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Relativity violations and beta decay

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In some quantum theories of gravity, deviations from the laws of relativity could be comparatively large while escaping detection to date. In the neutrino sector, precision experiments with beta decay yield new and improved constraints on these countershaded relativity violations. Existing data are used to extract bounds of 3×10^{-8} GeV on the magnitudes of two of the four possible coefficients, and estimates are provided of future attainable sensitivities in a variety of experiments.

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Investigating the behavior of neutrinos has yielded deep physical insights since Pauli predicted their existence in 1930 to rescue the conservation of energy in beta decay [1]. More recently, numerous experiments have accumulated compelling evidence for neutrino oscillations, thereby confirming the existence of physics beyond the standard model (SM) [2]. In this work, we investigate the prospects for using beta decay to search for another popular type of neutrino physics beyond the SM, namely violations of Lorentz symmetry in neutrino propagation. These violations could originate in a Planck-scale unification of quantum gravity with other forces such as string theory [3], and they are expected to be heavily suppressed in effective quantum field theory describing physics at accessible energy scales, typically by a factor involving the tiny ratio m_W/m_P of the electroweak to Planck scales.

The general framework describing deviations from Lorentz symmetry in realistic effective quantum field theory is the standard-model extension (SME) [4]. All quantum field operators for Lorentz violation involved in the propagation of neutrinos have been classified and enumerated [5]. Most of these can be studied using neutrino oscillations, which compare the way different neutrinos propagate and provide interferometric sensitivity to energy differences between neutrinos [6]. Some effects cannot be detected by neutrino oscillations because they are produced by "oscillation-free" operators that change all neutrino energies equally. Most of these can instead be studied by comparing neutrino propagation to other species, such as time-of-flight experiments matching the group velocity of neutrinos with that of photons. However, four oscillation-free operators leave unaffected the neutrino group velocity and so cannot be detected in this way. Instead, they must be accessed through physical processes that involve neutrino phase-space properties, such as quantum decays. These operators are rare examples of "countershaded" Lorentz violations [7]: relativity-violating effects that could be enormous compared to ones suppressed by the ratio m_W/m_P and that nonetheless could have escaped detection to date. These could provide an interesting path for building models with viable Lorentz violation obviating the typical requirement of a heavy suppression factor. The present work focuses on methods to constrain these unique effects.

The four countershaded neutrino operators are of renormalizable mass dimension $d = 3$, are odd under CPT, and are controlled by coefficients conventionally denoted $(a^{(3)}_{\rm of})_{jm}$, where j, m are angular-momentum quantum numbers with $j = 0, 1$. Conservation of energy and momentum is assured by taking these four coefficients to be constant as usual for couplings beyond the SM, so all physics other than Lorentz and CPT violation is conventional. Dimensional arguments suggest these coefficients are likely to dominate at accessible energies and can be measured sensitively in low-energy processes. Here, we demonstrate that experiments on beta decay, a well-studied low-energy process with the potential for precision measurement, provide high sensitivity to these oscillation-free effects. We use available data to improve the current limit on $(a^{(3)}_{\rm of})_{00}$ by over an order of magnitude and to obtain a first measurement of $(a^{(3)}_{\rm of})_{10}$. We show that targeted existing and forthcoming experiments can access all four coefficients and further improve sensitivities, thereby revealing an effect or substantially reducing the window for countershaded Lorentz violation.

For a beta decay involving an antineutrino of mass m_ν and 4-momentum $q^\alpha = (\omega, \boldsymbol{q})$, the antineutrino phase space can be written as $d^3q = f(\omega) d\omega d\Omega$, where the antineutrino function $f(\omega) \approx \omega^2 - \frac{1}{2} m_\nu^2 - 2\omega\delta\omega$ encodes the Lorentz-violating modifications

$$\delta\omega = -\sum_{jm} e^{im\omega_\oplus T_\oplus} \mathcal{N}_{jm} (a^{(3)}_{\rm of})_{jm} \tag{1}$$

