Light Sterile Neutrinos in the Supersymmetric $U(1)'$ Models and Axion Models

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Abstract

We propose the minimal supersymmetric sterile neutrino model (MSSNM) where the sterile neutrino masses are about 1 eV, while the active neutrino masses and the mixings among the active and sterile neutrinos are generated during late time phase transition. All the current experimental neutrino data include the LSND can be explained simultaneously, and the constraints on the sterile neutrinos from the big bang nucleosynthesis and large scale structure can be evaded. To realize the MSSNM naturally, we consider the supersymmetric intermediate-scale $U(1)'$ model, the low energy $U(1)'$ model with a secluded $U(1)'$-breaking sector, and the DFSZ and KSVZ axion models. In these models, the $\mu$ problem can be solved elegantly, and the 1 eV sterile neutrino masses can be generated via high-dimensional operators. For the low energy $U(1)'$ model with a secluded $U(1)'$-breaking sector, we also present a scenario in which the masses and mixings for the active and sterile neutrinos are all generated during late time phase transition.

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I. INTRODUCTION

There has been great progress in neutrino physics during last several years [1]. Solar neutrino [2] and atmospheric neutrino [3] experiments together with reactor neutrino [4, 5] experiments have established the oscillation solutions to the solar and atmospheric neutrino anomalies, which are consistent with three light active neutrino scheme. However, the LSND experiment found evidence for the oscillations $\bar{\nu}_\mu \to \bar{\nu}_e$ and $\nu_\mu \to \nu_e$ with an oscillation probability of around $3 \times 10^{-3}$ [6] and a $\Delta m^2 \sim 1$eV$^2$. The statistical evidence for the anti-neutrino oscillations is much stronger than that for the neutrino case, with some analyses finding a $5\sigma$ effect [7]. Although the other experiments eliminated a large fraction of the parameter space allowed by the LSND, they do not exclude the LSND result [8]. Hopefully, the Mini-BOONE experiment at Fermilab will settle this issue down in the near future [9].

If the LSND experiment is confirmed, to explain its data, we need introduce one or two sterile neutrinos, which are Standard Model (SM) singlets and can mix with the active neutrinos. Also, the masses for the sterile neutrinos should be in the eV range. Because of various possible mass hierarchies for the active and sterile neutrinos, there are three proposals: the 2+2 model [10], 3+1 model [11] as well as 3+2 model [12]. Of the three, 3+1 model seems less disfavored than the 2+2 model due to the null results of other oscillation experiments. And the 3+2 model [12] is apparently in better agreement with all the experimental data than the others.

There are two strong constraints on the sterile neutrino models which can explain the LSND result. The first constraint is that the big bang nucleosynthesis (BBN) allows about three effective light neutrinos ($N_{\nu}^{\text{eff}}$) in the equilibrium when the Universe temperature is around 1 MeV [13]. However, for above three proposals, the rapid active neutrinos-sterile neutrino(s) oscillations give the $N_{\nu}^{\text{eff}} = 4$ for the 3+1 and 2+2 models, and $N_{\nu}^{\text{eff}} = 5$ for the 3+2 model. The second constraint is the bound on the sum of all the neutrino masses in the equilibrium at the epoch of structure formation which corresponds to a temperature around an eV from the large scale structure surveys and WMAP [14, 15]. Suppose that the sterile neutrinos are in the equilibrium during the BBN epoch, the upper bound on the sum of all the neutrino masses is about 1.38 eV for the 3+1 and 2+2 models, and about 2.12 eV for the 3+2 model. These constraints are also very severe because for example the 3+2 model is close to be ruled out if we take the face value.
To avoid the BBN and large scale structure constraints and explain the LSND experiment, Chacko, Hall, Oliver and Perelstein proposed a class of models where the masses and mixings of the active and sterile neutrinos are generated during late time phase transition [16]. Based on the next to the minimal supersymmetric Standard Model (NMSSM), they introduced two SM singlet fields with vacuum expectation values (VEVs) in the 100 keV range so that at the BBN epoch the active as well as the sterile neutrinos are massless. Thus, there is no oscillation among them which can bring the sterile neutrinos into equilibrium. Since the sterile neutrinos decouple from Hubble expansion at very high temperatures, their abundance at the BBN epoch is suppressed leading to concordance with the BBN constraints. Moreover, the constraints on the sum of the neutrino masses from large scale structure surveys and WMAP are also easily avoided. The breaking of the global symmetries gives rise to a few Goldstone bosons, which can couple to both active and sterile neutrinos. These couplings are strong enough for the sterile neutrinos to disappear after they become non-relativistic, for example, by decaying into an active neutrino and a Goldstone boson. Thus, the relic abundance of the sterile neutrinos is low, and they do not significantly contribute to dark matter.

