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Constraining *CPT*-odd nonminimal interactions in the electroweak sector

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In this work, we propose two possibilities of *CPT*-odd and Lorentz-violating (LV) nonminimal couplings in the electroweak sector. These terms are gauge-invariant and couple a fixed 4-vector to the physical fields of the theory. After determining the LV contributions to the electroweak currents, we reassess the evaluation of the decay rate for the vector mediators W and Z. Using the experimental uncertainty in these decay rates, upper bounds of 1 part in 10^{-6} $(\mathrm{GeV})^{-1}$ and 10^{-5} $(\mathrm{GeV})^{-1}$ are imposed on the magnitude of the proposed nonmimal interactions.

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I. INTRODUCTION

Mechanisms of spontaneous Lorentz violation have been proposed in some candidate theories of quantum gravity. As a consequence, Lorentz-violating (LV) background tensors (generated as vacuum expectation values) are coupled to the physical fields of the standard model (SM). The most general effective theory considering the explicit breaking of Lorentz and *CPT* symmetry is the minimal standard model extension (mSME) [1], which is an extension of the $SU(3)\times SU(2)\times U(1)$ standard model, featuring terms breaking Lorentz and *CPT* symmetries in all of its sectors: lepton, quark, Yukawa, Higgs, and gauge. Investigation of Lorentz symmetry violation is a rich line of research, embracing developments in the electromagnetic sector [2,3], fermion sector [4], including photon-fermion interactions [5–7] and quantization techniques in the photon sector [8]. Such studies have scrutinized LV effects in very distinct physical systems, allowing to construct a precision program to determine to what extent the Lorentz covariance is preserved in nature (by means of tight upper bounds on the LV coefficients). Nonminimal LV interactions have been examined in an extension of the SME encompassing higher derivatives in both the gauge [9] and the fermion sector [10]. Some models containing higher-dimension operators [11,12] have also been proposed.

In the electroweak sector of the mSME [1], the SU(2) and U(1) gauge fields are properly coupled to LV fixed tensors in renormalizable dimension four terms. The mSME lepton sector is composed of a *CPT*-even and a *CPT*-odd term, that is,

$$\mathcal{L}_{\mathrm{lep}}^{\mathrm{even}}=(c_L)_{\mu\nu AB}\bar{L}_A\gamma^{\mu}iD^{\nu}L_B+(c_R)_{\mu\nu AB}\bar{R}_A\gamma^{\mu}iD^{\nu}R_B, \tag{1}$$

$$\mathcal{L}_{\mathrm{lepton}}^{\mathrm{odd}}=-(a_L)_{\mu AB}\bar{L}_A\gamma^{\mu}L_B-(a_R)_{\mu AB}\bar{R}_A\gamma^{\mu}R_B, \tag{2}$$

where $A,B=1,\,2,\,3$ are the lepton flavor labels. In the same way, the SU(2) and U(1) gauge sectors are modified by the *CPT*-even terms:

$$\mathcal{L}_{\mathrm{gauge}}^{\mathrm{even}}=-\frac{1}{2}(k_W)^{\mu\nu\alpha\beta}(W_{\mu\nu}^{a}W_{\alpha\beta}^{a})-\frac{1}{4}(k_B)_{\mu\nu\alpha\beta}B^{\mu\nu}B^{\alpha\beta}, \tag{3}$$

while the *CPT*-odd piece generates instabilities in the theory and are not considered. The k_W, k_B coefficients are real, dimensionless, and possess the same symmetries of the Riemann tensor. The pure *CPT*-even Higgs sector is also modified by the following term:

$$\mathcal{L}_{Higgs}^{\mathrm{even}}=\frac{1}{2}\Big[(k_{\phi\phi})_{\mu\nu}(D^{\mu}\phi^{a})^{\dagger}(D^{\nu}\phi^{a})+\mathrm{H.c.}-(k_{\phi B})_{\mu\nu}\phi^{a\dagger}\phi^{a}B^{\mu\nu}-(k_{\phi W})^{\mu\nu}(\phi^{\dagger}\times\phi)^{a}W_{\mu\nu}^{a}\Big], \tag{4}$$

while the Higgs *CPT*-odd has the form $i(k_{\phi})_{\mu}(\phi^{a})^{\dagger}(D^{\mu}\phi^{a})$, with $(k_{\phi})_{\mu}$ having dimension of mass.

