Time-reversal symmetry violation in several Lepton-Flavor-Violating processes

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Abstract: We compute a T-odd triple vector correlation for the $\mu \to e\gamma$ decay and the $\mu \to e$ conversion process. We find simple results in terms of the CP violating phases of the effective Hamiltonians. Then we focus on the minimal Left-Right symmetric extension of the Standard Model, which can lead to an appreciable correlation. We show that under rather general assumptions, this correlation can be used to discriminate between Parity or Charge-conjugation as the discrete Left-Right symmetry.
1 Introduction

Lepton Number Violating (LNV) and Lepton Flavor violating (LFV) processes are forbidden in the Standard Model (SM) and are thus a good probe of new physics. In principle new physics brings also new sources of CP violation.

In this paper we compute CP and T-odd triple vector correlations for the LFV $\mu \to e\gamma$ and $\mu \to e$ conversion processes, since much of the present and future experimental attention is devoted to these two processes. The MEG collaboration report the best experimental limit for the $\mu \to e\gamma$ process [1]

$$\text{Br}(\mu \to e\gamma) = \frac{\Gamma(\mu \to e\gamma)}{\Gamma(\mu \to e\nu\nu)} < 5.7 \times 10^{-13}$$

(1.1)

and the SINDRUM II collaboration gives the strongest limits for the $\mu \to e$ conversion process [2, 3], namely

$$\text{Br}(\mu + \text{Ti}(\text{Au}) \to e + \text{Ti}(\text{Au})) < 6.1(7) \times 10^{-13}.$$  

(1.2)
Upgrades of ongoing experiments have been considered with the final goal of achieving a sensitivity around $10^{-18} - 10^{-19}$ [4–7]. Given the current limits and the future improvements, there exist the possibility of having enough statistics to start probing CP violation beyond the SM in the next round of experiments. This is suggested and studied in [8, 9].

In this work we focus on quantities that test T violation –and therefore CP violation in a local, Lorentz invariant QFT– in the absence of final-state interactions. Among these quantities are triple vector correlations made up of the momenta or spins of the participating particles [10]. In [11], it is suggested that triplet vector correlations can be used to probe CP violation in the $\mu \rightarrow e$ conversion process. In this paper we present the first analytical computation for the correlation suggested in [11] for the $\mu \rightarrow e$ conversion process and we extend their work in two ways: first, we compute the correlation for the $\mu \rightarrow e\gamma$ decay and second we include the full set of effective operators responsible for the $\mu \rightarrow e$ conversion process.

In section 2 we introduce some theoretical tools for the $\mu \rightarrow e\gamma$ and the $\mu \rightarrow e$ conversion processes. In section 3 as an example of a theory that gives order one contribution to the triple vector correlation, we briefly introduce the minimal Left-Right (LR) symmetric extension of the SM. In section 4 and 5 we present the analytical computation of the triple correlation in the $\mu \rightarrow e\gamma$ and $\mu \rightarrow e$ conversion process respectively. Then in section 6 and for both processes, we study these correlations in the context of the minimal LR model, for both parity and charge-conjugation as the LR symmetries. Finally in section 7 we present our conclusions.

2 General theory

2.1 $\mu \rightarrow e\gamma$ process

The $\mu \rightarrow e\gamma$ decay is predicted to be negligible small in the SM with massive neutrinos, therefore if this process is seen it implies that new physics is behind it.

The effective Hamiltonian for the $\mu \rightarrow e\gamma$ process is given by

$$H_{eff} = \frac{4eG_F m_\mu}{\sqrt{2}} \bar{e}(p_e) \sigma_{\mu\nu} F_{\mu\nu} (A_L P_L + A_R P_R) \mu(p_\mu) + h.c.,$$

(2.1)

where $F_{\mu\nu}$ is the electromagnetic field strength for the photon field, $G_F$ is the Fermi constant, $P_{(R,L)} = \frac{1}{2}(1\pm\gamma_5)$, $m_\mu$ is the muon mass and $e$ and $\mu$ are the spinors for the electron and muon respectively. The coefficients $A_L$ and $A_R$ are calculated within a given physical model.

2.2 $\mu \rightarrow e$ conversion. Theory and Effective Hamiltonian

Theoretical studies of this process were performed in the past [12–15]. In [15] the outgoing electron coming from the conversion process, belongs to one of the states in the continuum energy spectrum for the Coulomb potential. As a matter of fact the outgoing electron must be treated as a plane wave. One way to argue this is by noticing that an electron in the continuum energy spectrum, is described by a Dirac spinor in the angular momentum basis
(an eigenstate of $\hat{J}_z^2, j_z, K$, and the Hamiltonian $H$). Experimentally, the detected electron has a definite 4-momentum implying that the outgoing electron must be a plane wave.

In this work we present a method for computing a triple vector correlation that tests T-violation in the $\mu \to e$ conversion process for various nuclei. We make use of the formalism developed in [16].

