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Lepton- and baryon-number violation in nuclear effective field theory

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Lepton- and baryon-number violation in nuclear effective field theory

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Abstract. Lepton and baryon numbers are accidental symmetries in the effective field theory (EFT) of the Standard Model. Their violation would lead to spectacular nuclear decays which can be described with nuclear EFTs. I discuss ways in which two such decays — the neutrinoless double-beta decay of heavy nuclei and the disappearance of the deuteron — provide complementary information to processes that take place outside the nuclear environment — neutrino masses, nucleon decay, and neutron-antineutron oscillations.

1. Accidental Symmetries

The beautiful views from Ischia seemed to me a good place to discuss the expanded horizons of the Standard Model (SM) viewed as an effective field theory (EFT)— the now popular SMEFT. At a momentum scale $M_{\rm EW} \sim 100$ GeV, the relevant degrees of freedom are apparently quarks, charged leptons, left-handed neutrinos, gauge bosons, and a Higgs boson with interactions constrained by Lorentz symmetry and the gauge group $\mathrm{SU}(3)_{\rm c}\times\mathrm{SU}(2)_{\rm L}\times\mathrm{U}(1)_{\rm Y}$. As we all know, originally Weinberg imposed [1] the restriction that interactions be represented by operators of canonical dimension no larger than four, which most people continue to call "the SM". But it is unlikely that no new — "beyond the SM" (BSM) — physics exists beyond $M_{\rm EW}$. Already in lectures on quantum field theory in the 1980s Weinberg articulated the SM as an EFT, following his formulation of the EFT framework in 1979 [2] and the immediate application to the SM [3].

What are the new horizons? Any new degrees of freedom manifest at a scale $M_{\rm hi} > M_{\rm EW}$ entail higher-dimensional operators when we describe processes with characteristic momenta $Q \lesssim M_{\rm hi}$. On the basis of naturalness [4, 5], one expects the corresponding strengths to be suppressed by powers of $M_{\rm hi}^{-1}$. This can be represented by an extension of the SM Lagrangian $\mathcal{L}_{\rm SM}$, schematically

$$\mathcal{L}_{\rm SMEFT} = \mathcal{L}_{\rm SM} + \sum_{d=5}^{\infty} \frac{1}{M_{\rm hi}^{d-4}} \sum_i c_i^{(d)} O_i^{(d)} , \tag{1}$$

where $O_i^{(d)}$ are various operators of dimension d with SM fields and symmetries, and $c_i^{(d)}$ are "Wilson" or "low-energy" coefficients (LECs). Effects of the $d \geq 5$ interactions on observables should then be suppressed by powers of $Q/M_{\rm hi}$ compared to the leading-order (LO) interactions in $\mathcal{L}_{\rm SM}$. The EFT framework not only incorporates the higher-energy physics but also explains why the interactions of "the SM" are dominant for $Q \ll M_{\rm hi}$. SM interactions generate the main regularities seen in nuclear structure, reactions, and decays. My talk in Ischia was about the tiny effects of some of the higher-dimensional operators.

Tiny effects stand out when they violate "accidental" symmetries, defined as symmetries that emerge up to some order in an EFT as a consequence of its particle content but are not symmetries of the underlying theory. An example [6] is isospin in nuclear EFTs [7], which is violated only at subleading orders even though light quarks are not close to being degenerate in mass. The accidental symmetries I tackle here are lepton (L) and baryon (B) numbers. As Weinberg discovered [8, 9] there is only one $d = 5$ operator, which violates L by two units ($|\Delta L| = 2$) and generates a Majorana neutrino mass. It might be no accident that neutrino mass was the first solid deviation from "the SM" found experimentally. Weinberg [8, 9] and others [10, 11, 12, 13] also pointed out that the most important violation of B should arise at $d = 6$, generating processes such as nucleon decay where B and L change by one unit ($\Delta B = \Delta L = \pm 1$). $\Delta B = \pm 2$ processes such as neutron-antineutron oscillation are qualitatively different and stem from $d = 9$ operators [14, 15, 16, 17]. The violation of these symmetries can in principle arise at different scales, and in the following $M_{\rm hi} \to M_{|\Delta L|=2}, M_{|\Delta B|=1}, M_{|\Delta B|=2}$. If these scales are very different, higher-dimensional operators could become more prominent.