Furthermore, using the idea of the late time phase transition, Mohapatra and Nasri considered the sterile neutrinos in the mirror matter models [17]. If the sterile neutrinos are the mirror neutrinos, they only need to generate the mixings between the active and sterile neutrinos (and not their masses) via the late time phase transition to avoid the constraints from the BBN, the large scale structure surveys and WMAP. An advantage of this model is that the contribution of the sterile neutrinos to the energy density of the universe at the BBN epoch is given by a free parameter unlike the model in Ref. [16]. Cosmological consequences of these two kinds of the models have also been discussed in Refs. [16, 17].

In this paper, we propose the minimal supersymmetric sterile neutrino model (MSSNM) with late time phase transition. The active neutrino masses and the mixings among the active and sterile neutrinos are generated during late time phase transition, while the 1 eV masses for the sterile neutrinos are introduced directly in the Lagrangian. We can also forbid the dangerous operators by introducing $Z_3 \times Z_2$ global symmetry. In the MSSNM, the current neutrino data from all the experiments include the LSND can be explained simultaneously, and one can automatically evade the constraints on sterile neutrinos from the BBN, large scale structure surveys and WMAP. However, there are two interesting questions in the
MSSNM: (1) how to produce the 1 eV masses for the sterile neutrinos? (2) how to solve the \( \mu \) problem because the supersymmetry breaking is mediated by gauge interactions? To realize the MSSNM naturally, we consider the supersymmetric intermediate-scale \( U(1)' \) model, the supersymmetric low energy \( U(1)' \) model with a secluded \( U(1)' \)-breaking sector, and the supersymmetric Dine–Fischler–Srednicki–Zhitnitskii (DFSZ) and Kim–Shifman–Vainshtein–Zakharov (KSVZ) axion models \([18, 19]\). In these models, the 1 eV masses for the sterile neutrinos can be obtained via the high-dimensional operators by integrating out the heavy fields, and the dimension-5 operators for the active neutrino masses and the mixings among the active and sterile neutrinos can also be generated by integrating out the heavy fields. Also, the \( \mu \) problem can be solved elegantly. Furthermore, for the low energy \( U(1)' \) model with a secluded \( U(1)' \)-breaking sector, we briefly present a scenario where the sterile neutrino masses are also generated during late time phase transition.

This paper is organized as follows: in Section II, we propose the MSSNM. We consider the supersymmetric intermediate-scale \( U(1)' \) model, the supersymmetric low energy \( U(1)' \) models with a secluded \( U(1)' \)-breaking sector, and the supersymmetric DFSZ and KSVZ axion models in Sections III, IV and V, respectively. Our discussions and conclusions are given in Section VI.

II. MINIMAL SUPERSYMMETRIC STERILE NEUTRINO MODEL WITH LATE TIME PHASE TRANSITION

We first specify our conventions. For the supersymmetric Standard Model, the SM fermions and Higgs fields are superfields belonging to chiral multiplets. The left-handed quark doublets, the right-handed up-type quarks, the right-handed down-type quarks, the left-handed lepton doublets, the right-handed neutrinos, the right-handed leptons, and one pair of Higgs doublets are denoted as \( Q_i, u^c_i, d^c_i, L_i, n_i, e^c_i, H_u \) and \( H_d \), respectively.

To construct the MSSNM with late time phase transition, we introduce one SM singlet field \( \phi \). The relevant superpotential is

\[
W_\nu = \lambda_{ij} L_i n_j H_u \frac{\phi}{M} + \frac{\kappa}{3} \phi^3 + m_{nij} n_i n_j , \tag{1}
\]

where \( \lambda_{ij} \) and \( \kappa \) are the Yukawa coupling constants, and \( m_{nij} \) are the masses for the sterile neutrinos that are around 1 eV and can be generated via high-dimensional operators.
after extra gauge or global symmetry breaking in the following model buildings. The non-renormalizable term (the first term) can be obtained at scale $M$ by integrating out the heavy fields where $M$ is around $10^{8-9}$ GeV.