We use the following representation for the $\gamma$ matrices

$$\gamma_0 = \beta = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \quad \gamma_i = \begin{pmatrix} 0 & \sigma_i \\ -\sigma_i & 0 \end{pmatrix},$$

and

$$\sigma_{\mu\nu} = \frac{i}{2} [\gamma_\mu, \gamma_\nu], \quad \gamma_5 = -i\gamma_1\gamma_2\gamma_3\gamma_0,$$

where the $\sigma'$s are the Pauli matrices.

The Dirac’s equation for the central field problem in polar coordinates is given by (the energy is given in units of the electron mass)

$$E\psi = H\psi = [-i\gamma_5\Sigma_r (\frac{\partial}{\partial r} + \frac{1}{r} - \frac{\beta}{r} K) + V + \beta]\psi,$$

where

$$\Sigma_r = \frac{1}{r} \sum_i \Sigma_i, \quad \Sigma_i = i\frac{1}{2} [\gamma_j, \gamma_k] \quad (\{i,j,k\} \text{ cyclic}),$$

$$K = \beta(\Sigma \cdot L + 1),$$

and $V$ is the Coulomb potential.

We write the wave function $\psi$ as [17]

$$\psi = \psi_\kappa^\mu = \begin{pmatrix} g_\kappa(r) \chi_\kappa^\mu \\ i f_\kappa(r) \chi_{-\kappa}^\mu \end{pmatrix},$$

$g_\kappa$ and $f_\kappa$ are the radial functions that obey the differential equations

$$\frac{dg_\kappa(r)}{dr} = -\kappa r^{-1} g_\kappa(r) + (E - V + 1) f_\kappa(r),$$

$$\frac{df_\kappa(r)}{dr} = \kappa r^{-1} f_\kappa(r) - (E - V - 1) g_\kappa(r).$$

In the high energy limit -all the masses are set to zero- and from eqs.(2.8) and (2.9), $f_\kappa(r)$ and $g_\kappa(r)$ satisfy

$$f_{-\kappa} = -g_\kappa, \quad g_{-\kappa} = f_\kappa.$$

We will use this result for the spinor describing the electrons coming from the conversion process.
The initial muon has the quantum numbers, \( \mu = \pm \frac{1}{2} \) and \( \kappa = -1 \) with a normalization
\[
\int d^3 x \psi^{(\mu)\dagger}_{1s}(\vec{x}) \psi^{(\mu)}_{1s}(\vec{x}) = 1. \tag{2.11}
\]

For the electrons in the continuum-energy states we use the same normalization considered in [15], namely
\[
\int d^3 x \psi^{(e)\dagger}_{\kappa,E}(\vec{x}) \psi^{(e)}_{\kappa,E}(\vec{x}) = 2\pi \delta_{\mu \mu'} \delta_{\kappa \kappa'} \delta(E - E'). \tag{2.12}
\]

In the conversion process the effective Hamiltonian is given by [15]
\[
H_{\text{eff}} = \frac{4G_F}{\sqrt{2}} (m_\mu A_1^\mu \bar{\mu} \sigma^{\mu\nu} P_L e F_{\mu\nu} + m_\mu A_1^\mu \bar{\mu} \sigma^{\mu\nu} P_R e F_{\mu\nu} + \text{h.c.})
+ \frac{G_F}{\sqrt{2}} \sum_{q=u,d,s} [(g_{LS(q)} \bar{e} P_R \mu + g_{RS(q)} \bar{e} P_L \mu) \bar{q} q + (g_{LP(q)} \bar{e} P_R \mu + g_{RP(q)} \bar{e} P_L \mu) \bar{q} \gamma_5 q
+ (g_{LV(q)} \bar{e} \gamma_\mu P_L \mu + g_{RV(q)} \bar{e} \gamma_\mu P_R \mu) \bar{q} \gamma_\mu q + (g_{LA(q)} \bar{e} \gamma_\mu P_L \mu + g_{RA(q)} \bar{e} \gamma_\mu P_R \mu) \bar{q} \gamma_\mu \gamma_5 q + \frac{1}{2} (g_{LT(q)} \bar{e} \sigma^{\mu\nu} P_R \mu + g_{RT(q)} \bar{e} \sigma^{\mu\nu} P_L \mu) \bar{q} \sigma_{\mu\nu} q + \text{h.c.}]. \tag{2.13}
\]

The nuclear form factors were calculated in [18]. The wave function for the muon and the electrons in the presence of a central field were obtained in [14, 15]. In particular in [15] updated data for the proton and neutron densities were used.

For \( r \to \infty \) the general solution \( \langle \psi_{as} \rangle \) for a Dirac particle in a Coulomb field at first order in \( H_{\text{eff}} \) is of the form [16]
\[
\psi_{as} = -i \sqrt{\frac{e^{ipr}}{|p|}} \sum_{\kappa \mu} e^{i\delta_\kappa} \langle \psi^{(e)\mu}_\kappa | H_{\text{eff}} | \psi^{(\mu)}_{1s} \rangle \left( \frac{\sqrt{E + \chi^\mu_{\kappa}(\hat{p})}}{\sqrt{E - \chi^\mu_{-\kappa}(\hat{p})}} \right) + \mathcal{O}(H_{\text{eff}}^2), \tag{2.14}
\]
where \( \hat{p} \) is in the direction of the outgoing electron. The phases \( e^{i\delta_\kappa} \) are the usual ones appearing in scattering problems in the presence of a Coulomb field and are given by
\[
\delta_\kappa = y \ln 2pr - \arg \Gamma(\gamma + iy) + \eta_\kappa - \frac{1}{2} \pi \gamma \tag{2.15}
\]
\[
y = \alpha Z E / p, \quad \gamma = \sqrt{\kappa^2 - \alpha^2 Z^2}, \quad e^{2ip\eta_\kappa} = -\frac{\kappa - iy/E}{\gamma + iy}. \tag{2.16}
\]

