Azimuthal asymmetry of recoil electrons in neutrino-electron elastic scattering as signature of neutrino nature

W. Sobków∗, A. Blaut

 Institute of Theoretical Physics, University of Wrocław, Pl. M. Born 9, PL-50-204 Wrocław, Poland

Abstract

In this paper, we show how a presence of the exotic scalar, tensor weak interactions in addition to the standard vector-axial (V-A) one may help to distinguish the Dirac from Majorana neutrinos in the elastic scattering of (anti)neutrino beam off the unpolarized electrons in the limit of vanishing (anti)neutrino mass. We assume that the incoming (anti)neutrino beam comes from the polarized muon decay at rest and is the left-right chiral mixture with assigned direction of the transversal spin polarization with respect to the production plane. We display that the azimuthal asymmetry in the angular distribution of recoil electrons is generated by the interference terms between the standard and exotic couplings, which are proportional to the transversal (anti)neutrino spin polarization and independent of the neutrino mass. This asymmetry for the Majorana neutrinos is distinct from the one for the Dirac neutrinos through the absence of interference between the standard and tensor couplings. Additionally, the interference term between the standard and scalar coupling of the only left chiral neutrinos, absent in the Dirac case, appears. We also indicate the possibility of utilizing the azimuthal asymmetry measurements to search for the new CP-violating phases. Our analysis is model-independent and consistent with the current upper limits on the non-standard couplings.

Keywords: neutrino nature, exotic couplings of right chiral neutrinos, neutrino-electron elastic scattering, transversal neutrino spin polarization

PACS: 13.15.+g, 14.60.St

1. Introduction

One of the fundamental problems in the neutrino physics is whether the neutrinos (ν)’s are the Dirac or Majorana fermions. The question of ν nature can

∗Corresponding author

Email addresses: wieslaw.sobkow@ift.uni.wroc.pl (W. Sobków), arkadiusz.blaut@ift.uni.wroc.pl (A. Blaut)
be probed in the context of non-vanishing $\nu$ mass and of standard vector-axial ($V - A$) weak interaction with only the left chiral (LCh) $\nu$'s, using purely leptonic processes such as the polarized muon decay at rest (PMDaR) or the neutrino-electron elastic scattering (NEES). There is an alternative way within the massless $\nu$ limit, when one admits the existence of exotic scalar (S), tensor (T), pseudoscalar (P) and $V + A$ weak interactions of the right chiral (RCh) $\nu$'s (right-handed helicity when $m_\nu \to 0$) in addition to the $V - A$ interaction of the LCh ones in the above processes. The appropriate tests involving the mass dependence have been proposed by Kayser and Langacker. It is also worthwhile remarking the others interesting ideas regarding the $\nu$ nature problem. One ought to emphasize that at present the neutrinoless double beta decay (NDBD) seems to be the best tool to investigate the $\nu$ nature problem. First tests concerning the problem of distinguishing between the Dirac and Majorana $\nu$'s in the limit of vanishing $\nu$ mass, when one departs from the $V - A$ interaction and one allows for the exotic S, T, P weak interactions in the NEES, have been reported by Rosen and Dass. The leptonic processes are also suitable to probe the time reversal violation (TRV) effects. It is relevant to point out that the existing data still leaves a small space for the exotic couplings of the interacting RCh $\nu$'s. It is noteworthy that the effects coming from the interacting $\nu$'s with right-handed chirality are also important for interpreting of results on the NDBD. Unfortunately, the proposed quantities in are composed of the squares of exotic couplings of the RCh $\nu$'s and at most of the interferences within exotic couplings, that are both very tiny. Furthermore, both transverse components of electron (positron) spin polarization and neutrino energy spectrum in the PMDaR contain only the interference terms between the standard $V$ and non-standard S, T couplings of LCh $\nu$'s. All the eventual interferences between the standard couplings of LCh and exotic couplings of RCh $\nu$'s vanish, because are proportional to a tiny $\nu$ mass and do not produce the effect. As the current experiments do not detect the RCh $\nu$'s, it seems meaningful to search for new tools including the linear terms from the exotic couplings that are independent of the $\nu$ mass, and obtained in model-independent way. It would enable to compare the predictions of various non-standard schemes with the experimental data, and look for the TRV effects. The suitable observables could be the $\nu$ quantities carrying information on the transversal components of (anti)neutrino spin polarization, both T-odd and T-even. Presently, such tests are still not available, because they require the observation of final $\nu$'s, the strong $\nu$ beam coming from the polarized source and the efficient $\nu$ polarimeters. However, it is worthy of indicating the potential possibilities of experiments of the $\nu$ polarimetry in the connection with the $\nu$ nature problem, the existence of interacting RCh $\nu$'s and the non-standard TRV phases predicted by many extensions of the SM. Let us recall that the SM can not be viewed as a ultimate theory, because it does not clarify the origin of parity violation at current energies, the observed baryon asymmetry of universe through a single CP-violating phase of the Cabibbo-Kobayashi-Maskawa quark-mixing matrix (CKM), the large...
hierarchy fermion masses, and others fundamental aspects. This situation led to the appearance of various non-standard gauge models including the Majorana $\nu$'s, exotic TRV interactions, mechanisms explaining the origin of fermion generations, masses, mixing and smallness of $\nu$ mass.

