A simple gauge theory with Majorana neutrinos

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Abstract

We construct a local gauge theory with Majorana neutrinos and study its implications. The theory is based on the group $SU(2) \times U(1) \times U(1)'$ with a mixing between the $U(1)$ and $U(1)'$ gauge bosons. Majorana neutrinos couple only to the gauged $U(1)'$. There is a small coupling of the Majorana neutrinos to $Z_\mu$ and a large coupling to the new gauge boson $X_\mu$. In addition the $X_\mu$ has small couplings to the Standard Model fermions which leads to new processes. We estimate limits for the coupling constants and compute differential cross sections.
I. INTRODUCTION

Majorana particles are introduced in many articles and for various reasons. Some articles study their effective couplings and predictions as is the case with the anapole form factor \cite{1, 2} and neutrinoless double beta decay \cite{3, 4}. In other cases Majorana states are introduced in the mass matrices, especially for neutrinos, where they influence the masses and the mixing of states \cite{5}.

In this article we study models where Majorana states are introduced in both the gauge interactions and the mass matrices. The model is a simple extension of the electroweak theory with an extra $U(1)^\prime$ group. The extended theory is $SU(2) \times U(1) \times U(1)^\prime$, where the $U(1)^\prime$ part of the Lagrangian is invariant under chiral gauge transformations and acts on right-handed Majorana neutrinos. The new fermions are singlets under the electroweak group and there is one state for each generation. We describe the mass and mixing of the gauge bosons and their interactions with the standard and new fermions. At the end, Majorana neutrinos couple to $Z_\mu$ and a new $X_\mu$ boson. The vertex of the neutrinos is axial-vector which restricts the spectra of the cross sections. In particular, we study the effects in $Z_\mu$ decays, $\nu_e + e^- \rightarrow \nu_e + e^-$ scattering and coherent scattering of neutrinos and antineutrinos on nuclei.

The mass mixing of the gauge bosons preserves the standard form of the electromagnetic and charged current interactions. Thus neutrino beams produced in particle and nuclear decays have a definite flavor. However, as the beams develop in time a Majorana component is generated through the fermion mass matrix and may reveal itself through new interactions. We adopt the point of view that the disappearance of particles in reactor experiments \cite{6, 7} are created by the above oscillations and compute possible cross sections. In addition, the LSND \cite{8} and the MiniBooNE \cite{9} reported an excess of electrons produced in a $\nu_\mu$ beam at a distance of $\leq 550$ m from the Beryllium target where the $\nu_\mu$’s are generated. They were attributed to quasi-elastic scattering \cite{9} of regenerated $\nu_e$’s. In section II we describe the theory and provide estimates for the coupling constants. The appearance of a Majorana component in oscillations is included in section III together with formulas for elastic scattering on atomic electrons. An appealing reaction is the coherent scattering on nuclei presented in section IV. The results are accompanied with numerical estimates for the reactions.
II. THE MODEL

Majorana neutrinos are standard states in many models. Frequently right-handed states are introduced in the fermion mass matrix as is the case in the seesaw mechanism.

The Majorana field can be written as

$$\Psi_M = \frac{1}{\sqrt{2}}(\Psi_R + \Psi^c_R)$$

with $$\Psi_R = \left(\frac{1 + \gamma^5}{2}\right)\Psi$$, $$\Psi^c = C\Psi_R$$ and $$C = i\gamma^2\gamma^0$$. This field cannot be subjected to a gauge transformation of the first kind

$$\Psi_M(x) \rightarrow \Psi'_M(x) = e^{i\epsilon_j(x)\lambda_j^j/2} \Psi_M(x)$$

because the Lagrangian

$$\mathcal{L} = \bar{\Psi}_M(x)i\gamma^\mu \partial_\mu \Psi_M(x)$$

transforms into

$$\mathcal{L}' = \bar{\Psi}_M(x)i\gamma^\mu \left(\partial_\mu + i\frac{\partial\epsilon_j(x)}{\partial x^\mu} \lambda_j^j/2\right) \Psi_M(x)$$

with the second term vanishing since $$\bar{\Psi}_M(x)\gamma^\mu \Psi_M(x) = 0$$. The situation is different when we start with the Lagrangian

$$\mathcal{L}_I = \bar{\Psi}_M(x)i\gamma^\mu \left(\partial_\mu + ig\gamma^5 \frac{\lambda_j^j}{2} X^j_\mu(x)\right) \Psi_M(x).$$