We consider states with \( \kappa = \pm 1 \), hence the only term relevant for our discussion is \( \eta_\kappa \)–the remaining ones will be just an overall phase in the solution \( \psi_{as} \).

Finally the total conversion rate per unit flux is 1
\[
\omega_{\text{conv.total}} = R^2 \int d\Omega \psi^{\dagger}_{as} \psi_{as} = 2\pi |\langle \psi^{(e)}_\mu | H_{\text{eff}} | \psi_{1s} \rangle|^2. \tag{2.17}
\]

\(^1\text{See appendix B for a more detailed discussion on this issue.}\)
3 The minimal Left-Right symmetric theory

As an example of a complete and predictive theory of lepton number violating phenomena, we consider the minimal LR symmetric extension of the SM [19–22]. In the minimal LR symmetric theory the scalar sector is [23–26]

\[
\Phi = \begin{pmatrix}
\phi_0^1 \\
\phi_1^1 \\
\phi_0^2 \\
\phi_1^2
\end{pmatrix}, \quad \Delta_{L,R} = \begin{pmatrix}
\delta_{L,R}^+ / \sqrt{2} \\
\delta_{L,R}^0 \\
\delta_{L,R}^+ / \sqrt{2}
\end{pmatrix},
\]

(3.1)

where \(\Phi\) is in the (2,2,0) representation of \(SU(2)_L \times SU(2)_R \times U(1)_{B-L}\) and the two scalar triplets \(\Delta_L\) and \(\Delta_R\), belong to the (3,1,2) and the (1,3,2) representations respectively.

The Yukawa interactions of leptons with the scalar triplets have the form

\[
\mathcal{L}_Y = \bar{L}_L Y Y \Phi \Phi \Phi + \bar{L}_L \bar{Y}_R Y \Delta_L \Delta_L L_L + \frac{1}{2} (L^T_L C \sigma_2 Y \Delta_L \Delta_L L_L + h.c.),
\]

(3.2)

where \(\tilde{\Phi} = \sigma_2 \phi_2^* \phi_2\), \(\sigma_2\) being the Pauli matrix, \(L_L\) is the lepton doublet of the standard model \((L^T_L = (\nu^L))\), \(L_R\) is its right-handed analogue \(L^T_R = (N^L)\) and \(N\) is the heavy Majorana neutrino.

In the mass eigenstate basis the flavor changing charged current Lagrangian is given by

\[
\mathcal{L}_{cc} = \frac{g}{\sqrt{2}} (\bar{\nu}_L V^\dagger_L W^L l_L + \bar{N}_R V^\dagger_R W^R l_R) + h.c.,
\]

(3.3)

\(V_R\) is the right-handed analogue of the PMNS mixing matrix \(V_L\). In general it has three different mixing angles and six arbitrary complex phases and we parametrize it as

\[
V_R = K_e \hat{V}_R K_N,
\]

with \(K_e = \text{diag}(e^{i \phi_e}, e^{i \phi_\mu}, e^{i \phi_\tau})\), \(K_N = \text{diag}(1, e^{i \phi_2}, e^{i \phi_3})\). The matrix \(\hat{V}_R\) has three mixing angles and the dirac phase \(\delta\). We choose for \(\hat{V}_R\) the standard form for the CKM matrix shown in the PDG [27].

The interaction terms of charged leptons with the doubly-charged scalars are

\[
\mathcal{L}_\Delta = \frac{1}{2} \bar{t}_F^R C Y^\prime \Delta_R \delta_{R}^{++} l_R + \frac{1}{2} \bar{t}_L^T C Y^\prime \Delta_L \delta_{L}^{++} l_L + h.c.,
\]

(3.4)

\[
Y^\prime_{\Delta_R} = \frac{g}{M_{W_R}} V_{R}^{*\dagger} m_N V_{R}^{\dagger},
\]

(3.5)

If charge conjugation (\(C\)) is the discrete LR symmetry and since the charged lepton masses are symmetric, the Yukawa couplings in (3.4) satisfy (for reviews on this topic see references [28–30])

\[
Y^\prime_{\Delta_L} = (Y^\prime_{\Delta_R})^*.
\]

(3.6)

For parity (\(P\)) and in the more interesting phenomenological situations, the charged lepton masses matrices are almost hermitian [31]. In [32] it was realized that it implies the near equality between the Yukawa couplings in eq. (3.4). More precisely

\[
Y^\prime_{\Delta_L} = Y^\prime_{\Delta_R} + \mathcal{O}(\tan 2\beta \sin \alpha),
\]

(3.7)
where $\beta$ is the usual ratio of the v.e.v’s $v_2/v_1$ and $\alpha$ is the spontaneous phase. In [33, 34] it is shown that $\tan 2\beta \sin \alpha \lesssim 2m_b/m_t$ so that the Yukawa coupling of the doubly charged scalar are nearly equal.

