Axion from pseudo Goldstone seesaw without Peccei-Quinn symmetry

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Introduction: The strong CP problem is a big challenge to the standard model (SM) of particle interactions. The most attractive solution to this problem is the anomalous Peccei-Quinn (PQ) global symmetry $U(1)_{PQ}$ proposed in 1977 [1]. After the PQ symmetry is spontaneously broken, a Goldstone boson can emerge as usual. Subsequently this Goldstone boson picks up a mass through the color anomaly [2–4]. Therefore, the Goldstone boson from the PQ symmetry breaking eventually becomes a pseudo Goldstone boson, named the axion [5, 6]. However, the original PQ model within a two Higgs doublet context was quickly ruled out by experimental data. Actually, the axion has not been observed in any experiments so far. This fact implies that the axion should interact with the SM particles at an extremely weak level [7–10]. To revive the PQ symmetry, Kim-Shifman-Vainshtein-Zakharov (KSVZ) [11–13] and Dine-Fischler-Srednicki-Zhitnitsky (DFSZ) [13, 14] published their invisible axion models during 1979 to 1981. In the KSVZ-type and DFSZ-type models [11–14], a gauge-singlet scalar drives the PQ symmetry to be spontaneously broken far far above the TeV scale. This opens a new window to explore the origin of axion at colliders.

Realistic models: We now begin to demonstrate our mechanism in details. The new scalars and fermions are summarised in Table I where $U(1)_p$ and $U(1)_X$ are two gauge symmetries and $Z_8$ is a discrete symmetry. The heavy scalars $\Delta$ and $\Omega$ cross the $U(1)_p$ and $U(1)_X$ symmetries and connect the two $U(1)_p$ Higgs scalars $\sigma, \xi$ to the two $U(1)_X$ Higgs scalars $\varphi, \eta$. The chiral fermions $\psi_{L,R}$ and $\chi_{L,R}$ are four color triplets with an electric charge $Q_e = -\frac{1}{3}$ or $+\frac{2}{3}$. Clearly our model is free of gauge anomaly.

For simplicity, we do not write down the full Lagrangian, which is assumed at renormalizable level. Because of our chosen charge assignments in Table II the Yukawa and mass terms only involving the new fermions should be nothing but

$$\mathcal{L} \supset -y_{\varphi \psi} \bar{\psi} \varphi \psi - y_{\chi \chi} \bar{\chi} \chi - m_{\psi_{L,R}} \bar{\psi}_{L,R} \psi_{L,R} + \text{H.c.} \quad (1)$$

As for the scalar potential, it should include the quartic terms as below,

$$V \supset \alpha_1 \Delta \sigma \varphi \varphi + \alpha_2 \Delta \xi \eta \eta + \beta_1 \Omega \sigma \varphi^* \varphi^* + \beta_2 \Omega \xi \eta^* \eta^* + \kappa (\sigma^* \xi)^2 + \text{H.c.} \quad (2)$$

In order to simplify the following demonstration, we conveniently integrate out the heavy scalars $\Delta$ and $\Omega$ from the above potential, i.e.

$$V \supset \frac{\alpha_1^2 \alpha_2}{M^2_{\Delta}} \sigma^* \xi (\varphi^* \eta)^2 + \frac{\beta_1^2 \beta_2}{M^2_{\Omega}} \sigma^* \xi (\varphi \eta^*)^2 + \kappa (\sigma^* \xi)^2 + \text{H.c.} \quad (3)$$
Here and thereafter the parameter \( \kappa \) is rotated to be real for convenience and without loss of generality.