This Lagrangian remains invariant under chirality gauge transformations

$$\Psi_M(x) \rightarrow \Psi'_M(x) = e^{i\gamma^5\epsilon_j(x)\lambda_j^j/2} \Psi_M(x)$$

provided the vector field $$X_\mu(x)$$ transforms as

$$X^j_\mu(x) \rightarrow X'^j_\mu(x) - \frac{i}{g} \frac{\partial \epsilon_j(x)}{\partial x^\mu}.$$ 

In this case the current associated with the transformation is

$$J^\mu = \bar{\Psi}_M(x)\gamma^\mu \gamma^5 \Psi_M(x).$$
is conserved before the introduction of fermion mass terms.

In this article we study a simple theory with Majorana neutrinos coupled to a new gauge field $X_\mu(x)$ belonging to the group $U(1)'$. As mentioned, the theory is based on the group $SU(2) \times U(1) \times U(1)'$ with the $U(1)$ and $U(1)'$ gauge fields mixing in the mass matrix. There is also the possibility of mixing in the kinetic terms which is not included in this article. In the following, we couple the covariant derivative to the fermion multiplets. Then replacing the initial gauge boson with physical fields we find a small component coupling Majorana fermions to $Z_\mu$ and in addition there are small couplings of the standard fermions to $X_\mu$.

The new covariant derivative is

$$D_\mu = \partial_\mu + igt^j \hat{W}^j_\mu + ig' \frac{Y}{2} \hat{B}_\mu + ig \frac{Q_X}{2} \hat{X}_\mu$$

(9)

with $t^j = \tau^j/2$ and $\tau^j$ are the Pauli matrices. The couplings $g, g', g_X$ and the charges $t^j, Y$ and $Q_X$ have their traditional meanings. Similar models were studied in the past [5, 10, 11] and also recently [12, 13].

We consider the masses of the gauge bosons generated through spontaneous symmetry breaking of $SU(2)$ by the Higgs doublet

$$\Phi(x) = \begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix}$$

(10)

acquiring a vacuum expectation value (VEV) $\langle \Phi \rangle = \begin{pmatrix} 0 \\ 1 \end{pmatrix} v/\sqrt{2}$. The singlet vector fields $\hat{B}_\mu(x)$ and $\hat{X}_\mu(x)$ couple to quantities of the $SU(2)$ group in exactly the same way, and thus in the mass generation remain massless. Only one neutral gauge boson receives a mass

$$\left( -\frac{1}{2} g \hat{W}^3_\mu + g' \frac{Y}{2} \hat{B}_\mu + g_X \frac{Q_X}{2} \hat{X}_\mu \right)^2 \frac{v^2}{2}.$$  

(11)

Among the massless bosons one corresponds to the photon. A mass for the third boson is generated when we break the symmetry between $U(1)$ and $U(1)'$. This happens, for example, when we introduce a new Higgs singlet coupled only to $\hat{X}_\mu$ and acquire a VEV $v_0/\sqrt{2}$. Its contribution to the mass matrix is $\frac{1}{4} g_X^2 Q_X^2 v_0^2/2$ which modifies the $(3,3)$ element of the mass matrix.
\[ M_{33}^2 = \frac{1}{4} \left( g_X^2 Q_X^2 + g_X^2 Q_X^2 \frac{v_0^2}{v^2} \right) \frac{v^2}{2} = \frac{1}{4} g_X^2 \bar{Q}_X^2 \frac{v^2}{2}. \]  