It is a remarkable feature of the minimal LR theory, that the TeV energy scale accessible at the LHC through the KS process [35] –and its associated LNV and LFV, predicts the rate for the neutrino-less double beta decay and low energy LFV. This deep connection and the related phenomenology are illustrated in [36, 37]. All these processes depend in a crucial way on the elements of the leptonic right-handed mixing matrix $V_R$, for which all its mixing angles, the Dirac phase and two Majorana phases can be determined at the LHC [32]. Useful information can also be obtained from EDM of the neutron and such [38–42]. This is deeply connected to the study of the strong CP parameter, which in the mLRSM turns out to be calculable [43–45].

The CMS collaboration [46] has reported an excess in the ee-channel for this process at $2.8\sigma$, but they claimed that this excess cannot be accommodated in the minimal version of the theory –assuming no mixing in the right-handed leptonic sector. Several works have been proposed [47–49] in order to explain this excess and the conclusion was that it would need a higher Left-Right symmetry breaking scale, or a more general mixing scenario with pseudo-Dirac heavy neutrinos. Recently in [50], it was realized that it is still possible to accommodate the CMS data in the minimal version of the theory ($g_L = g_R$), if the CP phases and non-degenerate masses are considered. Of course next run of the LHC will be crucial to establish or discard this excess.

4 Computation of a triple vector correlation in the $\mu \to e\gamma$ decay

$T$-odd asymmetries in the $\mu \to e\gamma$ were considered in the past. In [8, 9], it was shown that by studying the polarization of electron and the photon coming from the muon decay it is possible to extract the CP-violating phases from the experiment. The conclusion was that in order to extract the CP-violating phases both electron and photon polarizations must be measured.

In this paper instead, we present an alternative way of extracting the CP-violating phases of the effective Hamiltonian in the $\mu \to e\gamma$ decay. This is complementary to the work presented in [8, 9]. The novelty is that no measurements of the final photon polarizations are needed. We consider the T-violating triple vector product

$$\hat{s}_{\mu^+} \cdot (\hat{p}_{e^+} \times \hat{s}_{e^+}) = \cos \Phi \sin \theta_s,$$

(4.1)

where $\theta_s$ is the angle between the polarization’s direction ($\hat{s}_{e^+}$) of the positron and its momentum’s direction $\hat{p}_{e^+}$, $\Phi$ is the angle formed between $\hat{s}_{\mu^+}$ and the direction defined by $\hat{p}_{e^+} \times \hat{s}_{e^+}$ and $\Psi$ is the azimuthal angle. In Fig.1 the reference frame and setup are shown. Notice that this quantity changes sign under parity and naive time-reversal transformation $\hat{T}$ (i.e. $t \to -t$). For processes whose interactions are characterized by a small coupling, it can be shown at first order that the connected part of the S-matrix is hermitian [10]. Therefore, in this case the violation of the $\hat{T}$ symmetry amounts the violation of the time-reversal symmetry $T$. 

\[ -6 - \]
We define the triple vector correlation as

\[
\langle \hat{s}_{\mu^+} \cdot (\hat{p}_{e^+} \times \hat{s}_{e^+}) \rangle_{\Phi} = \frac{N(\cos \Phi > 0) - N(\cos \Phi < 0)}{N_{\text{total}}} = \frac{\int_{\Phi_0}^{\pi} d\Phi d\Gamma/d\Phi \cdot \text{sgn}(\hat{s}_{\mu^+} \cdot (\hat{p}_{e^+} \times \hat{s}_{e^+}))}{\Gamma_{\text{total}}},
\]

where \( N(\cos \Phi > 0) \) and \( N(\cos \Phi < 0) \) are the number of events satisfying \( \cos \Phi > 0 \) and \( \cos \Phi < 0 \) respectively.

The 4-momenta of the participating particles in the rest frame of the muon are given by

\[
p_{\mu^+} = (m_{\mu}, 0, 0, 0) \quad (4.3)
\]

\[
p_{e^+} = (E_{e^+}, |\vec{p}_{e^+}| \sin \theta_s, |\vec{p}_{e^+}| \cos \theta_s, 0) \quad (4.4)
\]

\[
p_{\gamma} = (E_{\gamma}, -|\vec{p}_{e^+}| \sin \theta_s, -|\vec{p}_{e^+}| \cos \theta_s, 0) \quad (4.5)
\]

where the mass of the positron has been neglected. The energy \( E_{e^+} \) of the positron and the energy \( E_\gamma \) of the photon are given by

\[
E_{e^+} \equiv E_{\gamma} = |\vec{p}_{e^+}| = \frac{m_{\mu}}{2}. \quad (4.6)
\]

From the effective Hamiltonian in eqn. (2.1) and eqns. (A.9), (A.10) and (A.11) in appendix A, a straightforward computation gives the following value for the correlation in
The main advantage of this quantity is that no measurements of the photon polarizations are needed.

