Linking axion-like dark matter, the XENON1T excess, inflation and the tiny active neutrino masses

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(Dated: July 16, 2020)

We consider a renormalizable theory, which successfully explains the number of Standard Model (SM) fermion families and whose non-SM scalar sector includes an axion dark matter as well as a field responsible for cosmological inflation. In such theory, the spontaneous breaking of the Peccei-Quinn (PQ) symmetry at a very large scale produces very large Majorana neutrino masses, thus allowing a natural explanation of the tiny active neutrino masses from a tree-level type I seesaw mechanism. In the theory under consideration, soft-breaking mass terms generate an axion mass consistent with the current experimental limits. Furthermore, the theory under consideration can also successfully accommodates the XENON1T excess provided that the PQ symmetry is spontaneously broken at the $10^{10}$ GeV scale.

I. INTRODUCTION

Currently, an axion is a very attractive subject in Particle Physics in both theoretical and experimental aspects \[1, 2\], thus providing a motivation to consider extensions of the Standard Model that include this particle in its field content. The axion is a CP-odd scalar field which arises in the solution of the strong-CP problem, and originally it is always a massless particle that prevent its possibility to be a dark matter candidate. In order to generate a mass for the axion, one can consider the implementation of radiative corrections as shown in \[3–5\] or gravitational effects \[6\]. In this paper, we present a new mechanism to generate the axion mass, which to the best of our knowledge has not been previously discussed in the literature.

It is interesting to note that nowadays the axion is widely considered as a candidate of dark matter (DM)\[7\]. The dark matter candidate is existed only in some beyond standard models. Among the SM extensions, the models based on the $SU(3)_C \times SU(3)_L \times U(1)_X$ gauge symmetries (called 3-3-1 models, for short) \[8, 9\] have several very interesting features, some of them being, the natural explanation of the number of SM fermion families, the electric charge quantization and the solution of the strong CP problem from the PQ symmetry \[10\], which is automatically fulfilled in the 3-3-1 models. In one of the 3-3-1 models, there exist both interesting features, namely the axion dark matter candidate and inflaton for Early Universe \[11, 12\]. In the above mentioned papers, the axion gets mass only from gravitational contribution.

It is worth mentioning that the CP-odd sector of the 3-3-1 model with axion has been considered in \[11–13\]. However, the rotation matrix that diagonalizes the squared mass matrix for the CP odd neutral scalar fields given in \[11\], which was obtained by Mathematica is not unitary, and thus such mixing matrix cannot be used for further studies such coupling constants, collider phenomenology, etc. The aim of this work is to reconsider the CP-odd scalar sector of the aforementioned model with the inclusion of the soft-breaking mass term, which in turn generates the mass of the axion. Such soft-breaking mass term was not considered in \[11\]. Our paper is organized as follows: in Section II we present brief review of the model. Section III is devoted to discrete and PQ symmetries needed for the axion existence. The Higgs potential and the resulting physical scalar spectrum is discussed in Section IV. We state our
conclusions in Section V.

II. REVIEW OF THE MODEL

The model under consideration is based on $SU(3)_C \times SU(3)_L \times U(1)_X \times Z_{11} \times Z_2$ and has the following fermion content:

$$
\psi_{aL} = (\nu_a, t_a, (\nu^\alpha_R)^0)^T \sim (1, 3, -1/3), \quad l_{aR} \sim (1, 1, -1), \quad N_{aR} \sim (1, 1, 0), \\
Q_{3L} = (u_3, d_3, U)^T_L \sim (3, 3, 1/3), \quad Q_{\alpha L} = (d_\alpha, -u_\alpha, D_\alpha)^T_L \sim (3, 3^*, 0), \\
u_{aR}, U_R \sim (3, 1, 2/3), \quad d_{aR}, D_{aR} \sim (3, 1, -1/3),
$$

(1)

where $\alpha = 1, 2$ and $a = \{3, \alpha\}$ are family indices. The $U$ and $D$ are exotic quarks with ordinary electric charges, whereas $N_{aR}$ are right-handed neutrinos.

