Neutrinos and Electromagnetic Gauge Invariance

F. Pisano*, J.A. Silva-Sobrinho

Instituto de Física Teórica, Universidade Estadual Paulista

Rua Pamplona, 145 – 01405-000 – São Paulo, SP

Brazil

M.D. Tonasse

Instituto de Física, Universidade do Estado do Rio de Janeiro,

Rua São Francisco Xavier, 524 – 20550-013 – Rio de Janeiro, RJ

Brazil

Abstract

It is discussed a recently proposed connection among $U(1)_{\text{em}}$ electromagnetic gauge invariance and the nature of the neutrino mass terms in the framework of $SU(3)_C \otimes G_W \otimes U(1)_N$, $G_W = SU(3)_L$, extensions of the Standard Model. The impossibility of that connection, also in the extended case $G_W = SU(4)_L$, is demonstrated.

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The neutrinos of the Standard Model are Weyl fermions. The pairs of states $(\nu_{lL}, \bar{\nu}_{lR})$, $l = \{e, \mu, \tau\}$, are the sole fermionic neutral non-massive degrees of freedom of the model. However, in a large class of extensions of the Standard Model the neutrino can be a massive fermion. A mass term for the neutrino can be either Dirac or Majorana, the last one being additive violating quantum numbers. If the theory does not contain pairs $(\nu_{lR}, \nu_{lL})$ a Dirac mass term is absent, but anyway a Majorana mass term is forbidden if the lepton number is a symmetry of nature.

Here we analyze a result contemplated by Özer [1] which can provide a connection between neutrino mass terms and the $U(1)_{em}$ electromagnetic gauge invariance. Such a connection comes out in a class of models with gauge symmetry

$$G_0 \equiv SU(3)_C \otimes G_W \otimes U(1)_N.$$  

(1)

where $G_W \equiv SU(3)_L, SU(4)_L$ are the $SU(2)_L$ weak isospin extended groups. Different representation contents are determined by embedding the electric charge operator

$$\frac{Q}{e} = \frac{1}{2} \left( \lambda_3^L + \xi \lambda_8^L + \zeta \lambda_{15}^L \right) + N$$  

(2)

in the neutral generators $\lambda_{3,8,15}^L$ of $G_W$. The parameters $\xi$ and $\zeta$ distinguish different embeddings, fermionic contents, and the respective flavordynamics. A salient feature in this class of models concerns the anomaly cancellation procedure which is carried out among fermion families, color degrees of freedom, and different group transformation properties under $G_0$ of the matter field multiplets. As a remarkable result, an explanation for the family problem arises.

The symmetry breaking pattern

$$G_0 \rightarrow SU(3)_C \times SU(2)_L \times U(1)_Y \rightarrow SU(3)_C \times U(1)_{em}$$  

(3)

and masses are achieved by introducing triplets and sextets of scalar fields when $G_W = SU(3)_L$, or 4-plets and 10-plets for $G_W = SU(4)_L$.

Let us consider the case in which $G_W = SU(3)_L$ and therefore $\zeta = 0$ (which we will call 331 models). In general, three triplets, $\phi^a$, and one sextet, $S$, constitute the sufficient
set of Higgs fields needed to break the original gauge symmetry and to generate masses. Substituting the Gell-Mann matrices

\[ \lambda_3^L = \text{diag} (+1, -1, 0), \quad \lambda_8^L = \frac{1}{\sqrt{3}} \text{diag} (+1, +1, -2) \]

in the Gell-Mann-Nishijima relation of Eq. (2) we obtain the electric charge of the fields contained in the triplets \( \phi^a \sim (1, 3, N_{\phi^a}) \), \( a = 1, 2, 3 \),

\[
Q(\phi^a) = \begin{pmatrix}
\frac{1}{2}(1 + \xi \sqrt{3}) + N_{\phi^a} \\
\frac{1}{2}(-1 + \xi \sqrt{3}) + N_{\phi^a} \\
\frac{1}{2}(-2 \xi \sqrt{3}) + N_{\phi^a}
\end{pmatrix}
\]

in unities of the proton charge. Performing the Kronecker product \( 3 \otimes 3 = 3^* \oplus 6 \) we obtain also the electric charges of the fields contained in the symmetric sextet \( S \sim (1, 6, N_S) \),