In this paper, we focus on the elastic scattering of electron $\nu$'s ($\overline{\nu}$'s) off the unpolarized electron target and show in model-independent way how the participation of the exotic S, T couplings of RCh $\nu$'s ($\overline{\nu}$'s) in addition to the standard couplings of LCh ones can be utilized to distinguish the Dirac from Majorana $\nu$'s, and to test the TRV in the limit of vanishing $\nu$ mass.

2. Elastic scattering of Dirac electron antineutrinos off unpolarized electrons

We assume that the incoming Dirac $\overline{\nu}_e$ beam comes from the decay of polarized negative muons at rest ($\mu^- \rightarrow e^- + \overline{\nu}_e + \nu_\mu$) and is a mixture of left-right chiral states with a fixed direction of the transversal spin polarization with respect to the production plane. LCh $\overline{\nu}_e$'s are mainly detected by the standard $V - A$ interaction and RCh ones are detected only by the exotic scalar, tensor interactions in the elastic scattering on the unpolarized electron target; $(\overline{\nu}_e + e^- \rightarrow \overline{\nu}_e + e^-)$. Our scenario admits also the detection of $\overline{\nu}_e$'s with left-handed chirality by the non-standard S and T interactions. The production plane for the $\overline{\nu}_e$ beam, shown in Fig. 1, is spanned by the unit vector $\hat{\eta}_\mu$ of the muon polarization and the $\overline{\nu}_e$ LAB momentum unit vector $\hat{q}$. In this plane, the vector $\hat{\eta}_\mu$ can be expressed, with respect to $\hat{q}$, as a sum of $(\hat{\eta}_\mu \cdot \hat{q})\hat{q}$ and $\eta_\mu^\perp = \hat{\eta}_\mu - (\hat{\eta}_\mu \cdot \hat{q})\hat{q}$. The reaction (detection) plane is spanned by the direction of the outgoing electron momentum $p_e$ and $\hat{q}$, Fig. 1. The amplitude for the $\overline{\nu}_e e^-$ scattering at low energies is as follows:

$$M_{\overline{\nu}_e e^-}^D = -\frac{G_F}{\sqrt{2}} \left\{ (\overline{\nu}_e \gamma^\alpha (e^L_A \gamma_5) u_e) (\overline{\nu}_e \gamma_\alpha (1 - \gamma_5) v_{\nu_e}) ight\} (1)$$

$$+ c_{S}^L(\overline{\nu}_e u_e)(\overline{\nu}_e (1 - \gamma_5) v_{\nu_e})$$

$$+ \frac{1}{2} c_T^L(\overline{\nu}_e \sigma^{\alpha\beta} u_e)(\overline{\nu}_e \sigma_{\alpha\beta} (1 - \gamma_5) v_{\nu_e})$$

$$+ c_{S}^L(\overline{\nu}_e u_e)(\overline{\nu}_e (1 + \gamma_5) v_{\nu_e})$$

$$+ \frac{1}{2} c_T^L(\overline{\nu}_e \sigma^{\alpha\beta} u_e)(\overline{\nu}_e \sigma_{\alpha\beta} (1 + \gamma_5) v_{\nu_e}) \right\},$$

where $G_F = 1.1663788(7) \times 10^{-5}$ GeV$^{-2}$ (0.6 ppm) [13] is the Fermi constant.

