

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 Kolck^{3,4}

¹*Theoretical Division, Los Alamos National Laboratory, Los Alamos, NM 87545, USA*

²*Nikhef, Theory Group, Science Park 105, 1098 XG, Amsterdam, The Netherlands*

³*Institut de Physique Nucléaire, CNRS/IN2P3, Université Paris-Sud, Université Paris-Saclay, 91406 Orsay, France*

⁴*Department 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 ($0\nu\beta\beta$) is the most sensitive laboratory probe of lepton number violation (LNV). In $0\nu\beta\beta$ 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 $0\nu\beta\beta$ 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].

$0\nu\beta\beta$ 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}Ge [12] and $T_{1/2} > 1.07 \times 10^{26}$ y for ^{136}Xe [10], with next-generation ton-scale experiments aiming at improvements in sensitivity by two orders of magnitude.

$0\nu\beta\beta$ 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 $0\nu\beta\beta$ through light Majorana-neutrino exchange [23, 24]. To interpret positive or null $0\nu\beta\beta$ 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 $0\nu\beta\beta$ induced by a light Majorana neutrino in the chiral EFT framework [27–29], which provides a systematic expansion of hadronic amplitudes

in p/Λ_χ , where $p \sim m_\pi \sim k_F \sim \mathcal{O}(100 \text{ MeV})$ and $\Lambda_\chi \sim 4\pi F_\pi \sim m_N \sim \mathcal{O}(1 \text{ GeV})$. The $0\nu\beta\beta$ transition operators were derived up to next-to-next-to-leading order (N^2LO) in Weinberg’s power-counting scheme [30, 31].

In this letter we demonstrate that Weinberg’s scheme for $0\nu\beta\beta$ 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 independent of the ultraviolet regulator, an additional $\Delta L = 2$ contact operator with coupling g_ν^{NN} has to be introduced to make the $nn \rightarrow ppee$ amplitude finite and regulator-independent. The finite part of g_ν^{NN} , which encodes hard-neutrino exchange, can be determined by (i) matching the chiral EFT $nn \rightarrow ppee$ 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_ν^{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_ν^{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

$$\mathcal{L}_{\Delta L=2} = -\frac{m_{\beta\beta}}{2} \nu_{eL}^T C \nu_{eL}, \quad (1)$$

where $C = i\gamma_2\gamma_0$ denotes the charge conjugation matrix.

The nuclear effective Hamiltonian can be written as

$$H_{\text{eff}} = H_{\text{strong}} + 2G_F^2 V_{ud}^2 m_{\beta\beta} \bar{e}_L C \bar{e}_L^T V_\nu, \quad (2)$$

in terms of the Fermi constant G_F and the V_{ud} element of the CKM matrix [32, 33]. The neutrino potential V_ν can be obtained from two-nucleon irreducible diagrams mediating $nn \rightarrow ppee$ to a given order in p/Λ_χ . Within Weinberg's power counting the only LO contribution [26] comes from the exchange of potential neutrinos, with $q^0 \ll |\mathbf{q}|$,

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

where $g_A \simeq 1.27$ is the nucleon axial coupling, m_π the pion mass, and \mathbf{q} the momentum transfer. N²LO terms arise from corrections to the single nucleon weak currents, irreducible one-loop diagrams, and contact interactions mediating $\pi\pi \rightarrow ee$, $n \rightarrow p\pi^+ee$, and $nn \rightarrow ppee$. In particular, the short-range potential includes a two-nucleon term [26]

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

where the LEC g_ν^{NN} is $\mathcal{O}((4\pi F_\pi)^{-2})$ 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 currents [38], due to a conflict between naive dimensional analysis and nonperturbative renormalization. We therefore investigate the scaling of g_ν^{NN} by studying the amplitude $\mathcal{A}(nn \rightarrow ppee) \equiv \mathcal{A}_{\Delta L=2}$ with strong interactions, H_{strong} , included nonperturbatively.