Setting the hypercharge \( Y = 1 \), the mass matrix attains the form

\[ M^2 = \frac{1}{4} \begin{pmatrix} g^2 & -gg' & -ggXQ_X \\ -gg' & g'^2 & g'gXQ_X \\ -ggXQ_X & g'gXQ_X & g_X^2 \bar{Q}_X^2 \end{pmatrix} \frac{v^2}{2} \]

with the basis of the gauge bosons being \( (\hat{W}_3^\mu, \hat{B}_\mu, \hat{X}_\mu)^T \). It is useful to notice that \( (s_W, c_W, 0)^T \) is still an eigenvector of the mass matrix with zero eigenvalue, where \( c_W = g/\sqrt{g^2 + g'^2} \) and \( s_W = g'/\sqrt{g^2 + g'^2} \). The other two eigenvectors are orthogonal and produce the mixing matrix

\[ U_w = \begin{pmatrix} s_W & c_W & 0 \\ c_W & -s_W & 0 \\ 0 & 0 & 1 \end{pmatrix}. \]

Then

\[ U_w^T M^2 U_w = \frac{1}{4} \begin{pmatrix} 0 & 0 & 0 \\ 0 & g^2 + g'^2 & -gXQ_X \sqrt{g^2 + g'^2} \\ 0 & -gXQ_X \sqrt{g^2 + g'^2} & g_X^2 \bar{Q}_X^2 \end{pmatrix} \frac{v^2}{2} \Rightarrow \begin{pmatrix} 0 & 0 & 0 \\ 0 & m_Z^2 & 0 \\ 0 & 0 & m_X^2 \end{pmatrix}. \]

As a final step we diagonalize the \( 2 \times 2 \) submatrix using the unitary matrix

\[ U_\alpha = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_\alpha & s_\alpha \\ 0 & -s_\alpha & c_\alpha \end{pmatrix}. \]

The steps so far show that the photon remains massless and has the structure of the electroweak theory

\[ A_\mu = s_W \hat{W}_3^\mu + c_W \hat{B}_\mu. \]
Thus electromagnetism is not affected. The other two physical fields are given in terms of
the original fields

\[ Z_\mu = c_\alpha c_W \hat{W}_3 - c_\alpha s_W \hat{B}_\mu - s_\alpha \hat{X}_\mu, \]
\[ X_\mu = s_\alpha c_W \hat{W}_3 - s_\alpha s_W \hat{B}_\mu + c_\alpha \hat{X}_\mu. \]

(18a)

To leading order they have the masses

\[ m_Z = \frac{1}{2} v \sqrt{g^2 + g'^2 + g_X^2}, \]

(19)

\[ m_X = \frac{1}{2} v_0 g_X Q_X, \]

(20)

and the mixing angle

\[ c_\alpha^2 = \frac{g^2 + g'^2}{g^2 + g'^2 + g_X^2}. \]

(21)

A light \( X_\mu \) boson appears when the VEV \( v_0 \) and/or the coupling constant \( g_X \) is small. A
limit on the mixing angle is obtained from the \( \rho \) parameter which is modified

\[ \frac{m_W^2}{m_Z^2} = \frac{g^2}{g^2 + g'^2 + g_X^2} = c_\alpha^2 c_W. \]

(22)

The experimental value for the ratio

\[ \rho = \frac{m_W^2}{m_Z^2 c_W^2} = c_\alpha^2 = 1.00037 \pm 0.00023 \]

restricts the mixing angle. To within two standard deviations

\[ 0.99991 \leq c_\alpha^2 \leq 1 \]

(23)

restricting \( s_\alpha^2 < 8 \times 10^{-5} \) or \( s_\alpha \leq 9 \times 10^{-3} \approx 0.01 \). A small mass of the order of MeV
is accommodated when the VEV \( v_0 \sim 1 \text{ GeV} \).

