The KSVZ Axion and Pseudo-Nambu-Goldstone Boson Models for the XENON1T Excess

Tianjun Li
CAS Key Laboratory of Theoretical Physics, Institute of Theoretical Physics,
Chinese Academy of Sciences, Beijing 100190, China and
School of Physical Sciences, University of Chinese Academy of Sciences, Beijing, 100049, China

The XENON1T excess can be explained by the Axion Like Particle (ALP) dark matter with mass around 2.5 keV. However, there are three problems needed to be solved: suppressing the coupling \( g_{a\gamma} \) between the ALP and photon, and generating the proper coupling \( g_{ae} \) between the ALP and electron as well as the correct ALP mass. We propose three models to solve these problems. In our models, the \( g_{ae} \) couplings are produced by integrating out the vector-like leptons, but we need some fine-tunings to obtain the ALP mass. Similarly, one can study the DFSZ axion model. In the Z\( _8 \) and \( U(1)_X \) models with approximate Pseudo-Nambu-Goldstone Bosons (PNGBs), the coupling \( g_{a\gamma} \) is suppressed due to \( SU(3)_C \times U(1)_{EM} \) anomaly free. In the \( Z_8 \) model, the PNGB mass can be generated naturally at the keV scale via the dimension-8 operator. To solve the PNGB quality problem in the \( Z_8 \) model, we embed it into the model with \( U(1)_X \) gauge symmetry.

Introduction.— Using the low-energy electronic recoil data with an exposure of 0.65 ton-years, the XENON Collaboration recently reported the results for new physics search \([1]\). They have observed 285 events over an expected background of 232±15 events, and found an excess for the electron recoil energies below 7 keV, rising towards lower energies and prominent between 2 and 3 keV. Also, they showed that the solar axion and the solar neutrino magnetic moment hypotheses no longer have the substantial statistical significance, and their significance levels are respectively reduced to 2.1\( \sigma \) and 0.9\( \sigma \). This excess has been studied extensively via solar axion, Axion Like Particles (ALPs), the non-standard neutrino-electron interactions with light mediators, and dark photon, etc \([2, 50]\).

It is well-known that the Peccei-Quinn mechanism \([51, 52]\) provides a natural solution to the strong CP problem in the Quantum Chromodynamics (QCD), and predicts a light Pseudo-Nambu-Goldstone Boson (PNGB), dubbed as axion \( a \) from QCD anomalous \( U(1)_{PQ} \) global symmetry breaking. The electroweak axion \( 51 \) \( 54 \) was ruled out by the \( K \rightarrow \pi a \) and \( J/\Psi \rightarrow a\gamma \) experiments. And there are two viable invisible axion models: the Dine-Fischler-Srednicki-Zhitnitsky (DFSZ) model \([55, 56]\) and Kim-Shifman-Vainshtein-Zakharov (KSVZ) model \([57, 58]\) with \( U(1)_{PQ} \) symmetry breaking scale from about 10\( ^{10} \) GeV to 10\( ^{12} \) GeV. Interestingly, the ALPs, which are the generalizations of axion, may be intrinsic the structure of string theory. The ALP dark matter can explain the XENON1T excess via the electron absorption \([3, 18]\), and let us study its properties before our model building. The Lagrangian between axion and photon/fermions is

\[
\mathcal{L}_a^{\text{int}} \supset \frac{\alpha_{EM}}{8\pi} \frac{C_{a\gamma}}{f_a} a F \tilde{F} + C_{af} \frac{\partial \phi}{2f_a} f \gamma^\mu \gamma_5 \tilde{f} ,
\]

where \( \alpha_{EM} \) is structure constant, \( f_a \) is the axion decay constant, and \( C_{a\gamma} \) and \( C_{af} \) are the couplings. The above Lagrangian can be rewritten as

\[
\mathcal{L}_a^{\text{int}} \supset \frac{1}{4} g_{a\gamma} a F \tilde{F} - i g_{af} a f \gamma^5 \tilde{f} ,
\]

where

\[
 g_{a\gamma} = \frac{\alpha_{EM} C_{a\gamma}}{2\pi} , \quad g_{af} = \frac{C_{af} m_f}{f_a} .
\]