LV studies in the electroweak sector were initially developed in connection with meson decays $(\pi^{-}\rightarrow\mu^{-}+\bar{\nu}_{\mu})$, where the LV effects were considered at the level of the Feynman propagator of the W boson [13], $\langle W^{\mu\dagger}W^{\nu}\rangle=-i(g^{\mu\nu}+\chi^{\mu\nu})/M_W^2$, with contributions coming from the Higgs (ϕ) and the W sectors: $\chi^{\mu\nu}=k_{\phi\phi}^{\mu\nu}-\frac{i}{2g}k_{\phi W}^{\mu\nu}+k_W^{\alpha\mu\beta\nu}p_{\alpha}p_{\beta}$. Comparison with experimental data led to upper bounds of 1 part in 10^4. Contributions of the k_W coefficients, presented in Eq. (3), to the W propagator, $\langle W^{\mu\dagger}W^{\nu}\rangle=-i(g^{\mu\nu}+\chi^{\mu\nu})/M_W^2$, jointly with contributions stemming from the Higgs sector, $k_{\phi\phi},k_{\phi W}$, see Eq. (4), were more explicitly considered in Ref. [14], with implications on the allowed nuclear decays and forbidden β decays. This framework was also used: (i) to reinterpret experiments dedicated to searching for preferred directions in forbidden β-decays, implying upper bounds as tight as 10^{-8} on the LV parameters [15]; (ii) to constrain β decay rate asymmetries to the level of 1 part in 10^6 [16]; (iii) to study isotopes that undergo orbital electron capture [17]; (iv) to analyze LV effects on the kaon decay and evaluate asymmetries in the respective lifetime [18]. Another interesting study considered LV coefficients $(c_L)_{\mu\nu AB}$ of the lepton sector (1), with the same flavor $(A=B)$,

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$$\mathcal{L}_{\text{lepton}} = c_{\alpha\beta}[i\bar{\psi}\gamma^{\alpha}\partial^{\beta}\psi + i\bar{\psi}_{\nu}\gamma^{\alpha}\partial^{\beta}\psi_{\nu} + \bar{\psi}_{(L)}\gamma^{\alpha}W^{\beta(-)}\psi_{\nu(L)} + \bar{\psi}_{\nu(L)}\gamma^{\alpha}W^{\beta(+)}\psi_{(L)}], \tag{5}$$

where $\psi, \psi_{\nu}, \psi_{(L)}$ represent leptons, neutrinos and left-handed leptons (of a given flavor), to examine direct effects on the pion-decay rate [19], attaining upper bounds of the level of 10^{-4}. Some works also examined the possibility of LV electroweak terms to make feasible forbidden processes ($Z_0 \to \gamma + \gamma$) [20] or modify reactions such as $\gamma + e \to W + \nu_e$, $\gamma + \gamma \to W + W$ [21]. Lepton flavor violating decays triggered by renormalizable and non-renormalizable (dimension five) terms belonging to the Higgs sector were recently considered as well [22]. Tree-level Z-boson contributions to the polarized Möller scattering were carried out, allowing us to improve k_W upper bounds by two orders of magnitude [23]. Lorentz violation influence on neutrino oscillations was also probed using a distinct framework [24].

A dimension five LV nonminimal coupling (NMC), representing unusual interactions between fermions and photons, $g^{\nu}\bar{\psi}\gamma^{\mu}\tilde{F}_{\mu\nu}\psi$, was first introduced by means of the derivative, $D_{\mu} = \partial_{\mu} + ieA_{\mu} + i\frac{\lambda}{2}\epsilon_{\mu\lambda\alpha\beta}V^{\lambda}F^{\alpha\beta}$, in the Dirac equation [25], where V^{μ} can be identified with the Carroll-Field-Jackiw four-vector. Such a coupling has been addressed in numerous aspects [26], including the radiative generation of CPT-odd LV terms [27], topological phases [28], and generation of electric dipole moment [29]. Dimension-five CPT-even NMCs were also proposed in the context of the Dirac equation [30], with MDM and EDM experimental measurements being used to state upper bounds at the level of 1 part in $10^{20}\,(\text{eV})^{-1}$ and $10^{24}\,(\text{eV})^{-1}$, respectively. A systematic investigation on NMCs of dimension five and six was recently proposed in Ref. [31].

Nonminimal interactions have been a topical issue in the latest years, mainly in the fermion and electromagnetic sectors. However, a NMC in the lepton electroweak sector of the SM has not been proposed yet. In this work, we introduce two possibilities of CPT-odd LV nonminimal interactions in the electroweak sector, the first one being proposed in the $U(1)_Y$ sector of the GSW model, while the second is considered in its $SU(2)_L$ sector, both as extensions of the covariant derivative. Knowing the interaction Lagrangian, we evaluate the LV corrections to the decay rates of the following mediators: $Z_0 \to \bar{l} + l$ and $W^{-} \to l + \bar{\nu}_l$, attaining upper bounds as tight as $10^{-6}\,(\text{GeV})^{-1}$.