\[
Q(S) = \begin{pmatrix}
1 + \xi \sqrt{3} + N_S & \frac{\xi}{\sqrt{3}} + N_S & \frac{1}{2}(1 - \xi \sqrt{3}) + N_S \\
\frac{\xi}{\sqrt{3}} + N_S & -1 + \xi \sqrt{3} + N_S & -\frac{1}{2}(1 + \xi \sqrt{3}) + N_S \\
\frac{1}{2}(1 - \xi \sqrt{3}) + N_S & -\frac{1}{2}(1 + \xi \sqrt{3}) + N_S & -\frac{2 \xi}{\sqrt{3}} + N_S
\end{pmatrix}
\]

where

\[
N_S = 2 N_{\phi^a},
\]

so, we have only three parameters, namely the embedding parameter \( \xi \) and the U(1)\(_N\) charges \( N_{\phi^a} \) and \( N_S \), in order to know which components can develop vacuum expectation values.

Mass matrices for the electroweak neutral gauge bosons decoupling from gluons can be constructed starting with the Lagrangian

\[
\mathcal{L}_{\phi^a, S} = (\mathcal{D}_\mu \phi^a_1)\dagger (\mathcal{D}^\mu \phi^a_1) + [(\mathcal{D}_\mu S_{mn})\dagger (\mathcal{D}^\mu S_{mn})],
\]

with the gauge covariant derivatives

\[
\mathcal{D}_\mu \phi^a_1 = \partial_\mu \phi^a_i - i g_L \left( \vec{W}_\mu \cdot \vec{\lambda}^L / 2 \right)_i^j \phi^a_j - i g_N N_{\phi^a} B_\mu \phi^a_i
\]

\[
\mathcal{D}_\mu S_{mn} = \partial_\mu S_{mn} - i g_L \left[ (\vec{W}_\mu \cdot \vec{\lambda}^L / 2)_m^k S_{kn} + (\vec{W}_\mu \cdot \vec{\lambda}^L / 2)_n^k S_{km} \right] - i g_N N_S B_\mu S_{mn}
\]
where \( \{ \vec{W}_\mu \} \) and \( B_\mu \) are the octet and a singlet of gauge bosons associated to \( SU(3)_L \) and \( U(1)_N \) groups and we have denoted the respective gauge coupling constants by \( g_L \) and \( g_N \). The diagonalization of the neutral gauge boson mass matrix includes the photon as a linear combination of the \( W_{\mu}^3 \), \( W_{\mu}^8 \) and \( B_\mu \) states. If we want to preserve the \( U(1)_{\text{em}} \) local invariance, the inclusion of a mass term for the photon

\[
\frac{1}{2} M^2 A^\mu A^\mu
\]

in the QED Lagrangian is not allowed since

\[
A^\mu A^\mu \to [A^\mu - \partial_\mu \alpha(x)][A^\mu - \partial^\mu \alpha(x)] \neq A^\mu A^\mu,
\]

then we are restricted to a singular mass matrix for the electroweak gauge bosons. Hence, we demand that the contributions to the photon mass in the covariant derivatives coming from any scalar field vanish. According to Ref. [1] this requirement leads to the following conditions:

if \( \langle S_{11} \rangle \neq 0 \) and \( N_S \neq 0 \), then \( \langle S_{22} \rangle = 0 \) and \( 1 + \frac{\xi}{\sqrt{3}} + N_S = 0 \), \( \text{(10)} \)

if \( \langle S_{22} \rangle \neq 0 \) and \( N_S \neq 0 \), then \( \langle S_{11} \rangle = 0 \) and \( -1 + \frac{\xi}{\sqrt{3}} + N_S = 0 \), \( \text{(11)} \)

if \( N_S = 0 \), then \( \langle S_{11} \rangle = \langle S_{22} \rangle = 0 \). \( \text{(12)} \)