One probably has noticed that the Yukawa couplings in Eq. (1) can arise from a two Higgs doublet context. In this case, the new \( U(1)_P \) gauge symmetry should be replaced by the SM \( SU(2)_L \times U(1)_Y \) electroweak symmetries. Accordingly, the new colored fermions \( \psi_{L,R} \) and \( \chi_{L,R} \) respectively are the SM up-type and down-type quarks, while the new Higgs singlets \( \sigma \) and \( \xi \) are the neutral components of the two Higgs doublets. In addition, the heavy scalars \( \Delta \) and \( \Omega \) come from two heavy isodoublets to construct the scalar potential (2) and then mediate the effective operators (3). Another alternative scheme is to abandon the \( U(1)_X \) gauge symmetry and the Higgs singlet \( \eta \). We then can replace \( \eta \) by \( \varphi^* \) as well as \( \eta^* \) by \( \varphi \) in the potentials (2) and (3). For simplicity we do not give the details of these variant models which will be studied elsewhere.

Recognize that any known axion models must contain a PQ global symmetry \( U(1)_{PQ} \) with the \( U(1)_{PQ} - SU(3)_c - SU(3)_c \) anomaly. However, it is impossible to define any anomalous global symmetries in our models because of the terms in Eq. (2) or Eq. (3). Instead, our models respect a \( Z_8 - SU(3)_c - SU(3)_c \) anomaly as well as a \( Z_8 - U(1)_Y - U(1)_Y \) anomaly.

### Pseudo Goldstone bosons: When the Higgs scalars \( \sigma, \xi, \varphi \) and \( \eta \) develop the nonzero vacuum expectation values (VEVs) \( v_{\sigma,\xi,\varphi,\eta} \), they can be expressed by

\[
\begin{align*}
\sigma &= \frac{1}{\sqrt{2}} (v_\sigma + h_\sigma) \exp(iG_\sigma/v_\sigma), \\
\xi &= \frac{1}{\sqrt{2}} (v_\xi + h_\xi) \exp(iG_\xi/v_\xi), \\
\varphi &= \frac{1}{\sqrt{2}} (v_\varphi + h_\varphi) \exp(iG_\varphi/v_\varphi), \\
\eta &= \frac{1}{\sqrt{2}} (v_\eta + h_\eta) \exp(iG_\eta/v_\eta),
\end{align*}
\]

with \( h_{\sigma,\xi,\varphi,\eta} \) being the Higgs bosons and \( G_{\sigma,\xi,\varphi,\eta} \) being the Goldstone bosons. We then insert these expressions in the effective potential (3). Through the series expansions, we can read the mass terms as follows,

\[
V \supset -\frac{\alpha v_\sigma v_\xi v_\varphi v_\eta}{8M_\Delta^2} \left[ \left( \frac{G_\xi}{v_\xi} - \frac{G_\sigma}{v_\sigma} \right) + 2 \left( \frac{G_\eta}{v_\eta} - \frac{G_\varphi}{v_\varphi} \right) \right]^2
-\frac{\beta v_\sigma v_\xi v_\varphi v_\eta}{8M_\Omega^2} \left[ \left( \frac{G_\xi}{v_\xi} - \frac{G_\sigma}{v_\sigma} \right) - 2 \left( \frac{G_\eta}{v_\eta} - \frac{G_\varphi}{v_\varphi} \right) \right]^2
-\kappa v_\sigma^2 v_\xi^2 \left( \frac{G_\xi}{v_\xi} - \frac{G_\sigma}{v_\sigma} \right)^2
\quad \text{with} \quad \alpha = \text{Re} (\alpha_1^* \alpha_2) , \quad \beta = \text{Re} (\beta_1^* \beta_2).
\]

The two \( U(1)_P \) Goldstone bosons \( G_\sigma \) and \( G_\xi \) have the following orthogonal transformation,

\[
G_\sigma = (v_\xi G_\xi + v_\varphi G_\varphi) / \sqrt{v_\sigma^2 + v_\xi^2} , \quad G_\xi = (v_\eta G_\xi - v_\varphi G_\varphi) / \sqrt{v_\sigma^2 + v_\xi^2}.
\]