II. BASICS ABOUT THE GSW MODEL

In the Glashow-Salam-Weinberg electroweak model (GSW), with a $SU(2)_L \times U(1)_Y$ gauge structure spontaneously broken via the Higgs mechanism, the vector bosons, $W^{\pm}$, Z^0 and γ are mediator of the interactions, being introduced via minimal coupling to the matter fields. In this theory, left-handed leptons (L_l) are represented by isodoublets

$$L_l = \begin{bmatrix} \psi_{\nu_l} \\ \psi_l \end{bmatrix}_L = \frac{1-\gamma_5}{2}\begin{bmatrix} \psi_{\nu_l} \\ \psi_l \end{bmatrix}, \tag{6}$$

while right-handed leptons (R_l) are isosinglets,

$$R_l = (\psi_l)_R = \left(\frac{1+\gamma_5}{2}\right)\psi_l, \tag{7}$$

and $l = 1, 2, 3$ is the lepton flavor label: $\psi_l = (e, \mu, \tau)$. The part of the electroweak Lagrangian, in which the leptons interact directly with the gauge fields, is $\mathcal{L}_{\text{EW}} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{lepton}}$, where

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}\mathbf{W}_{\mu\nu}\cdot\mathbf{W}^{\mu\nu} - \frac{1}{4}B_{\mu\nu}B^{\mu\nu}, \tag{8}$$

$$\mathcal{L}_{\text{lepton}} = \bar{L}_l\gamma^{\mu}iD_{\mu}L_l + \bar{R}_l\gamma^{\mu}iD_{\mu}R_l, \tag{9}$$

with $\mathbf{W}_{\mu} = (W^1_{\mu}, W^2_{\mu}, W^3_{\mu})$ being a four-vector gauge field which is a three-vector in isospin space, and B_{μ} a gauge four-vector field, whose field strengths are $B_{\mu\nu} = \partial_{\mu}B_{\nu} - \partial_{\nu}B_{\mu}$ and

$$\mathbf{W}_{\mu\nu} = \partial_{\mu}\mathbf{W}_{\nu} - \partial_{\nu}\mathbf{W}_{\mu} + g\mathbf{W}_{\mu}\times\mathbf{W}_{\nu}. \tag{10}$$

The covariant derivative involves both gauge fields,

$$D_{\mu} = \partial_{\mu} - ig\mathbf{T}\cdot\mathbf{W}_{\mu} - i\frac{g'}{2}YB_{\mu}. \tag{11}$$

Here, $\mathbf{T} = (T_1, T_2, T_3)$ stands for the generators of the group $SU(2)_L$, and Y is the generator of $U(1)_Y$ group, fulfilling $[T_i, T_j] = i\varepsilon_{ijk}T_k$ and $[T_i, Y] = 0$. Furthermore, $Y = -1$ or $Y = -2$ for left-handed and right-handed leptons, respectively. The lepton Lagrangian (9) can be written as $\mathcal{L} = i\bar{L}_l\gamma^{\mu}\partial_{\mu}L_l + i\bar{R}_l\gamma^{\mu}\partial_{\mu}R_l + \mathcal{L}^{(l)}_{\text{int}}$, with the interaction part given as

$$\mathcal{L}^{(l)}_{\text{int}} = \tilde{g}(J^{(l)\alpha}_{+}W^{(+)}_{\alpha} + J^{(l)\alpha}_{-}W^{(-)}_{\alpha} + J^{(l)\alpha}_{0}Z_{\alpha}) - eJ^{(l)\alpha}_{\text{EM}}A_{\alpha}, \tag{12}$$

where $\tilde{g} = g/(2\sqrt{2})$ and there appear charged currents, $J^{(l)\alpha}_{+}, J^{(l)\alpha}_{-}$, a neutral current, $J^{(l)\alpha}_{0}$, and the electromagnetic current, $J^{(l)\alpha}_{\text{EM}}$, given as

$$J^{(l)\alpha}_{+} = 2\bar{L}_l\gamma^{\alpha}T_{+}L_l = \bar{\psi}_{\nu_l}\gamma^{\alpha}(1-\gamma_5)\psi_l, \tag{13}$$

$$J^{(l)\alpha}_{-} = 2\bar{L}_l\gamma^{\alpha}T_{-}L_l = \bar{\psi}_l\gamma^{\alpha}(1-\gamma_5)\psi_{\nu_l}, \tag{14}$$

$$J^{(l)\alpha}_{0} = (\sqrt{2}\cos\theta)^{-1}[\bar{\psi}_{\nu_l}\gamma^{\alpha}(1-\gamma_5)\psi_{\nu_l} - \bar{\psi}_l\gamma^{\alpha}(g'_V - g'_A\gamma_5)\psi_l], \tag{15}$$

$$J^{(l)\alpha}_{\text{EM}} = -\left[\bar{L}_l\gamma^{\alpha}\left(\frac{g'\cos\theta}{2} - g\sin\theta T_3\right)L_l + g'\cos\theta\bar{R}_l\gamma^{\alpha}R_l\right]. \tag{16}$$