In a \( G_0 \) model with \( G_W = SU(3)_L \) the most general gauge boson mass matrix is

\[
\frac{1}{2} M^2 = \frac{g_L^2}{4} \begin{pmatrix} m_{11} & m_{12} & m_{13} \\ m_{12} & m_{22} & m_{23} \\ m_{13} & m_{23} & m_{33} \end{pmatrix}
\] \( \text{(13)} \)

where

\[
m_{11} = \langle \phi_i^a \rangle^2 + \langle \phi_j^a \rangle^2 + 2 \left( 2 \langle S_{11} \rangle^2 + \langle S_{13} \rangle^2 + 2 \langle S_{22} \rangle^2 + \langle S_{23} \rangle^2 \right) \]

\[
m_{22} = \frac{1}{3} \left[ \langle \phi_i^a \rangle^2 + \langle \phi_j^a \rangle^2 + 4 \langle \phi_k^a \rangle^2 \right]
\]

\[
m_{12} = m_{13} = m_{23}
\]
\[ \rho = (t \text{ with form under } 331 \text{ as where } a \text{ to } l \text{ gives mass to neutrinos. The gauge boson mass matrix elements are whose neutral components can develop a vacuum expectation value. A nonvanishing } \rho \xi a L. \text{ The scalar sector is the set of three SU}(3) \equiv m = 1, \rho^0 = 2^{23}, \xi a L. \text{ Leptons transform under 331 as } \Psi_{aL} = (\nu_a, \ l_a, \ l_a^c)^T \sim (1, 3, 0) \text{ (20)} \]

where \( a = 1, 2, 3 \) is the family index and \( l_a^c \) is the charge conjugate field corresponding to \( l_a \). The scalar sector is the set of three SU(3) triplets \( \eta = (\eta^0, \ \eta^-, \ \eta^+)^T \sim (1, 3, 0) \), \( \rho = (\rho^+, \ \rho^0, \ \rho^{++})^T \sim (1, 3, +1) \), \( \chi = (\chi^-, \ \chi^-, \ \chi^0)^T \sim (1, 3, -1) \) and the sextet

\[ S = \begin{pmatrix} \sigma^0_1 & h_2^- & h_1^+ \\ h_2^- & H_1^- & \sigma^0_2 \\ h_1^+ & \sigma^0_2 & H_2^{++} \end{pmatrix} \sim (1, 6, 0) \text{ (21)} \]

whose neutral components can develop a vacuum expectation value. A nonvanishing \( \langle \sigma^0_1 \rangle \) gives mass to neutrinos. The gauge boson mass matrix elements are

\[ m_{11} = \langle \eta \rangle^2 + \langle \rho \rangle^2 + 4\langle \sigma_1 \rangle^2 + 2\langle \sigma_2 \rangle^2 \text{ (22)} \]
\[ m_{22} = \frac{1}{3} \left( \langle \eta \rangle^2 + \langle \rho \rangle^2 + 4\langle \chi \rangle^2 + 4\langle \sigma_1 \rangle^2 + 2\langle \sigma_2 \rangle^2 \right) \text{ (23)} \]
\[ m_{33} = 4t^2 \left( \langle \rho \rangle^2 + \langle \chi \rangle^2 \right) \text{ (24)} \]
\[ m_{12} = \frac{1}{\sqrt{3}} \left( \langle \eta \rangle^2 - \langle \rho \rangle^2 + 4\langle \sigma_1 \rangle^2 + 2\langle \sigma_2 \rangle^2 \right) \text{ (25)} \]
\[ m_{13} = -2t\langle \rho \rangle^2 \text{ (26)} \]
\[ m_{23} = \frac{2}{\sqrt{3}} t \left( \langle \rho \rangle^2 + 2\langle \chi \rangle^2 \right) \text{ (27)} \]
This is a singular matrix even if $\langle S_{11} \rangle = \langle \sigma_1^0 \rangle \neq 0$.