arising from the modified antineutrino dispersion relation [5] $\omega = |\boldsymbol{q}| + m_\nu^2/|\boldsymbol{q}| + \delta\omega$. In Eq. (1), the sidereal time $T_\oplus$ controls the harmonic variation of the antineutrino function in the laboratory produced by the Earth's sidereal rotation at frequency $\omega_\oplus \simeq 2\pi/(23\,{\rm h}\,56\,{\rm min})$. The factors $\mathcal{N}_{jm}$ contain information about the direction of propagation of the antineutrinos, expressed relative to the canonical Sun-centered frame [8]. Denoting the electron mass by m_e and its 4-momentum by $p^\alpha = (E, \boldsymbol{p})$, the differential spectrum for a single beta decay is given by $d\Gamma/dT = C(T) \int d\Omega f(T_0 - T)$, where $C(T)$ is a function of the electron kinetic energy $T = E - m_e$ and T_0 is the conventional end-point energy for $m_\nu = 0$. For simplicity, we assume measurable Lorentz violation is limited to the neutrino sector. This choice is compatible with existing constraints on other species, including restrictions from electroweak symmetry [9], and with observability requirements imposed by standard field redefinitions [4]. Early theoretical works considering neutrino Lorentz violation in beta decay include Refs. [10]. Lorentz-violating effects arising from weak interactions in beta decay without neutrino Lorentz violation are studied in Ref. [11], while CPT violation without Lorentz violation in double beta decay is considered in Ref. [12].

First, consider precision experiments designed to detect neutrino mass directly by studying the end point of tritium beta decay. Recent experimental measurements using tritium have been performed in Troitsk [13–15] and Mainz [16], while the next-generation Karlsruhe Tritium Neutrino experiment (KATRIN) [17] is expected to begin taking data shortly. In these experiments, beta-decay electrons are guided by a magnetic field from the decay region to an electrostatic filter, where electrons with energies near the end point $T_0 \simeq 18.6$ keV are selected. The electrons guided to the filter emerge from the decay region within a solid angle $\Delta\Omega$ measured about the z axis along the magnetic field, which we take as horizontal, with the acceptance cone having angular aperture θ_0. For this configuration, the orientation factors $\mathcal{N}_{jm}$ in Eq. (1) become

$$\mathcal{N}_{jm} = \sum_{m'm''} Y_{jm'}(\theta,\phi) d^{(j)}_{m'm''}(-\pi/2) e^{-im''\xi} d^{(j)}_{m''m}(-\chi), \quad (2)$$

where $Y_{jm'}(\theta,\phi)$ are spherical harmonics in the laboratory frame, $d^{(j)}_{m'm''}$ are the little Wigner matrices, ξ is the angle of the magnetic field at the source measured east of local north, and χ is the colatitude of the laboratory.

The selection of electrons lying within $\Delta\Omega$ introduces experimental sensitivity to direction-dependent effects arising from neutrino Lorentz violation and hence provides access in principle to all four coefficients $(a^{(3)}_{\rm of})_{jm}$. For $m = \pm 1$, the phase factor in Eq. (1) varies sinusoidally in the sidereal time $T_\oplus$, so time stamps for the data permit a search for sidereal variations and hence measurements of $(a^{(3)}_{\rm of})_{11}$ and $(a^{(3)}_{\rm of})_{1-1} \equiv -(a^{(3)}_{\rm of})^*_{11}$ or, equivalently, of ${\rm Re}(a^{(3)}_{\rm of})_{11}$ and ${\rm Im}(a^{(3)}_{\rm of})_{11}$. For $m = 0$ there is no time dependence, but the Lorentz violation modifies the shape of the differential beta spectrum and so a study of the time-averaged spectral shape instead can enable measurements of $(a^{(3)}_{\rm of})_{00}$ and $(a^{(3)}_{\rm of})_{10}$.

In this work, we use published results from the Troitsk and Mainz experiments to place conservative constraints on the coefficients $(a^{(3)}_{\rm of})_{j0}$ and others, and we estimate sensitivities attainable in principle from the unpublished raw data in these experiments and in KATRIN. In practice, data for the tritium end-point spectrum are available only over a small energy range $\Delta T_c = T_0 - T_c$ from some cutoff energy T_c to T_0, so the Lorentz-violating modifications to the spectral shape are only partly observable. For this energy range, the decay rate takes the form

$$\frac{d\Gamma}{dT} \approx B + C\Big[(\Delta T + k(T_\oplus))^2 - \frac{1}{2}m_\nu^2\Big], \quad (3)$$

where B is the experimental background rate, C is approximately constant, and $\Delta T = T_0 - T$. The function $k(T_\oplus)$ contains the SME coefficients,

$$\begin{aligned} k(T_\oplus) = {} & \frac{1}{\sqrt{4\pi}}(a^{(3)}_{\rm of})_{00} - \sqrt{\frac{3}{4\pi}}\cos^2\frac{1}{2}\theta_0 \sin\chi\cos\xi (a^{(3)}_{\rm of})_{10} \\ & - \sqrt{\frac{3}{2\pi}}\cos^2\frac{1}{2}\theta_0\Big[\sin\xi\,{\rm Im}((a^{(3)}_{\rm of})_{11}e^{i\omega_\oplus T_\oplus}) \\ & + \cos\chi\cos\xi\,{\rm Re}((a^{(3)}_{\rm of})^*_{11}e^{-i\omega_\oplus T_\oplus})\Big]. \end{aligned} \quad (4)$$