Here, θ is the weak mixing angle, and g, g' are the coupling constants, and the vector-axial interaction is controlled by $g'_A = 1$, $g'_V = 1$–$4\sin^2\theta$. In the electroweak theory, the photon field (A_{μ}) and the neutral intermediate boson (Z_{μ})

are given by combinations of the fields W_μ^3 and B_μ, that is, $A_\mu = B_\mu \cos\theta + W_\mu^3 \sin\theta$, $Z_\mu = -B_\mu \sin\theta + W_\mu^3 \cos\theta$. The inverse relations are also well known, $B_\mu = \cos\theta A_\mu - \sin\theta Z_\mu$, $W_\mu^3 = \sin\theta A_\mu + \cos\theta Z_\mu$. The generators and isovector can be also written as $\mathbf{T}=(T_+, T_3, T_-)$, $\mathbf{W}_\alpha = (W_\alpha^{(+)}/\sqrt{2}, W_\mu^3, W_\alpha^{(-)}/\sqrt{2})$, where $T_\pm = \sigma_x/2 \pm i(\sigma_y/2)$, $T_3 = \sigma_z/2$, and $W_\mu^{(\pm)} = \frac{1}{\sqrt{2}}(W_\mu^1 \mp i W_\mu^2)$, and σ_x, σ_y, σ_z are the Pauli matrices.

III. A NONMINIMAL COUPLING IN THE U$(1)_Y$ SECTOR OF THE GSW MODEL

We have already mentioned how LV terms are inserted in the mSME electroweak sector. Another route to consider Lorentz violation involves higher dimensional, nonrenormalizable NM (nonminimal) operators. Gauge invariant NM interactions in the electroweak sector can be proposed in the context of the covariant derivative (11). A first possibility, in the $U(1)_Y$ sector of the GSW model, is the NM derivative

$$D_\mu = \partial_\mu - ig\mathbf{T}\cdot\mathbf{W}_\mu - i\frac{g'}{2}YB_\mu + ig_2' Y B_{\mu\nu} C^\nu, \tag{17}$$

where C^ν is a fixed 4-vector that establishes a preferred direction in spacetime and violates Lorentz symmetry. Replacing such a derivative in Lagrangian (9), the nonminimal coupling yields additional electromagnetic and neutral LV interactions,

$$\mathcal{L}_{LV(1)} = J^{(l)\nu}_{EM(LV)} A_\nu + J^{(l)\nu}_{0(LV)} Z_\nu, \tag{18}$$

given explicitly as

$$\begin{aligned} J^{(l)\nu}_{EM(LV)} = {} & \frac{g_2'}{2}\cos\theta[\bar\psi_{\nu_l}\gamma^\mu(1-\gamma_5)\psi_{\nu_l}]C^\nu\partial_\mu \\ & - \frac{g_2'}{2}\cos\theta[\bar\psi_{\nu_l}\gamma^\nu(1-\gamma_5)\psi_{\nu_l}]C^\mu\partial_\mu \\ & + g_2'\cos\theta[j_1^\mu C^\nu\partial_\mu] - g_2'\cos\theta[j_1^\nu C^\mu\partial_\mu], \end{aligned} \tag{19}$$

$$\begin{aligned} J^{(l)\nu}_{0(LV)} = {} & -\frac{g_2'}{2}\sin\theta[\bar\psi_{\nu_l}\gamma^\mu(1-\gamma_5)\psi_{\nu_l}]C^\nu\partial_\mu \\ & + \frac{g_2'}{2}\sin\theta[\bar\psi_{\nu_l}\gamma^\nu(1-\gamma_5)\psi_{\nu_l}]C^\mu\partial_\mu \\ & - g_2'\sin\theta[j_1^\mu C^\nu\partial_\mu] + g_2'\sin\theta[j_1^\nu C^\mu\partial_\mu], \end{aligned} \tag{20}$$

with $j_1^\mu(x) = \bar\psi_l(x)\gamma^\mu(3+\gamma_5)\psi_l(x)/2$. These expressions are useful to show the processes that are directly affected, at tree-level, by the nonminimal derivative (17). We now examine the effect of this nonminimal coupling on the decay of the Z_0 mediator in a pair lepton and antilepton, $Z_0 \to \bar l + l$, evaluating the contributions implied to the decay rate. The total neutral current, $(J_0^{(l)\mu} + J_{0(LV)}^{(l)\mu})Z_\mu$, that contributes for this process is