The coupling constants are denoted with the superscripts $L$ and $R$ as $c^L_V, c^L_A, c^R_L, c^{R,L}_T$ respectively to the incoming $\overline{\nu}_e$ of left- and right-handed chirality. Because we take into account the TRV, all the coupling constants are complex. Calculations are carried out with use of the covariant density matrix for the polarized initial $\overline{\nu}_e$. The formula for the projector $\Lambda^{(s)}_{\overline{\nu}_e}$ in the massless $\overline{\nu}_e$ limit is given by:

$$\lim_{m_{\overline{\nu}_e} \rightarrow 0} \Lambda^{(s)}_{\overline{\nu}_e} = \frac{1}{2} \left\{ (q^\mu \gamma_\mu) [1 - \gamma_5(\hat{\eta}_\mu \cdot \hat{q}) - \gamma_5 S^\perp \cdot \gamma] \right\},$$

(2)
Figure 1: Production plane of the $\nu_e$ beam is spanned by the vectors $\hat{\eta}_\mu$ and $\hat{q}$ for $\mu^- \rightarrow e^- + \pi^0 + \nu_\mu$. Reaction plane is spanned by the vectors $\hat{p}_e$ and $\hat{q}$ for $\pi^- + e^- \rightarrow \nu_\mu + e^-$. $\hat{\eta}_\nu$ is expressed, with respect to $\hat{q}$, as a sum of $(\hat{\eta}_\nu \cdot \hat{q})\hat{q}$ and $\eta^\perp_\nu$. 

4
where \( \hat{\eta_\nu} \) is the unit 3-vector of \( \nu_e \) spin polarization in its rest frame; \( (\hat{\eta_\nu} \cdot \hat{q})\hat{q} \) is the longitudinal component of \( \nu_e \) spin polarization; \( \eta^{\perp}_{\nu}=\hat{\eta_\nu}-(\hat{\eta_\nu} \cdot \hat{q})\hat{q} \) is the transversal component of \( \nu_e \) spin polarization; \( S^\perp=(0,\eta^{\perp}_{\nu}) \).

We see that in spite of the singularities \( m_\nu^{-1} \) in the Lorentz boosted spin polarization 4-vector of massive \( \nu_e \) \( S' \) (in the laboratory frame), the projector \( \Lambda^{(s)}_s \) including \( \eta^{\perp}_{\nu} \) remains finite \[16\]. One should notice that the last term in \( \Lambda^{(s)}_s \) has different \( \gamma \)-matrix structure from that of the longitudinal polarization contribution. This term will generate the non-vanishing interferences between the standard and exotic couplings in the differential cross section for the \( \nu_e e^- \) scattering.

Using the current data for the muon decay at rest \[17\], we calculate the upper limit on the magnitude of \( \eta^{\perp}_{\nu} \) and lower bound for \( (\hat{\eta_\nu} \cdot \hat{q}) \), \[18\]:

\[
|\eta^{\perp}_{\nu}| = 2 \sqrt{Q_L^s(1 - Q_L^s)} \leq 0.537, \quad (3)
\]
\[
|\hat{\eta_\nu} \cdot \hat{q}| = |2Q_L^s - 1| \geq 0.843,
\]
\[
Q_L^s = 1 - \frac{1}{4}(m^S_{L,R} |^2 + m^S_{LL} |^2) - 3|m^T_{LR} |^2 \geq 0.922.
\]