In summary we find that given a source of polarized anti-muons, by measuring the 3-momentum $\vec{p}_{e^+}$ of the outgoing positron and its polarization $\vec{s}_{e^+}$, the asymmetry shown in eqn. $(4.7)$ is sensitive to the CP-violating phases of the effective Hamiltonian shown in $(2.1)$. In [51–55] it is shown that measurements of the polarization of electrons coming from the muon decay are feasible. We assume a 100 % polarized muon flux so that our results must be trivially rescaled by the actual polarization of the initial muons.

5 Computation of a triple vector correlation in the $\mu \to e$ conversion process

Following the same lines of the last section, we define an asymmetry given by comparing the number of events with $\vec{s}_{\mu} \cdot (\vec{p}_{e^+} \times \vec{s}_{e^+}) > 0$ with the ones satisfying $\vec{s}_{\mu} \cdot (\vec{p}_{e^+} \times \vec{s}_{e^+}) < 0$ in the $\mu \to e$ conversion process.

Once again we define the triple vector correlation as

$$\langle \vec{s}_{\mu} \cdot (\vec{p}_{e^+} \times \vec{s}_{e^+}) \rangle_\Phi \equiv \frac{N(\cos \Phi > 0) - N(\cos \Phi < 0)}{N_{\text{total}}} = \frac{\omega_{\text{conv}}(\cos \Phi > 0) - \omega_{\text{conv}}(\cos \Phi < 0)}{\omega_{\text{conv, total}}} \quad (5.1)$$

and as previously, $\Phi$ is the angle between the plane formed by the vectors $\vec{p}_{e^+}$ and $\vec{s}_e$ and the polarization of the muon $\vec{s}_\mu$. We used the same coordinate system shown in Fig.1 but clearly there is no photon coming from the muon decay.

A direct computation of the asymmetry shown in eqn. $(5.1)$ gives

$$\langle \vec{s}_{\mu} \cdot (\vec{p}_{e^+} \times \vec{s}_{e^+}) \rangle_\Phi = \frac{1}{2} \sin \theta_e \Re m(C_L C_R^*) + \mathcal{O}(\alpha Z) + \mathcal{O}(\frac{m_e}{E_e}). \quad (5.2)$$

where

$$C_R \equiv DA_R + S^{(p)}(g_{LS}^{(p)} + g_{LV}^{(p)}) + S^{(n)}(g_{LS}^{(n)} + g_{LV}^{(n)}) \quad (5.3)$$

$$C_L \equiv DA_L + S^{(p)}(g_{RS}^{(p)} + g_{RV}^{(p)}) + S^{(n)}(g_{RS}^{(n)} + g_{RV}^{(n)}) \quad (5.4)$$

and

$$g_{LS,RS}^{(p)} \equiv \sum_q G^{(q,p)}_L g_{LS,RS(q)}, \quad \tilde{g}_{LS,RS}^{(n)} \equiv \sum_q G^{(q,n)}_L g_{LS,RS(q)} \quad (5.5)$$

$$g_{LV,RV}^{(p)} = 2g_{LV,RV(u)} + g_{LV,RV(d)}, \quad \tilde{g}_{LV,RV}^{(n)} \equiv g_{LV,RV(u)} + 2g_{LV,RV(d)}. \quad (5.6)$$

$D$, $S^{(n,p)}$ are nuclear constants already calculated and tabulated in [15] for various elements. This expression is valid for non-relativistic muons. We also drop terms of the order $\alpha Z$ and terms of the order $m_e/E_e$. In practice equation $(5.2)$ must be multiplied by the polarization of the initial muons, which is of the order of 15% in the conversion process [56].
6 Triplet vector correlation in the minimal Left-Right theory

As a concrete example of a theory beyond the SM that gives order one values for the T-odd triple vector correlation \[11\], we consider the minimal LR symmetric model. In what follows we analyze separately the contributions to the asymmetries (4.7) and (5.2) in the case of \( P \) and \( C \) as the LR symmetries. In \[11\] it is found that this contribution can be of order one, since there are new contributions coming from interactions of charged leptons with the singly-charged and doubly-charged scalar fields.

6.1 \( \mu \to e\gamma \) decay

In this section and for the \( \mu \to e\gamma \) decay, we study the contributions to the triple vector correlation for both Parity and Charge Conjugation as the LR symmetry.