L and B violation lead to new, distinctive nuclear processes, which are crucial in unraveling the structure of the corresponding BSM operators. Practically the only way to test the existence of the $d = 5$ operator is to measure the neutrinoless double-beta ($0\nu 2\beta$) decay of heavy nuclei, which proceeds dominantly through the transformation of two neutrons into two protons and two electrons ($nn \to ppee$). $|\Delta B| = 1, 2$ operators lead, in addition to free proton decay and neutron-antineutron oscillation, also to the decay of otherwise stable nuclei such as the deuteron ($d \to NX$ and $d \to X$, where X is not a nucleon N). The relative importance of one- and two-nucleon mechanisms depends on the structure of the BSM operators. Thus, the nuclear environment offers information complementary to free processes. In fact, we demonstrated several years ago (for reviews, see Refs. [18, 19]) how the eventual measurement of light nuclear electric dipole moments will allow inferences about the structure of time-reversal-violating interactions from "the SM" and from $d = 6$ BSM operators.

We are now taking the first steps to do the same for L and B violation, as I report below for $|\Delta L| = 2$, $|\Delta B| = 1$, and $|\Delta B| = 2$ in Secs. 3, 4, and 5, respectively. The Chiral EFT (ChEFT) [7] that enables this approach is first briefly recalled in Sec. 2, while some final thoughts are offered in Sec. 6.

2. Chiral EFT

Below the nonperturbative scale $M_{\rm QCD} \sim 1$ GeV, QCD is best described by EFTs with hadronic degrees of freedom. In order to draw conclusions from nuclear observables, we need an EFT at the momentum scale $M_{\rm nuc} \sim 100$ MeV characteristic of most nuclei. Since this scale is not very different from the pion mass $m_\pi \simeq 140$ MeV, it is useful to retain pions as explicit degrees of freedom. ChEFT [7] is based on the approximate $\mathrm{SU}(2)_{\rm L} \otimes \mathrm{SU}(2)_{\rm R}$ chiral symmetry of QCD, which is broken spontaneously to the isospin subgroup $\mathrm{SU}(2)_{\rm L+R}$ and explicitly by the u- and d-quark masses. Pions are pseudo-Goldstone bosons with interactions containing either derivatives or powers of m_π^2, which are proportional to powers of $F_\pi^{-1} \sim 2\pi M_{\rm QCD}^{-1}$, where $F_\pi \simeq 185$ MeV is the pion decay constant. Since nucleons have a mass $m_N \simeq 940$ MeV $\sim M_{\rm QCD}$, they are nonrelativistic in the regime of validity of the theory.

The chiral Lagrangian

$$\mathcal{L}_{\rm ChEFT} = \mathcal{L}_{B,L} + \mathcal{L}_{|\Delta L|=2} + \mathcal{L}_{|\Delta B|=1} + \mathcal{L}_{|\Delta B|=2} + \dots \qquad (2)$$

is constructed from all terms involving isodoublet nucleon ($N = (p\ n)^T$), antinucleon ($N^c = (p^c\ n^c)^T$), and isotriplet pion (π_k) fields, which transform under symmetries in the same way as