2.1. Azimuthal distribution of recoil electrons in case of Dirac electron antineutrinos

The differential cross section for the scattering of Dirac \( \nu_e \)'s off the unpolarized electrons, when \( m_\nu \rightarrow 0 \), takes the form:

\[
\frac{d^2\sigma}{dy_e d\phi_e} = \left( \frac{d^2\sigma}{dy_e d\phi_e} \right)_{(V-A)} + \left( \frac{d^2\sigma}{dy_e d\phi_e} \right)_{(S,T)} + \left( \frac{d^2\sigma}{dy_e d\phi_e} \right)_{(AT)},
\]

\[
\left( \frac{d^2\sigma}{dy_e d\phi_e} \right)_{(V-A)} = B \left\{ 1 + \hat{\eta_\nu} \cdot \hat{q} \right\} \left[ |c^L_e|^2 + |c^R_e|^2 + |c^A_e|^2 \right],
\]

\[
\left( \frac{d^2\sigma}{dy_e d\phi_e} \right)_{(S,T)} = B \left\{ 1 + \hat{\eta_\nu} \cdot \hat{q} \right\} \left[ \frac{1}{2} y_e \left( y_e + 2 \frac{m_e}{E_e} \right) |c^S_e|^2 \right.
\]
\[
+ \left( 2 - y_e^2 - \frac{m_e}{E_e} y_e \right) |c^R_e|^2 - y_e(y_e - 2) \text{Re}(c^S_e c^R_e) \right]
\]
\[
+ (1 + \hat{\eta_\nu} \cdot \hat{q}) \left\{ \frac{1}{2} y_e \left( y_e + 2 \frac{m_e}{E_e} \right) |c^L_e|^2 \right.
\]
\[
+ \left( 2 - y_e^2 - \frac{m_e}{E_e} y_e \right) |c^L_e|^2 - y_e(y_e - 2) \text{Re}(c^S_e c^L_e) \right\},
\]

\[
\left( \frac{d^2\sigma}{dy_e d\phi_e} \right)_{(V-S)} = B \left\{ -4 \sqrt{y_e(y_e + 2 \frac{m_e}{E_e})} \right\} \eta^{\perp}_{\nu} \cdot (\hat{p}_e \times \hat{q}) \text{Im}(c^L_e c^R_e) \}
\]
\[
\left( \frac{d^2 \sigma}{dy_e d\phi_e} \right)_{(AT)}^{(V,S)} + \left( \frac{d^2 \sigma}{dy_e d\phi_e} \right)_{(AT)}^{(R,"'S)} = -B \left\{ -2 \sqrt{y_e} \left( 1 + \frac{m_e}{E_{\nu}} \right) \left[ \gamma_T \cdot (\hat{p}_e \times \hat{q}) \right] \right. \\
+ \left( \frac{d^2 \sigma}{dy_e d\phi_e} \right)_{(AT)}^{(R,"'S)} \right\},
\]

where \( \phi \) is the angle between \( \gamma_T \) and \( \gamma_L \) only; \( \phi_0 = \phi - \phi_e \) is the angle between \( \hat{p}_e \) and \( \gamma_L \). \[ \begin{align*}
\left( \frac{d^2 \sigma}{dy_e d\phi_e} \right)_{(V,S)}^{(R,"'S)} &= -B \left[ \gamma_T \right] \left[ \gamma_L \right] \\
&= \left\{ -4|c_T|^2|c_R|^2 \cos(\phi + \beta_{SV} - \phi_e) + 2|c_T|^2|c_R|^2 \cos(\phi + \beta_{TA} - \phi_e) \right\},
\end{align*} \]