We work at LO in chiral EFT, and focus on the scattering of two neutrons to two protons in the 1S_0 wave, where H_{strong} has short-range and Yukawa components,

$$V_0(\mathbf{q}) = \tilde{C} + V_\pi(\mathbf{q}), \quad V_\pi(\mathbf{q}) = -\frac{g_A^2}{4F_\pi^2} \frac{m_\pi^2}{\mathbf{q}^2 + m_\pi^2}, \quad (5)$$

with $\tilde{C} \sim \mathcal{O}(F_\pi^{-2}, m_\pi^2 F_\pi^{-4})$ [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 $\mathcal{A}_{\Delta L=2}$ from the exchange of a light neutrino ($\mathcal{A}_{\Delta L=2}^{(\nu)}$) are shown in Fig. 1. The blue ellipse denotes the iteration of the Yukawa potential $V_\pi(\mathbf{q})$. The diagrams in the second and third rows include an infinite number of bubbles, dressed with iterations of V_π . Without loss of generality for our arguments, we use the kinematics $n(\mathbf{p}) n(-\mathbf{p}) \rightarrow p(\mathbf{p}') p(-\mathbf{p}') e(\mathbf{p}_{e1} = 0) e(\mathbf{p}_{e2} = 0)$, with $|\mathbf{p}| = 1$ MeV and correspondingly $|\mathbf{p}'| = 38$ MeV.

$\mathcal{A}_{\Delta L=2}^{(\nu)}$ can be expressed in terms of the Yukawa “in” and “out” wavefunctions $\chi_{\mathbf{p}}^\pm(\mathbf{r})$ and the propagators

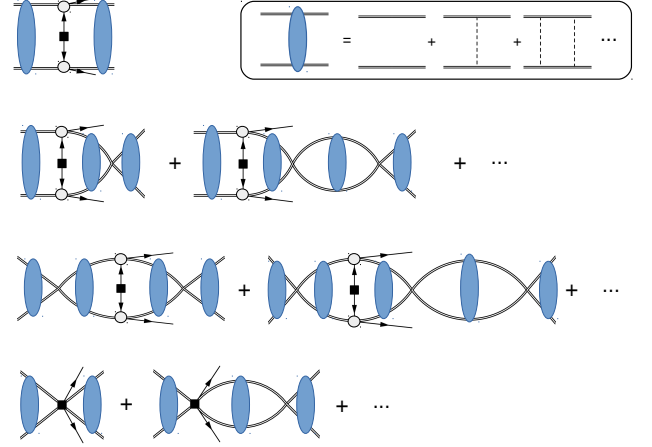


FIG. 1. Diagrammatic representation of LO contributions to $nn \rightarrow ppee$. Double, dashed, and plain lines denote nucleons, pions, and leptons, respectively. Gray circles denote the nucleon axial and vector currents, and the black square an insertion of $m_{\beta\beta}$. The blue ellipse represents iteration of V_π . In the counterterm amplitude (fourth line) the black square represents g_ν^{NN} . The \dots in the second to fourth lines denote diagrams with arbitrary numbers of bubble insertions.

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

$$\mathcal{A}_{\Delta L=2}^{(\nu)} = \mathcal{A}_A + K_{E'} \mathcal{A}_B + \bar{\mathcal{A}}_B K_E + K_{E'} \mathcal{A}_C K_E, \quad (6)$$

$$K_E = \frac{\chi_{\mathbf{p}}^+(\mathbf{0}) \tilde{C}}{1 - \tilde{C} G_E^+(\mathbf{0}, \mathbf{0})},$$

where \mathcal{A}_A , \mathcal{A}_B , and \mathcal{A}_C denote the first diagram in the first, second, and third rows of Fig. 1, respectively (without the wavefunctions at $\mathbf{0}$, in the case of \mathcal{A}_B and \mathcal{A}_C). $\bar{\mathcal{A}}_B$ is similar to \mathcal{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 $\mathcal{A}_{\Delta L=2}^{(\nu)}$. $\chi_{\mathbf{p}}^+(\mathbf{0})$ is finite and the divergence in $G_E^+(\mathbf{0}, \mathbf{0})$ is absorbed by \tilde{C}^{-1} , so that K_E is finite and scheme-independent [34]. We note that:

(i) All diagrams in \mathcal{A}_A are finite. The tree level is finite and each V_π iteration improves the convergence by bringing in a factor of $d^3\mathbf{k}/(\mathbf{k}^2)^2$, where one \mathbf{k}^2 comes from the pion propagator and the other from the two-nucleon propagator.

(ii) All the diagrams in \mathcal{A}_B and $\bar{\mathcal{A}}_B$ are finite. The first loop goes as $d^3\mathbf{k}/(\mathbf{k}^2)^2$, while V_π insertions further improve the convergence.