SMEFT interactions. The SM interactions are well known [7, 20],

$$
\begin{aligned}
\mathcal{L}_{B,L} &= N^\dagger \left(i\partial_0 + \frac{\nabla^2}{2m_N}\right) N + N^{c\dagger}\left(i\partial_0 + \frac{\nabla^2}{2m_N}\right) N^c - \frac{1}{2}\pi_k\left(\partial^2 + m_\pi^2\right)\pi_k \\
&+\frac{g_A}{F_\pi}\left(N^\dagger\sigma_k\tau_l N + N^{c\dagger}\sigma_k\tau_l^T N^c\right)\nabla_k\pi_l \\
&-\sqrt{2}G_F V_{ud}\left\{\left(\bar{e}_L\gamma^\mu\nu_{eL}\right)\left[F_\pi\partial_\mu\pi^- + N^\dagger\tau_+\left(\delta_{\mu 0} - g_A\delta_{\mu k}\sigma_k\right)N\right] + \text{H.c.}\right\} \\
&-\sum_{s=0,1}\left(C_{0s} + D_{2s}m_\pi^2\right)\left(N^T P_k^{(s)} N\right)^\dagger\left(N^T P_k^{(s)} N\right) \\
&+\sum_{s=0,1}\frac{C_{2s}}{8}\left\{\left(N^T P_k^{(s)} N\right)^\dagger\left[N^T P_k^{(s)}(\vec{\nabla} - \overleftarrow{\nabla})^2 N\right] + \text{H.c.}\right\} \\
&-H_0\left(N^{cT}\tau_2 Y_{kl} N\right)^\dagger\left(N^{cT}\tau_2 Y_{kl} N\right) + \ldots\,, \qquad (3)
\end{aligned}
$$

where σ_k (τ_k) are Pauli spin (isospin) matrices (with subscript $\pm$ denoting raising/lowering operators), $P_k^{(1)} = \sigma_2\sigma_k\tau_2/\sqrt{8}$ ($P_k^{(0)} = \tau_2\tau_k\sigma_2/\sqrt{8}$) is the projector onto the NN $^3S_1$ ($^1S_0$) state, and $Y_{kl} = \sigma_2\sigma_k\tau_2\tau_l/2$ is the projector onto the $N\bar{N}$ isospin-triplet $^3S_1$ state. Here $g_A \simeq 1.27$ is the nucleon axial-vector coupling, G_F is the Fermi coupling constant, and V_{ud} is the ud element of the Cabibbo-Kobayashi-Maskawa (CKM) matrix, while C_{0s}, C_{2s}, D_{2s} are real LECs determined from nucleon-nucleon (NN) scattering and H_0 is a complex LEC contributing to nucleon-antinucleon ($N\bar{N}$) annihilation. The nonrelativistic character of (anti)nucleons guarantees that interactions with more than $2A$ baryon fields do not contribute to the A-baryon system. Here and below, terms that are not important for the remainder of the paper are buried in the "...".

The nonperturbative nature of the NN system means that renormalization is quite different [21] than in a purely perturbative context. Naturalness implies [22] significant changes with respect to the scaling of the LECs in the mesonic and single-nucleon sectors. The unnaturally large values of the NN scattering lengths require fine tuning and resummation of C_{0s} contact interactions, which run with the regulator cutoff Λ or, equivalently, with the renormalization scale μ. The finite part of the LEC $C_{01} = \mathcal{O}(4\pi/m_N\kappa)$, for example, is responsible for the deuteron binding momentum $\kappa = \sqrt{m_N B_d} \simeq 45$ MeV, with $B_d \simeq 2.225$ MeV the binding energy. For small momenta, pions can be treated perturbatively together with other contact interactions in a Q/M_{NN} expansion, where $M_{NN} \equiv 4\pi F_\pi^2/g_A^2 m_N \simeq 300$ MeV [23, 24]. With derivative interactions such as the one with LEC $C_{21} = \mathcal{O}(4\pi/m_N\kappa^2 M_{NN})$, pions account for the energy dependence of NN amplitudes, for example at NLO the $^3S_1$ np effective range $r_{np} \simeq 1.75$ fm. As Q increases, pions become nonperturbative in the lowest partial waves [25] and renormalization requires additional LECs at LO in spin-triplet waves [26, 27]. In $^1S_0$, an $m_\pi^2 \ln\Lambda$ (or alternatively $m_\pi^2 \ln\mu$) dependence appears in the amplitude at the order one-pion exchange enters, which demands D_{20} at the same order [28, 29]. Higher orders remain perturbative and the expansion for observables is now in $Q/M_{\rm QCD}$.