is the ratio of the kinetic energy of the recoil electron \( T_e \) to the incoming antineutrino energy \( E_\nu \): \( B \equiv \left( E_\nu m_e/4\pi^2 \right) \left( G_F^2/2 \right) \), \( \theta_e \) is the angle between \( \hat{p}_e \) and \( \hat{q} \) (recoil electron scattering angle); \( m_e \) is the electron mass; \( \phi_e \) is the angle between the production plane and the reaction plane (azimuthal angle of outgoing electron momentum). We see that the interference terms, Eqs. \( \eta_{V,L} \), between the standard couplings in the NEES; \( c_T \) and \( c_R \) couplings in the massless \( \nu_\tau \) limit. There are no interferences between \( c_T^{L,R} \) and \( c_R^{L,R} \) couplings of the LCh \( \nu_\tau \)'s for \( m_\tau \to 0 \). It can be noticed that the interferences include only the contributions from the transverse components of the \( \nu_\tau \) spin polarization, both \( T \)-even and \( T \)-odd:

Using the experimental values of standard couplings: \( c_T^{L,R} = 1 \pm (-0.04 \pm 0.015), c_R^{L,R} = 1 \pm (-0.507 \pm 0.014) \), we find the upper limits on the exotic couplings in the NEES: \( |c_T^{L,R}| \leq 0.501, |c_R^{L,R}| \leq 0.486, |c_T^{L,R}| \leq 0.003, |c_R^{L,R}| \leq 0.489, \) simultaneously shifting the standard values to new ones \( c_T^{L,R} = 0.945, c_R^{L,R} = 0.479 \), which still lie within the experimental bars. If one uses the above limits, one gets the same magnitude of total cross section as in the standard case. Now, we calculate the upper limits on the azimuthal asymmetry between the \( (0, \pi) \) and \( (\pi, 2\pi) \) angles (up-down asymmetry) using the upper limit on \( |\eta_T^{L,R}| \) and lower
bound on \((\tilde{\eta}_\nu \cdot \mathbf{q})\):

\[
A(\phi_{e'}) = \frac{\int_{\phi_{e}'}^{\phi_{e}'} d\sigma d\phi_e - \int_{\phi_{e}'}^{\phi_{e}'+2\pi} d\sigma d\phi_e}{\int_{\phi_{e}'}^{\phi_{e}'+2\pi} d\sigma d\phi_e + \int_{\phi_{e}'}^{\phi_{e}'+2\pi} d\sigma d\phi_e}.
\]

(11)

The differential cross section is integrated over \(y_e \in [0, 1/(1 + (m_e/2E_\nu))]\). We get for the case of TRV and time reversal conservation (TRC) with \(E_\nu = 50 \text{ MeV}\), respectively:

\[
A_D(\phi_{e'})^{T-vio} = 0.024 \cos(\phi - \phi_{e'}) \text{ for } \beta_{VS} = \frac{\pi}{2}, \beta_{AT} = \frac{\pi}{2},
\]

(12)

\[
A_D(\phi_{e'})^{T-cons} = -0.024 \sin(\phi - \phi_{e'}) \text{ for } \beta_{VS} = 0, \beta_{AT} = 0.
\]

(13)

3. Elastic scattering of Majorana electron neutrinos off unpolarized electrons

The amplitude for the elastic scattering of the Majorana electron neutrinos (\(\nu_e\)'s) on the unpolarized electrons at low energies has the form:

\[
M_{\nu_e e^-}^M = \frac{G_F}{\sqrt{2}} \left\{ (-2)(\mathbf{\tau}_{e'}\gamma^\alpha (c_V - c_A\gamma_5)u_e)(\mathbf{\bar{\tau}}_{\nu_e}\gamma^\alpha\gamma_5u_{\nu_e})
+ 2c_P^R(\mathbf{\bar{u}}_{\nu_e}u_e)(\mathbf{\bar{u}}_{\nu_e}(1 + \gamma_5)u_{\nu_e})
+ 2c_L^S(\mathbf{\bar{u}}_{\nu_e}u_e)(\mathbf{\bar{u}}_{\nu_e}(1 - \gamma_5)u_{\nu_e}) \right\}.
\]

(14)

We see that the above matrix element in the neutrino part does not contain the contribution from the V and T interactions in contrast to the Dirac case, where both terms partake. In addition, the A and S contributions are multiplied by factor 2. This arises from the fact that the Majorana neutrino is described by the self-conjugate field. Absence of the index L for \(c_V, c_A\) couplings means that both LCh and RCh \(\nu_e\)'s may participate in the standard A interaction of Majorana \(\nu_e\)'s. In consequence, the new term with the interference between \(c_V\) and \(c_L^S\) couplings in the differential cross section appears. Such interference vanishes in the Dirac case. All the couplings are assumed to be complex as for the Dirac case. The others assumptions concerning the production of \(\nu_e\) beam and the detection of \(\nu_e\)'s by the interaction with the unpolarized electron target are the same as in the Dirac scenario.