The scalar potential contains terms as

$$V(\eta, \rho, \chi, S) = ... + f_1 \eta^\dagger S \eta^* + f_2 \det S + f_3 \varepsilon_{ijk}(S \eta^*)_i \rho_j \chi_k + f_4 \varepsilon_{i\delta j \kappa \ell} S_{il} S_{jm} \rho_k \chi_n.$$  \hspace{1cm} (28)

If we wish to avoid tree level neutrino masses, $\langle \sigma_1^0 \rangle = 0$, we have a fine-tuning among the $f_1, ..., 4$ coupling constants. In order to avoid fine tuning between the coupling constants of the most general Higgs potential $V(\eta, \rho, \chi, S)$ and maintain lepton and baryon number conservation the following discrete symmetries must be introduced if $\langle \sigma_1^0 \rangle = 0$ [3]:

$$Q_{1L} \rightarrow -Q_{1L}, \quad \eta \rightarrow -\eta,$$

$$Q_{jL} \rightarrow -iQ_{jL}, \quad \rho, \chi \rightarrow i\rho, \chi,$$

$$\Psi_{1L} \rightarrow i\Psi_{1L}, \quad S \rightarrow -S,$$

$$u_{jR} \rightarrow u_{jR}, \quad J_{1R} \rightarrow iJ_{1R},$$

$$d_{jR} \rightarrow id_{jR}, \quad J_{2,3R} \rightarrow J_{2,3R},$$

where $Q_{1L}$, $Q_{jL}$ ($j = 2, 3$) are SU(3)$_L$ triplets of quarks, and $J_{1,2,3R}$ are the corresponding right-handed exotic quark singlets [3]. The leptonic Yukawa couplings involving the Higgs sextet $S \sim (1, 6, 0)$ allowed by the gauge invariance have the general form

$$\mathcal{L}_{l,S} = -\frac{1}{2} \sum_{l,m} G_{lm} \bar{\Psi}_{il} \Psi_{jm} S_{ij} + \text{H.c.}$$  \hspace{1cm} (30)

where $i, j$ denote SU(3) indices, $l, m = e, \mu, \tau$, and $G_{lm} = G_{ml}$. In the $\xi = -\sqrt{3}$ model these couplings for the neutrinos are

$$\mathcal{L}_{\nu,S} = -\frac{1}{2} \sum_{l,m} G_{lm} (\bar{\nu}_{iR} \nu_{mL} S_{11} + \text{H.c.}).$$  \hspace{1cm} (31)

Then Majorana mass terms are not allowed if $\langle S_{11} \rangle = \langle \sigma_1^0 \rangle = 0$. Notwithstanding such condition is entirely independent and does not affect the mass spectrum pattern of the gauge sector.

In the $\xi = 1/\sqrt{3}$ model [3], with a symmetric sextet $S \sim (1, 6, +2/3)$, the Yukawa couplings are
\[ \mathcal{L}_{\nu, S} = -\frac{1}{2} \sum_{l,m} G_{lm} \left( \bar{\nu}_{lR} \nu_{mL} S_{11} + \bar{\nu}_{lR} \nu_{mL} S_{12} + \bar{\nu}_{lR} \nu_{mL} S_{13} + \bar{\nu}_{lR} \nu_{mL} S_{22} + \text{H.c.} \right). \] (32)

The sufficient scalar sector of the \( \xi = 1/\sqrt{3} \) model contains the set of three triplets \( \eta = (\eta^0, \eta_1^-, \eta_2^+) \sim (1, 3, -2/3)^T \), \( \rho = (\rho^+, \rho_1^0, \rho_2^0) \sim (1, 3, 1/3)^T \), and \( \chi = (\chi^+, \chi_1^0, \chi_2^0) \sim (1, 3, 1/3)^T \) with the same transformation properties as the \( \rho \) multiplet. The corresponding gauge boson mass matrix given in Eq. (13) including the sextet

\[ S = \begin{pmatrix} S_0^0 & S_1^0 & S_1^- \\ S_1^0 & S_2^0 & S_2^- \\ S_1^- & S_2^- & S_-^-- \end{pmatrix} \sim (1, 6, +2/3) \] (33)