In contrast, the (complex) $N\bar{N}$ scattering lengths are apparently not particularly large on the scale set by M_{NN} [30, 31]. At low momenta, all interactions can be treated in perturbation theory, starting with $H_0 = \mathcal{O}(4\pi/m_N M_{NN})$ which is then determined in terms of the isospin-triplet $^3S_1$ scattering length $a_{\bar{n}p} \simeq 0.44 + i0.96$ fm [31]. As Q increases one expects that, like in NN, pions become nonperturbative in the lowest waves together with contact interactions needed for renormalization [32].

With these interactions, one can build NN and $N\bar{N}$ wavefunctions and address processes driven by strong and electroweak interactions [7], or by BSM physics.

3. Neutrinoless double-beta decay

Weinberg's $d = 5$ L-violating operator [8, 9] gives rise at low energies to [33, 34]

$$\begin{aligned}\mathcal{L}_{|\Delta L|=2} &= -\frac{m_{\beta\beta}}{2}\,\nu_{eL}^T C\nu_{eL} \\ &\quad - \left(2\sqrt{2}\,G_F V_{ud}\right)^2 m_{\beta\beta}\, g_\nu^{\rm NN}\left[\left(\bar{e}_L C\bar{e}_L^T\right)\left(N^T P_+^{(0)} N\right)\left(N^T P_-^{(0)} N\right)^\dagger + \text{H.c.}\right] + \ldots \quad (4)\end{aligned}$$

The first term, where $C = i\gamma_2\gamma_0$ is the charge conjugation matrix, combines neutrino masses and elements of the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix into the effective neutrino mass $m_{\beta\beta} = \mathcal{O}(M_{\rm EW}^2/M_{|\Delta L|=2})$. Together with the SM coupling of the neutrino to the nucleon, it gives rise to neutrino exchange at distances $\gtrsim M_{\rm QCD}^{-1}$. In the dominant NN channel, $^1S_0$, the transition operator for $nn \to ppee$ is proportional to

$$V_{\nu\rm L}^{(^1S_0)}(\vec{q}) = \frac{\tau_+^{(1)}\tau_+^{(2)}}{\vec{q}^{\,2}}\left[1 + 2g_A^2 + \frac{g_A^2 m_\pi^4}{(\vec{q}^{\,2} + m_\pi^2)^2}\right], \tag{5}$$

where I neglected the tiny $m_{\beta\beta}$ compared to the magnitude of the momentum $\vec{q}$ transferred between nucleons labeled by superscripts. The second term in Eq. (4) is the simplest L-violating contact two-electron/four-nucleon interaction, with the nucleon pair in the $^1S_0$ channel. It can be thought of as representing Majorana neutrino exchange at distances $\lesssim M_{\rm QCD}^{-1}$. With factors of $m_{\beta\beta}$, G_F, and V_{ud} factored out, its LEC, denoted as $g_\nu^{\rm NN}$, depends on the details of the QCD dynamics at short distances and is poorly known.

For the calculation of $0\nu2\beta$ decay, we need to estimate the importance of $g_\nu^{\rm NN}$. Naturalness implies that the order in the EFT expansion where $g_\nu^{\rm NN}$ enters is no higher than the order where long-range neutrino exchange generates essential regulator dependence ("divergence" in old parlance). The surprise [33, 34] is that $g_\nu^{\rm NN}$ is needed at LO! The average of the transition operator (5) with the LO $^1S_0$ wavefunction has an $\ln\Lambda$ (or alternatively $\ln\mu$) dependence, which can only be absorbed by $g_\nu^{\rm NN}$. Fortunately there is no renormalization argument for new L-violating short-range interactions at NLO [34].