Parity as the LR symmetry: in \[57\] the authors presented a complete study of the contributions to several LFV processes in the context of the minimal LR extension of the SM. It was found in their work that \( A_L \) and \( A_R \) in eqn. (2.1) are explicitly given by

\[
A_R = \frac{1}{16\pi^2} \sum_n (V_R^\dagger n)(V_R)_{np} \left[ \frac{M_W^2}{M_{W_R}^2} S_3(X_n) - \frac{X_n M_W^2}{3 m_{\delta^+_R}} \right],
\]

\[
A_L = \frac{1}{16\pi^2} \sum_n (V_R^\dagger n)(V_R)_{np} X_n \left[ -\frac{1}{3} \frac{M_W^2}{m_{\delta^+_R}^2} - \frac{1}{24} \frac{M_W^2}{M_{H^+_1}^2} \right] + O(\tan 2\beta \sin \alpha),
\]

\[
X_n = \left( \frac{m_N}{M_{W_R}} \right)^2, \quad S_3(x) = -\frac{x}{8} \frac{1+2x}{(1-x)^2} + \frac{3x^2}{4(1-x)^2} \frac{x}{(1-x)^2} (1-x + \log x) + 1.
\]

where \( N_n \) are the heavy neutrino mass eigenstates and \( n = 1, 2, 3 \).

Notice that the function \( S_3 \) is always small as far as \( m_N \) is not much bigger than \( M_{W_R} \), so that the term with the loop function \( S_3 \) can neglected for a wide range of the heavy neutrino masses (see figure 2). In this case the contribution to the correlation defined in (4.7) is suppressed. Assuming \( m_{\delta^+_R}^2 \gtrsim m_{\delta^+_R}^2 \ll M_{W_R} \), the correlation also vanishes – as already emphasized in \[11\]. Finally we neglect the contribution of the charged Higgs \( H^+_1 \) since its mass cannot be lower than \( (15-20) \) TeV \[33, 58\]. This poses no problem for the theory, since its mass emerges at the large scale of symmetry breaking \[23, 59\]. The gauge boson and doubly-charged scalar masses can be obtained at the LHC through the so called KS process and the decays of the doubly charged scalars \[35\]. Let us just mention that all the mixing angles and the Dirac phase in \( V_R \) can be in principle obtained at the LHC \[32\], as well as two of the Majorana phases. This is example shows vividly the complementary role played by the high and low energy experiments in the establishment of the LR theory \[36, 60–65\].

For the sake of illustration, imagine that type II see-saw is the dominant source of neutrino masses i.e. \( \frac{M_N}{(\Delta_R)} = \frac{M_L}{(\Delta_L)} \) and \( V_L = V_R \). In this case it is possible to show \[36\]

\[
\frac{m_{N2}^2 - m_{N1}^2}{m_{N3}^2 - m_{N1}^2} = \frac{m_{\nu2}^2 - m_{\nu1}^2}{m_{\nu3}^2 - m_{\nu1}^2} \approx \pm 0.03,
\]

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where the ± corresponds to normal/inverted hierarchy respectively. In figure 3 (top) we show the value of the triple vector correlation correlation given in (4.7) in the \((m_0, \delta)\)-plane, where \(m_0\) is the lightest neutrino mass and \(\delta\) is the Dirac phase present in \(\tilde{V}_R\). From Fig.3 (top) we see that the value of the correlation (4.7) goes from \(10^{-6}\) to \(10^{-5}\) for \(m_0 < 1\) TeV and \(0 < \delta < 2\pi\).

Clearly the triple vector correlation can be bigger for general values of neutrino masses and mixings. However from eqns. (3.7), the contribution to the triple vector correlation shown in (4.7) is bounded to be less \(10^{-2}\) since \(\tan 2\beta \sin \alpha < 10^{-2}\) from the quark masses [33, 34, 45].

**Charge conjugation as the LR symmetry**: in this case from eqn. (3.6) we have that

\[
A_R = \frac{1}{16\pi^2} \sum_n (V_R^\dagger e_n V_R)_{\mu \nu} \left[ \frac{M_W^2}{M_W^2} S_3(X_n) - \frac{X_n}{3 \left( M_W^2 \right)^2} \right],
\]

\[
A_L = \frac{1}{16\pi^2} \sum_n (V_R^\dagger e_n V_R)_{\mu \nu} \left[ -\frac{1}{3} \frac{M_W^2}{M_W^2} + \frac{1}{24} \frac{M_W^2}{M_W^2} \right].
\]

Notice that some of the external phases appearing in \(V_R\) do not cancel in (4.7) and the triple vector correlation is proportional to \(\sin 2(\phi_\mu - \phi_e)\). In this case the triple vector correlation is not suppressed by the small \(\theta_{13}\) mixing-angle. In Fig.3 (bottom) we show the value of the triple vector correlation in the \((m_0, \delta)\)-plane. We take \(\phi_\mu - \phi_e = 0\) for \(\phi_\mu - \phi_e = \pi/4\) it will reach in maximum value of around 0.5 in almost all the parameter space. We do it in both normal and inverted neutrino mass hierarchies.

We can see from Fig.3 (bottom) that \(C\) as the LR symmetry gives larger contributions to the triple vector correlation than the \(P\) case, because for parity the triple vector correlation is suppressed due to the near equality between the Yukawa couplings in eqn (3.4).
Figure 3. (Top) Contour plots illustrating the value of the asymmetry defined in (4.7) as a function of the lightest neutrino mass $m_{0}$ and the Dirac phase $\delta$ for $P$ as the LR symmetry. (Bottom) Contour plots illustrating the value of the asymmetry defined in (4.7) as a function of the lightest neutrino mass $m_{0}$ and the Dirac phase $\delta$ (assuming $\phi_{u} - \phi_{e} = 0$) for $C$ as the LR symmetry. (Left) Normal hierarchy for neutrino masses. (Right) Inverse hierarchy for neutrino masses. We take the gauge boson mass $m_{W_{R}} = 3.5$ TeV, the heaviest right-handed neutrino mass $m_{\text{heaviest}} = 1$ TeV and the masses for the doubly charged scalars of 1 TeV. The mixing angles are $\theta_{12} = 35^\circ$, $\theta_{23} = 45^\circ$, $\theta_{13} = 7^\circ$.