After the electroweak symmetry breaking, the renormalizable effective Lagrangian for the neutrino sector is

$$-\mathcal{L}_\nu = (g_{ij}\nu_i\phi + m_{nij}n_in_j \pm \text{H.C.}) + V(\phi),$$

where $g_{ij} = \lambda_{ij} < H_u^0 > / M$, $\nu_i$ is the left-handed neutrino, and the scalar potential is $V(\phi) = -\mu^2|\phi|^2 + \kappa^2|\phi|^4$. Here, we have assumed that the supersymmetry breaking effects produce the negative soft mass-squared for $\phi$. Thus, there is one global $U(1)$ symmetry under which $\nu_i$ and $\phi$ have opposite charges with the same magnitude while $n_i$ is neutral. When $\phi$ acquires VEV, this global $U(1)$ symmetry is broken. And then, the active neutrinos obtain masses, and there is one pseudo-Goldstone boson which has diagonal couplings to the neutrinos in the mass basis.

Similar to the discussions in Ref. [16], we can avoid the constraints from the BBN, the large scale structure surveys and WMAP. First, because the total energy density in radiation at the time of BBN does not differ significantly from the SM prediction [13], we require that the "hidden sector" fields $n_i$ and $\phi$ not be in thermal equilibrium with the "observable sector" fields ($\nu_i$ and $\gamma$, etc) before and during the BBN. To be concrete, we require that the two sectors decouple at a certain temperature $T_0 > 1$ GeV, and do not recouple until the temperature of the observable sector drops below $T_W \sim 1$ MeV, the temperature at which the weak interactions decouple. Thus, we have

$$g_{ij} \lesssim 10^{-5}, \, g_{ij}\kappa \lesssim 10^{-10} r^{-1},$$

where $r$ is the ratio of temperature of the hidden sector to that of the observable sector at the time of BBN. The energy density in the hidden sector is suppressed by a factor of $r^4$ compared to that of observable sector from the naive estimation, so, $r \lesssim 0.3$ is enough for one to avoid the BBN constraints. Moreover, because at least one active neutrino has mass around 0.05 eV, we obtain that $f \gtrsim 10$ keV. And to avoid producing sterile neutrinos by oscillations prior to the decouplings of weak interactions, we obtain that $f \lesssim r$ MeV, and then $g_{ij} \gtrsim r^{-1} 10^{-7}$ and $\kappa \lesssim 10^{-3}$. In short, from the BBN constraints, we have

$$10 \text{ keV} \lesssim f \lesssim r \text{ MeV}, \, r^{-1} 10^{-7} \lesssim g_{ij} \lesssim 10^{-5}, \, g_{ij}\kappa \lesssim 10^{-10} r^{-1}.$$
Second, the constraints on the sum of neutrino masses from the large scale structure surveys and WMAP are automatically evaded in above model, and do not lead to extra limits on $f$. The above lower bound on $g_{ij}$ implies that the reactions $\nu \bar{\nu} \leftrightarrow n \bar{n}$, $\phi \bar{\phi}$ become unfrozen before the sterile neutrinos become non-relativistic. These reactions thermalize the hidden sector fields with the active neutrinos. The density of the thermal sterile neutrino with mass $m_s$ at temperatures $T < m_s$ is suppressed by a Boltzmann factor $e^{-m_s/T}$, and the excess sterile neutrinos disappear either via a decay process $n \rightarrow \nu \phi$, or via an annihilation process $n \bar{n} \rightarrow \nu \bar{\nu}$. Thus, the massive sterile neutrinos will not give a significant contribution to dark matter. And only the sum of the masses of the active neutrinos and the Goldstone boson has to satisfy the constraints in Ref. [15].

In our model, the supersymmetry breaking must be mediated via the gauge interactions. And $\phi$ only feels the supersymmetry breaking via its coupling to $L_i$ and $n_i$ so that its supersymmetry breaking soft mass is around 100 keV. However, the $\mu$ problem becomes a severe problem because of the gauge mediated supersymmetry breaking.

In addition, in the MSSNM, we must highly suppress some other renormalizable operators in the superpotential which are allowed by the gauge symmetry, for example, $\phi^2$, $\phi^2 n_i$ and $H_u L_i \phi$, etc. To achieve this, we introduce a global $Z_3 \times Z_2$ discrete symmetry. Under the $Z_3$ symmetry, the particles in the MSSNM transform as

\[
(Q_i, n_i, e_i^c) \longrightarrow (Q_i, n_i, e_i^c) , \ (u_i^c, H_d) \longrightarrow e^{-i2\pi/3} (u_i^c, H_d) , \\
(\phi, d_i^c, L_i, H_u) \longrightarrow e^{i2\pi/3} (\phi, d_i^c, L_i, H_u). \tag{5}
\]

And under the $Z_2$ symmetry, the particles in the MSSNM transform as

\[
(Q_i, u_i^c, d_i^c, L_i, n_i, e_i^c) \longrightarrow - (Q_i, u_i^c, d_i^c, L_i, n_i, e_i^c) \\
(\phi, H_u, H_d) \longrightarrow (\phi, H_u, H_d). \tag{6}
\]

Note that the $\mu H_u H_d$ term in the superpotential is allowed by the $Z_3 \times Z_2$ discrete symmetry.