In hindsight this should not have come as a surprise: the situation here is similar to the need for D_{20} along with one-pion exchange [28, 29]. The same sensitivity to short-range physics appears in electromagnetic interactions [35, 36], since the Coulomb potential also has a $\vec{q}^{\,-2}$ decrease at large momenta. Renormalization similarly demands at the same order a combination of two possible isospin-violating NN contact interactions, which can be determined from the charge-independence-breaking (CIB) scattering length. Due to isospin, the structure of one of these interactions is the same as the $g_\nu^{\rm NN}$ term [33, 34]. If the coefficient of the other interaction were known, $g_\nu^{\rm NN}$ could be extracted. This conclusion is backed up by an analysis of QCD in the limit of large number of colors [37]. Efforts are underway to obtain a phenomenological estimate of $g_\nu^{\rm NN}$ [38] and to calculate it directly from lattice QCD [39, 40]. This is essential to draw quantitative conclusions about $m_{\beta\beta}$ from an eventual signal of $0\nu2\beta$ decay.

4. $\Delta B = 1$ deuteron decay

The 12 independent $|\Delta B| = 1$ operators involving u and d quarks (and one lepton) at $d = 6$ [8, 9, 10, 11, 12, 13] generate a chiral Lagrangian [41]

$$\mathcal{L}_{|\Delta B|=1} = \frac{i}{2}N^\dagger\left(\bar{\alpha}_0 + \bar{\alpha}_1\tau^3\right)N \; - i\bar{C}_0\left(N^T P_k^{(1)} N\right)^\dagger\left(N^T P_k^{(1)} N\right) + \ldots, \tag{6}$$

with real LECs $\bar{\alpha}_{0,1} = \mathcal{O}(M_{\rm QCD}^5 c_i^{(6)2}/(4\pi M_{|\Delta B|=1})^4)$ and $\bar{C}_0 = \mathcal{O}(M_{NN}^4 c_i^{(6)2}/M_{|\Delta B|=1}^4\kappa^2)$. The non-Hermiticity of the Lagrangian reflects nucleon and nuclear decay into an antilepton and mesons with momenta beyond the regime of validity of the EFT.

In fact, the LECs $\bar{\alpha}_{0,1}$ determine the proton and neutron decay widths,

$$\Gamma_p = \bar{\alpha}_0 + \bar{\alpha}_1\,, \qquad \Gamma_n = \bar{\alpha}_0 - \bar{\alpha}_1\,, \tag{7}$$

up to $\mathcal{O}(Q^2/M^2_{\rm QCD})$ corrections [41]. From the expected size of the LECs, $M^4_{|\Delta B|=1}\Gamma_{p,n}/c_i^{(6)2} \sim 10^{-4}\,\mathrm{GeV}^5$, to be compared to $\simeq 3\cdot 10^{-4}\,\mathrm{GeV}^5$ obtained for the decay to a meson and a positron or antimuon in lattice QCD [42].

The deuteron width for $d \to NX$ decay can be calculated in ChEFT with perturbative pions. One finds [41] it to be dominated by the sum of proton and neutron widths with a correction at relative $\mathcal{O}(\kappa M_{NN}/M^2_{\rm QCD})$,

$$\Gamma_d = \Gamma_p + \Gamma_n - \frac{\kappa}{\pi}(\kappa-\mu)^2\,\bar{C}_0(\mu)\,. \tag{8}$$

The contribution from the LEC $\bar{C}_0$ is the first nuclear correction to the deuteron decay rate. It is independent of the renormalization scale μ, as can be seen from the renormalization-group equation for $\bar{C}_0(\mu)$ [41]. It is presently unknown, but it is expected to be of the order of a few percent. The dominance of nucleon decay rates agrees with earlier, model-based calculations [43, 44], although our estimate of corrections is smaller likely because of our explicit implementation of chiral symmetry.