We propose that Majorana neutrinos couple to the \( \hat{X}_\mu \) boson with a coupling invariant
under chiral gauge transformations of the group \( U(1)' \). To this end we invert equations (18a)–
(18b) and obtain the initial fields in terms of the physical states

\[ \hat{W}_3 = s_W A_\mu + c_\alpha c_W Z_\mu + s_\alpha c_W X_\mu, \]
\[ \hat{B}_\mu = c_W A_\mu - c_\alpha s_W Z_\mu - s_\alpha s_W X_\mu, \]
\[ \hat{X}_\mu = -s_\alpha Z_\mu + c_\alpha X_\mu. \]

(25a)

(25b)

(25c)
Substituting these original fields into (11) we can select the generators that couple to each physical boson. Then taking the matrix elements between all fermion states one obtains the couplings. We give the couplings of $X_\mu$ to leptons

$$\mathcal{L}_{NC}(X_\mu) = i \left\{ \frac{gs_\alpha}{4c_W} \left[ \bar{u}_\nu \gamma^\mu (1 - \gamma^5) u_\nu - (1 - 4s^2_W)u_e \gamma^\mu u_e + \bar{u}_e \gamma^\mu \gamma^5 u_e \right] + g_X c_\alpha \bar{\Psi}_L \gamma^\mu \frac{Q_X}{2} \Psi_L + g_X c_\alpha \bar{\Psi}_R \gamma^\mu \frac{Q_X}{2} \Psi_R \right\} X_\mu. \quad (26)$$

The last term with $\Psi_R$ is a coupling of Majorana particles. The complete coupling of the Majoranas $\Psi_M$ is

$$\mathcal{L}_{NC}^M = ig_X Q_X \bar{\Psi}_M \gamma^\mu \gamma^5 \Psi_M X_\mu = -ig_X Q_X \bar{\Psi}_M \gamma^\mu \gamma^5 \Psi_M (s_\alpha Z_\mu - c_\alpha X_\mu). \quad (27)$$

The mixing of the gauge bosons introduced a small coupling of the Majorana to the $Z_\mu$ boson and a large coupling to $X_\mu$. The coupling to $Z_\mu$ brings a new decay. The experimental results cannot eliminate a partial decay width of less than 0.5 MeV for each channel. This gives another bound

$$g_X Q_X s_\alpha < 0.06 \quad (28)$$

which is weaker than the bound for $s_\alpha$ from the $\rho$-parameter. In figures 2–4 we will use much smaller values for $g_X Q_X$ and $s_\alpha$ which still give sizable cross sections. The couplings to quarks will be discussed later.

**III. APPEARANCE OF A MAJORANA COMPONENT THROUGH OSCILLATIONS**

Neutrino and antineutrino beams in laboratories are produced in particle or nuclear decays via charged currents. The produced neutrinos are always left-handed and antineutrinos are right-handed. In theories with right-handed neutrinos Majorana couplings are introduced in the mass matrices with the effect that over one oscillation length a Majorana component is generated on the beam. This phenomenon may take place in the reactor experiments where the observed fluxes of antineutrinos at a distance of $\sim 100$ meters are smaller than those expected from the energy produced by the reactors.
For instance, an initial $\nu_e$ beam will develop a $\Psi_k = (N_{R_k} + N^C_{R_k})$ with the probability $[16–18]$

$$P_{\bar{\nu}_e \rightarrow \Psi_k} = -4 \sum_{i<j} U^*_{R_k,i} U_{e,i} U_{R_k,j} U^*_{e,j} \sin^2 \left( \frac{E_i - E_j}{2} \right),$$

(29)

where $U_{R_k,i}$ connects the Majorana state $R_k$ to the state $i$ and $E_i = (p^2 + m_i^2)^{1/2}$. When the difference in energies $E_i - E_j$ is of the order keV or MeV the oscillation length is small. We adopt the point of view that the reduction of neutrino fluxes of the order of 5% originates from an oscillation into Majorana states and describe experiments.

A reaction to search for the Majorana state is

$$\Psi_M(k) + e^- (p) \rightarrow \Psi_M(k') + e^- (p').$$

(30)

This is a very clean channel where the experiments observe the energy and direction of the recoiling electrons. The amplitude for this process represented by the diagram in Fig. is

$$M = -g g_X Q_X \frac{c_s}{4c_W} \bar{\Psi}_M(k') \gamma^\mu \gamma^5 \Psi_M(k) \frac{1}{2m_e^2 - 2E'_e m_e - m_X^2} \bar{u}(p') \gamma_\mu (g_V - g_A \gamma^5) u(p),$$