The result (3) reveals that $k(T_\oplus)$ acts to shift the tritium end-point spectrum along the energy axis without changing its shape, independent of the value of m_ν^2. Sensitivity to this effect therefore requires experimental access to absolute energy measurements. The shift can be positive or negative and depends in part on the location and orientation of the experiment and on the sidereal time. The coefficients ${\rm Re}(a^{(3)}_{\rm of})_{11}$ and ${\rm Im}(a^{(3)}_{\rm of})_{11}$ induce harmonic oscillations of the spectrum along the energy axis at frequency $\omega_\oplus$, while $(a^{(3)}_{\rm of})_{00}$ and $(a^{(3)}_{\rm of})_{10}$ shift the location of the end point relative to the usual case. For data collected over a long period the harmonic oscillations average away, and so only the coefficients $(a^{(3)}_{\rm of})_{00}$ and $(a^{(3)}_{\rm of})_{10}$ produce observable effects. For simplicity in what follows, we take only one nonzero coefficient at a time, noting that fortuitous cancellations cannot simultaneously occur in experiments with different values of χ, ξ, θ_0.

Conservative constraints on the coefficients $(a^{(3)}_{\rm of})_{j0}$ can be placed using published results and the time-averaged form of Eq. (3). Consider first the Troitsk experiment [13–15]. The experiment is located at a colatitude $\chi \simeq 35°$, has decay-pipe axis pointing barely west of north with $\xi \simeq -5°$, and has acceptance-cone aperture $\theta_0 \simeq 20°$ [18]. The averaged end-point energy measured in this experiment is 18576 eV, which cannot be taken as an absolute energy measurement [15] but lies within about 2 eV of the expected end point. Taking ± 5 eV as the upper limit of a possible constant shift yields the constraints $|(a^{(3)}_{\rm of})_{j0}| < 2 \times 10^{-8}$ GeV for both $j = 0$ and $j = 1$.

The Mainz experiment reports a series of 12 measurements of the end-point energy under different experimental conditions, denoted Q1–Q12 [16]. The apparatus was located at colatitude $\chi \simeq 40°$ and the axis of the decay pipe had orientation $\xi \simeq -65°$ relative to local north [19]. The theoretical maximum value for the electron kinetic energy for this experiment is $T_0 = 18574.3 \pm 1.7$ eV [16]. For definiteness consider run Q12, which involves a wider acceptance-cone aperture $\theta_0 \simeq 62°$ and has measured end point 18576.6 ± 0.2 eV. Interpreting the difference between the measured and theoretical end points via Eq. (3) gives $(a^{(3)}_{\rm of})_{00} = 8.2 \pm 6.1 \times 10^{-9}$ GeV and $(a^{(3)}_{\rm of})_{10} = -2.4 \pm 1.8 \times 10^{-8}$ GeV. Constraints at similar levels can be obtained from the other runs. More conservatively, we can infer that a constant shift of ± 5 eV would be observable, which for this experiment leads to the constraints $|(a^{(3)}_{\rm of})_{00}| < 2 \times 10^{-8}$ GeV and $|(a^{(3)}_{\rm of})_{10}| < 5 \times 10^{-8}$ GeV.

Taken together, the above results permit us safely to conclude that $|(a^{(3)}_{\rm of})_{j0}| \lesssim 3 \times 10^{-8}$ GeV for both $j = 0$ and $j = 1$ at better than a 90% confidence level. Despite their conservative nature, these constraints significantly improve on existing limits [20], which have been extracted from studies of threshold effects. Threshold effects have been investigated only in the purely isotropic model, for which $\mathring{a}^{(3)} \equiv (a^{(3)}_{\rm of})_{00}/\sqrt{4\pi}$ is the sole nonzero coefficient for $d = 3$. The best existing constraint $|\mathring{a}^{(3)}| < 1.9 \times 10^{-7}$ GeV is obtained using IceCube data [5], so the result for $(a^{(3)}_{\rm of})_{00}$ presented here represents an improvement of more than an order of magnitude. Our constraint on $(a^{(3)}_{\rm of})_{10}$ is the first in the literature.