$$\begin{aligned} & = -\frac{g}{4\cos\theta}\bar\psi_l(x)\gamma^\mu(g_V' - \gamma_5)\psi_l(x)Z_\mu(x) \\ & \quad - g_2'\sin\theta[j_1^\eta C^\mu\partial_\eta Z_\mu(x)] + g_2'\sin\theta[j_1^\mu C^\lambda\partial_\lambda Z_\mu(x)], \end{aligned} \tag{21}$$

where the first term is the usual Lorentz invariant contribution, the second and third terms stem from Eq. (20). Expression (21) shows how the NMC (17) affects the vertex of the neutral interaction. We point out that the LV terms regarded in the nonminimal coupling (17) do not modify the pure bilinear fermion Lagrangian, as it happens in the case of the fermion sector of the SME, for which a cross section evaluation is discussed in Ref. [32], taking into account some modifications amounted to the Feynman propagator, flux factor, as far as suitable spinor redefinitions. In the present NMC, however, the basic alteration is not in the pure bilinear sector, but in the vertex of the neutral interaction, so that the proceedings of Ref. [32] cannot be directly applied here.

The scattering matrix for such a process is

$$S = -i\int d^4x (J_0^{(l)\mu} + J_{0(LV)}^{(l)\mu})Z_\mu = S_0 + S_{LV(1)} + S_{LV(2)}, \tag{22}$$

where the zero order and first order contributions in the LV parameters are

$$S_0 = i\frac{g}{4\cos\theta}\int d^4x \bar\psi_l(x)\gamma^\mu(g_V' - \gamma_5)\psi_l(x)Z_\mu(x), \tag{23}$$

$$S_{LV(1)} = ig_2'\sin\theta\int d^4x[j_1^\eta(x)C^\mu\partial_\eta Z_\mu(x)], \tag{24}$$

$$S_{LV(2)} = -ig_2'\sin\theta\int d^4x[j_1^\mu(x)C^\lambda\partial_\lambda Z_\mu(x)]. \tag{25}$$

In order to evaluate these elements, we propose plane wave expansions, $Z_\mu^0(x) = N_k\varepsilon_\mu(k,\lambda)\exp(-ik\cdot x)$, $\psi_l(x) = N_q u_l(q,s)\exp(-iq\cdot x)$, $\psi_{\bar l}(x) = N_{q'}v(q',s')\exp(iq'\cdot x)$, where k, q, q' stand for the 4-momentum of the Z^0 boson and the emerging leptons, respectively, and $N_q = (2Vq_0)^{-1/2}$. With these expressions, we obtain

$$S_0 = i\frac{g}{4\cos\theta}(2\pi)^4\frac{\delta^4(q+q'-k)}{[8V^3q_0q_0'k_0]^{1/2}}M_0, \tag{26}$$

$$S_{LV(a)} = \frac{ig_2'\sin\theta}{2}(2\pi)^4\frac{\delta^4(q+q'-k)}{[8V^3q_0q_0'k_0]^{1/2}}M_{LV(a)}, \tag{27}$$

with $a = 1, 2$ representing the two LV contributions, which involves

$$M_0 = \varepsilon_\mu(k,\lambda)\bar u_l(q,s)\gamma^\mu(g_V' - \gamma_5)v(q',s'), \tag{28}$$

$$M_{LV(1)} = C^\mu k_\eta \varepsilon_\mu(k,\lambda) j^\eta_{qq'}, \tag{29}$$

$$M_{LV(2)} = -C^\lambda k_\lambda \varepsilon_\mu(k,\lambda) j^\mu_{qq'}, \tag{30}$$

and $j^\mu_{qq'} = \bar u_l(q,s)\gamma^\mu(3+\gamma_5)v(q',s')$. The decay rate for the process $Z_0 \to \bar l + l$ is given as usually evaluated, that is,

$$\Gamma_{ll} = \frac{1}{T} V \int \frac{d^3 q}{(2\pi)^3} V \int \frac{d^3 q'}{(2\pi)^3} \frac{1}{3} \sum_{\lambda} \sum_{s,s'} |S|^2, \quad (31)$$

where S is given in (22), implying

$$|S|^2 = S_0 S_0^\dagger + S_0 S_{LV(1)}^\dagger + S_{LV(1)} S_0^\dagger + S_0 S_{LV(2)}^\dagger + S_{LV(2)} S_0^\dagger, \quad (32)$$

in first order in the LV parameters. Substituting Eq. (32) in Eq. (31), we achieve

$$\Gamma_{ll} = \Gamma_{S_0 S_0^\dagger} + \Gamma_{S_0 S_{LV(1)}^\dagger} + \Gamma_{S_{LV(1)} S_0^\dagger} + \Gamma_{S_0 S_{LV(2)}^\dagger} + \Gamma_{S_{LV(2)} S_0^\dagger}. \quad (33)$$