The best fit for the XENON1T excess gives \([18]\)

\[
m_a = 2.5 \text{ keV} , \quad g_{ae} = 2.5 \times 10^{-14} .
\]

In particular, the cooling constraint \( g_{ae} < 2.5 \times 10^{-13} \) can be satisfied \([3, 4]\). The stronger constraint on the decay width for the axion decay into diphoton arises from the observation of the cosmic X-ray backgroud (CXB) gives \([50]\)

\[
\frac{C_{a\gamma}}{g_{ae}} \lesssim 2.9 \times 10^{-3} \left( \frac{2.5 \text{ keV}}{m_a} \right)^{3/2} \left( \frac{2.5 \times 10^{-14}}{g_{ae}} \right) .
\]

And then we obtain

\[
C_{a\gamma} \lesssim 2.9 \times 10^{-3} \left( \frac{f_a}{2 \times 10^{10} \text{ GeV}} \right) .
\]

For the QCD axion models, we have

\[
C_{a\gamma} = \frac{E}{N} - 1.92(4) ,
\]

where \( E \) and \( N \) are respectively the electromagnetic and QCD anomaly factors, and 1.92(4) is generated by the
mixing of the axion with the QCD mesons below the confinement scale.

Next, let us discuss the properties of the ALP dark matter particle, which can explain the XENON1T excess. First, we shall show $f_a \simeq 2 \times 10^{10}$ GeV later, and then the traditional QCD axion will have a mass around 2.85 $\times 10^{-4}$ eV. Thus, the ALP dark matter particle cannot be the traditional QCD axion. Second, from Eq. (5), we obtain $C_{a\gamma} \lesssim 2.9 \times 10^{-3}$. In general, there exists about 0.1% fine-tuning for Eq. (5), and the natural solution to it is that both the first term and the second term on the right-hand side vanish: the first condition implies that we do not have $[U(1)_{EM}]^2 U(1)_{PQ}$ anomaly, while the second condition means no mixing between axion and QCD mesons and thus we do not have $[SU(3)]^2 U(1)_{PQ}$ anomaly. Therefore, the ALP dark matter particle, which can explain the XENON1T excess, might arise from breaking of a $SU(3) \times U(1)_{EM}$ anomaly free $U(1)_{X}$ symmetry (or its discrete subgroup) and is a PNGB.

In short, to explain the XENON1T excess via a PNGB dark matter, we need to address three problems: how to suppress the coupling $g_{a\gamma}$, and how to generate the coupling $g_{ac}$ as well as the correct ALP mass. We shall propose three models to solve these problems: the KSVZ axion model with $U(1)_{PQ}$ symmetry, the model with $Z_8$ discrete symmetry, and the model with $U(1)_{X}$ gauge symmetry. In our models, assuming that the right-handed electron is charged under $U(1)_{PQ}$, $Z_8$, and $U(1)_{X}$ symmetries, we can produce the $g_{ac}$ couplings by integrating out the vector-like leptons. In the KSVZ axion model, the coupling $g_{a\gamma}$ can be suppressed by choosing proper sets of vector-like fermions. And with some fine-tuning, we can obtain the correct axion mass from high-dimensional operators via quantum gravity effects. Similarly, one can study the DFSZ model, where the coupling $g_{ac}$ is present and thus we do not need to generate it. In the $Z_8$ and $U(1)_{X}$ models, we do not have $SU(3)_{C} \times U(1)_{EM}$ anomaly, so the coupling $g_{a\gamma}$ is suppressed. In the $SU(3)_{C} \times SU(2)_{L} \times U(1)_{Y}$ model, we obtain the decay constant around $2 \times 10^{10}$ GeV for the best fit. The correct PNGB $a$ mass around 2.5 keV can be generated from dimension-8 operator naturally. We also show that a has a lifetime long enough to be a dark matter candidate. Moreover, the PNGB dark matter density around the observed value can be generated via the misalignment mechanism, while its thermal density is negligible. Furthermore, to solve the PNGB quality problem via quantum gravity effects in the $Z_8$ model, we embed it into a $SU(3)_{C} \times SU(2)_{L} \times U(1)_{Y} \times U(1)_{X}$ model. The $U(1)_{X}$ gauge symmetry is broken down to a $Z_8$ discrete symmetry around the string scale $10^{17}$ GeV, and then the $Z_3$ model can be realized.