3.1. Azimuthal distribution of recoil electrons in case of Majorana neutrinos

The azimuthal distribution of the recoil electrons for the scattering of Majorana \(\nu_e\)'s on the unpolarized electrons with \(m_\nu \rightarrow 0\) has the form:

\[
\frac{d^2\sigma}{dy_e d\phi_e} = \left( \frac{d^2\sigma}{dy_e d\phi_e} \right)_{(V-A)} + \left( \frac{d^2\sigma}{dy_e d\phi_e} \right)_{(S_L,S_R)}
+ \left( \frac{d^2\sigma}{dy_e d\phi_e} \right)_{(VS_R)} + \left( \frac{d^2\sigma}{dy_e d\phi_e} \right)_{(VS_L)}.
\]

(15)
The interference contribution can be written down as follows:

\[
\left( \frac{d^2 \sigma}{dy_e d\phi_e} \right)_{(V-A)} = B \left\{ |c_V - c_A|^2 (2 + (1 - \hat{q}_V \cdot \hat{q}) (y_e - 2) y_e) + |c_V + c_A|^2 (2 + (1 + \hat{q}_V \cdot \hat{q}) (y_e - 2) y_e) - \frac{2 m_e y_e}{E_\nu} (|c_V|^2 - |c_A|^2) \right\},
\]

\[
\left( \frac{d^2 \sigma}{dy_e d\phi_e} \right)_{(S_L, S_R)} = B \left\{ 2 y_e \left( y_e + \frac{2 m_e}{E_\nu} \right) \left[ (1 + \hat{q}_V \cdot \hat{q}) |c_R|^2 + (1 - \hat{q}_V \cdot \hat{q}) |c_L|^2 \right] \right\},
\]

\[
\left( \frac{d^2 \sigma}{dy_e d\phi_e} \right)_{(V S_R)} = B \left\{ 8 \sqrt{y_e (y_e + \frac{2 m_e}{E_\nu})} \left[ - \eta^\perp_{\nu} \cdot (\hat{p}_e \times \hat{q}) \text{Im}(c_V c_R^{R*}) + (\eta^\perp_{\nu} \cdot \hat{p}_e) \text{Re}(c_V c_R^{R*}) \right] \right\},
\]

\[
\left( \frac{d^2 \sigma}{dy_e d\phi_e} \right)_{(V S_L)} = B \left\{ 8 \sqrt{y_e (y_e + \frac{2 m_e}{E_\nu})} \left[ \eta^\perp_{\nu} \cdot (\hat{p}_e \times \hat{q}) \text{Im}(c_V c_L^{L*}) + (\eta^\perp_{\nu} \cdot \hat{p}_e) \text{Re}(c_V c_L^{L*}) \right] \right\}.
\]

The interference contribution can be written down as follows:

\[
\left( \frac{d^2 \sigma}{dy_e d\phi_e} \right)_{(V S_R)} + \left( \frac{d^2 \sigma}{dy_e d\phi_e} \right)_{(V S_L)} = 8 B |\eta^\perp_{\nu}| \sqrt{\frac{m_e}{E_\nu}} y_e \left[ (2 - (2 + \frac{m_e}{E_\nu}) y_e \right],
\]