Similarly, the two \( U(1)_X \) Goldstone bosons \( G_\varphi \) and \( G_\eta \) have

\[
G_\varphi = (v_\eta G_\eta + v_\varphi G_\varphi) / \sqrt{v_\varphi^2 + v_\eta^2} , \quad G_\eta = (v_\varphi G_\eta - v_\varphi G_\varphi) / \sqrt{v_\varphi^2 + v_\eta^2} .
\]

Under the new base of \( G_p, G_{\sigma \xi}, G_X \) and \( G_{\varphi \eta} \), we can

| Scalars & Fermions | \( \sigma \) | \( \xi \) | \( \varphi \) | \( \eta \) | \( \Delta \) | \( \Omega \) | \( \psi_L \) | \( \psi_R \) | \( \chi_L \) | \( \chi_R \) |
|-------------------|-------------|-------------|-------------|-------------|---------|---------|---------|---------|---------|---------|
| spin              | 0           | 0           | 0           | 0           | 0       | \( \frac{1}{2} \) | \( \frac{1}{2} \) | \( \frac{1}{2} \) | \( \frac{1}{2} \) |
| \( U(1)_P \)     | +1          | +1          | 0           | 0           | -1      | -1      | 1       | 0       | 0       | 1       |
| \( U(1)_X \)     | 0           | 0           | +1          | +1          | -2      | +2      | 0       | 0       | 0       | 0       |
| \( Z_8 \)        | -i          | +i          | \( e^{\pi i/4} \) | \( e^{3\pi i/4} \) | +       | -       | -i      | +       | +       | +i      |
rewrite the mass terms \( [\mathbf{3}] \) to be

\[
V \supset \frac{1}{2} \left[ G_{\sigma \xi} \ G_{\varphi \eta} \right] \left[
\begin{array}{c}
2b v_p^2 \\
2d v_p v_X
\end{array}
\right] \left[
\begin{array}{c}
G_{\sigma \xi} \\
G_{\varphi \eta}
\end{array}
\right] \text{ with }
\]

\[
v_p = \sqrt{v_\sigma^2 + v_\xi}, \quad v_X = \sqrt{v_\eta^2 + v_\xi},
\]

\[
b = -\kappa - \frac{v_\varphi^2 v_\eta^2 a}{8M_\Delta^2 v_\sigma v_\xi} - \frac{v_\eta v_\xi}{8M_\Omega^2 v_\sigma v_\xi},
\]

\[
c = -\frac{v_\varphi v_\eta}{2M_\Delta} - \frac{v_\sigma v_\xi}{2M_\Omega},
\]

\[
d = -\frac{v_\varphi v_\eta}{4M_\Delta^2} + \frac{v_\sigma v_\xi}{4M_\Omega^2}.
\]

This means the \( U(1)_P \) Goldstone \( G_P \) and the \( U(1)_X \) Goldstone \( G_X \) would be eaten by the longitudinal components of the \( U(1)_P \) gauge boson \( P_\mu \) and the \( U(1)_X \) gauge boson \( X_\mu \), respectively. As for the other \( U(1)_P \) Goldstone \( G_{\sigma \xi} \) and the other \( U(1)_X \) Goldstone \( G_{\varphi \eta} \), they are two pseudo Goldstone bosons and their mass eigenstates are

\[
A = G_{\sigma \xi} \cos \theta - G_{\varphi \eta} \sin \theta \quad \text{with}
\]

\[
M_A^2 = bv_p^2 + cv_X^2 + \sqrt{(bv_p^2 - cv_X^2)^2 + 4d^2 v_p^2 v_X^2},
\]

\[
a = G_{\sigma \xi} \sin \theta + G_{\varphi \eta} \cos \theta \quad \text{with}
\]

\[
m_a^2 = bv_p^2 + cv_X^2 - \sqrt{(bv_p^2 - cv_X^2)^2 + 4d^2 v_p^2 v_X^2}.
\]