The bottom line is that if a value for the triple vector correlation bigger than $10^{-2}$ is measured, it can only be the consequence of $C$ as the LR symmetry.

6.2 $\mu \rightarrow e$ conversion process

In this section we consider the triple vector correlation for the $\mu \rightarrow e$ conversion process in the context of the minimal LR symmetric extension of the SM.

In [57] it was shown that the contribution of the doubly-charged scalar may dominate due to a logarithmic enhancement. In this case only $\tilde{g}_{L_{V_{R}}}^{(p)}$ contributes to the triple vector correlation. More precisely eq. (5.2) is of the form

$$\langle \vec{s}_{\mu} \cdot (\vec{p}_{e} \times \vec{s}_{e}) \rangle = \frac{\sin \theta_{s} \Im (\tilde{g}_{L_{V_{R}}}^{(p)} \tilde{g}_{R_{V}}^{(p)})}{2 |\tilde{g}_{L_{V}}^{(p)}|^2 + |\tilde{g}_{R_{V}}^{(p)}|^2} = \frac{\sin \theta_{s} \Im (F_{L}^{(\gamma)} F_{R}^{(\gamma)^*})}{2 |F_{L}^{(\gamma)}|^2 + |F_{R}^{(\gamma)}|^2},$$

(6.7)
where $F_L^{(\gamma)}$ and $F_R^{(\gamma)}$ are defined in [57]. For the logarithmic enhanced terms this functions have the same flavor structure of the coefficients $A_L$ and $A_R$ defined previously. Therefore the same conclusion obtained in the $\mu \to e\gamma$ case, holds for the $\mu \to e$ conversion process as well.

7 Conclusions

We derived analytical expressions for a T-odd triple vector correlation in the $\mu \to e\gamma$ decay and the $\mu \to e$ conversion process and found simple results in terms of the CP-violating phases of the effective Hamiltonians. The expression obtained in the $\mu \to e$ conversion omits relativistic corrections for the muons, but is otherwise complete. For the $\mu \to e\gamma$ decay we conclude that in order to extract the CP violating phases of the theory from the experiment, no measurements of the photon polarizations are needed.

Then as an example of a theory that leads order one values for the triple vector correlation we consider the TeV scale, minimal Left-Right symmetric extension of the SM. Remarkably, due to the relation between left and right Yukawa couplings in (3.4) –see also eqs. (3.6) and (3.7)– this triple vector correlation can be used to discriminate between charge-conjugation or parity as the Left-Right symmetry. More precisely, a value for the triple vector correlation bigger than $10^{-2}$ can only be the consequence of charge-conjugation as the Left-Right symmetry.

A Kinematics of the $\mu \to e\gamma$ process

In this appendix we give some kinematical tools we used for the computation in the $\mu \to e\gamma$ decay.

Suppose we want to find the spinor $u(p)$, which represent a fermion with 4-momentum $p$ and suppose also that this fermion is polarized in a given arbitrary direction $\hat{n}$. Then as a solution of the free Dirac’s equation $u(p)$ satisfies

$$\left(p - m\right)u(p) = 0.$$  \hfill (A.1)

The solution to this equation is of the form (taking only the positive square root)

$$u(p) = \left(\frac{\sqrt{p \cdot \sigma} \xi}{\sqrt{p \cdot \bar{\sigma} \xi}}\right),$$  \hfill (A.2)

where $\xi$ is a two-component Pauli spinor, normalized such that

$$\xi^\dagger \xi = 1,$$  \hfill (A.3)

and

$$p \cdot \sigma = p^0 \sigma^0 - p^i \sigma^i,$$  \hfill (A.4)

$$p \cdot \bar{\sigma} = p^0 \sigma^0 + p^i \sigma^i.$$  \hfill (A.5)
The point is that $\xi$ transforms under rotations as an ordinary two-component spinor of the rotation group and therefore it determines the spin orientation of the fermion. In the rest frame of a massive fermion its spin is given by:

$$ \vec{s} = |\vec{s}|(\sin\Phi \cos\Psi, \sin\Phi \sin\Psi, \cos\Phi) \quad (A.6) $$

and if we want to find a fermion whose spin is given by (A.6), we can perform a rotation around the x-axis by an angle $-\Phi$ and the around the z-axis by an angle $-(\frac{\pi}{2} - \Psi)$.