(iii) The first two-loop diagram in \mathcal{A}_C has a logarithmic divergence, which stems from an insertion of the most

singular component of the neutrino potential, namely

$$\tilde{V}_\nu(\mathbf{q}) = \tau^{(1)+}\tau^{(2)+} \frac{1}{\mathbf{q}^2} \left(1 - \frac{2}{3} g_A^2 \boldsymbol{\sigma}^{(1)} \cdot \boldsymbol{\sigma}^{(2)} \right). \quad (7)$$

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

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

$$\begin{aligned} \mathcal{A}_C^{(\text{div})} &= - \left(\frac{m_N}{4\pi} \right)^2 (1 + 2g_A^2) \left[\Delta + L_{\mathbf{p},\mathbf{p}'}(\mu) \right], \quad (8) \\ L_{\mathbf{p},\mathbf{p}'}(\mu) &= \frac{1}{2} \left(\log \frac{\mu^2}{-(|\mathbf{p}| + |\mathbf{p}'|)^2 + i0^+} + 1 \right), \end{aligned}$$

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

$$\mathcal{A}_{\Delta L=2}^{(NN)} = K_{E'} \frac{2g_\nu^{NN}}{\tilde{C}^2} K_E, \quad (9)$$

and we can renormalize $\mathcal{A}_{\Delta L=2}$ by replacing $\mathcal{A}_C \rightarrow \mathcal{A}_C + 2g_\nu^{NN}/\tilde{C}^2$ in Eq. (6). In the $\overline{\text{MS}}$ scheme,

$$\mathcal{A}_C \rightarrow \delta\mathcal{A}_C + \left(\frac{m_N}{4\pi} \right)^2 \left[2\tilde{g}_\nu^{NN}(\mu) - (1 + 2g_A^2) L_{\mathbf{p},\mathbf{p}'}(\mu) \right] \quad (10)$$

after defining the dimensionless coupling

$$\tilde{g}_\nu^{NN} = \left(\frac{4\pi}{m_N \tilde{C}} \right)^2 g_\nu^{NN}. \quad (11)$$

This coupling obeys the renormalization-group equation (RGE)

$$\mu \frac{d\tilde{g}_\nu^{NN}}{d\mu} = \frac{1}{2} (1 + 2g_A^2), \quad (12)$$

confirming that $\tilde{g}_\nu^{NN} \sim \mathcal{O}(1)$. Since $\tilde{C}(\mu = m_\pi) \approx -0.9/F_\pi^2$, we find that $g_\nu^{NN} \sim \mathcal{O}(F_\pi^{-2})$ instead of $\mathcal{O}((4\pi F_\pi)^{-2})$. A similar enhancement also occurs in four-nucleon couplings induced by higher-dimensional LNV operators. Treating V_π as a subleading correction [35, 39] is equivalent to working to LO in pionless EFT, and does not affect our conclusions about the importance of g_ν^{NN} [26]. Details on how to obtain $\delta\mathcal{A}_C$ will be provided in future work [40].

$\mathcal{A}_{\Delta L=2}$ 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 with a smeared δ -function,

$$\tilde{C} \delta^{(3)}(\mathbf{r}) \rightarrow \frac{\tilde{C}(R_S)}{(\sqrt{\pi} R_S)^3} \exp\left(-\frac{r^2}{R_S^2}\right) \equiv \tilde{C}(R_S) \delta_{R_S}^{(3)}(\mathbf{r}), \quad (13)$$

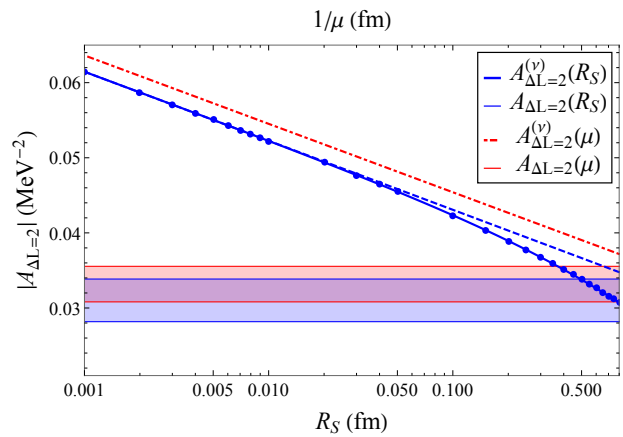


FIG. 2. Matrix element $\mathcal{A}_{\Delta L=2}^{(\nu)}$ for $|\mathbf{p}| = 1$ MeV and $|\mathbf{p}'| = 38$ MeV, as a function of R_S . The dashed line shows a fit to $a + b \log R_S$, which captures the small R_S behavior. The solid line corresponds to a fit that includes $\mathcal{O}(R_S, R_S \log R_S)$ power corrections. The dash-dotted line shows $\mathcal{A}_{\Delta L=2}^{(\nu)}$ in $\overline{\text{MS}}$ as a function of $1/\mu$. The horizontal bands represent the total amplitude $\mathcal{A}_{\Delta L=2}$ with $g_\nu^{NN} = (C_1 + C_2)/2$, as discussed in the main text.