has the elements

\[ m_{11} = \langle \eta \rangle^2 + \langle \rho_1 \rangle^2 + \langle \chi_1 \rangle^2 + 4\langle S_1 \rangle^2 + \langle S_3 \rangle^2 \] (34)
\[ m_{22} = \frac{1}{3} \left[ m_{11} + 4(\langle \rho_2 \rangle^2 + \langle \chi_2 \rangle^2 + 2\langle S_2 \rangle^2) + 3\langle S_3 \rangle^2 \right] \] (35)
\[ m_{33} = \frac{4}{9} t^2 \left[ m_{11} + 3\langle \eta \rangle^2 + \langle \rho_2 \rangle^2 + 8\langle S_2 \rangle^2 + \langle \chi_2 \rangle^2 + 3\langle S_3 \rangle^2 \right] \] (36)
\[ m_{12} = \frac{1}{\sqrt{3}} (\langle \eta \rangle^2 - \langle \rho_1 \rangle^2 - \langle \chi_1 \rangle^2 + \langle S_1 \rangle^2 - \langle S_3 \rangle^2) \] (37)
\[ m_{13} = \frac{4}{3} t (-\langle \eta \rangle^2 - \langle \rho_1 \rangle^2 - \langle \chi_1 \rangle^2 + 2\langle S_1 \rangle^2 - 2\langle S_3 \rangle^2) \] (38)
\[ m_{23} = \frac{4}{3\sqrt{3}} t (-\langle \eta \rangle^2 + \langle \rho_1 \rangle^2 - \langle \rho_2 \rangle^2 - \langle \chi_2 \rangle^2 + 2\langle S_1 \rangle^2 + 4\langle S_2 \rangle^2 + 2\langle S_3 \rangle^2). \] (39)

This matrix is singular only if hold the conditions \( \langle S_{11} \rangle = \langle S_{12} \rangle = \langle S_{22} \rangle = \langle \chi_1 \rangle = 0 \), so neutrinos are massless at tree level. Another possible Yukawa couplings in this model are

\[ \mathcal{L}_{l, \eta} = -\frac{1}{2} \sum_{l,m} \epsilon^{ijk} h_{lm} \Psi_{il} C^{-1} \Psi_{jm} \eta_k^* \] (40)

with \( C \) being the charge conjugation operator and \( h_{lm} = -h_{ml} \) which implies an antisymmetric \( 3 \times 3 \) mass matrix for the neutrinos. Hence, there are one massless and two mass degenerate neutrino states, at least at tree level.

For the \( \xi = -1/\sqrt{3} \), \( \zeta = -2\sqrt{6}/3 \) model with \( G_W = SU(4)_L \) the neutrino Yukawa couplings are

\[ \mathcal{L}_{\nu, H} = -\frac{1}{2} \sum_{l,m} G_{lm} (\bar{\nu}^c_{lR} \nu_{mL} H_{11} + \bar{\nu}_{lR} \nu_{mL} H_{13} + \bar{\nu}^c_{lR} \nu_{mL} H_{13} + \bar{\nu}_{lR} \nu_{mL} H_{33} + \text{H.c.}) \] (41)
where the $H_{ij}$, $i, j = 1, \ldots, 4$ are the elements of the $SU(4)$ symmetric $10$-plet

$$H = \begin{pmatrix} H_{1}^{0} & H_{1}^{+} & H_{2}^{0} & H_{2}^{-} \\ H_{1}^{+} & H_{1}^{++} & H_{3}^{+} & H_{3}^{0} \\ H_{2}^{0} & H_{3}^{+} & H_{4}^{0} & H_{4}^{-} \\ H_{2}^{-} & H_{3}^{0} & H_{4}^{-} & H_{2}^{-} \end{pmatrix} \sim (1, 10, 0). \quad (42)$$

The vacuum structure $\langle H_{3}^{0} \rangle \neq 0$, $\langle H_{1,2,4}^{0} \rangle = 0$ is sufficient for giving a finite mass to the charged leptons but neutrinos remain massless at tree level.

The mass matrix elements of the gauge bosons in the model with $G_{W} = SU(4)$, up to a factor $g_{SU(4)}^{2}/4$, are

$$m_{11} = \langle \eta \rangle^{2} + \langle \eta' \rangle^{2} + \langle \rho \rangle^{2} + 2 \left( 2\langle H_{1} \rangle^{2} + \langle H_{2} \rangle^{2} + \langle H_{3} \rangle^{2} \right) \quad (43)$$

$$m_{22} = \frac{1}{3} \left[ m_{11} + 4 \left( \langle \eta_{2} \rangle^{2} + \langle \eta'_{2} \rangle^{2} + 4\langle H_{4} \rangle^{2} \right) \right] \quad (44)$$