The first term, $\Gamma_{S_0 S_0^\dagger}$, is the decay rate for the Lorentz invariant usual process $Z_0 \to \bar{l} + l$. In this evaluation, $\Gamma_{S_0 S_{LV(1)}^\dagger} = 0$, $\Gamma_{S_{LV(1)} S_0^\dagger} = 0$, as a consequence of the current conservation, due to the presence of the momentum k_α in Eqs. (29), (30). The LV contribution is associated with $\Gamma_{S_0 S_{LV(2)}^\dagger}$ and $\Gamma_{S_{LV(2)} S_0^\dagger}$, so that the total decay rate, $\Gamma = \Gamma_{S_0 S_0^\dagger} + \Gamma_{S_0 S_{LV(2)}^\dagger} + \Gamma_{S_{LV(2)} S_0^\dagger}$, is

$$\begin{aligned}\Gamma_{ll} = \frac{g^2 (8 M_Z)}{1536\pi \cos^2\theta} \Bigg\{ & \left[(g_V'^2 + 1) - 6 g_V'^2 \frac{m_l^2}{M_Z^2} \right] \\ & - \frac{g_2' \sin 2\theta}{g} (C \cdot k) \left[(3 g_V' - 2) - 27 g_V' \frac{m_l^2}{M_Z^2} \right] \Bigg\} \\ & \times \Theta(M_Z - 2 m_l). \end{aligned} \quad (34)$$

We now use $k^2 = M_Z^2$ and $C \cdot k = C_0 M_Z$. As the Z_0 mass ($M_Z = 9.1 \times 10^{10}$ eV) is much larger than lepton masses, we can neglect the mass ratios for the electron, muon and tau ($m_e^2/M_Z^2 \simeq 2 \times 10^{-11}$, $m_\mu^2/M_Z^2 \simeq 10^{-6}$, $m_\tau^2/M_Z^2 \simeq 4 \times 10^{-4}$), which are smaller than the experimental uncertainty in decay rate measurements. Thus, the result is written as

$$\Gamma_{ll} = \frac{g^2 (g_V'^2 + 1) M_Z}{192\pi \cos^2\theta} [1 - 8 \times |g_2' C_0| M_Z] \times \Theta(M_Z - 2 m_l), \quad (35)$$

with the LV contribution appearing as a direct correction to the usual decay rate. We have used $g = e/\sin\theta$, $g_V' = 1–4\sin^2\theta$, $\sin^2\theta = 0.23$. In accordance with Ref. [33], the Z_0 decay rate (considering lepton universality) is $\Gamma_{ll} = (83.985 \pm 0.086)$ MeV, or $\Gamma_{ll} = 83.985(1 \pm 0.001)$ MeV, so that the experimental uncertainty is of 1 part in 10^3. We thus impose $8|g_2' C_0| M_Z < 1.0 \times 10^{-3}$, which leads to the upper bound $|g_2' C_0| < 1.3 \times 10^{-15}$ $(\mathrm{eV})^{-1}$, that is,

$$|g_2' C_0| < 1.3 \times 10^{-6} \ (\mathrm{GeV})^{-1}. \quad (36)$$

IV. A NONMINIMAL COUPLING IN THE SU(2)$_L$ SECTOR OF THE GSW MODEL

Analogously to the previous case, a gauge invariant nonminimal interaction in the $SU(2)_L$ sector of the GSW model can be proposed as

$$D_\mu = \partial_\mu - i g \mathbf{T} \cdot \mathbf{W}_\mu - i \frac{g'}{2} Y B_\mu + i g_3' (\mathbf{T} \cdot \mathbf{W}_{\mu\nu}) V^\nu, \quad (37)$$

where V^ν is a fixed 4-vector that establishes a preferred direction in spacetime and violates Lorentz symmetry. The interaction term, $\bar{L}_l \gamma^\mu i (i g_3' \mathbf{T} \cdot \mathbf{W}_{\mu\nu} V^\nu) L_l$ embraces the following interactions at tree-level,

$$\mathcal{L}_{LV(2)} = \mathcal{J}_{+(LV)}^{(l)\nu} W_\nu^{(+)} + \mathcal{J}_{-(LV)}^{(l)\nu} W_\nu^{(-)} + \mathcal{J}_{0(LV)}^{(l)\nu} Z_\nu, \quad (38)$$

involving the vector bosons, in which the related currents read

$$\mathcal{J}_{+(LV)}^{(l)\nu} = \tilde{g}_3' [-\bar{\psi}_{\nu_l} \gamma^\mu (1 - \gamma_5) \psi_l V^\nu \partial_\mu + \bar{\psi}_{\nu_l} \gamma^\nu (1 - \gamma_5) \psi_l V^\mu \partial_\mu], \quad (39)$$

