardi, and Y. Nir, Phys. Rept. 466, 105 (2008), arXiv:0802.2962 [hep-ph].
[6] A. Gando et al. (KamLAND-Zen), Phys. Rev. Lett. 110, 062502 (2013), arXiv:1211.3863 [hep-ex].
[7] M. Agostini et al. (GERDA), Phys. Rev. Lett. 111, 122503 (2013), arXiv:1307.4720 [nucl-ex].
[8] J. B. Albert et al. (EXO-200), Nature 510, 229 (2014), arXiv:1402.6956 [nucl-ex].
[9] S. Andringa et al. (SNO+), Adv. High Energy Phys. 2016, 6194250 (2016), arXiv:1508.05759 [physics.ins-det].
[10] A. Gando et al. (KamLAND-Zen), Phys. Rev. Lett. 111, 082503 (2013), [Addendum: Phys. Rev. Lett.117,no.10,109903(2016)], arXiv:1605.02889 [hep-ex].
[11] S. R. Elliott et al. (2016) arXiv:1610.01210 [nucl-ex].
[12] M. Agostini et al., Nature 544, 47 (2017), arXiv:1703.00570 [nucl-ex].
[13] J. B. Albert et al. (EXO), Phys. Rev. Lett. 120, 072701 (2018), arXiv:1707.08707 [hep-ex].
[14] C. Alchino et al. (CUORE), (2017), arXiv:1710.07988 [nucl-ex].
[15] O. Azzolini et al. (CUPID-0), (2018), arXiv:1802.07791 [nucl-ex].
[16] S. Weinberg, Phys. Rev. Lett. 43, 1566 (1979).
[17] K. S. Babu and C. N. Leung, Nucl. Phys. B619, 667 (2001), arXiv:hep-ph/0106054 [hep-ph].
[18] G. Prezeau, M. Ramsey-Musolf, and P. Vogel, Phys. Rev. D68, 034016 (2003), arXiv:hep-ph/0303205 [hep-ph].
[19] A. de Gouvêa and J. Jenkins, Phys. Rev. D77, 013008 (2008), arXiv:0708.1344 [hep-ph].
[20] L. Lehman, Phys. Rev. D90, 125023 (2014).
[21] M. L. Graesser, JHEP 08, 099 (2017), arXiv:1606.04549 [hep-ph].
[22] V. Cirigliano, W. Dekens, J. de Vries, M. L. Graesser, and E. Mereghetti, JHEP 12, 082 (2017), arXiv:1708.09390 [hep-ph].
[23] S. M. Bilenky and C. Giunti, Phys. Rev. D 39, 1311 (1989).
[24] S. M. Bilenky and S. T. Petcov, Phys. Rev. D 37, 1299 (1988).
[25] S. M. Bilenky and S. T. Petcov, Phys. Rev. D 40, 1299 (1989).
[26] V. Cirigliano, W. Dekens, E. Mereghetti, and A. Walker-Loud, (2017), arXiv:1710.01729 [hep-ph].
[27] P. F. Bedaque and U. van Kolck, Ann. Rev. Nucl. Part. Sci. 52, 339 (2002), arXiv:nucl-th/0203055 [nucl-th].
[28] E. Epelbaum, H.-W. Hammer, and U.-G. Meißner, Rev. Mod. Phys. 81, 1773 (2009), arXiv:0811.1338 [nucl-th].
[29] R. Machleidt and D. R. Entem, Phys. Rev. C 74, 054003 (2001), arXiv:nucl-th/0104030 [nucl-th].
[30] M. Kobayashi and T. Maskawa, Prog. Theor. Phys. 49, 652 (1973).
[31] D. B. Kaplan, M. J. Savage, and M. B. Wise, Nucl. Phys. B478, 629 (1996), arXiv:nucl-th/9605002 [nucl-th].
[32] S. R. Beane, P. F. Bedaque, M. J. Savage, and U. van Kolck, Nucl. Phys. A700, 377 (2002), arXiv:nucl-th/0104030 [nucl-th].
[33] A. Nogga, R. G. E. Timmermans, and U. van Kolck, Phys. Rev. C72, 054006 (2005), arXiv:nucl-th/0506005 [nucl-th].
[34] B. Long and C.-J. Yang, Phys. Rev. C 86, 024001 (2012), arXiv:1202.4053 [nucl-th].
[35] M. Pavón Valderrama and D. R. Phillips, Phys. Rev. Lett. 114, 082502 (2015), arXiv:1407.0437 [nucl-th].
[36] D. B. Kaplan, M. J. Savage, and M. B. Wise, Nucl. Phys. B534, 329 (1998), arXiv:nucl-th/9802075 [nucl-th].
[37] V. Cirigliano, W. Dekens, J. de Vries, M. L. Graesser, E. Mereghetti, S. Pastore, and U. van Kolck, “Neutrinoless double-beta decay in chiral effective field theory,” (in preparation, 2018).
[38] N. Barnea, L. Contessi, D. Gazit, F. Pedreira, and U. van Kolck, Phys. Rev. Lett. 114, 052501 (2015), arXiv:1311.4966 [nucl-th].
[39] A. Nicholson, E. Berkowitz, C. C. Chang, M. A. Clark, B. Joo, T. Kurth, E. Rinaldi, B. Tiburzi, P. Vranas, and A. Walker-Loud, in Proceedings, 34th International Symposium on Lattice Field Theory (Lattice 2016): Southampton, UK, July 24-30, 2016 (2016) arXiv:1608.04793 [hep-lat].
[40] P. E. Shanahan, B. C. Tiburzi, M. L. Wagman, F. Winter, E. Chang, Z. Davoudi, W. Detmold, K. Orginos, and M. J. Savage, Phys. Rev. Lett. 119, 062003 (2017), arXiv:1701.03456 [hep-lat].
[41] B. Moussallam, Nucl. Phys. B504, 381 (1997), arXiv:hep-ph/9701400 [hep-ph].
[42] M. Piarulli, L. Girlanda, R. Schiavilla, A. Kievsky, A. Lovato, L. E. Marcucci, S. C. Pieper, M. Viviani, and R. B. Wiringa, Phys. Rev. C94, 054007 (2016), arXiv:1606.06335 [nucl-th].
[43] M. Piarulli, L. Girlanda, R. Schiavilla, A. Kievsky, A. Lovato, L. E. Marcucci, S. C. Pieper, M. Viviani, and R. B. Wiringa, Phys. Rev. C94, 054007 (2016), arXiv:1606.06335 [nucl-th].
[44] P. Reinert, H. Krebs, and E. Epelbaum, (2017), arXiv:1711.08821 [nucl-th].
[45] R. B. Wiringa, V. G. J. Stoks, and R. Schiavilla, Phys. Rev. C51, 38 (1995), arXiv:nucl-th/9408016 [nucl-th].
[46] S. Pastore, J. Carlson, V. Cirigliano, W. Dekens, E. Mereghetti, and R. B. Wiringa, Phys. Rev. C97, 014606 (2018), arXiv:1710.05026 [nucl-th].
[47] S. C. Pieper, AIP Conference Proceedings 1011, 143 (2008).
[48] F. ˇSimkovic, A. Faessler, V. Rodin, P. Vogel, and J. Engel, Phys. Rev. C77, 045503 (2008), arXiv:0710.2055 [nucl-th].
[49] J. Menéndez, A. Poves, E. Caurier, and F. Nowacki, Nucl. Phys. A818, 139 (2009), arXiv:0801.3760 [nucl-th].