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Universal seesaw from left–right and Peccei–Quinn symmetry breaking

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ABSTRACT

To generate the lepton and quark masses in the left–right symmetric models, we can consider a universal seesaw scenario by integrating out heavy fermion singlets which have the Yukawa couplings with the fermion and Higgs doublets. The universal seesaw scenario can also accommodate the leptogenesis with Majorana or Dirac neutrinos. We show that the fermion singlets can obtain their heavy masses from the Peccei–Quinn symmetry breaking.

1. Introduction

In the $SU(3)_c \times SU(2)_L \times U(1)_Y$ standard model (SM), the charged fermion singlets have Yukawa interactions with the fermion and Higgs doublets so that the quarks and the charged leptons can obtain their Dirac masses after the Higgs doublet develops its vacuum expectation value (VEV). However, we cannot get the neutrino masses in this way because the right-handed neutrinos are absent in the SM. To naturally generate the neutrino masses, which are far lower than the charged fermion masses, we can consider the seesaw [1] extension of the SM by introducing the right-handed neutrino singlets with heavy Majorana masses [1] and/or the left-handed Higgs triplet(s) with small VEV(s) [2].

The SM and its seesaw extension can be embedded in the $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ left–right symmetric models [3], where the left-(right-)handed fermions are placed in $SU(2)_L$ [$SU(2)_R$] doublets. For the breakdown of $SU(2)_R \times U(1)_{B-L}$ to $U(1)_Y$, the simplest choice is to introduce a right-handed Higgs doublet and its left-handed partner. The fermion doublets associated with the additional fermion singlets can have the Yukawa couplings with the Higgs doublets [4–7]. Subsequently we can integrate out the fermion singlets to generate the masses of the quarks, the charged leptons, and the neutral neutrinos. In this scenario, the seesaw is a universal origin of the fermion masses. The universal seesaw can also accommodate the leptogenesis [8] mechanism with Majorana [8] or Dirac [9] neutrinos to explain the matter–antimatter asymmetry in the universe. The key point of the universal seesaw is the existence of the heavy fermion singlets, including the color singlets for generating the lepton masses and the color triplets for generating the quark masses. The heavy masses of the fermion singlets can be simply input by hand as they are allowed by the gauge symmetry. A more interesting possibility is that the fermion singlets obtain their masses through certain spontaneous symmetry breaking. For example, the original work on the universal seesaw introduced a new $SU(3)$ gauge symmetry [4].

In this Letter we shall consider a spontaneous Peccei–Quinn [10] (PQ) symmetry breaking to generate the heavy masses of the fermion singlets for the universal seesaw. Specifically we shall impose a global symmetry under which the fermion doublets have no charges while the left- and right-handed fermion singlets carry equal but opposite charges. Accordingly, the left- and right-handed Higgs doublets also carry equal but opposite charges due to their Yukawa couplings with the fermion singlets. In this context, the left–right symmetry should be the charge-conjugation. The spontaneous symmetry breaking of the global symmetry is driven by a complex scalar singlet. The fermion singlets can obtain the heavy masses through their Yukawa couplings with the complex scalar singlet. The Nambu–Goldstone boson (NGB), associated with the global symmetry breaking, can become a pseudo NGB (pNGB) as it picks up a tiny mass through the color anomaly [11]. Since the global symmetry is mediated to the SM quarks at tree and two-loop level, it can be identified with the PQ symmetry to solve the strong CP problem. Consequently, the pNGB should be an invisible [12,13] axion [10,14].

2. The model

The Higgs scalars include two left-handed doublets $\phi_{L_{1,2}}(\mathbf{1}, \mathbf{2}, \mathbf{1}, -1)$, two right-handed doublet $\phi_{R_{1,2}}(\mathbf{1}, \mathbf{1}, \mathbf{2}, -1)$, a real singlet