Dedicated analyses of the raw data by the Troitsk, Mainz, and KATRIN collaborations could improve on these constraints. For the Troitsk and Mainz experiments, sensitivities of $|(a^{(3)}_{\rm of})_{j0}| \lesssim 10^{-9}$ GeV or better appear attainable by focusing attention on absolute energies. For KATRIN [17], the apparatus is located at colatitude $\chi \simeq 41°$, has orientation relative to local north of $\xi \simeq 16°$, and has acceptance-cone aperture $\theta_0 \simeq 51°$. With 30 days of data, statistical confidence levels suggest a reach about 2 orders of magnitude beyond the new constraints reported above. In all these experiments, first measurements of ${\rm Re}(a^{(3)}_{\rm of})_{11}$ and ${\rm Im}(a^{(3)}_{\rm of})_{11}$ could be achieved by binning in sidereal time and fitting to Eq. (3).

We remark in passing that tritium end-point experiments also have sensitivity to the full spectrum of neutrino SME coefficients. Competitive constraints arise for effects that dominate at low energies and hence for operators of low d. For example, the $d = 2$ coefficients $(c^{(2)}_{\rm eff})^{ab}_{1m}$, where $a, b = e, \mu, \tau$ are flavor indices and $m = 0, \pm 1$, control helicity-flipping CPT-even Lorentz violation [5]. Calculation reveals that linear combinations $c^{(2)}_m$ of these coefficients act to shift the squared mass in tritium beta decay, so the end-point spectrum is governed by an effective squared mass $m^2_{\rm eff} = m^2_\nu + k_m c^{(2)}_m$, where k_m depends on χ, ξ, θ_0, and also on $T_\oplus$ for $m = \pm 1$ [21]. For instance, for the Troitsk experiment $k_0 \simeq 0.5$, and the reported measurement $m^2_{\rm eff} = -0.67 \pm 2.53$ eV^2 translates into the bound $c^{(2)}_0 < 4 \times 10^{-18}$ GeV^2, independent of m^2_ν. Using the standard mixing parameters [2] $\theta_{12} = 0.59$, $\theta_{13} = 0.16$, $\theta_{23} = 0.79$, assuming zero conventional CP phase δ, and taking only one coefficient $(c^{(2)}_{\rm eff})^{ab}_{10}$ at a time, this bound translates into six constraints in units of 10^{-17} GeV^2: $-1 < (c^{(2)}_{\rm eff})^{ee}_{10}$, $-2 < (c^{(2)}_{\rm eff})^{\mu\mu}_{10}$, $-2 < (c^{(2)}_{\rm eff})^{\tau\tau}_{10}$, ${\rm Re}(c^{(2)}_{\rm eff})^{e\mu}_{10} < 1$, $-3 < {\rm Re}(c^{(2)}_{\rm eff})^{e\tau}_{10}$ and ${\rm Re}(c^{(2)}_{\rm eff})^{\mu\tau}_{10} < 1$. Similarly, for the Mainz experiment $k_0 \simeq 0.2$, and the reported combined limit $m^2_\nu = -0.6 \pm 3.0$ eV^2 yields the constraint $c^{(2)}_0 < 1 \times 10^{-17}$ GeV^2, which gives the six constraints $-2 < (c^{(2)}_{\rm eff})^{ee}_{10}$, $-4 < (c^{(2)}_{\rm eff})^{\mu\mu}_{10}$, $-5 < (c^{(2)}_{\rm eff})^{\tau\tau}_{10}$, ${\rm Re}(c^{(2)}_{\rm eff})^{e\mu}_{10} < 3$, $-8 < {\rm Re}(c^{(2)}_{\rm eff})^{e\tau}_{10}$ and ${\rm Re}(c^{(2)}_{\rm eff})^{\mu\tau}_{10} < 3$, all in units of 10^{-17} GeV^2. These are first results for $(c^{(2)}_{\rm eff})^{ab}_{10}$ for all flavors ab except $e\mu$. For KATRIN $k_0 \simeq 0.5$, and a sensitivity of 0.04 eV^2 corresponds to a competitive reach of $c^{(2)}_0 < 8 \times 10^{-20}$ GeV^2, offering an improvement of about 2 orders of magnitude. Sidereal analyses for these experiments could yield m^2_ν-independent two-sided constraints on the coefficients $(c^{(2)}_{\rm eff})^{ab}_{1\pm1}$.