$$\mathcal{J}_{-(LV)}^{(l)\nu} = \tilde{g}_3' [-\bar{\psi}_l \gamma^\mu (1 - \gamma_5) \psi_{\nu_l} V^\nu \partial_\mu + \bar{\psi}_l \gamma^\nu (1 - \gamma_5) \psi_{\nu_l} V^\mu \partial_\mu], \quad (40)$$

$$\begin{aligned}\mathcal{J}_{0(LV)}^{(l)\nu} = & -\frac{g_3' \cos\theta}{4} \{ \bar{\psi}_{\nu_l} \gamma^\mu (1 - \gamma_5) \psi_{\nu_l} V^\nu \partial_\mu \\ & - \bar{\psi}_l \gamma^\mu (1 - \gamma_5) \psi_l V^\nu \partial_\mu - \bar{\psi}_{\nu_l} \gamma^\nu (1 - \gamma_5) \psi_{\nu_l} V^\mu \partial_\mu \\ & + \bar{\psi}_l \gamma^\nu (1 - \gamma_5) \psi_l V^\mu \partial_\mu \}, \end{aligned} \quad (41)$$

with $\tilde{g}_3' = g_3'/(2\sqrt{2})$. The current, $J_{-(LV)}^{(l)\mu}$, given by Eq. (40), affects the processes mediated by the W^- particle, including the decay $W^- \to l + \bar{\nu}_l$. The total electroweak current that contributes to this process is

$$\begin{aligned}(J_-^{(l)\mu} + \mathcal{J}_{-(LV)}^{(l)\mu}) W_\mu^{(-)} = & [\tilde{g} j_2^\mu W_\mu^{(-)}(x) - \tilde{g}_3' j_2^\eta V^\mu \partial_\eta W_\mu^{(-)}(x) \\ & + \tilde{g}_3' j_2^\mu V^\lambda \partial_\lambda W_\mu^{(-)}(x)], \end{aligned} \quad (42)$$

where $j_2^\mu(x) = \bar{\psi}_l(x) \gamma^\mu (1 - \gamma_5) \psi_{\nu_l}$, and the first term is the usual Lorentz invariant contribution. The scattering matrix for the process $(W^- \to l + \bar{\nu}_l)$, at leading order, can be written as

$$\mathcal{S} = -i \int d^4 x (J_-^{(l)\mu} + \mathcal{J}_{-(LV)}^{(l)\mu}) W_\mu^{(-)}, \quad (43)$$

that implies $\mathcal{S} = \mathcal{S}_0 + \mathcal{S}_{LV(1)} + \mathcal{S}_{LV(2)}$, with

$$\mathcal{S}_0 = -i \frac{g}{2\sqrt{2}} \int d^4 x [j_2^\mu(x) W_\mu^{(-)}(x)], \quad (44)$$

$$\mathcal{S}_{LV(1)} = i \frac{g_3'}{2\sqrt{2}} \int d^4 x [j_2^\eta(x) V^\mu \partial_\eta W_\mu^{(-)}(x)], \quad (45)$$

$$\mathcal{S}_{LV(2)} = -i \frac{g_3'}{2\sqrt{2}} \int d^4 x [j_2^\mu(x) V^\lambda \partial_\lambda W_\mu^{(-)}(x)]. \quad (46)$$

Following the same steps of the previous calculation, we obtain the decay rate for the usual Lorentz invariant process $(W^- \to l + \bar{\nu}_l)$:

$$\Gamma_{\mathcal{S}_0 \mathcal{S}_0^\dagger} = \frac{g^2}{48\pi} M_W \left(1 - \frac{m_l^2}{M_W^2}\right)^2 \left(1 + \frac{m_l^2}{2 M_W^2}\right) \Theta(M_W - m_l), \quad (47)$$

where M_W, m_l stand for the W^- boson and lepton masses. As it occurs in the previous case, the quantities $\Gamma_{\mathcal{S}_0\mathcal{S}^\dagger_{LV(1)}}$, $\Gamma_{\mathcal{S}_{LV(1)}\mathcal{S}^\dagger_0}$ also vanish. The terms, $\Gamma_{\mathcal{S}_0\mathcal{S}^\dagger_{LV(2)}}$, $\Gamma_{\mathcal{S}_{LV(2)}\mathcal{S}^\dagger_0}$ are computed, leading to the following decay rate:

$$\Gamma = \left[\frac{g^2}{48\pi}M_W + (g_3'V_0)\frac{10gM_W^2}{384\pi}\right]\Theta(M_W - m_l), \qquad (48)$$

where $V\cdot k = V_0 M_W$ for the rest frame of the W^- mediator, and we have neglected the contributions in $m_l^2/M_W^2, m_l^4/M_W^4$. This result can be also expressed as

$$\Gamma = \frac{g^2}{48\pi}M_W\left[1 + (g_3'V_0)\frac{5M_W}{4g}\right]\Theta(M_W - m_l). \qquad (49)$$