To solve the $\mu$ problem and generate the 1 eV sterile neutrino masses naturally via the high-dimensional operators, we shall consider the extra gauge symmetry or global symmetry. In the following Sections, we will show that the MSSNM can be realized elegantly in the intermediate-scale $U(1)'$ model, the low energy $U(1)'$ model with a secluded $U(1)'$-breaking sector, and the axion models. In these models, we do not need to introduce above $Z_3 \times Z_2$ discrete symmetry.
TABLE I: The $U(1)'$ charges of the relevant particles in the intermediate-scale $U(1)'$ model. Here $Q_S$ and $Q_L$ are the $U(1)'$ charges for $S$ and $L_i$, respectively.

| Field | $n_i$ | $S'$ | $\phi$ | $H_u$ | $H_d$ | $X$ | $\overline{X}$ | $X'$ | $\overline{X}'$ |
|-------|-------|------|--------|-------|-------|-----|-------------|------|-------------|
| Charge | $-3Q_S/2$ | $-3Q_S$ | $-Q_S/3$ | $-Q_L + 17Q_S/6$ | $Q_L - 29Q_S/6$ | $-17Q_S/6$ | $11Q_S/6$ | $-2Q_S/3$ | $2Q_S/3$ |

III. INTERMEDIATE-SCALE $U(1)'$ MODEL

We first consider the intermediate-scale $U(1)'$ model where the $\mu$ problem can be solved simultaneously. To break the $U(1)'$ gauge symmetry, we introduce two SM singlet Higgs fields $S$ and $S'$ with $U(1)'$ charges $Q_S$ and $-3Q_S$, respectively. After the supersymmetry is broken, the Higgs potential for $S$ and $S'$ is

$$V(S, S') = m_S^2 |S|^2 + m_{S'}^2 |S'|^2 + \frac{1}{2} g_Z^2 Q_S^2 \left(|S|^2 - 3|S'|^2\right)^2,$$

where $m_S^2$ and $m_{S'}^2$ are the supersymmetry breaking soft masses for $S$ and $S'$, respectively. To break the intermediate-scale $U(1)'$ gauge symmetry, we assume that the sum of the supersymmetry breaking soft masses for $S$ and $S'$ is negative, i.e., $m_S^2 + m_{S'}^2 < 0$. Then, there is a runaway direction along the D-flat direction $\sqrt{3} \langle S \rangle = |\langle S' \rangle|$. However, the potential can be stabilized by the loop corrections or higher-dimensional operators. Thus, the $S$ and $S'$ fields can acquire the intermediate-scale VEVs. For example, suppose that there is a high-dimensional operator in the superpotential

$$W \supset S^3 S' / M_{Pl},$$

where $M_{Pl}$ is the Planck scale, we obtain that the $S$ and $S'$ fields can have the VEVs around $10^{10}$ GeV. So, the $U(1)'$ gauge symmetry is broken at intermediate scale elegantly.

To realize the MSSNM naturally, we introduce four SM singlet fields $X$, $\overline{X}$, $X'$, and $\overline{X}'$. The $U(1)'$ charges for the relevant particles are given in Table I.

The relevant superpotential is

$$W = y_i L_i H_u X + y'_j n_j \phi \overline{X} + y_X S X \overline{X} + y_{X'} S' X \phi + y_{\overline{X}} \overline{X}' \phi^2 + M_{X'} X \overline{X}'$$

$$+ \tilde{\lambda}_{ij} \frac{S^3}{M_{Pl}^2} n_i n_j + h \frac{S^2}{M_{Pl}} H_u H_d,$$

where $m_{X'}$ is the singlet mass parameter, $M_{Pl}$ is the Planck scale, and $h$ is the Higgs coupling constant.
where $y_i, y'_j, y_X, y'_{X'}$, $\bar{\lambda}_{ij}$ and $h$ are Yukawa couplings, and $M_{X'}$ is the vector-like mass for $X'$ and $X'$. We assume that $M_{X'}$ is about $10^{14}$ GeV, which can be generated from the non-renormalizable operators after the gauge symmetry in the Grand Unified Theory (GUT) is broken because $M_{X'} \sim M^2_{\text{GUT}}/M_{\text{Pl}}$ where $M_{\text{GUT}}$ is the GUT scale.