and obtain $\psi_{\mathbf{p}}^-(\mathbf{r})$ and $\psi_{\mathbf{p}}^+(\mathbf{r})$ by solving the Schrödinger equation. We determine $\tilde{C}(R_S)$ by requiring that the 1S_0 scattering length be reproduced ($\tilde{C} \approx -0.4/F_\pi^2$ at $R_S = 0.8$ fm). We find that $1/\tilde{C}(R_S)$ has linear ($1/R_S$) and logarithmic divergences [35] and that the 1S_0 phase shifts at nonzero momentum are indeed R_S -independent.

We then compute

$$\mathcal{A}_{\Delta L=2}^{(\nu)} = - \int d^3\mathbf{r} \psi_{\mathbf{p}}^-(\mathbf{r})^* V_{\nu,0}(\mathbf{r}) \psi_{\mathbf{p}}^+(\mathbf{r}), \quad (14)$$

where $V_{\nu,0}(\mathbf{r})$ is obtained by Fourier-transforming the 1S_0 projection of Eq. (3). In Fig. 2 we plot $\mathcal{A}_{\Delta L=2}^{(\nu)}$ 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 $V_{\nu,CT}(\mathbf{r}) = -2g_\nu^{NN}(R_S)\delta_{R_S}^{(3)}(\mathbf{r})$. The corresponding amplitude,

$$\mathcal{A}_{\Delta L=2}^{(NN)} = - \int d^3\mathbf{r} \psi_{\mathbf{p}}^-(\mathbf{r})^* V_{\nu,CT}(\mathbf{r}) \psi_{\mathbf{p}}^+(\mathbf{r}), \quad (15)$$

is also regulator-dependent. As expected from Eq. (9), we find its leading divergent behavior to be well reproduced by $1/\tilde{C}(R_S)^2$. We can then make $\mathcal{A}_{\Delta L=2} = \mathcal{A}_{\Delta L=2}^{(\nu)} + \mathcal{A}_{\Delta L=2}^{(NN)}$ finite for $R_S \rightarrow 0$ and R_S -independent by choosing $\tilde{g}_\nu^{NN}(R_S) = -(a/2)(1 + 2g_A^2) \log R_S + b + cR_S + \dots$, with the coefficient of the logarithm quite close to the $\overline{\text{MS}}$ expectation $a = 1$.

Relating g_ν^{NN} to electromagnetic isospin violation: The finite part of g_ν^{NN} can be obtained by matching the chiral EFT amplitude to a lattice QCD calculation

performed at the same kinematic point, as it is done in the strong-interacting sector [41]. First lattice results related 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 $0\nu\beta\beta$ 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 electromagnetic current with exchange of hard virtual photons [44, 45]. In the $\Delta L = 2$ case, up to a proportionality factor, the *same* operators are generated by the insertion of two weak currents with exchange of hard neutrinos. This comes about because the neutrino propagator and weak vertices combine to give a massless gauge-boson propagator in Feynman gauge, multiplied by $8G_F^2 V_{ud}^2 m_{\beta\beta} \bar{e}_L e_L^c$ [26]. The LECs needed for $0\nu\beta\beta$ are therefore related to the LECs associated with the isospin $I = 2$ component of the product of two electromagnetic currents, which belongs to the $5_L \times 1_R$ irreducible representation of chiral $SU(2)_L \times SU(2)_R$.