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$\sigma(\mathbf{1},\mathbf{1},\mathbf{1},0)$ and a complex singlet $\xi(\mathbf{1},\mathbf{1},\mathbf{1},0)$. In the fermion sector, besides the usual quark and lepton doublets, $q_L(\mathbf{3},\mathbf{2},\mathbf{1},\frac{1}{3})$, $q_R(\mathbf{3},\mathbf{1},\mathbf{2},\frac{1}{3})$, $l_L(\mathbf{1},\mathbf{2},\mathbf{1},-1)$, and $l_R(\mathbf{1},\mathbf{1},\mathbf{2},-1)$, there are four types of fermion singlets, $D_{L,R}(\mathbf{3},\mathbf{1},\mathbf{1},-\frac{2}{3})$, $U_{L,R}(\mathbf{3},\mathbf{1},\mathbf{1},\frac{4}{3})$, $E_{L,R}(\mathbf{1},\mathbf{1},\mathbf{1},-2)$, and $N_{L,R}(\mathbf{1},\mathbf{1},\mathbf{1},0)$. The discrete left–right symmetry is taken to be the charge-conjugation, which yields

$$\begin{aligned}&\sigma\leftrightarrow-\sigma, && \xi\leftrightarrow\xi, && \phi_{L_i}\leftrightarrow\phi^*_{R_i},\\ &q_L\leftrightarrow q^c_R, && D_L\leftrightarrow D^c_R, && U_L\leftrightarrow U^c_R,\\ &l_L\leftrightarrow l^c_R, && E_L\leftrightarrow E^c_R, && N_L\leftrightarrow N^c_R.\end{aligned}\tag{1}$$

We further impose a global symmetry under which the fields carry the following quantum numbers,

$$\begin{aligned}&0 \quad \text{for } \sigma, q_L, q^c_R, l_L, l^c_R,\\ &1 \quad \text{for } D_L, D^c_R, U_L, U^c_R, \phi_{L_1}, \phi^*_{R_1}, \phi^*_{L_2}, \phi_{R_2},\\ &2 \quad \text{for } \xi,\\ &x \quad \text{for } E_L, E^c_R \quad (x=\pm1),\\ &y \quad \text{for } N_L, N^c_R \quad (y=\pm1).\end{aligned}\tag{2}$$

Clearly, the above assignment is consistent with the discrete left–right symmetry (1).

We write down the full scalar potential as below,

$$\begin{aligned}V={}&\frac{1}{2}\mu^2_\sigma\sigma^2+\frac{1}{4}\lambda_\sigma\sigma^4+\mu^2_\xi|\xi|^2+\lambda_\xi|\xi|^4\\ &+\rho^2_i\big(|\phi_{L_i}|^2+|\phi_{R_i}|^2\big)+\kappa_i\big(|\phi_{L_i}|^4+|\phi_{R_i}|^4\big)\\ &+\alpha_{12}\big(|\phi_{L_1}|^2|\phi_{L_2}|^2+|\phi_{R_1}|^2|\phi_{R_2}|^2\big)\\ &+\alpha'_{12}\big(\phi^\dagger_{L_1}\phi_{L_2}\phi^\dagger_{L_2}\phi_{L_1}+\phi^\dagger_{R_1}\phi_{R_2}\phi^\dagger_{R_2}\phi_{R_1}\big)\\ &+\alpha''_{12}\big(\phi^T_{L_1}\tau_2\phi_{L_2}\phi^\dagger_{L_2}\tau_2\phi^*_{L_1}+\phi^T_{R_1}\tau_2\phi_{R_2}\phi^\dagger_{R_2}\tau_2\phi^*_{R_1}\big)\\ &+\beta_{ij}|\phi_{L_i}|^2|\phi_{R_j}|^2+\big(\beta'_{12}\phi^\dagger_{L_1}\phi_{L_2}\phi^\dagger_{R_1}\phi_{R_2}\\ &+\beta''_{12}\phi^T_{L_1}\tau_2\phi_{L_2}\phi^\dagger_{R_1}\tau_2\phi^*_{R_2}+\text{H.c.}\big)\\ &+\gamma\sigma^2|\xi|^2+\eta_i\sigma\big(|\phi_{L_i}|^2-|\phi_{R_i}|^2\big)\\ &+\varepsilon_i\sigma^2\big(|\phi_{L_i}|^2+|\phi_{R_i}|^2\big)+\zeta_i|\xi|^2\big(|\phi_{L_i}|^2+|\phi_{R_i}|^2\big)\\ &+\big[\omega\xi\big(\phi^\dagger_{L_1}\phi_{L_2}+\phi^T_{R_1}\phi^*_{R_2}\big)+\text{H.c.}\big].\end{aligned}\tag{3}$$