The KSVZ Axion Model.— First, we construct the KSVZ axion model which can explain the XENON1T excess. We introduce the vector-like fermions $(XQ_{1i}, XQ_{2i})$, $(XU_{1i}, XU_{2i})$, $(XD_{1j}, XD_{2j})$, $(XL_{1i}, XL_{2i})$, and $(XE_{1i}, XE_{2i})$, as well as a SM singlet axion field $a$. For simplicity, we assume the vector-like fermions have $U(1)_{PQ}$ charge $+1$, while $S$ has $U(1)_{PQ}$ charge $-2$. These particles and their quantum numbers under the $SU(3)_{C} \times SU(2)_{L} \times U(1)_{Y} \times U(1)_{PQ}$ symmetry and global symmetries are summarized in Table I.

The Lagrangian is given by:

$$-\mathcal{L} = -m_a^2 |S|^2 + \lambda_S |S|^4 + (y_{XQ})^i_j XQ_{1i}XQ_{2j}^c + (y_{XU})^i_j XU_{1i}XU_{2j}^c + (y_{XD})^i_j XD_{1i}XD_{2j}^c + (y_{XL})^i_j XL_{1i}XL_{2j}^c + (y_{XE})^i_j XE_{1i}XE_{2j}^c + \text{H.C.}) \quad (7)$$

To have small $C_{a\gamma}$, we need to find the sets of vector-like fermions which gives $E/N$ close to 1.92(4). Because the contribution to the electromagnetic anomaly factor from $(XL_{1i}^c, XL_{2i})$ is the same as the $(XE_{1i}^c, XE_{2i})$, we do not consider $(XE_{1i}^c, XE_{2i})$ for simplicity. Of course, any $(XL_{1i}^c, XL_{2i}^c)$ can be replaced by a $(XE_{1i}^c, XE_{2i})$ in the following discussions. For $n$ pairs of $(Q_{XQ_i}^c, XQ_{2i}^c)$, $m$ pairs of $(U_{XU_i}^c, XU_{2i}^c)$, $k$ pairs of $(D_{XD_j}^c, XD_{2j}^c)$, and $l$ pairs of $(L_{XL_i}^c, XL_{2i}^c)$, we can obtain the condition $C_{a\gamma} \simeq 0$

$$\frac{10n + 8m + 2k + 6l}{6n + 3m + 3k} \approx 1.92(4) \quad (8)$$

It is not difficult to find the approximate solution to the above equation, for example, $\frac{10n + 8m + 2k + 6l}{6n + 3m + 3k} = 2$ for $n = m = 0$, $k = 6$, and $l = 4$. In addition, assuming that the right-handed electron and muon are charged under $U(1)_{PQ}$ symmetry and introducing the vector-like fermions $(XL_{1i}, XL_{2i})$ and $(XL_{1j}, XL_{2j})$, we can generate the coupling $g_{ac}$ as we discuss in the following $Z_8$ and $U(1)_{X}$ models. If the QCD axion only obtains mass via instanton effect, its mass will be too small since the decay constant is around $10^{10}$ GeV as in the following discussions. Therefore, the key question is how to generate the correct axion mass around 2.5 keV. As we know, the global $U(1)_{PQ}$ symmetry can be broken by the quantum gravity effects. To be concrete, we consider the following effective operator with dimension $d = 2m + n$ that violates the PQ symmetry by $n$ units:

$$V \supset \lambda_n m_n |S|^{2m} (e^{-i\delta_n} S^n + e^{i\delta_n} S^{n*})$$

$$\approx m_a^2 f_a^2 \left( \frac{\theta^2}{2} - \frac{\theta}{n} \tan \delta_n \right),$$

| $XQ_{1i}$ | $(3, 2, 1/6, 1)$ | $XQ_{2i}$ | $(3, 2, -1/6, 1)$ |
| $XU_{1i}$ | $(3, 1, 2/3, 1)$ | $XU_{2i}$ | $(3, 1, -2/3, 1)$ |
| $XD_{1j}$ | $(3, 1, -1/3, 1)$ | $XD_{2j}$ | $(3, 1, 1/3, 1)$ |
| $XL_{1i}$ | $(1, 2, -1/2, 1)$ | $XL_{2i}$ | $(1, 2, 1/2, 1)$ |
| $XE_{1i}$ | $(1, 1, -1, 1)$ | $XE_{2i}$ | $(1, 1, 1, 1)$ |
| $S$ | $(1, 1, 0, -2)$ |
where we have expanded for $\theta = \frac{2\pi}{3} \ll 1$ by neglecting an irrelevant constant. Here, $M_{P_1}$ is the reduced Planck scale, $\lambda_m^m$ is real and $\delta_m^m$ the phase of the coupling, $S = \frac{1}{\sqrt{2}} (f_a + s) e^{ia_f}$, and $m_a^2 = \frac{\lambda_a^m}{2} f_a (\sqrt{2} M_{P_1})^{-d-4} \cos \delta_m^m$. In particular, the linear term or tadpole term will shift the QCD vacuum from $\langle \theta \rangle = 0$. Therefore, if we have multiple high-dimensional operators, we can find the fine-tuned solution where the sum of the linear terms is zero or so small that the solution to the strong CP problem can be preserved. And the condition is

$$\sum_{m,n} \frac{\tan \delta_m^m}{n} \simeq 0 \ .$$

(9)

Also, the axion mass is given by

$$m_a = \sqrt{\sum_{m,n} \frac{\lambda_m^m f_a^2}{2}} \left( \frac{f_a}{\sqrt{2} M_{P_1}} \right)^{d-4} \cos \delta_m^m \ .$$

(10)

Therefore, with some fine-tuning, we have shown that the KSVZ axion model can explain the XENON1T excess. Similarly, one can study the DFSZ model, where the coupling $g_{ae}$ is present. Then we do not need to generate it.

The $Z_8$ Model.-- We shall propose a $SU(3)_C \times SU(2)_L \times U(1)_Y \times Z_8$ model where $Z_8$ is a global discrete symmetry. First, let us explain our convention, which is the same as the supersymmetric Standard Model (SM). The SM quark doublets, right-handed up-type quarks, right-handed down-type quarks, lepton doublets, right-handed charged leptons, right-handed neutrinos, and the SM Higgs doublet are denoted as $Q_i$, $U^c_i$, $D^c_i$, $L_i$, $E^c_i$, $N^c_i$, and $H$, respectively. We shall construct the models where the masses and mixings for the SM quarks and neutrinos are generated in a traditional way. Thus, $Q_i$, $U^c_i$, $D^c_i$, $L_i$, $N^c_i$, and $H$ are not charged under $Z_8$ discrete symmetry. Also, we assume that the $Z_8$ quantum numbers for right-handed electron $E^c_\mu$, muon $\nu_\mu$, and tau $\nu_\tau$ are $+1$, $-1$, and 0, respectively. To break the $Z_8$ gauge symmetry and have a approximate PNGB, we introduce a SM singlet scalar $S$ with charge $-1$ under $Z_8$. Moreover, to generate the electron and muon Yukawa couplings, we introduce two pairs of vector-like fermions ($XL_1$, $XL^c_1$) and ($XL_2$, $XL^c_2$). These particles and their quantum numbers under the $SU(3)_C \times SU(2)_L \times U(1)_Y \times Z_8$ gauge and discrete symmetries are summarized in Table II.