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

$$\begin{aligned} O_1 &= \bar{N} \mathcal{Q}_L N \bar{N} \mathcal{Q}_L N - \frac{\text{Tr}[\mathcal{Q}_L^2]}{6} \bar{N} \boldsymbol{\tau} N \cdot \bar{N} \boldsymbol{\tau} N + \{L \leftrightarrow R\}, \\ O_2 &= 2 \left(\bar{N} \mathcal{Q}_L N \bar{N} \mathcal{Q}_R N - \frac{\text{Tr}[\mathcal{Q}_L \mathcal{Q}_R]}{6} \bar{N} \boldsymbol{\tau} N \cdot \bar{N} \boldsymbol{\tau} N \right), \end{aligned} \quad (16)$$

where $\mathcal{Q}_L = u^\dagger Q_L u$, $\mathcal{Q}_R = u Q_R u^\dagger$, $u = \exp(i\boldsymbol{\tau} \cdot \boldsymbol{\pi}/(2F_\pi))$, and $Q_{L,R}$ are ‘‘spurions’’ transforming under the chiral group as $Q_L \rightarrow L Q_L L^\dagger$, $Q_R \rightarrow R Q_R R^\dagger$. In the electromagnetic case $Q_L = Q_R = \tau_3/2$, while in $0\nu\beta\beta$ $Q_L = \tau^+$, $Q_R = 0$. In our conventions O_1 enters the $\Delta L = 2$ Lagrangian with coefficient $2G_F^2 V_{ud}^2 m_{\beta\beta} g_\nu^{NN}$. Defining the electromagnetic LECs multiplying $O_{1,2}$ as $e^2 C_{1,2}/4$, chiral symmetry dictates $g_\nu^{NN} = C_1$.

In the electromagnetic case, O_1 and O_2 only differ at the multipion level, and an isospin-breaking two-nucleon observable, such as the $I = 2$ combination of scattering lengths $a_{\text{CIB}} = (a_{nn} + a_{pp})/2 - a_{np}$, only constrains the sum $C_1 + C_2$. Extracting this combination from data provides a rough estimate of g_ν^{NN} under the assumption $C_1 \sim C_2$. As in the $\Delta L = 2$ case, we introduce the dimensionless couplings $\tilde{C}_i \equiv [4\pi/(m_N \tilde{C})]^2 C_i$ and compute the scattering lengths $a_{pp,nn,np}$ including the leading sources of isospin breaking – the Coulomb potential and pion mass splitting – and $\tilde{C}_1 + \tilde{C}_2$. Similarly to the $\Delta L = 2$ case, we find that $\tilde{C}_1 + \tilde{C}_2$ needs to be promoted to LO and obeys the RGE

$$\mu \frac{d}{d\mu} \frac{\tilde{C}_1 + \tilde{C}_2}{2} = \frac{1}{2} \left(1 + g_A^2 \frac{m_{\pi^+}^2 - m_{\pi^0}^2}{e^2 F_\pi^2} \right), \quad (17)$$

while, of course, \tilde{C}_1 has the same RGE as \tilde{g}_ν^{NN} . By fitting to a_{CIB} using $a_{np} = -23.7$ fm, $a_{nn} = -18.9$ fm, and $a_{pp} = -7.8$ fm, we find $(\tilde{C}_1 + \tilde{C}_2)/2 \approx 2.5$ at $\mu = m_\pi$ in

the $\overline{\text{MS}}$ scheme. Using instead the R_S scheme, we find $(\tilde{C}_1 + \tilde{C}_2)/2 \approx 2.0$ at $R_S = 0.5$ fm.¹ This estimate, based on data and chiral symmetry, again confirms that $g_\nu^{NN} \sim \mathcal{O}(F_\pi^{-2})$.

Numerical impact: To roughly estimate the impact of the contact term, we assume for concreteness $C_1 = C_2$ and hence $g_\nu^{NN} = (C_1 + C_2)/2$ at some \bar{R}_S or $\bar{\mu}^{-1}$ in the range 0.002 – 0.8 fm. The total two-nucleon amplitude $\mathcal{A}_{\Delta L=2} = \mathcal{A}_{\Delta L=2}^{(\nu)} + \mathcal{A}_{\Delta L=2}^{(NN)}$ then becomes independent of the regulator, as illustrated in Fig. 2, where the widths of the horizontal bands reflect the ambiguity in the choice of the point \bar{R}_S or $\bar{\mu}$ where $C_1 = C_2$ is assumed. (They do not account for the uncontrolled error of the assumption itself.) The relative size of the two components depends on R_S , with $\mathcal{A}_{\Delta L=2}^{(NN)}/\mathcal{A}_{\Delta L=2}^{(\nu)} \sim 30\%$ at $R_S \sim 0.1$ fm, decreasing to $\sim 10\%$ at $R_S \sim 0.6$ fm. More insight can be obtained by plotting the matrix-element densities ρ_ν and ρ_{NN} defined as

$$\mathcal{A}_{\Delta L=2}^{(\nu)} = \int dr \rho_\nu(r), \quad \mathcal{A}_{\Delta L=2}^{(NN)} = \int dr \rho_{NN}(r). \quad (18)$$

Figure 3 (top panel) shows that $\rho_{NN}(r)$ is concentrated at smaller distances than $\rho_\nu(r)$, and its contribution to the amplitude is thus partially diluted.