$$m_{33} = \frac{1}{6} \left[ m_{11} + \langle \eta_{2} \rangle^{2} + \langle \eta'_{2} \rangle^{2} + 9\langle \chi \rangle^{2} + 6 \left( \langle H_{2} \rangle^{2} + \langle H_{3} \rangle^{2} \right) + 4\langle H_{4} \rangle^{2} \right] \quad (45)$$

$$m_{44} = 4t^{2} \left( \langle \rho \rangle^{2} + \langle \chi \rangle^{2} \right) \quad (46)$$

$$m_{12} = \frac{1}{\sqrt{3}} \left[ \langle \eta_{1} \rangle^{2} + \langle \eta'_{1} \rangle^{2} - \langle \rho \rangle^{2} + 2 \left( 2\langle H_{1} \rangle^{2} - \langle H_{2} \rangle^{2} - \langle H_{3} \rangle^{2} \right) \right] \quad (47)$$

$$m_{13} = \frac{1}{\sqrt{6}} \left[ \langle \eta_{1} \rangle^{2} + \langle \eta'_{1} \rangle^{2} - \langle \rho \rangle^{2} + 4 \left( \langle H_{1} \rangle^{2} + \langle H_{2} \rangle^{2} + \langle H_{3} \rangle^{2} \right) \right] \quad (48)$$

$$m_{14} = -2t\langle \rho \rangle^{2} \quad (49)$$

$$m_{23} = \frac{1}{3\sqrt{2}} \left[ \langle \eta_{1} \rangle^{2} + \langle \eta'_{1} \rangle^{2} - 2 \left( \langle \eta_{2} \rangle^{2} + \langle \eta'_{2} \rangle^{2} \right) + \langle \rho \rangle^{2} + 4 \left( \langle H_{1} \rangle^{2} - \langle H_{2} \rangle^{2} - \langle H_{3} \rangle^{2} - 2\langle H_{4} \rangle^{2} \right) \right] \quad (50)$$

$$m_{24} = \frac{2}{\sqrt{3}} t \langle \rho \rangle^{2} \quad (51)$$

$$m_{34} = \frac{2}{\sqrt{6}} t \left( \langle \rho \rangle^{2} + 3\langle \chi \rangle^{2} \right) \quad (52)$$

where $\eta \sim (1, 4, 0)$, $\eta' \sim (1, 4, 0)$, $\rho \sim (1, 4, +1)$, and $\chi \sim (1, 4, -1)$. This is the most general $G_{W} = SU(4)_{L}$ electroweak gauge boson mass matrix. As it can be checked, it is a singular matrix for any $\langle H_{i}^{0} \rangle \neq 0$. None of the $H$ matrix elements imposes any constraint on the neutrino Yukawa couplings of Eq. (11).

We outline our results. The constraint given in Eq. (12) is not realized in the $G_{0}$ models. This is true also in the 331 model of Ref. [6] which is equivalent to the model of Ref. [2].
In fact these two representations are identical since we can map from one representation to another by a unitary transformation \[ U \]. It can be checked that condition in Eq. (11) does not hold in the \( \xi = 1/\sqrt{3} \) model since for the nonvanishing vacuum expectation value of any neutral component of the sextet with \( N_S = 2/3 \) the electromagnetic gauge invariance is not preserved. As we have shown, also in the \( G_W = SU(4)_L \) extensions, there is no connection among neutrino masses and electromagnetic gauge invariance. However, a remarkable fact is that a possible constraint on neutrino masses in the \( \xi = -\sqrt{3} \) model can be arisen from the scalar potential through the set of discrete symmetries of Eqs. (29) in order to avoid a fine tuning among the scalar potential coupling constants which is consistent for \( \langle \sigma_1 \rangle = 0 \).

For concluding we wish to point out that in a large class of gauge models which contain the right-handed neutrino the electric charge quantization can be obtained only if the neutrino is a Majorana particle \[ 8 \]. This fact is true even for the Standard Model enlarged for containing a right-handed neutrino. Thus in the \( G_0 \) models, since the conditions of Eq. (12) are not true, there is no contradiction with this fact.
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