where \(\alpha_{V S_R} \equiv \alpha_V - \alpha_S^{R}, \alpha_{V S_L} \equiv \alpha_V - \alpha_S^{L}\) are the relative phases between the \(c_V, c_R\) and \(c_V, c_L\) couplings, respectively. Now, we calculate the upper limits on the azimuthal asymmetry between \((0, \pi)\) and \((\pi, 2\pi)\) angles for the TRV and TRC, using the same limits as for the Dirac case and \(E_\nu = 50 \text{ MeV}\):

\[
A_M(\phi)_{T-violi} \leq -0.062 \cos(\phi - \phi_{e'}) \text{ for } \alpha_{V S_R} = \frac{\pi}{2}, \alpha_{V S_L} = \frac{3\pi}{2}, \quad (21)
\]

\[
A_M(\phi)_{T-conse} \leq 0.062 \sin(\phi - \phi_{e'}) \text{ for } \alpha_{V S_R} = 0, \alpha_{V S_L} = 0. \quad (22)
\]

We see that the possible effect of up-down azimuthal asymmetry for the Majorana \(\nu_e\)'s is larger than in the Dirac case. Moreover, there is a different dependence on the angle \(\phi_{e'}\) in the case of the TRV and TRC, similarly as for the Dirac neutrinos, so the precise measurement of maximal azimuthal asymmetry would answer the question of whether the TRV takes place.

4. Conclusions

We have shown that there is the distinction between the Dirac and Majorana \(\nu\)'s in the limit of vanishing \(\nu\) mass, when the incoming \(\nu\) beam is the mixture
of LCh and RCh states, and has the fixed direction of transversal component of the (anti)neutrino spin polarization with respect to the production plane. If the $\nu$ beam comes from the polarized source (e.g. PMDaR), where the exotic S, T interactions produce the $\nu$’s with the right-handed chirality, while the V-A interaction generates the LCh $\nu$’s, the (anti)neutrino polarization vector may acquire the transversal component (both T-even and T-odd), which is left invariant under Lorentz boost. Next, this left-right chiral mixture is scattered off the unpolarized electrons in the presence of both standard and exotic interactions. The precise measurement of the azimuthal asymmetry of recoil electrons, generated by the interference terms between the standard and exotic couplings, proportional to $\eta_+^\perp$, would allow to distinguish between the Dirac and Majorana $\nu$’s, and test the TRV. For the Majorana $\nu$’s, the upper limit on the expected magnitude of up-down azimuthal asymmetry is larger than for the Dirac case.

According to the SM, the angular distribution of recoil electrons should be azimuthally symmetric in the massless $\nu$ limit, and then there is no difference between the Dirac and Majorana $\nu$’s. The basic difference between the both cases follows from the absence of interference terms between the standard and exotic tensor interactions in the differential cross section for the Majorana $\nu$’s. The additional distinction arises from the occurrence of interference between the standard and S couplings of the LCh Majorana $\nu$’s. This type of interference annihilates for the Dirac $\nu$’s.

It is also important to note that the eventual effects connected with the neutrino mass and mixing for the tests with a near detector are inessential, see e.g. [19]. It is relevant to stress that the azimuthal asymmetry measurements require the very intense polarized (anti)neutrino sources and large unpolarized (polarized) target of electrons (or nucleons), and also long duration of experiment. To make the above tests feasible, the low-threshold, real-time detectors should measure both the polar angle and azimuthal angle of outgoing electron momentum with a high resolution. It is necessary to point out that there is a real interest in the development of low-threshold technology in the context of dark matter searches and the study of neutrino interactions. The silicon cryogenic detectors, the high purity germanium detectors [20], the semiconductor detectors [21] and the bolometers [22] are worth mentioning. The two experiments aiming at the measurement of recoil electron scattering angle and of azimuthal angle, i.e. Hellaz [23] and Heron [24], have also been proposed. Recently, the interesting proposal for particle detection based on the infrared quantum counter concept has emerged [27]. Our studies are reported in hope that it may encourage the neutrino collaborations (e.g. KARMEN, PSI, TRIUMF, BooNE, Borexino, Super-Kamiokande) working with the polarized muon decay, other artificial polarized $\nu$ sources and neutrino beams to realize the measurements of the azimuthal asymmetry of recoil electrons. It seems to be a real challenge, but new tests using the neutrino polarimeters could shed much more light on the $\nu$ nature, detect the existence of the exotic couplings of interacting RCh $\nu$’s and the non-standard phases of TRV than the present measurements based on the electron (positron) observables and energy spectrum of $\nu$’s. Finally, it is worthy of pointing out the fact that in the massless $\nu$ limit, there are physical
and in principle observable effects, coming from the mixture between the LCh and RCh \( \nu \) states (left- and right-handed helicity components) in the spin \( 1/2 \) quantum state, when the exotic S, T interactions coupling these two types of states exist in the weak processes of \( \nu \) production and detection.