(31)

where $g_V = -(1 - 4s_W^2)$ and $g_A = 1$. In the scattering the recoiling electrons emerge close to the direction of the incident beam (very small scattering angle, $\theta \sim 2^{\circ}/\sqrt{E_\nu}$, with $E_\nu$ in GeV [19]). The cross section differential in the electron energy has the form

$$\frac{d\sigma}{dE_e'} = \frac{g^2 g'^2 Q_X^2}{32\pi} \left( \frac{c_s}{c_W} \right)^2 (g^2 + g^2_A) \frac{m_e}{(m_X^2 + 2E_e' m_e)^2} \left[ 1 + \left( 1 - \frac{E_e'}{E_\nu} \right)^2 \right],$$

(32)

where $m_e$ is the mass of the electron, $E'_e$ is the energy of the recoiling electron. The cross section decreases as $E'_e \rightarrow E_\nu$ but not as fast as the background. The background originate from the reaction of the electroweak theory

$$\bar{\nu}_e + e^- \rightarrow \bar{\nu}_e + e^-$$

(33)

whose cross section is

$$\frac{d\sigma}{dE'_e} = \frac{G^2_F m_e}{4\pi} \left[ 1 + 2.25 \left( 1 - \frac{E'_e}{E_\nu} \right)^2 \right],$$

(34)
where we have replaced the weak couplings by their limits. We plot in Fig. 2 the two cross sections (32) and (34) as functions of \( x = E'_\nu/E_\nu \) and for various values of the couplings. The new reaction can be identified either as an increase above the Standard Model cross section or as a modification of the distribution of the recoil energy. Figures 3 and 4 show the dependences of the cross section (32) on the \( X \) boson mass and the neutrino incident energy \( E_\nu \), respectively. We see that process (30) competes with the Standard Model background at \( g_X Q_X = 5 \times 10^{-5} \) and \( s_\alpha = 10^{-4} \) when the boson mass is of the order of 10 MeV.

IV. COHERENT SCATTERING ON NUCLEI

The mixing of the gauge bosons brings about couplings of \( X_\mu \) to ordinary leptons and quarks. The couplings to quarks are the same as those occurring in the standard model for \( Z_\mu \), but now weighted by an overall factor \( c_\alpha \) for \( Z_\mu \) and the factor \( s_\alpha \) for \( X_\mu \). Specifically, the couplings to up- and down-quarks are

\[
\mathcal{L}'(X_\mu) = i \frac{g s_\alpha}{4 c_W} \left\{ \bar{u} \gamma^\mu \left[ \left( 1 - \frac{8}{3} s_W^2 \right) - \gamma^5 \right] u \right. \\
\left. + \bar{d} \gamma^\mu \left[ \left( -1 + \frac{4}{3} s_W^2 \right) + \gamma^5 \right] d \right\} X_\mu. \tag{35}
\]

The cross section differential in the recoiling kinetic energy, \( T = E'_N - M_N \), is given by

\[
\frac{d\sigma}{dT} = \frac{\lambda^2 M_N}{\pi} \left[ 2 - \frac{2T}{E_\nu} + \left( \frac{T}{E_\nu} \right)^2 \right] \left( (1 - 4s_W^2)Z - N \right)^2, \tag{36}
\]

where \( M_N \) is the mass of the nucleus, \( E'_N \) is the energy of the recoiling nucleus, \( Z \) and \( N \) are the numbers of protons and neutrons in the target, respectively,

\[
\lambda = gg_X \frac{Q_X}{2 c_W} \frac{c_\alpha s_\alpha}{m_X^2 + 2 TM_N} \frac{1}{M_N}. 
\]

The cross section (36) as a function of \( T \) for the \(^{23}\text{Na}\) and \(^{133}\text{Cs}\) nuclei at \( E_\nu = 10 \text{ MeV} \), \( g_X Q_X = 5 \times 10^{-5} \) and \( s_\alpha = 10^{-4} \) is shown in Fig. 6. For \(^{127}\text{I}\), this will be of the same order...
as for $^{133}$Cs since the numbers of neutrons in both nuclei are close to each other. Note that the cross section is almost independent on $E_\nu$ as long as $T \ll E_\nu$.