The allowed Yukawa interactions should be

$$\begin{aligned}\mathcal{L}_Y={}&-y_D\big(\bar{q}_L\tilde{\phi}_{L_2}D_R+\bar{q}^c_R\tilde{\phi}^*_{R_2}D^c_L\big)-h_D\xi\bar{D}_LD_R\\ &-y_U\big(\bar{q}_L\phi_{L_1}U_R+\bar{q}^c_R\phi^*_{R_1}U^c_L\big)-h_U\xi\bar{U}_LU_R\\ &-y_E\big(\bar{l}_L\tilde{\phi}_{L_x}E_R+\bar{l}^c_R\tilde{\phi}^*_{R_x}E^c_L\big)-h_E\xi_x\bar{E}_LE_R\\ &-y_{N_R}\big(\bar{l}_L\phi_{L_y}N_R+\bar{l}^c_R\phi^*_{R_y}N^c_L\big)\\ &-y_{N_L}\big(\bar{l}_L\phi_{L_y}N^c_L+\bar{l}^c_R\phi^*_{R_y}N_R\big)\\ &-\left[f^D_N\bar{N}_LN_R+\frac{1}{2}f^M_N\big(\bar{N}_LN^c_L+\bar{N}^c_RN_R\big)\right]\xi_y\\ &+\text{H.c.},\end{aligned}\tag{4}$$

where according to the global symmetry (2), the Higgs doublets ϕ_{L_x}, ϕ_{R_x}, ϕ_{L_y}, ϕ_{R_y} and singlets ξ_x, ξ_y are specified by

$$\begin{aligned}&\phi_{L_x}(\phi_{R_x})=\begin{cases}\phi_{L_2}(\phi_{R_2}) & \text{for } x=+1,\\ \phi_{L_1}(\phi_{R_1}) & \text{for } x=-1,\end{cases}\\ &\xi_x=\begin{cases}\xi & \text{for } x=+1,\\ \xi^* & \text{for } x=-1,\end{cases}\end{aligned}\tag{5}$$

and

$$\begin{aligned}&\phi_{L_y}(\phi_{R_y})=\begin{cases}\phi_{L_1}(\phi_{R_1}) & \text{for } y=+1,\\ \phi_{L_2}(\phi_{R_2}) & \text{for } y=-1,\end{cases}\\ &\xi_y=\begin{cases}\xi & \text{for } y=+1,\\ \xi^* & \text{for } y=-1.\end{cases}\end{aligned}\tag{6}$$

The parity-odd singlet σ is responsible for the spontaneous D-parity violation [15] to guarantee the different VEVs of the left- and right-handed Higgs doublets [7]. We may simply replace the spontaneous D-parity violation by softly breaking the discrete left–right symmetry. The right-handed Higgs doublet ϕ_{R_1} and its left-handed partner ϕ_{L_1} will drive the left–right and electroweak symmetry breaking, respectively. Through the ω-term in the potential (3), the other Higgs doublets ϕ_{R_2} and ϕ_{L_2} will also acquire their VEVs,

$$\langle\phi_{R_2}\rangle\simeq-\frac{\omega\langle\xi\rangle\langle\phi_{R_1}\rangle}{M^2_{\phi_{R_2}}},\qquad\langle\phi_{L_2}\rangle\simeq-\frac{\omega\langle\xi\rangle\langle\phi_{L_1}\rangle}{M^2_{\phi_{L_2}}},\tag{7}$$

like the type-II seesaw [2]. Note that the left-handed VEVs is constrained by [16]

$$\sqrt{\langle\phi_{L_1}\rangle^2+\langle\phi_{L_2}\rangle^2}\simeq174\ \text{GeV}.\tag{8}$$

The Higgs singlet ξ accounts for the global symmetry breaking, after which the fermion singlets can obtain their heavy masses,

$$\begin{aligned}&M_D=h_D\langle\xi\rangle, && M_U=h_U\langle\xi\rangle, && M_E=h_E\langle\xi\rangle,\\ &M^D_N=f^D_N\langle\xi\rangle, && M^M_N=f^M_N\langle\xi\rangle.\end{aligned}\tag{9}$$

We emphasize that with the Yukawa couplings (4), the strong CP phase will not vanish at tree level. The universal seesaw models can provide a solution to the strong CP problem without an axion [17] if the left–right symmetry is not the charge-conjugation but the parity [18]. However, the parity as the left–right symmetry is inconsistent with the global symmetry (2), which is essential for the mass generation of the fermion singlets.