5. $\Delta B = 2$ deuteron decay

There are four independent $d = 9$ $|\Delta B| = 2$ operators involving u and d quarks [14, 15, 16, 17], three of which are in the same chiral irreducible representation, one in another [45]. Their descendants in ChEFT are [46, 20]

$$\begin{aligned}\mathcal{L}_{|\Delta B|=2} \;=\;& -\delta m\left(n^{c\dagger}n + \text{H.c.}\right) - \bar{B}_0\left[\left(N^T P_k^{(1)} N\right)^\dagger \left(N^{cT}\tau_2 Y_{k-}N\right) - \text{H.c.}\right] \\ & - i\bar{C}_0\left(N^T P_k^{(1)} N\right)^\dagger\left(N^T P_k^{(1)} N\right) + \dots\,,\end{aligned} \tag{9}$$

where $\delta m = \mathcal{O}(c_i^{(9)} M^6_{\rm QCD}/(4\pi)^4 M^5_{|\Delta B|=2})$, $\bar{B}_0 = \mathcal{O}(c_i^{(9)} M^6_{\rm QCD}/(4\pi)^3\kappa M^2_{NN} M^5_{|\Delta B|=2})$, and $\bar{C}_0 = \mathcal{O}(\bar{B}_0^2 M^2_{\rm QCD}/(4\pi)^2)$.

Due to their chiral properties, only three of the four $|\Delta B| = 2$ operators contribute at lowest orders to the n-$\bar{n}$ oscillation time. Up to $\mathcal{O}(Q^2/M^2_{\rm QCD})$ corrections [46, 20],

$$\tau_{n\bar{n}} = (\delta m)^{-1}\,. \tag{10}$$

For the first lattice QCD simulations of δm, see Ref. [47, 48]. The fourth $|\Delta B| = 2$ operator can only contribute to $\tau^{-1}_{n\bar{n}}$ together with additional isospin violation, leading to a suppression with respect to δm by $\alpha_{\rm em}/4\pi \sim \mathcal{O}(Q^3/M^3_{\rm QCD})$, where $\alpha_{\rm em}$ is the fine-structure constant.

The deuteron decay rate for $d \to X$ can be calculated just like in the $\Delta B = 1$ case. For the three operators contributing to δm, at LO the deuteron decay rate is determined by $\tau_{n\bar{n}}$, but there are corrections at $\mathcal{O}(Q/M_{\rm NN})$ [20],

$$\Gamma_d = -\frac{m_N}{\kappa\tau^2_{n\bar{n}}}\,\mathrm{Im}\,a_{\bar{n}p}\left[1+\kappa\left(r_{np}+2\,\mathrm{Re}\,a_{\bar{n}p}-\frac{g_A^2 m_N}{3\pi F_\pi^2}\,\frac{2-2\xi-5\xi^2+6\xi^3}{1+2\xi}-\frac{\tau_{n\bar{n}}(\kappa-\mu)\,\bar{B}_0(\mu)}{\sqrt{2}\pi\,\mathrm{Im}\,a_{\bar{n}p}}\right)\right], \tag{11}$$

where $\xi = \kappa/m_\pi \simeq 0.32$. The LO result agrees with models [49, 50, 51, 52] for a zero-range potential. The NLO correction is again independent of μ on account of the running of $\bar{B}_0(\mu)$ [20], but this LEC is unknown. With an estimate based on naturalness, the

NLO correction is roughly 50% and dominated by the effective range. As a consequence, $R_d \equiv \Gamma_d^{-1}/\tau_{n\bar{n}}^2 = (1.1 \pm 0.3) \cdot 10^{22}\,\mathrm{s}^{-1}$. Taking the largest value of R_d, the Sudbury Neutrino Observatory limit on Γ_d^{-1} [53] gives $\tau_{n\bar{n}} > 1.6 \cdot 10^8\,\mathrm{s}$, which is a factor of about two stronger than the direct Institut Laue-Langevin limit [54]. The central value for R_d is smaller by a factor $\simeq 2 - 2.5$ than Refs. [51, 55], suggesting either deficiencies in the potential models or a failure of the perturbative-pion expansion. The latter scheme has been tested with the electromagnetic form factors of the deuteron [56, 57, 58, 59], but not yet in $N\bar{N}$ scattering. $N\bar{N}$ scattering certainly deserves more EFT scrutiny.