Here the mixing angle \( \theta \) is determined by

\[
\tan 2\theta = \frac{2d v_p v_X}{cv_X^2 - bv_p^2}.
\]

We expect the pseudo Goldstone \( A \) much heavier than the other pseudo Goldstone \( a \). For this purpose, we can take

\[
bcv_p^2 \gg cv_X^2, dv_p v_X,
\]

\[
A \simeq G_{\sigma \xi} + \frac{dv_X}{bv_p} G_{\varphi \eta},
\]

\[
a \simeq G_{\varphi \eta} - \frac{dv_X}{bw_p} G_{\sigma \xi},
\]

\[
m_a^2 \simeq 2c \left( \frac{v_X^2}{b} \right) v_X^2.
\]

In this case the \( U(1)_P \) pseudo Goldstone \( G_{\sigma \xi} \) and the \( U(1)_X \) pseudo Goldstone \( G_{\varphi \eta} \) respectively dominate the heavy pseudo Goldstone \( A \) and the light pseudo Goldstone \( a \).

Remarkably the above mechanism for generating the light pseudo Goldstone \( a \) has an essence of seesaw, like the type-I/III+II seesaw mechanism for neutrino masses [21][30]. In this sense, we entitle this mechanism to 'pseudo Goldstone seesaw'.

**Axion:** Through Eq. (1), the chiral fermions \( \psi_{L,R} \) and \( X_{L,R} \) can acquire the following mass terms,

\[
\mathcal{L} \supset -\frac{1}{\sqrt{2}} g_{\psi \varphi} \bar{\psi}_L \psi_R e^{iG_{\psi \varphi} / v_\psi} - \frac{1}{\sqrt{2}} g_{\psi \chi} \bar{\psi}_L \chi_R e^{-iG_{\psi \chi} / v_\chi}
\]

\[
- m_{\psi \chi} \bar{\psi}_R \chi_L + \text{H.c.}.
\]

We then can remove the Goldstone bosons \( G_\sigma \) and \( G_\xi \) from the above masses by redefining the fermions, i.e.

\[
\psi_L e^{-iG_{\psi \varphi} / v_\psi} \rightarrow \psi_L, \quad \chi_R e^{-iG_{\psi \chi} / v_\chi} \rightarrow \chi_R.
\]

As a result, the Goldstone bosons \( G_\sigma \) and \( G_\xi \) should appear in the kinetic terms, i.e.

\[
\mathcal{L} \supset -\frac{1}{\sqrt{2}} \left( \partial_\mu G_\sigma \right) \bar{\psi}_L \gamma_\mu \psi_L - \frac{1}{\sqrt{2}} \left( \partial_\mu G_\xi \right) \bar{\psi}_R \gamma_\mu \chi_R
\]

\[
= \frac{G_\sigma}{v_\sigma} \partial_\mu \left( \bar{\psi}_L e^{iG_{\psi \varphi} / v_\psi} \right) + \frac{G_\xi}{v_\xi} \partial_\mu \left( \bar{\psi}_R e^{-iG_{\psi \chi} / v_\chi} \right)
\]

\[
- \left( \frac{G_\sigma}{v_\sigma} - \frac{G_\xi}{v_\xi} \right) \left( \frac{\alpha_v}{8\pi} G_{\mu \nu} \tilde{G}^{\mu \nu} + Q_v^2 \frac{\alpha_{\tilde{F}}}{8\pi} F_{\mu \nu} \tilde{F}^{\mu \nu} \right). \quad \quad \quad (22)
\]

Here \( G_{\mu \nu} \) and \( F_{\mu \nu} \), respectively are the gluon and photon filed strengths while the tilde notation is for the dual. In the second step the total derivative terms have been abandoned and then in the last step the axial triangle anomaly has been adopted.