3. Universal seesaw

By integrating out the charged fermion singlets, $D_{L,R}$, $U_{L,R}$, and $E_{L,R}$, we can have the following dimension-5 operators,

$$\begin{aligned}\mathcal{O}_5\supset{}&y_D\frac{1}{M_D}y^T_D\bar{q}_L\tilde{\phi}_{L_2}\tilde{\phi}^\dagger_{R_2}q_R\\ &+y_U\frac{1}{M_U}y^T_U\bar{q}_L\phi_{L_1}\phi^\dagger_{R_1}q_R\\ &+y_E\frac{1}{M_E}y^T_E\bar{l}_L\tilde{\phi}_{L_x}\tilde{\phi}^\dagger_{R_x}l_R+\text{H.c.}\end{aligned}\tag{10}$$

After the left–right symmetry breaking, the right-handed charged fermions can obtain their Yukawa couplings with the left-handed fermion and Higgs doublets. Such Yukawa couplings can give the Dirac masses of the charged fermions after the electroweak symmetry breaking. Note that for generating the top mass $m_t=171.2$ GeV and the bottom mass $m_b=4.20$ GeV [16], we should expect the VEVs $\langle\phi_{L_1}\rangle=\mathcal{O}(100\ \text{GeV})$ and $\langle\phi_{L_2}\rangle=\mathcal{O}(1\text{–}10\ \text{GeV})$.

In the neutrino sector, the neutral fermion singlets $N_{L,R}$ can form the Majorana or pseudo-Dirac fields, depending on the size of M^D_N and M^M_N. In the pseudo-Dirac case with $M^D_N\gg M^M_N$, the induced dimension-5 operators are

$$\mathcal{O}_5\supset\left[y_{N_R}\frac{1}{M^D_N}y^T_{N_R}+y_{N_L}\frac{1}{(M^D_N)^T}y^T_{N_L}\right]\bar{l}_L\phi_{L_y}\phi^\dagger_{R_y}l_R+\text{H.c.},\tag{11}$$

which can account for the light Dirac neutrinos. In the Majorana case with $M_N^M \gg M_N^D$, we obtain

$$\mathcal{O}_5 \supset \left(y_{N_R}\frac{1}{M_N^M}y_{N_L}^T + y_{N_L}\frac{1}{M_N^M}y_{N_R}^T\right)\bar{l}_L\phi_{L_y}\phi_{R_y}^\dagger l_R + \frac{1}{2}\left(y_{N_R}\frac{1}{M_N^M}y_{N_R}^T + y_{N_L}\frac{1}{M_N^M}y_{N_L}^T\right) \times \left(\bar{l}_L\phi_{L_y}\phi_{L_y}^T l_L^c + \bar{l}_R^c\phi_{R_y}^*\phi_{R_y}^\dagger l_R\right) + \text{H.c.} \tag{12}$$

Clearly, the right-handed neutrinos can get a Majorana mass term from the operators involving the right-handed doublets. Their Yukawa couplings to the left-handed lepton and Higgs doublets can be induced by the operators involving the left- and right-handed doublets. This means we can realize the type-I seesaw [1]. Furthermore, the operators involving the left-handed doublets will give an additional contribution to the neutrino masses, playing a role of the type-II seesaw [2].