Returning to countershading studies, neutron decay [22] offers another interesting experimental option. Precision experiments investigating the beta spectrum from neutron decay typically have lesser sensitivity but could detect the full spectral distortion instead of only the end-point shift. This includes energy regions where the coefficients $(a^{(3)}_{\rm of})_{jm}$ induce maximal deviation from the conventional spectrum, which can lie comparatively far from the end point $T_0 \simeq 0.78$ MeV [23].

Consider, for example, experiments that measure only T. The differential beta spectrum $d\Gamma/dT$ must then be constructed by integrating over the lepton directions of travel, so only the isotropic coefficient $\mathring{a}^{(3)}$ can play a role. We find $d\Gamma/dT \propto F(Z,T)|\boldsymbol{p}|(T + m_e)(\Delta T + \mathring{a}^{(3)})^2$, where $F(Z,T)$ is the Fermi function. The conventional spectrum peaks at $T \simeq 0.25$ MeV, while the residual Lorentz-violating spectrum has a maximum at $T_m \simeq 0.41$ MeV. The ratio R of the residual to the conventional spectra at T_m is $R \simeq 5 \times 10^3 \mathring{a}^{(3)}$/GeV. An experiment with a plausible sensitivity of 0.1% in this energy region would thus have an estimated reach of $|\mathring{a}^{(3)}| < 2 \times 10^{-7}$ GeV, comparable to the constraint from IceCube [5].

Some neutron-decay experiments are designed to measure the antineutrino-electron correlation parameter a associated with the angle between the phase velocities of the two emitted leptons. Denoting by N_+ the number of parallel leptons at energy T and by N_- the number of antiparallel ones, a standard observable is the asymmetry $a_{\rm exp} = (N_+ - N_-)/(N_+ + N_-)$, which in the absence of Lorentz violation is $a_{\rm exp} = a|\boldsymbol{p}|/E$. The correction involves all $j = 1$ coefficients $(a^{(3)}_{\rm of})_{jm}$, including a signal from sidereal variations associated with $(a^{(3)}_{\rm of})_{11}$. Assuming a plausibly attainable 0.1% measurement of $a_{\rm exp}$ near 0.35 MeV for an experiment located at $\chi = 45°$ and taking $a = -0.103$ [2], we find estimated reaches of $|(a^{(3)}_{\rm of})_{10}| < 5 \times 10^{-8}$ GeV and $|{\rm Re}(a^{(3)}_{\rm of})_{11}|$, $|{\rm Im}(a^{(3)}_{\rm of})_{11}| < 4 \times 10^{-8}$ GeV.

Another class of experiments seeks to measure the correlation parameter B controlling the angle between the neutron spin and the antineutrino phase velocity. This requires sensitivity to the neutron polarization and reconstruction of the antineutrino emission. A standard observable is $B_{\rm exp} = (N_{--} - N_{++})/(N_{--} + N_{++})$, where N_{--} and N_{++} count events with both the electron and the proton emitted against and along the direction of the neutron spin, respectively. The Lorentz-violating correction to the asymmetry involves all $j = 1$ coefficients $(a_{\rm of}^{(3)})_{jm}$, along with a dependence on a, B, and the spin-electron correlation parameter A. For a plausible 0.1% measurement of $B_{\rm exp}$ near 0.35 MeV at $\chi = 45°$ and taking $a = -0.103$, $B = 0.9807$, $A = -0.1176$ [2], we obtain estimated attainable sensitivities $|(a_{\rm of}^{(3)})_{10}| < 2 \times 10^{-6}$ GeV and $|{\rm Re}(a_{\rm of}^{(3)})_{11}|$, $|{\rm Im}(a_{\rm of}^{(3)})_{11}| < 1 \times 10^{-6}$ GeV.

Single beta decays of other nuclei and particles such as pions [24], kaons, or muons are also worth studying and are likely to yield similar limits on $(a_{\rm of}^{(3)})_{jm}$. A qualitatively different option is provided by double beta decay. This is a second-order weak process, so a lesser sensitivity is generically to be expected. However, precision experiments searching for the neutrinoless mode typically also generate a large sample of two-antineutrino events. The corresponding statistical reach can be significant, and all four $(a_{\rm of}^{(3)})_{jm}$ appear in angular correlations if the electron directions can be reconstructed.