By inserting Eq. (10) to Eq. (22) and then considering Eqs. (14-15) and (18-19), we can derive the couplings of the heavy pseudo Goldstone \( A \) and the light pseudo Goldstone \( a \) to the gluon pairs \( G\tilde{G} \) and the photon pairs \( FF \), i.e.

\[
\mathcal{L} \supset -\frac{G_{\sigma \xi}}{v_\sigma v_\xi / v_p} \left( \frac{\alpha_v}{8\pi} G\tilde{G} + Q_v^2 \frac{\alpha_{\tilde{F}}}{8\pi} FF \right)
\]

\[
= -A \cos \theta + a \sin \theta \left( \frac{\alpha_v}{8\pi} G\tilde{G} + Q_v^2 \frac{\alpha_{\tilde{F}}}{8\pi} FF \right)
\]

\[
\simeq \left( \frac{v_p}{v_\sigma v_\xi} A + \frac{dv_X}{bw_p v_\xi} \right) \left( \frac{\alpha_v}{8\pi} G\tilde{G} + Q_v^2 \frac{\alpha_{\tilde{F}}}{8\pi} FF \right). \quad \quad \quad (23)
\]

Therefore, the light pseudo Goldstone \( a \) potentially can serve as the role of axion, depending on its mass \( m_a \) and decay constant \( f_a \[10][31] \). This can be easily achieved.

From Eqs. (19) and (23), we can read

\[
m_a = \sqrt{2\epsilon} \left( \frac{v_\sigma^2 + v_\xi^2}{b} \right), \quad \epsilon \equiv 1 - \frac{d^2}{bc},
\]

\[
f_a = \frac{bv_\sigma v_\xi}{d} \sqrt{v_\sigma^2 + v_\xi^2}.
\]

When the VEVs \( v_{\sigma, \xi, \varphi, \eta} \) are expected at the TeV scale, the \( f_a - m_a \) values can fulfil the experimental constraints \[10][31] \) for a proper choice of the dimensionless parameters \( b, c \) and \( d \). However, unless a large cancellation is
accepted to make the induced parameter $\epsilon$ small enough, the $f_a - m_a$ values should deviate from those in the KSVZ and DFSZ models [10, 31].

In order to realize a light mass $m_a$ and a large decay constant $f_a$ as naturally as possible, we can arrange a D-parity for the scalars carrying the $U(1)_X$ charges, i.e.

$$\Delta \xrightarrow{D} \Omega , \; \varphi \xrightarrow{D} e^{i\pi/4} \varphi^* , \; \eta \xrightarrow{D} e^{i3\pi/4} \eta^* . \tag{25}$$

This D-parity is spontaneously violated by a real gauge-singlet scalar [32], i.e.

$$\zeta \xleftarrow{D} -\zeta \text{ with } \zeta = v_\zeta + h_\zeta . \tag{26}$$

Here $v_\zeta$ is the VEV and $h_\zeta$ is the Higgs boson. The heavy scalars $\Delta$ and $\Omega$ thus can be expected to have a tiny mass split,

$$V \supset \mu \zeta (\Delta^* - \Omega^*) \Omega$$

$$\Rightarrow M_\Delta^2 - M_\Omega^2 = 2\mu v_\zeta \ll M^2 \equiv \frac{1}{2} (M_\Delta^2 + M_\Omega^2) . \tag{27}$$

Meanwhile, the couplings $\alpha_{1,2}$ and $\beta_{1,2}$ in the potential [2] should be constrained by

$$\alpha_1 = i\beta_1 , \quad \alpha_2 = -i\beta_2 \Rightarrow \alpha_1^* \alpha_2 = -\beta_1^* \beta_2 . \tag{28}$$

Accordingly the key parameters $b, c, d$ in Eq. [24] can be determined by

$$b \simeq -\kappa - \frac{\alpha v_\varphi^2 v_\eta^2}{8 M^2 v_\varphi v_\xi} , \quad c \simeq \frac{\alpha \mu v_\xi v_\varphi v_\zeta}{M^4} , \quad d \simeq -\frac{\alpha v_\varphi v_\eta}{2MF^2} . \tag{29}$$