The scalar potential in our model is given by

$$V = -m_3 S \bar{S}^2 - m_2 |H|^2 + \lambda_S |S|^4 + \lambda_H |S|^2 |H|^2 + \lambda_H |H|^4 + \frac{y}{M_{P_1}^4} |S|^8 + \frac{1}{M_{P_1}^4} \left( y^8 S^8 + \text{H.C.} \right) \ .$$

(11)

For simplicity, we assume $y > |y^8|$ so that the potential is stabilized. From the the dimension-8 operator $y^8 S^8 / M_{P_1}^4$, we obtain the mass of the PNGB $a$ is at the order of $(|S|^8 / M_{P_1}^4)$.

\begin{table}[h]
\begin{center}
| $Q_i$ | (3, 2, 1/6, 0) | $U^c_i$ | (3, 1, -2/3, 0) |
| $D^c_i$ | (3, 1, 1/3, 0) | $L_i$ | (1, 2, 1/2, 0) |
| $S$ | (1, 1, 1, 1) | $E^c_i$ | (1, 1, 1, -1) |
| $E^c_i$ | (1, 1, 1, 0) | $N^c_i$ | (1, 1, 0, 0) |
| $XL_1$ | (1, 2, -1/2, -1) | $XL^c_1$ | (1, 2, 1/2, 1) |
| $XL_2$ | (1, 2, -1/2, 1) | $XL^c_2$ | (1, 2, 1/2, -1) |
| $H$ | (1, 2, -1/2, 1) | $S$ | (1, 1, 0, -1) |
\end{center}
\end{table}

Table II. The particles and their quantum numbers under the $SU(3)_C \times SU(2)_L \times U(1)_Y \times Z_8$ gauge and discrete symmetries.

The Lagrangian for the Yukawa couplings and vector-like fermion masses is

$$\mathcal{L} = y_{ij}^L Q_i U^c_j \overline{T} + y_{ij}^D Q_i D^c_j H + y_{ij}^E Q_i E^c_j H + y_{ij}^L L_i N^c_j \overline{T} + y_{ij}^X L_i E^c_j H + y_{ij}^Z L_i X L^c_j \overline{T} + y_{ij}^S L_i X L^c_j \overline{T} + M_N^N N^c_i N^c_j + M^X \nu L_i X L^c_j \overline{T} + M_z^L L_i X L^c_j H + \text{H.C.} \ .$$

(12)

Using $y_{ij}^L L_i N^c_j \overline{T}$ and $M_N^N N^c_i N^c_j$ terms, we can generate the neutrino masses and mixings via Type I seesaw mechanism. For simplicity, we choose $y_1^2 \neq 0$ and $y_2^5 = 0$, while $y_3^2 = y_4^2 = y_5^2 = 0$. After integrating out the vector-like particles ($XL_1$, $XL^c_1$) and ($XL_2$, $XL^c_2$), we obtain

$$-\mathcal{L} \supset -\frac{1}{M_1^4} y_{1L}^1 y_{1L}^1 S L_1 E^c_1 H - \frac{1}{M_2^4} y_{1L}^2 y_{1L}^2 S L_2 E^c_2 H + \text{H.C.} \ .$$

(13)

Thus, we obtain

$$f_a \equiv \langle S \rangle = \frac{m_e}{g_{ae}} = 2 \times 10^{10} \text{ GeV} \times \left( \frac{2.5 \times 10^{-14}}{g_{ae}} \right) \ .$$

Therefore, for the best fit, we have $f_a = 2 \times 10^{10}$ GeV. And then the mass of the PNGB $a$ is around keV scale from the dimension-8 operator $y^8 S^8 / M_{P_1}^4$ in Eq. (11), and we can indeed take it as 2.5 keV.