The scalar sector of the model is composed of three $SU(3)_L$ scalar triplets and one $SU(3)_L$ singlet scalar field. They have the following transformations under the $SU(3)_C \times SU(3)_L \times U(1)_X$ symmetry:

$$
\chi = (\chi_1^+, \chi_2^-, \chi_3^0)^T \sim (1, 3, -1/3), \quad \eta = (\eta_1^0, \eta_2^-, \eta_3^1)^T \sim (1, 3, -1/3), \\
\rho = (\rho_1^+, \rho_2^-, \rho_3^0)^T \sim (1, 3, 2/3), \quad \phi \sim (1, 1, 0).
$$

(2)

To generate masses for gauge bosons and fermions, the scalar fields should acquire vacuum expectation values (VEVs). These fields can be expanded around the minimum as follows

$$
\chi = \left( \begin{array}{c} \frac{1}{\sqrt{2}}(R_1^3 + iI_1^3) \\ \chi^- \\ \frac{1}{\sqrt{2}}(v_\chi + R_3^3 + iI_3^3) \end{array} \right), \quad \eta = \left( \begin{array}{c} \frac{1}{\sqrt{2}}(v_\eta + R_1^3 + iI_1^3) \\ \eta^- \\ \frac{1}{\sqrt{2}}(R_\eta^3 + iI_\eta^3) \end{array} \right), \quad \rho = \left( \begin{array}{c} \frac{1}{\sqrt{2}}(v_\rho + R_\rho + iI_\rho) \\ \rho_1^+ \\ \rho_3^0 \end{array} \right),
$$

(3)

$$
\phi = \frac{1}{\sqrt{2}}(v_\phi + R_\phi + iI_\phi).
$$

(4)

Note that since $\phi$ carries non-zero PQ charge (as shown below), it has to be complex as shown in (4). The VEV $v_\chi$ is responsible for the first stage of gauge symmetry breaking, whereas $v_\eta, v_\rho$ trigger the second stage of electroweak symmetry breaking.

III. DISCRETE AND PECCIEI-QUINN SYMMETRIES

In order to keep intact the physics results of the Ref.[11] the Lagrangian of the model must be invariant by the discrete symmetries $Z_{11} \times Z_2$ which are summarised in Table I. Here we have used a notation $\omega_k \equiv e^{2i\pi k/5}, k = 0, \pm1 \cdots \pm5$.

| $SU(3)_C$ | $SU(3)_L$ | $U_{aR}$ | $d_{aR}$ | $U_{3R}$ | $D_{1R}$ | $\psi_{aL}$ | $l_{aR}$ | $N_{aR}$ | $\eta$ | $\chi$ | $\rho$ | $\phi$ |
|---------|---------|------|------|------|------|------|------|------|------|------|------|------|
| $Q_{1L}$ | $Q_{3L}$ | $u_{aR}$ | $d_{aR}$ | $\nu_{aR}$ | $t_{aR}$ | $l_{aR}$ | $N_{aR}$ | $\eta$ | $\chi$ | $\rho$ | $\phi$ |
| $SU(3)_C$ | $SU(3)_L$ | $U_{aR}$ | $d_{aR}$ | $\nu_{aR}$ | $t_{aR}$ | $l_{aR}$ | $N_{aR}$ | $\eta$ | $\chi$ | $\rho$ | $\phi$ |
| $U(1)_X$ | $Z_{11}$ | $Z_2$ | $\omega^+_{11}$ | $\omega_0$ | $\omega_1$ | $\omega_2$ | $\omega_3$ | $\omega_4$ | $\omega_5$ | $\omega^{-1}_{11}$ | $\omega^{-1}_1$ | $\omega^{-1}_2$ | $\omega^{-1}_3$ | $\omega^{-1}_4$ | $\omega^{-1}_5$ |

Table I: Quantum numbers of fields in the model

Let us discuss about the PQ symmetry. These discrete symmetries yield the following Yukawa couplings

$$
-\mathcal{L}^Y = y_{1i} \bar{Q}_{3L} U_{3R} X + y_{2j} \bar{Q}_{iL} D_{1R} X^* + y_{3k} \bar{Q}_{3L} u_{aR} \eta^a + y_{4l} \bar{Q}_{iL} d_{aR} \eta^* + y_{5m} \bar{Q}_{3L} d_{aR} \rho + y_{6n} \bar{Q}_{iL} u_{aR} \rho^* + g_{ab} \bar{\psi}_{aL} e_{bR} + (y^D_{\nu})_{ab} \bar{\psi}_{aL} \eta N_{bR} + (y_N)_{ab} \bar{\phi} N_{aR}^C N_{bR} + \text{H.c.}
$$

(5)
where unlike Ref. [11], the neutrino Yukawa term $\bar{\psi}_{ab}^{(p)} v_{ij} (\bar{\psi}_{aL})_j (\psi_{bL})^c_i (\rho^*)_k$ is not present since that term is not invariant under the $Z_2$ symmetry. It is worth mentioning that the Yukawa interactions given above do not allow for terms which interchange $\chi \leftrightarrow \eta$, since they do not respect the $Z_{11} \times Z_2$ symmetry.