Even if the VEVs $v_{\sigma, \xi, \varphi, \eta, \zeta}$ and the cubic coupling $\mu$ are all taken at the TeV scale, the dimensionless parameter $b$ can be dominated by the quartic coupling $\kappa$ to guarantee the heavy pseudo Goldstone $A$, while the other two dimensionless parameters $c$ and $d$ can be highly suppressed by the larger mass $M$ or the smaller coupling $\alpha$ to guarantee the light axion $a$.

In Fig. 1 we show the origin of the axion-gluon-gluon coupling $aG\tilde{G}$ and the axion-photon-photon coupling $aF\tilde{F}$. Both of the colored fermions $\psi$ and $\chi$ participate in generating the axion-gluon-gluon and axion-photon-photon couplings. Benefitted from the $Z_3 - SU(3)_c - SU(3)_c$ anomaly, the axion $a$ can leave a non-cancelled coupling to the gluon pairs, although the $Z_3$ discrete symmetry has forbidden the anomalous PQ symmetry. In this sense, we would like to emphasize that the axion in our mechanism is induced without the introduction of PQ symmetry.

**Phenomenology:** Our model can provide rich phenomena at colliders. Firstly, the TeV scale breaking of the $U(1)_X$ gauge symmetry motivates the new gauge boson $X_\mu$ in the range from the weak scale to the TeV scale. This new gauge boson can become a dark photon because the $U(1)_X - U(1)_Y$ kinetic mixing is always available. Such dark photon can be expected to find at colliders. For example, the $e^+ e^- \rightarrow q\bar{q}X$ channel could be observed at the CEPC [33, 34]. Alternatively the dark photon can be produced at the LHC by $gg^*\tilde{g}$ fusion through the new fermion loops. We then can find the axion as a missing energy in the dark photon decays, i.e. $X \rightarrow h_\varphi, \eta \rightarrow a + a$ and $h_\varphi, \eta \rightarrow a + + a$. Due to the Higgs portal interactions, the axion could also significantly contribute to the invisible decays of the SM Higgs boson. Moreover, the new color-triplet fermions $\psi$ and $\chi$ can naturally acquire the masses of the order of TeV through their Yukawa couplings with the $U(1)_Y$ Higgs scalars $\sigma$ and $\xi$, which develops the TeV scale VEVs. These new colored fermions can mix with the SM quarks as their electric charge is $-\frac{1}{3}$ or $+\frac{2}{3}$. Through their single or pair production, these new fermions can be verified at the LHC [35, 37]. As an example, the search is sensitive to both charged current and neutral current processes, $pp \rightarrow \psi(\chi) q \rightarrow Wq\bar{q}'$ and $pp \rightarrow \psi(\chi) q \rightarrow Zq\bar{q}'$ with a leptonic decay of the vector gauge boson.

**Conclusion:** We have demonstrated that the pseudo Goldstone seesaw mechanism can account for the ori-
gin of axion without resorting to the anomalous PQ global symmetry. After the spontaneous breaking of the $U(1)_P \times U(1)_X$ gauge symmetries and the $Z_8$ discrete symmetry, the $U(1)_P$ pseudo Goldstone and the $U(1)_X$ pseudo Goldstone respectively obtain a heavy mass and an ultralight mass, besides a tiny mass mixing. The heavy $U(1)_P$ pseudo Goldstone rather than the ultralight $U(1)_X$ pseudo Goldstone couples to the new color-triplet fermions. As a result, the axion is dominated by the ultralight $U(1)_X$ pseudo Goldstone and has an extremely weak interaction with the SM particles. This novel scenario allows all of the new symmetries to be spontaneously broken at the TeV scale and hence opens a new window to explore the axion at colliders.

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