We have performed a similar analysis for $A = 6, 12$ nuclei, using Variational Monte Carlo nuclear wavefunctions [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 $0\nu\beta\beta$ 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 $V_{\nu,CT}(\mathbf{r})$ with AV18’s short-range CIB potential.

For $\Delta I = 0$ transitions such as the ${}^6\text{He} \rightarrow {}^6\text{Be}$ shown in Fig. 3 (middle panel), we find $\mathcal{A}_{\Delta L=2}^{(NN)}/\mathcal{A}_{\Delta L=2}^{(\nu)} \sim 10\%$, similarly to the $nn \rightarrow ppee$ case. In realistic $0\nu\beta\beta$ transitions, however, the total nuclear isospin changes by two units, $\Delta I = 2$. This implies the presence of a node in $\rho_\nu(r)$ due to the orthogonality of the initial and final spatial wavefunctions. The resulting partial cancellation between the regions with $r \lesssim 2$ fm and $r \gtrsim 2$ fm [51] leads to a relative enhancement of the short-range contribution, as illustrated in Fig. 3 (bottom panel) for ${}^{12}\text{Be} \rightarrow {}^{12}\text{C}$. Numerically we find $\mathcal{A}_{\Delta L=2}^{(NN)}/\mathcal{A}_{\Delta L=2}^{(\nu)} \sim 25\%$ (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 components of the CIB interaction induce effects of similar size.

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

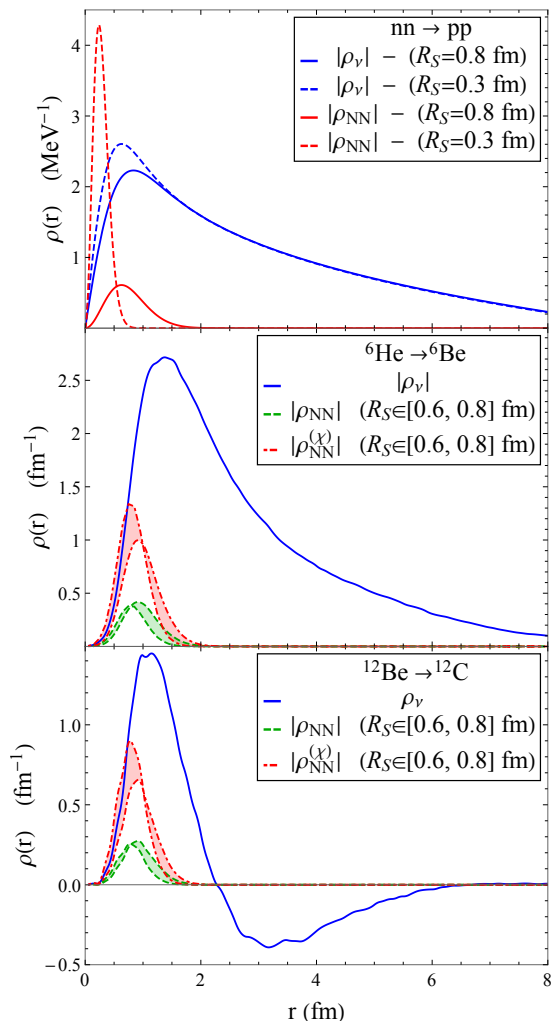


FIG. 3. $\rho_\nu(r)$ and $\rho_{NN}(r)$ for the $nn \rightarrow ppee$ process (top), and for nuclear transitions with $A = 6$ (middle) and $A = 12$ (bottom). In the middle and bottom panels the green (ρ_{NN}) and red ($\rho_{NN}^{(\chi)}$) bands correspond to $g_\nu^{NN} = (C_1 + C_2)/2$ extracted from 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 $^{12}\text{Be} \rightarrow ^{12}\text{C}$ and experimentally relevant transitions to be of comparable size.

Conclusion: The above arguments suggest that the new short-range $\Delta L = 2$ 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_ν^{NN} from lattice QCD and an assessment of the impact of the short-range potential in nuclei of experimental interest.

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