For the two-antineutrino mode, we limit attention here to an analysis using only the differential spectrum $d\Gamma/dK$ of the summed energies $K = T_1 + T_2$ of the two emitted electrons. This involves integration over the neutrino directions, which implies only isotropic effects are observable and hence that the residual spectrum depends only on $\mathring{a}^{(3)}$. Defining $\Delta K = K_0 - K$ where K_0 is the maximum kinetic energy available in the decay, the sum electron spectrum $d\Gamma/dK$ is modified by the factor $\Delta K^5 \to (\Delta K + 2\mathring{a}^{(3)})^5$, revealing a distortion of the spectral shape. Consider ^{136}Xe, for example, for which $K_0 \simeq 2.46$ MeV and the maximum of the conventional sum electron spectrum occurs at energy $K \simeq 0.86$ MeV. The residual spectrum has maximum at $K_m \simeq 1.02$ MeV, so the ratio of the residual and conventional spectra at K_m is $R \simeq 7 \times 10^3 \mathring{a}^{(3)}/$GeV. An experiment with a plausible precision of 0.1% near K_m would hence have an estimated reach of $|\mathring{a}^{(3)}| < 2 \times 10^{-7}$ GeV.

In the neutrinoless mode of double beta decay, the neutrino is virtual and must have Majorana couplings to generate a nonzero amplitude. The coefficients $(a_{\rm of}^{(3)})_{jm}$ are Dirac couplings, so their contribution to the amplitude must be suppressed by some other Majorana coupling and competitive sensitivities are unlikely. However, numerous Lorentz-violating Majorana operators exist [5]. Here, we outline a few implications. The usual half life is $T_{1/2} = (G|M|^2 m_\nu^2)^{-1}$, where G is a known function of the nuclear radius R and other quantities, and M is the nuclear matrix element. Calculation shows leading-order effects can arise only from CPT-odd Majorana operators controlled by the SME coefficients $\hat{g}_{M+}^{\alpha\beta}$ [5]. Suppressing the orientation dependence and denoting the dimensionless effective coefficient by g, we find the half-life corrections include the replacement $m_\nu^2 \to m_\nu^2 + m_\nu g/R + (g/R)^2$. This reveals the striking feature that Lorentz-violating neutrinoless double beta decay can occur even for negligible m_ν [25], with the role of the Majorana mass played by g. Experiments placing an upper bound on m_ν^2 can therefore also report a constraint on g. For example, the current limit on the neutrinoless double beta decay of ^{136}Xe [26] corresponds to the constraint $|g| \lesssim 10^{-9}$. Note that experiments with different locations and orientations generically have different sensitivities due to the directional and sidereal dependences, as do experiments with different isotopes due to the R dependence [27].

The results reported here demonstrate that studies of beta decay can achieve interesting sensitivities to countershaded neutrino Lorentz violation. For gravitational Lorentz violation, countershading has recently been eliminated to the keV scale [20]. Here, the constraints obtained lie at the eV scale, improving by more than an order of magnitude the existing bound on one effect and setting a first bound on another. Models involving these neutrino operators at scales comparable to m_e are now excluded, while the new constraints must be satisfied even by models with effects at neutrino-oscillation scales.

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[1] W. Pauli, *Offener Brief an die Gruppe der Radioaktiven bei der Gauvereins-Tagung zu Tübingen* (1930).
[2] J. Beringer *et al.*, Phys. Rev. D **86**, 010001 (2012).
[3] V. A. Kostelecký and S. Samuel, Phys. Rev. D **39**, 683 (1989); V. A. Kostelecký and R. Potting, Nucl. Phys. **B359**, 545 (1991); Phys. Rev. D **51**, 3923 (1995).
[4] D. Colladay and V. A. Kostelecký, Phys. Rev. D **55**, 6760 (1997); **58**, 116002 (1998); V. A. Kostelecký, Phys. Rev. D **69**, 105009 (2004).
[5] V. A. Kostelecký and M. Mewes, Phys. Rev. D **85**, 096005 (2012).
[6] L. B. Auerbach *et al.*, Phys. Rev. D **72**, 076004 (2005); P. Adamson *et al.*, Phys. Rev. Lett. **101**, 151601 (2008); **105**,