After integrating out the electron and muon, we obtain the effective Lagrangian between the PNGB $a$ and photon

$$\mathcal{L}_{\text{eff}} = \frac{\alpha_{em} m_a^2}{48 \pi f_a} \left( \frac{1}{m_e^2} - \frac{1}{m_\mu^2} \right) a F_{\mu\nu} \tilde{F}^{\mu\nu} \ .$$

(14)

And then we get

$$C_{a\gamma} = \frac{1}{6} \left( \frac{m_a^2}{m_e^2} - \frac{m_a^2}{m_\mu^2} \right) \simeq 4.17 \times 10^{-8} \ ,$$

(15)

which is much smaller than $2.9 \times 10^{-3}$ and is negligible. The PNGB $a$ can decay into two photons via the above effective interaction, and the decay rate is

$$\Gamma_{a \rightarrow \gamma\gamma} \simeq \frac{\alpha_{em} m_a^2}{9216 \pi^3} \frac{m_a^7}{f_a^2} \simeq 4.17 \times 10^{-57} \text{ GeV} \left( \frac{m_a}{2 \text{ keV}} \right)^7 \left( \frac{2 \times 10^{10} \text{ GeV}}{f_a} \right)^2 \ .$$
To solve the PNGB quality problem, we propose the
since we can fine-tune some parameters in our models.
In our model, the relativistic PNGBs can be pro-
duced from the scatterings between electron/muon and
the Higgs bosons in the thermal bath. The resulting abundance is [61]
\[
\Omega_{a}^{(th)}h^2 \approx 3.28 \times 10^{-4} \left( \frac{T_R}{3 \times 10^5 \text{ GeV}} \right) \left( \frac{m_a}{2.5 \text{ keV}} \right) \\
\times \left( \frac{2 \times 10^{10} \text{ GeV}}{f_a} \right)^2,
\]
where \( T_R \) is the reheating temperature. Thus, the ther-
mal relic density of \( a \) is negligible.

The PNGB \( a \) can be produced by the misalignment
mechanism as well. When the Hubble parameter is
smaller than the mass of \( a \), it begins to oscillate around
its potential minimum. The temperature \( T_{\text{osc}} \) at the on-
set of the PNGB oscillation is [3 61]
\[
T_{\text{osc}} \approx 1.12 \times 10^6 \text{ GeV} \left( \frac{m_a}{2.5 \text{ keV}} \right)^{1/2}.
\]
For the temperature higher than \( T_{\text{osc}} \), the PNGB field
\( a \) has a field value which is not the potential minimum
in general. We define the initial oscillation amplitude as
\( \phi_{\text{initial}} \equiv \theta_{\text{mis}}/a \) with \( \theta_{\text{mis}} \) the misalignment angle, and
obtain the oscillation energy of the PNGB \( a \) [3 61]
\[
\Omega_{a}^{(\text{osc})}h^2 \approx 0.1 \left( \frac{\theta_s}{4} \right)^2 \left( \frac{f_a}{2 \times 10^{10} \text{ GeV}} \right)^2 \\
\times \begin{cases} 
\frac{T_{\phi}}{(10^6 \text{ GeV})} & \text{for } T_R \lesssim T_{\text{osc}} \\
\frac{m_a}{2.5 \text{ keV}} & \text{for } T_R \gtrsim T_{\text{osc}}
\end{cases}
\]
Thus, to realize the observed dark matter relic density, we
need large initial misalignment angle. We consider
the reheating temperature is higher than the oscillation
temperature, and the 3\( Z \) symmetry breaking is after
inflation. Thus, the decays of the topological defects such
as cosmic string and domain wall might contribute to the relic
density of the PNGB \( a \) as well.