Considering that the experimental uncertainty in the measures of this decay is at the level of $\sim 4.0\times 10^{-2}$, and using $g = e/\sin\theta, \sin^2\theta = 0.23$, we impose $7(g_3'V_0)M_W < 4.0\times 10^{-2}$, yielding $|g_3'V_0| < 7\times 10^{-14}\ (\mathrm{eV})^{-1}$, or

$$|g_3'V_0| < 7\times 10^{-5}\ (\mathrm{GeV})^{-1}. \qquad (50)$$

As the current (39), involving the mediator W^+, is analogue to the current (40), we conclude that these latter developments equally hold to the decay $W^+ \to \bar{l} + \nu_l$, which becomes constrained by a bound similar to Eq. (50).

V. CONCLUSION AND FINAL REMARKS

We have explicitly computed the corrections implied by two *CPT*-odd nonminimal electroweak couplings to the decay rates of the processes, $Z_0 \to \bar{l} + l$ and $W^- \to l + \bar{\nu}_l$ ($W^+ \to \bar{l} + \nu_l$). Regarding the experimental imprecision in the measurements, upper limits were imposed on the magnitude of the LV nonminimal coupling at the level of $10^{-6}\ (\mathrm{GeV})^{-1}$ and $10^{-5}\ (\mathrm{GeV})^{-1}$.

In LV theories the background components are considered fixed in the Sun's frame, in such a way there appear sidereal variations in the Earth frame [30,34], being necessary to translate the bounds from the Earth lab, at the colatitude χ, rotating around the Earth's axis with angular velocity Ω, to the Sun's frame. For experiments up to a few weeks long, for a rank-1 tensor, A_μ, it holds $A^T_\mu = \mathcal{R}_{\mu\alpha}A_\alpha$, where the label T indicates the quantity measured in the Sun's frame, and $\mathcal{R}_{0i} = \mathcal{R}_{i0} = 0$ and $\mathcal{R}_{00} = 1$. Thus, $A^T_0 = A_0$, so that the upper bounds (36), (50) could be equally written in the Sun's frame. However, the situation is not so simple, as pointed out in Ref. [13] (for pion decays), once the decay rates (35), (49) were carried out in the rest frame of the decaying vector bosons, not in the Lab (Earth) frame, where the measurements are performed. In order to take into account this point, one option is to translate the upper bounds (36), associated with an evaluation at the vector boson rest frame, directly to the Sun's frame, with the boost

$$C^0 = \gamma_z(C^0_T + \alpha^i C^i_T), \qquad (51)$$

where $\gamma_z = \gamma(\mathrm{v}_z)$ is the Lorentz factor, v_z is the boson velocity in the Sun's frame, $\alpha^i = \mathrm{v}^i_z/c$. The data about width decays were attained in the LEP accelerator [33], constructed to work with center-of-mass energy around 91 GeV, reaching 161 GeV in 2000. As the Z_0 mass is close to the center-of-mass energy, it happens that the Lorentz factor is nearly 1 ($\gamma_z \gtrsim 1$), and not larger than 2, which also implies a not (meaningful) relativistic velocity (v^i_z). In this case, the upper bounds (36), (50) can be read in the Sun's frame as

$$|g_2'(C^0_T + \alpha^i C^i_T)| \lesssim 1\times 10^{-6}\ (\mathrm{GeV})^{-1}, \qquad (52)$$

$$|g_3'(V^0_T + \alpha^i V^i_T)| \lesssim 1\times 10^{-5}\ (\mathrm{GeV})^{-1}. \qquad (53)$$

If the case the center-of-mass energy is really close to the boson mass ($\gamma_z \simeq 1$), these bounds simplify to

$$|g_2'C^0_T| < 1\times 10^{-6}, |g_3'V^0_T| < 1\times 10^{-5}, \qquad (54)$$

measured in $(\mathrm{GeV})^{-1}$.

Another possibility is to write the results (35), (49) in the Lab frame, in which $C\cdot k = M_z(C_0 - \gamma_z\alpha^i C^i)$, $V\cdot k = M_z(V_0 - \gamma_z\alpha^i V^i)$, procedure that, for $\gamma_z \simeq 1$, yields $|g_2'(C_0 - \alpha^i C^i)| \simeq |g_2'C_0)| \lesssim 1\times 10^{-6}\ (\mathrm{GeV})^{-1}$, $|g_3'(V_0 - \alpha^i V^i)| \simeq |g_3'V_0| \lesssim 1\times 10^{-5}\ (\mathrm{GeV})^{-1}$. In such a situation these latter bounds can be read in the Sun frame, leading to the results (54).

Other impacts of these NMC can be investigated in electroweak phenomenan, including di¤erential decay rates of polarized processes, which could, in principle, yield improved upper bounds.

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