V. SUMMARY

We investigated a model in which Majorana neutrinos couple directly to a gauge boson. For this purpose we enlarged the group of the electroweak theory to the group $SU(2) \times U(1) \times U(1)'$ with couplings for the Majorana neutrinos in the group $U(1)'$. In addition, we introduced a single Higgs particle $\chi(x)$, which mixes the $U(1)$ and $U(1)'$ gauge bosons in the mass matrix, and derived formulas for the masses and mixing angles.

The theory preserves the good properties of the electroweak theory and makes several interesting new predictions.

1. The electromagnetic interactions are not affected by the $U(1)'$ gauge boson.

2. The Majorana field $\Psi_M$ being a singlet under $SU(2)$ does not initially have a coupling to the charged gauge bosons. However, a new Yukawa coupling to the singlet Higgs is allowed

$$g'' \bar{\Psi}_M \chi(x) \Psi_M$$

whose VEV $v_0$ generates a Majorana term in the fermion mass matrix. As usually, the $(1,1)$ matrix element is zero but there are Dirac and Majorana matrix elements breaking the $SU(2)$ symmetry softly, thus producing the oscillations discussed in section III. Starting with a $\nu_\mu$ beam a small component of $\Psi_M$ is generated downstream at distances of a few hundred meters provided energy differences, like $(E_3 - E_2)$, are large. Elastic scattering on atomic electrons is possible but it will have a sharp peak in the forward direction. Other processes with more complicated diagrams [24–26] are also possible which have different angular distributions. They can be enhanced by the small masses in the propagators. They will be the subject for further investigation.

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Appendix A: The amplitude for neutrino–nucleus scattering

The amplitude for the process shown in Fig. 5 is

\[ M_c = -g g_X \frac{c_A s_A}{4 c_W} \overline{\Psi}_M \gamma^\mu \gamma^5 Q_X \frac{i}{q^2 - m_X^2} \bar{u}(p') \gamma_\mu u(p) \left[ (1 - 4 s_W^2) Z - N \right], \]  

(A1)

where \( q^2 = (p - p')^2 = 2 M_N^2 - 2(p \cdot p') = 2 M_N^2 - 2 M_N E_N' = 2 M_N (M_N - E_N'). \)

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FIG. 1. Neutral current Majorana neutrino–electron scattering with $X - Z$ mixing.
FIG. 2. The cross section as a function of the ratio $x = E'_e/E_\nu$ at two values of the X boson mass, 10 MeV and 15 MeV. The incident neutrino energy is fixed to be 100 MeV. The other parameters are $g_X Q_X = 5 \times 10^{-5}$, $s_\alpha = 10^{-4}$, $c_\alpha = \sqrt{1 - s_\alpha^2}$. The Standard Model (SM) cross section is also shown.
FIG. 3. Dependence of the cross section \((32)\) on the \(X\) boson mass at two values of the incident neutrino energy, 10 MeV and 1000 MeV. The ratio \(x = E'_\nu/E_\nu = 0.9\). The other parameters are \(g_X Q_X = 5 \times 10^{-5}, s_\alpha = 10^{-4}, c_\alpha = \sqrt{1 - s_\alpha^2}\).
FIG. 4. Dependence of the cross section (32) on the incident neutrino energy at two values of the X boson mass, 10 MeV and 100 MeV. The ratio $x = E'_\nu / E_\nu = 0.9$. The other parameters are $g_X Q_X = 5 \times 10^{-5}$, $s_\alpha = 10^{-4}$, $c_\alpha = \sqrt{1 - s_\alpha^2}$.
FIG. 5. Coherent elastic Majorana neutrino–nucleus scattering with $X - Z$ mixing.
FIG. 6. Dependence of the cross section (36) on the kinetic energy of the recoiling nuclei at the incident neutrino energy $E_\nu = 10$ MeV and two values of the $X$ boson mass, 10 MeV and 100 MeV. The other parameters are $g_X Q_X = 5 \times 10^{-5}$, $s_\alpha = 10^{-4}$, $c_\alpha = \sqrt{1 - s_\alpha^2}$. 