The \( U(1)_X \) Model.— In the above model, 3\( Z \) is a
discrete symmetry, and can be broken via the quantum
gravity effects. Thus, the above discussions might
not be valid in general if we consider quantum gravity
corrections, which is called the PNGB quality prob-
lem. Because we do not solve the strong CP problem,
in principle we are fine with quantum gravity corrections
since we can fine-tune some parameters in our models.
To solve the PNGB quality problem, we propose the
\( SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X \) model where the
\( U(1)_X \) gauge symmetry is broken down to the 3\( Z \) dis-
crete symmetry around the string scale 10\(^{17} \) GeV. In ad-
dition to the particles in the 3\( Z \) model, we shall intro-
duce two pairs of vector-like particles \( (X_{E1}, X_{E1}^c) \) and
\( (X_{E2}, X_{E2}^c) \), as well as a SM singlet Higgs scalar field \( T \)
with \( U(1)_X \) charge 8. The particles and their quantum
numbers under the \( SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X \)
gauge symmetry are given in Table III. And one can eas-
ily show that our model is anomaly free.

The scalar potential is given by
\[
V = -m_S^2 |S|^2 - m_T^2 |T|^2 - m_{T'}^2 |T'|^2 + \lambda_S |S|^4 + \lambda_T |T|^4 + \lambda_{ST} |S|^2 |T|^2 + \lambda_{ST'} |S|^2 |T'|^2 \\
+ \frac{y}{M^6_{Pl}} |T|^2 |S|^8 + \frac{1}{M^8_{Pl}} (y' TS^8 + \text{H.C.}) .
\]
To stabilize the potential after \( U(1)_X \) gauge symmetry
breaking, we require
\[
\frac{y}{M^6_{Pl}} |\langle T \rangle|^2 > \frac{1}{M^8_{Pl}} |y' \langle T \rangle| .
\]

Table III. The particles and their quantum numbers under the
\( SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X \) gauge symmetry.

| \( Q_i \) | \( (3.2.1/6.0) \) | \( U_{i}^{c} \) | \( (3.1.2/3.0) \) |
| --- | --- | --- |
| \( D_{L}^{c} \) | \( (3.1.1/3.0) \) | \( L_{i} \) | \( (1.2.1/2.0) \) |
| \( E_{i}^{c} \) | \( (1.1.1.1) \) | \( E_{i} \) | \( (1.1.1.1) \) |
| \( E_{i} \) | \( (1.1.1.0) \) | \( N^{c}_{i} \) | \( (1.1.0.0) \) |
| \( X_{L_{i}} \) | \( (1.2.1/2.1) \) | \( X_{L_{i}} \) | \( (1.2.1/2.1) \) |
| \( X_{L_{2}} \) | \( (1.2.1/2.1) \) | \( X_{L_{2}} \) | \( (1.2.1/2.1) \) |
| \( X_{E_{i}} \) | \( (1.1.1.0) \) | \( X_{E_{i}} \) | \( (1.1.1.0) \) |
| \( X_{E_{2}} \) | \( (1.1.1.0) \) | \( X_{E_{2}} \) | \( (1.1.1.0) \) |
| \( H \) | \( (1.2.1/2.1) \) | \( S \) | \( (1.1.0.1) \) |
| \( T \) | \( (1.1.0.8) \) |
functions for \((XE_1, XE_2')\) and \((XE_2, XE_2')\) are highly suppressed on the 3-brane at \(y = 0\), and then the Yukawa couplings \(y_{\alpha E}^{XE}\) will be very small. The rest discussions are similar to the above Section, so we shall not repeat it here. In short, we can solve the PNGB quality problem in the \(SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X\) model.

**Conclusion.**—We proposed three models to explain the XENON1T excess. In our models, the \(g_{\alpha \gamma}\) couplings are generated by integrating out the vector-like leptons, and the correct PNGB mass arises from high-dimensional operators. In the KSVZ axion model, the coupling \(g_{\alpha \gamma}\) can be suppressed by choosing proper sets of vector-like fermions, but we need some fine-tuning to obtain the ALP mass. In the \(Z_8\) model, the coupling \(g_{\alpha \gamma}\) is suppressed due to \(SU(3)_C \times U(1)_{EM}\) anomaly free, and the PNGB mass can be generated naturally at the keV scale via the dimension-8 operator. To solve the PNGB quality problem in the \(Z_8\) model, we embedded it into the model with \(U(1)_X\) gauge symmetry.

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