The contribution of the fourth $|\Delta B| = 2$ operator to Γ_d is not suppressed like to $\tau_{n\bar{n}}$. If the corresponding coefficient happens to be much larger than the others, then at LO the deuteron decay width becomes [20]

$$\Gamma_d = -\frac{\kappa}{\pi}(\kappa - \mu)^2\, \bar{C}_0(\mu) = -4\frac{\kappa^3}{m_N}\,\mathrm{Im}\, a_{np} \tag{12}$$

in terms of the imaginary part of the ${}^3S_1$ np scattering length a_{np}. In this case Γ_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 $d \to X$ and free n-$\bar{n}$ oscillation are both observed, one might be able to infer the dominant source(s) of $|\Delta B| = 2$ violation.

6. Conclusion

Each topic above illustrated an aspect of what Chiral EFT brings to nuclear physics. When discussing $0\nu2\beta$ decay, we see the crucial role played by renormalization: it ensures that observables are independent of the arbitrary regularization procedure, which is one of the ways model assumptions can creep in at low energies. Deuteron decay into one nucleon is an example of how one can quantify departures to known model results, in this case the difference to free-nucleon decay, thanks to a controlled expansion. Deuteron decay into mesons is an instance where the constraints of chiral symmetry lead to relations between in-vacuum and in-medium quantities that might shed light onto the structure of underlying BSM mechanisms.

Now what? All the considerations I presented concerned two-nucleon (sub)processes. Clearly the task ahead is to consider heavier nuclei. For that, one needs consistent wavefunctions, which unfortunately exist only at low orders and for light nuclei. Most calculations have been performed within hybrid schemes where phenomenological wavefunctions are used in either unrenormalized "*ab initio*" or mean-field approaches [7]. Still, such schemes can provide order-of-magnitude estimates.

In the case of $0\nu2\beta$ decay, the impact of the new short-range $|\Delta L| = 2$ mechanism has already been gauged. One should emphasize that this mechanism is not equivalent to strong-interaction correlations that exist beyond the nuclear mean field and are automatically included in *ab initio* calculations. Instead, it represents physics at the QCD scale that needs to be added to any type of nuclear calculation. Assuming the unknown LEC to be of the same size as the one determined from NN CIB, the short-range operator is found to indeed give a contribution to the $0\nu2\beta$ amplitude in nuclei that is comparable (albeit somewhat smaller) to that of the long-range mechanism [33, 34, 60, 61, 62, 63]. In the longer term one of course would like to use wavefunctions free of model assumptions, just like the $|\Delta L| = 2$ mechanism itself.

In the case of deuteron decay, we exploited the small binding of the deuteron to provide analytical results based on an expansion in Q/M_{NN} that includes a perturbative treatment of pion exchange. The same framework can be used for other light nuclei, where the binding energy per particle is still relatively small [7]. In denser nuclei, such as ^{16}O, this expansion is less likely to be valid. If that is the case, we need to resum Q/M_{NN} corrections by treating pion exchange nonperturbatively, being left with an expansion in Q/M_{QCD}. While this complicates matters, it only partially affects power-counting estimates, as contributions from $Q \sim M_{NN}$ are transferred from LECs to explicit pion exchange. For $|\Delta B| = 1$, the decay is still power-counted to be dominated by free-nucleon decay, with corrections from $\bar{C}_0$ and from one-pion exchange expected to appear at relative $\mathcal{O}(M_{NN}^2/M_{\rm QCD}^2) \sim 0.1$ [41]. For $|\Delta B| = 2$, the infrared enhancement by κ^{-1}, which increases the overall sensitivity of Γ_d to $\tau_{n\bar{n}}$, should become less pronounced. Still, 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 $\mathcal{O}(Q^2/M_{\rm QCD}^2)$, or beyond [20]. Thus, we expect in both cases the qualitative features of the deuteron results — dominance of single-nucleon decay and partial discrimination among $|\Delta B| = 2$ operators — to survive. But explicit calculations are, of course, needed.

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