References

[1] S. L. Glashow, Nucl. Phys. 22, 579 (1961); S. Weinberg, Phys. Rev. Lett. 19, 1264 (1967); A. Salam, in: N. Svartholm (Ed.), Elementary Particle Theory, Almqvist and Wiksells, Stockholm, 1969; R. P. Feynman, M. Gell-Mann, Phys. Rev. 109, 193 (1958); E. C. G. Sudarshan, R. E. Marshak, Phys. Rev. 109, 1860 (1958).

[2] B. Kayser, R. E. Shrock, Phys. Lett. B 112, 137 (1982).

[3] P. Langacker, D. London, Phys. Rev. D 39, 266 (1989).

[4] V. B. Semikoz, Nucl. Phys. B 498, 39 (1997).

[5] S. Pastor, J. Segura, V. B. Semikoz, J. W. F. Valle, Phys. Rev. D 59, 013004 (1998).

[6] J. Barranco et al., Phys. Lett. B 739, 343 2014.

[7] D. Singh, N. Mobed and G. Papini, Phys. Rev. Lett. 97, 041101 (2006).

[8] T. D. Gutierrez, Phys. Rev. Lett. 96, 121802 (2006).

[9] M. Doi et al., Phys. Lett. B 103, 219 (1981); W. C. Haxton et al., Phys. Rev. Lett. 47, 153 (1981); H. Ejiri, J. Phys. Soc. Jpn. 74, 2101 (2005).

[10] S. P. Rosen, Phys. Rev. Lett. 48, 852 (1982).

[11] G. V. Dass, Phys. Rev. D 32, 1239 (1985).

[12] D. Bogdan, A. Faessler and A. Petrovici, Neutrino masses and right-handed current in the neutrinoless double beta decay 76Ge 76Se + 2e*1, Volume 150, Issues 1-3, 3 January 1985, Pages 29-34.

[13] A. Riotto and M. Trodden, Annu. Rev. Nucl. Part. Sci. 49, 35 (1999).

[14] M. Kobayashi and T. Maskawa, Prog. Theor. Phys. 49, 652 (1973).

[15] D. M. Webber et al., Phys. Rev. Lett. 106, 041803 (2011).

[16] L. Michel and A. S. Wightman, Phys. Rev. 98, 1190 (1955).

[17] K.A. Olive et al. (Particle Data Group), Chin. Phys. C 38, 090001 (2014).

[18] W. Fetscher, Phys. Rev. D 49, 5945 (1994).
[19] W. Sobków, S. Ciechanowicz and M. Misiaszek, Phys. Lett. B 713, 258 (2012).

[20] B.S. Neganov et al., hep-ex/0105083.

[21] C. E. Aalseth et al., Phys. Rev. Lett. 106, 131301 (2011).

[22] C. Enss, Cryogenic Particle Detection (Springer-Verlag Berlin Heidelberg, 2005).

[23] F. Arzarello et al.: Report No. CERN-LAA/94-19, College de France LPC/94-28, 1994. J. Seguinot et al.: Report No. LPC 95 08, College de France, Laboratoire de Physique Corpusculaire, 1995.

[24] R. E. Lanou et al.: The Heron project, Abstracts of Papers of the American Chemical Society 2(217), 021-NUCL 1999.

[25] A. F. Borghesani et al., arXiv: 1506.07987 [physics.ins-det].