Assuming fermions of opposite chiralities have opposite PQ charges and $X_d = X_d^c = 1$ we summarise PQ charges of fermions in Table II.

| $X_{PQ}$ | $u_{aL}$ | $d_{aL}$ | $U_L$ | $D_L$ | $L_a$ | $L_a$ | $\nu_{aL}$ | $\nu_{aR}$ | $N_{aR}$ |
|----------|----------|----------|-------|-------|-------|-------|-----------|-----------|---------|
| -1       | 1        | 1        | 1     | 1     | 1     | 1     | 1         | -1        | 1       |

Table II: PQ charges of fermions in the model

It is now clear that the entire Lagrangian of the model is $U_{PQ}(1)$ invariant, providing a natural solution to the strong-CP problem.

The following remarks are in order

1. The discrete symmetry $Z_N$ can naturally be accommodated when it has enough number of fields in its spectrum.

2. If the $Z_{11} \times Z_2$ symmetry is imposed, then the PQ symmetry automatically appears in the model. Therefore, the CP problem can be solved by the dynamical properties of the axion field, which is a physical scalar field belonging to the CP odd neutral scalar sector.

3. From the last two terms of Eq. (5), it follows that the tiny masses for the light active neutrinos are generated from a type I seesaw mechanism mediated by very heavy right handed Majorana neutrino getting masses at the scale $v_\phi$ of spontaneous breaking of the PQ symmetry.

IV. HIGGS POTENTIAL

The model scalar potential takes the form:

$$V = \mu_3^2 \phi^* \phi + \mu_4^2 (\chi^*)^2 \chi + \mu_5^2 \rho^* \rho + \mu_6^2 \eta^* \eta + \lambda_1 (\chi^*)^2 \chi + \lambda_2 (\eta^*)^2 \eta + \lambda_3 (\rho^*)^2 \rho + \lambda_4 (\eta^*) \eta (\eta^*) \eta + \lambda_5 (\chi^*) \chi (\rho^*) \rho + \lambda_6 (\eta^*) \eta (\rho^*) \rho + \lambda_7 (\chi^*) \chi (\eta^*) \eta + \lambda_8 (\chi^*) \chi (\rho^*) \rho + \lambda_9 (\eta^*) \eta (\rho^*) \rho + \lambda_{10} (\phi^*)^2 \phi + \lambda_{11} (\phi^*) \phi (\eta^*) \eta + \lambda_{12} (\phi^*) \phi (\rho^*) \rho + \lambda_{13} (\phi^*) \phi (\eta^*) \eta + \lambda_{14} \epsilon^{ijk} \eta \eta \phi \phi + V_{soft} + H.c,$$

where $V_{soft}$ is a $Z_2 \times Z_{11}$ soft-breaking mass term introduced to generate the axion mass. This soft-breaking mass term is given by:

$$V_{soft} = \frac{\mu_5^2}{2} (\phi^* \phi + \phi \phi^*) = \frac{\mu_5^2}{2} \left[(v_\phi + R_{\phi})^2 - R_{\phi}^2 \right].$$

(7)

The VEV $v_\phi$ is responsible for PQ symmetry breaking resulting in existence of invisible axion due to very high scale around $10^{10} - 10^{11}$ GeV. Then SU($3)_L \times U(1)_X$ breaks to the SM group by $v_\chi$ and two others $v_{\rho}, v_\eta$ are needed for the usual U$(1)_Q$ symmetry. Hence

$$v_\phi \gg v_\chi \gg v_{\rho}, v_\eta.$$  

(8)

The constraint conditions of such scalar potential were analyzed in Ref. [11].