151601 (2010); **85**, 031101 (2012); Phys. Rev. D **82**, 112003 (2010); A. A. Aguilar-Arevalo *et al.*, Phys. Lett. B **718**, 1303 (2013); T. Katori, Mod. Phys. Lett. A **27**, 1230024 (2012); Y. Abe *et al.*, Phys. Rev. D **86**, 112009 (2012); B. Rebel and S. Mufson, Astropart. Phys. **48**, 78 (2013); V. A. Kostelecký and M. Mewes, Phys. Rev. D **70**, 076002 (2004); J. S. Díaz, V. A. Kostelecký, and M. Mewes, Phys. Rev. D **80**, 076007 (2009); arXiv:1308.6344; J. S. Díaz, T. Katori, J. Spitz, and J. M. Conrad, arXiv:1307.5789.

[7] V. A. Kostelecký and J. D. Tasson, Phys. Rev. Lett. **102**, 010402 (2009).

[8] R. Bluhm, V. A. Kostelecký, C. Lane, and N. Russell, Phys. Rev. D **68**, 125008 (2003); R. Bluhm, V. A. Kostelecký, C. Lane, and N. Russell, Phys. Rev. Lett. **88**, 090801 (2002); V. A. Kostelecký and M. Mewes, Phys. Rev. D **66**, 056005 (2002).

[9] V. A. Kostelecký and M. Mewes, Phys. Rev. D **69**, 016005 (2004).

[10] J. M. Carmona and J. L. Cortés, Phys. Lett. B **494**, 75 (2000); E. Di Grezia, S. Esposito, and G. Salesi, Mod. Phys. Lett. A **21**, 349 (2006); A. E. Bernardini and O. Bertolami, Phys. Rev. D **77**, 085032 (2008).

[11] J. P. Noordmans, H. W. Wilschut, and R. G. E. Timmermans, Phys. Rev. C **87**, 055502 (2013); H. W. Wilschut *et al.*, arXiv:1303.6419.

[12] G. Barenboim, J. F. Beacom, L. Borissov, and B. Kayser, Phys. Lett. B **537**, 227 (2002).

[13] V. M. Lobashev *et al.*, Phys. Lett. B **460**, 227 (1999).

[14] V. N. Aseev *et al.*, Phys. Rev. D **84**, 112003 (2011).

[15] V. N. Aseev *et al.*, Phys. At. Nucl. **75**, 464 (2012).

[16] C. Kraus *et al.*, Eur. Phys. J. C **40**, 447 (2005).

[17] J. Angrik *et al.*, Report No. FZKA-7090, 2005.

[18] V. S. Pantuev (private communication).

[19] C. Weinheimer (private communication).

[20] V. A. Kostelecký and N. Russell, Rev. Mod. Phys. **83**, 11 (2011) [2013 edition: arXiv:0801.0287v6].

[21] Notice $c_m^{(2)}$ can generate seemingly unphysical values $m_{\rm eff}^2 < 0$ like those for tachyonic neutrinos discussed in A. Chodos, A. I. Hauser, and V. A. Kostelecký, Phys. Lett. **150B**, 431 (1985).

[22] A recent review is J. S. Nico and W. M. Snow, Annu. Rev. Nucl. Part. Sci. **55**, 27 (2005).

[23] The spectral distortion affects the neutron lifetime, which can vary with sidereal time, neutron polarization, and momentum. For large $(a_{\rm of}^{(3)})_{jm}$, this could conceivably play a role in the neutron lifetime puzzle discussed, e.g., in S. Paul, Nucl. Instrum. Methods Phys. Res., Sect. A **611**, 157 (2009).

[24] See, e.g., the recent discussion for pions in B. Altschul, Phys. Rev. D **87**, 096004 (2013).

[25] The possibility of a neutrinoless mode for double beta decay with $m_\nu \to 0$ was originally noted in the context of superweak interactions by G. Feinberg and M. Goldhaber, Proc. Natl. Acad. Sci. U.S.A. **45**, 1301 (1959); B. Pontecorvo, Phys. Lett. **26B**, 630 (1968).

[26] M. Auger *et al.*, Phys. Rev. Lett. **109**, 032505 (2012).

[27] These features could provide a possible path for reconciling the null searches for neutrinoless double beta decay with the signal from ^{76}Ge claimed in H. V. Klapdor-Kleingrothaus and I. V. Krivosheina, Mod. Phys. Lett. A **21**, 1547 (2006).