4. Peccei–Quinn symmetry

After the global symmetry breaking, the complex scalar singlet ξ can be described by

$$\xi = \frac{1}{\sqrt{2}}(f+\rho)\exp\left(i\frac{a}{f}\right) \tag{13}$$

with a being a NGB. It is easy to derive the tree-level couplings of the NGB to the colored fermion singlets,

$$\mathcal{L} \supset \frac{1}{2f}(\partial_\mu a)\sum_Q \bar{Q}\gamma^\mu\gamma_5 Q, \tag{14}$$

where Q denotes D and U. Because of the $Q-q$ mixing, the NGB can also couple to the SM quarks at tree level,

$$\mathcal{L} \supset -\frac{1}{2f}(\partial_\mu a)\sum_{Q,q}\bar{q}\gamma^\mu\gamma_5 q\left(y_Q\frac{\langle\phi_L\rangle\langle\phi_R\rangle}{M_Q^2}y_Q^T\right) \sim \frac{1}{2f}(\partial_\mu a)\sum_{Q,q}\bar{q}\gamma^\mu\gamma_5 q\left[\mathcal{O}\left(\frac{m_q}{M_Q}\right)\right]. \tag{15}$$

In addition, the coupling (14) associated with the quark–gluon interaction will induce

$$\mathcal{L} \supset \frac{1}{2f}(\partial_\mu a)\sum_{Q,q}\bar{q}\gamma^\mu\gamma_5 q\left[\frac{\alpha_s^2}{\pi^2}\ln\left(\frac{M_Q}{m_q}\right)\right] \tag{16}$$

at two-loop order, as in the Kim–Shifman–Vainshtein–Zakharov [12] (KSVZ) model. For $M_Q \gg m_q$, the loop contribution (16) will dominate over the tree-level contribution (15). The NGB should be a pNGB as it picks up a tiny mass [14,19],

$$m_a^2 = N^2\frac{Z}{(1+Z)^2}\frac{f_\pi^2}{f^2}m_\pi^2\left[\frac{\alpha_s^2}{\pi^2}\ln\left(\frac{M_Q}{m_q}\right)\right]^2 \tag{17}$$

with $N=3$ for three families of the SM quarks while $Z \simeq m_u/m_d$.

Since the pNGB couples to the axial vector current of the SM quarks, it should be the axion while the global symmetry should be the PQ symmetry. For convenience, we express the axion mass as

$$m_a = \frac{\sqrt{Z}}{(1+Z)}\frac{f_\pi}{f_a}m_\pi \simeq 6.2\ \mu\text{eV}\left(\frac{10^{12}\ \text{GeV}}{f_a}\right), \tag{18}$$

where the axion decay constant f_a can be derived from Eq. (17). The PQ symmetry breaking scale should be high enough to fulfill the theoretical and experimental constraints [20]. For example, the PQ symmetry breaking may happen before inflation to avoid the cosmological domain wall problem. With an appropriate PQ symmetry breaking scale, the axion can act as the dark matter.

5. Leptogenesis

After the left–right symmetry breaking, the leptogenesis can be applied to explain the matter–antimatter asymmetry in the universe. If the neutral fermion singlets and hence the neutrinos form the pseudo-Dirac fields, the decays of the neutral fermion singlets can produce a lepton asymmetry stored in the left-handed lepton doublets and an equal but opposite lepton asymmetry stored in the right-handed neutrinos. Since the effective Yukawa interactions between the left- and right-handed neutrinos are extremely weak, they will go into equilibrium at a very low temperature where the sphaleron [21] action is not active. Therefore, the sphaleron process can partially transfer the left-handed lepton asymmetry to a baryon asymmetry. This leptogenesis scenario with Dirac neutrinos is titled as neutrinogenesis [9]. In the other case with the neutral fermion singlets being the Majorana fields, we can have the conventional leptogenesis with Majorana neutrinos.

6. Conclusion

In this Letter we connected the universal seesaw scenario, where not only the neutral neutrinos but also the charged fermions obtain their masses through the seesaw mechanism, to the PQ symmetry for solving the strong CP problem. In our model, the fermion singlets, including the color triplets for generating the quark masses and the color singlets for generating the lepton masses, have the Yukawa couplings to a complex scalar singlet. The scalar singlet acquires a large VEV to spontaneously break the global PQ symmetry. So, the fermion singlets can obtain their heavy masses for the realization of the universal seesaw. The colored fermion singlets mediate the PQ symmetry to the SM quarks at tree and two-loop level so that the axion can pick up a tiny mass through the color anomaly. We thus naturally related the PQ symmetry to the neutrino mass-generation [22]. Our model also accommodates the leptogenesis with Majorana or Dirac neutrinos.

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