From an analysis of the scalar potential, we find that the physical CP odd neutral scalar mass eigenstates are:

$$
\begin{pmatrix}
G_Z^3 \\
A_5 \\
G_Z \\
a
\end{pmatrix} =
\begin{pmatrix}
\cos \theta_3 & -\sin \theta_3 & 0 & 0 \\
\sin \theta_3 \cos \phi & \cos \theta_3 \cos \phi & -\sin \theta_3 \cos \phi & \sin \theta_3 \\
-\sin \theta_3 \sin \phi & \sin \theta_3 \cos \phi & \cos \theta_3 & 0 \\
-\sin \theta_3 \sin \phi \cos \phi & -\sin \theta_3 \cos \phi \cos \phi & -\sin \theta_3 \cos \phi \cos \phi & -\sin \theta_3 \cos \phi \cos \phi
\end{pmatrix}
\begin{pmatrix}
I_3^3 \\
I_5^3 \\
I_3 \\
I_5
\end{pmatrix}.
$$

(9)
The scalar spectrum of the model has the following features:

\[
\tan \alpha = \frac{v_\eta}{v_\rho}, \quad \tan \theta_3 = \frac{v_\eta}{v_\chi},
\]

\[
\tan \theta_\phi \approx \frac{v_\eta}{v_\phi}, \quad \tan \theta_4 \approx \tan \alpha.
\]

The mixing angles in the CP odd scalar sector depend on the ratio of \(v_\eta\) to \(v_\rho\), \(v_\chi\) and \(v_\phi\).

By rotating the scalar fields to the physical basis, we find that the scalar fields of the model under consideration can be expressed as follows:

\[
\chi \simeq \begin{pmatrix} \varphi^0 \\ G_Y- \\ \sqrt{2} (v_\chi + H_1 + iG_{Z1}) \end{pmatrix}, \quad \eta \simeq \begin{pmatrix} \frac{1}{\sqrt{2}} (v_\eta + h_5 + iA_5) \\ H^-_1 \\ G_{X0} \end{pmatrix}, \quad \rho \simeq \begin{pmatrix} 0 \\ G_{W+} \\ H_2^+ \end{pmatrix},
\]

\[
\phi \simeq \frac{1}{\sqrt{2}} (v_\phi + \Phi + ia).
\]

The scalar spectrum of the model has the following features:

1. In the charged scalar sector, one has four singly charged fields: two of them are massless being the Goldstones bosons eaten by the longitudinal components of the bilepton gauge bosons \(W^\pm\) and \(Y^\pm\), whereas the physical charged scalar fields have masses at the TeV scale.

2. In the CP-odd scalar sector, there are six fields: two massless Goldstones eaten by the longitudinal components of the \(Z\) and \(Z'\) gauge bosons and one massless is part of Goldstones boson associated to the longitudinal components of the neutral bilepton \(G_{X^0}\) gauge boson. One massive CP-odd field has the same mass as another in CP-even sector. Hence the above pair is a complex bilepton scalar field denoted by \(\varphi^0\). Another field associated with the singlet \(\phi\) is identified to the axion \(a\). Originally, this axion is massless, but it becomes massive thanks to the softly breaking mass term. The axion acquires a mass \(m_a\) equal to the soft-breaking mass parameter \(\mu_a\). Finally, we have only one massive CP-odd scalar field \(A_4\).

3. In the CP-even scalar sector, there are six fields. One massless field is part of \(G_{X^0}\), another massive is associated to \(\varphi^0\). One heavy field associated with the imaginary part of the singlet \(\phi\) is identified with the inflaton \(\Phi\) and one SM-like Higgs boson \(h\) with mass \(126\) GeV. For a detailed study of how the field \(\Phi\) drives inflation in 3-3-1 models, see for instance Ref [12, 14]. The remaining two fields are one heavy with mass at TeV scale and another with mass at EW scale. Besides that, the inflaton \(\Phi\) has mass around \(10^{11}\) GeV, thanks to its very large vacuum expectation value and then it provides \(10^{11}\) GeV scale masses for the right handed Majorana neutrinos, thus allowing the implementation of a type I seesaw mechanism that produces the tiny light active neutrino masses. Furthermore, such \(10^{11}\) GeV scale right handed Majorana neutrino masses are consistent with the constraints arising from leptogenesis, as shown in [15]. Furthermore, the light active neutrino mass matrix arising from the type I seesaw mechanism in our model has the form:

\[
M_\nu = M^D_\nu M^{-1}_N \left(M^D_\nu\right)^T, \quad M^D_\nu = y^D_\nu \frac{v_\eta}{\sqrt{2}}, \quad M_N = y_N \frac{v_\phi}{\sqrt{2}},
\]

which implies that the smallness of the active neutrino masses is a consequence of their inverse scaling with the very large scale of breaking of the Peccei-Quinn symmetry.