It is straightforward to show that the desired spinor is

$$ \xi^n = \begin{pmatrix} e^{-i\frac{\Psi}{2}} \cos\frac{\Phi}{2} \\ e^{i\frac{\Psi}{2}} \sin\frac{\Phi}{2} \end{pmatrix}. \quad (A.7) $$

Alternatively one can find the same result by requiring $\xi^n$ to be an eigenvector of $\vec{\sigma} \cdot \hat{n}$, where $\hat{n}$ is a unitary vector in the direction of $\vec{s}$.

How about the term $\sqrt{p \cdot \sigma}$ and $\sqrt{p \cdot \bar{\sigma}}$?, if we assume that the fermion is at rest. We have that $\sqrt{p \cdot \sigma} = \sqrt{p \cdot \bar{\sigma}} = \sqrt{m}$. Hence for a fermion at rest and with its spin oriented as (A.6) $u(p)$ has the form

$$ u(p) = \sqrt{m} \mu \left( \begin{array}{c} \xi^n \\ \xi^n \end{array} \right). \quad (A.8) $$

For an anti-muon we use $v(p)$ instead $u(p)$. In this case $v(p)$ is given by

$$ v(p) = \begin{pmatrix} \sqrt{p \cdot \sigma} \cdot \eta \\ -\sqrt{p \cdot \bar{\sigma}} \cdot \eta \end{pmatrix} \quad (A.9) $$

where $\eta \dagger \eta = 1$.

For the electron and for the reference frame shown in Fig.1 we find

$$ u_e^+(p_e^+) = \sqrt{\frac{|p_e^+|^2}{2}} \begin{pmatrix} -2e^{i\frac{\phi}{2}} \sin\frac{\theta_s}{2} \\ 2ie^{-i\frac{\phi}{2}} \sin\frac{\theta_s}{2} \\ 2ie^{i\frac{\phi}{2}} \cos\frac{\theta_s}{2} \\ -2e^{-i\frac{\phi}{2}} \cos\frac{\theta_s}{2} \end{pmatrix} \quad (A.10) $$

The photon has two possible polarizations along the direction of motion and in the particular frame we are considering in Fig.1 its polarization vector is given by,

$$ e_\pm^+(p_\gamma) = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \pm i \cos \theta_s \\ \pm i \sin \theta_s \\ 1 \end{pmatrix} \quad (A.11) $$

where we can explicitly see that when $\theta_s = 0$, the photon can only have a polarization $\pm 1$ along the y-axis.
In this appendix we briefly comment about the amplitude of the $\mu \rightarrow e$ conversion process and the Born's approximation we used.

In computing the $\mu \rightarrow e$ conversion process, one usually assumes the so called Born’s approximation for the outgoing electrons. This approximation has two meanings: one is computing the conversion rate to a given order in some small coupling; and the other is the assumption that electrons coming from the conversion process are plane waves. The point is that we can do better and have a complete control of both approximations at the same time. More precisely, for the relativistic one-electron atom and in the limit of big $r$ ($r >> r_0$, where $V(r \geq r_0) = 0$), the solution of the Dirac’s equation at first order in the perturbation $H_{\text{eff}}$ is of the form [16]

$$
\psi_{\text{as}} = -i \frac{\sqrt{\pi}}{|p|} e^{i p r} \sum_{\kappa \mu} e^{i \delta_{\kappa \mu}} \langle \psi_{\kappa} | \psi_{\mu} \rangle \left( \frac{\sqrt{E + 1} \chi_{\kappa}^{\mu}(\hat{p})}{-\sqrt{E - 1} \chi_{\mu}^{\kappa}(\hat{p})} \right) + O(H_{\text{eff}}^2),
$$

(B.1)

where $\psi_i$ is any stationary state of the Coulomb potential, $\psi_{\kappa}^\mu$ is one of the continuum energy solutions and $H_{\text{eff}}$ is the effective Hamiltonian for the $\mu \rightarrow e$ conversion process. Furthermore it can be shown that $\psi_{\text{as}}$ is an eigenfunction of $\vec{\alpha} \cdot \vec{p} + \beta$ with eigenvalue $E$ so that $\psi_{\text{as}}$ describes, indeed, a plane wave [16].

In the high energy limit –neglecting the electron mass– the solution $\psi_{\text{as}}$ simplifies to

$$
\psi_{\text{as}} = -i \frac{\sqrt{\pi}}{r} e^{i p r} \sum_{\kappa \mu} e^{i \delta_{\kappa \mu}} \langle \psi_{\kappa}^{\mu} | \psi_{\mu} \rangle \left( \frac{\chi_{\kappa}^{\mu}(\hat{p})}{-\chi_{\mu}^{\kappa}(\hat{p})} \right).
$$

(B.2)

Finally if we are interested in computing the total conversion amplitude per unit flux (for a detector placed at fixed radius $r = R$) the total conversion rate is given by

$$
\omega_{\text{conv,total}} = R^2 \int d\Omega \psi_{\text{as}}^\dagger \psi_{\text{as}} = 2\pi |\langle \psi_{\text{as}}^{\mu} | H_{\text{eff}} | \psi_{\text{as}} \rangle|^2.
$$

(B.3)

We can absorb the $\sqrt{2\pi}$ factor into the normalization of the wave function $\psi_{\kappa}^{\mu}$, as shown in eq. (2.12).

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