4. After including the axion field \(a(x)\), the total Lagrangian as,

\[
L_{total} = L_{SM} + \left(\bar{\theta} + \frac{a(x)}{f_a}\right) \frac{g^2}{32\pi^2} G^{\mu\nu} \tilde{G}_{\mu\nu} + \text{kinetic + interactions},
\]

where the axion decay constant \(f_a = \frac{\alpha}{\sqrt{2}}\), is related to the magnitude of the VEV that breaks the \(U(1)_{PQ}\) symmetry. Hence, the CP violating term \(G\tilde{G}\) is proportional to \((\bar{\theta} + \frac{a(x)}{f_a})\), when the axion field is redefined, \(a(x) \to a(x) - (a(x))\), thus implying that the CP violating term \(G\tilde{G}\) is no longer present in the Lagrangian, so that the Strong CP problem is solved.
5. Another important point is that the invisible axion of the model has a small component of $I_X^3$ which couples only to exotic quarks at tree level via the following Yukawa interaction

$$L_{(aq'q')} = -\frac{i}{\sqrt{2}}(\sin \theta_3 \sin \theta_\phi \cos \theta_4) \left[ h U_L U_R - h_{\alpha\beta} D_{\alpha L} D_{\beta R} \right] a + H.c.$$  

6. Eq. (9) leads to $I_\rho = -\sin \theta_4 \cos \theta_\phi A_\rho + \cos \theta_4 G_Z - \sin \theta_4 \sin \theta_\phi a$. Hence from Eq. (5) we obtain that the electron-positron-axion interaction has the form:

$$L_{ae} = -\frac{i g_{ae}}{\sqrt{2}} \sin \theta_\rho \sin \theta_4 a \bar{\gamma}_5 e = -i \frac{\sqrt{2} m_e}{\eta} \sin \theta_\phi \sin \theta_4 a \bar{\gamma}_5 e \equiv -i g_{ae} a \bar{\gamma}_5 e,$$

where $m_e = 0.488$ MeV is the electron mass at the $M_Z$ scale. Considering, $v_\rho = v_\eta = \frac{v}{\sqrt{2}} \approx 174$ GeV, which corresponds to $\theta_4 \approx \frac{\pi}{2}$, we get the following value for the electron-positron-axion coupling:

$$g_{ae} = 2.5 \times 10^{-14},$$

which is consistent with its corresponding experimental value $g_{ae} = 2.5 \times 10^{-14}$ arising from the XENON1T excess [16].

V. CONCLUSIONS

In this paper we have presented a new mechanism to generate the mass of axion by introducing a softly breaking mass term in the scalar potential in the 3-3-1 models, where the VEV of the added scalar singlet $\phi \sim (1, 1, 0)$ is much larger than the scale of breaking of the $SU(3)_L \times U(1)_X$ gauge symmetry. The model under consideration contains three important features: Dark Matter axion, inflation and type I seesaw mechanism to generate the light active neutrino masses. In such model, the axion is massless at the tree level but acquires a mass consistent with the current experimental limits, thanks to a soft-breaking mass term in the scalar potential. Such axion can play the role of cold dark matter. Furthermore, very large right handed Majorana neutrino masses are generated from the spontaneous breaking of the PQ symmetry at the energy scale around $10^{10} - 10^{11}$ GeV, thus allowing the implementation of tree level type I seesaw mechanism that produces the tiny active neutrino masses. In addition, from the analysis of the model scalar potential we find that in the CP-odd neutral scalar sector, there are six fields: one massive CP-odd field $A_\rho$, two massless Goldstones eaten by the longitudinal components of the $Z$ and $Z'$ and one massless which is part of Goldstones boson associated to the longitudinal component of the neutral bilepton $G_{X0}$ gauge boson. One massive CP-odd field has the same mass of another one in the CP-even sector. Hence the above pair is a complex bilepton scalar field denoted by $\phi^0$. Finally, another field associated to the neutral scalar singlet $\phi$ is identified with axion $a$ which becomes massive due to a soft-breaking mass term. Such axion is also very useful for successfully accommodating the XENON1T excess provided that the PQ symmetry is spontaneously broken at the $10^{10}$ GeV scale.

Acknowledgments

H. N. L acknowledges the financial support of the Vietnam Academy of Science and Technology under grant NVCC 05.03/20-20. A.E.C.H received funding from Chilean grant Fondecyt No. 1170803.

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