After the $U(1)'$ gauge symmetry breaking, we obtain the superpotential

$$W = \lambda_{ij} L_i n_j H_u \frac{\phi}{M_X} + \frac{\kappa}{3} \phi^3 + m_{nij} n_i n_j + \mu H_u H_d,$$

where

$$\lambda_{ij} = y_i y'_j, \quad M_X = y_X \langle S \rangle, \quad \kappa = 3 y_{X'} y'_{X'} \frac{\langle S \rangle}{M_{X'}}.$$

To produce the suitable sterile neutrino masses $m_{nij}$, we need $\bar{\lambda}_{ij} \sim 10^{-3}$. Such a value for $\bar{\lambda}_{ij}$ could be generated if the corresponding operators for sterile neutrino masses in Eq. (9) were themselves due to the high-dimensional operators involving additional fields with VEVs close to $M_{\text{Pl}}$, e.g., associated with an anomalous $U(1)'$ gauge symmetry [20].

We do not consider the $U(1)'$ anomaly cancellation in this paper because the anomaly free $U(1)'$ models can be constructed easily by introducing SM vector-like fields if one follows the procedures in Refs. [21, 22].

IV. LOW ENERGY $U(1)'$ MODELS WITH A SECLUDED $U(1)'$-BREAKING SECTOR

In the low energy supersymmetric $U(1)'$ models, the $\mu$ problem can be solved elegantly with an effective $\mu$ parameter generated by the VEV of the SM singlet field $S$ which breaks the $U(1)'$ symmetry. And the Minimal Supersymmetric Standard Model (MSSM) upper bound of $M_Z$ on the tree-level mass of the corresponding lightest MSSM Higgs scalar is relaxed because of the Yukawa term $h S H_d H_u$ in the superpotential [23] and the $U(1)'$ D-term [24]. More generally, for specific $U(1)'$ charge assignments for the ordinary and exotic fields one can simultaneously ensure the absence of anomalies; that all fields of the low energy effective theory are chiral, avoiding a generalized $\mu$ problem; and the absence of dimension-4 proton decay operators [21].
There are stringent limits from direct searches at the Tevatron\cite{25} and from indirect precision tests at the $Z$-pole, at LEP 2, and from weak neutral current experiments\cite{26}. The constraints depend on the particular $Z'$ couplings, but in typical models one requires $M_{Z'} > (500 - 800)$ GeV and the $Z - Z'$ mixing angle $\alpha_{Z-Z'}$ to be smaller than a few $\times 10^{-3}$. To explain the $Z - Z'$ mass hierarchy, Erler, Langacker and Li proposed a supersymmetric model with a string-motivated secluded $U(1)'$-breaking sector, where the squark and slepton spectra can mimic those of the MSSM, the electroweak symmetry breaking is driven by relatively large $A$ terms, and a large $Z'$ mass can be generated by the VEVs of additional SM singlet fields that are charged under the $U(1)'$\cite{27}. The phenomenological consequences of the low energy $U(1)'$ models, especially the models with a secluded $U(1)'$-breaking sector, have been studied extensively\cite{22, 28}.

First, let us briefly review the supersymmetric $U(1)'$ model with a secluded $U(1)'$-breaking sector\cite{27}. There are one pair of Higgs doublets $H_u$ and $H_d$, and four SM singlets, $S$, $S_1$, $S_2$, and $S_3$. The $U(1)'$ charges for the Higgs fields satisfy

$$Q_S = -Q_{S_1} = -Q_{S_2} = \frac{1}{2} Q_{S_3}, \quad Q_{H_d} + Q_{H_u} + Q_S = 0. \quad (13)$$

The superpotential for the Higgs fields is

$$W_H = h S H_d H_u + \lambda S_1 S_2 S_3, \quad (14)$$

where the Yukawa couplings $h$ and $\lambda$ are respectively associated with the effective $\mu$ term and with the runaway direction. The corresponding $F$-term scalar potential is

$$V_F = h^2 \left( |H_d|^2 |H_u|^2 + |S|^2 |H_d|^2 + |S|^2 |H_u|^2 \right) + \lambda^2 \left( |S_1|^2 |S_2|^2 + |S_2|^2 |S_3|^2 + |S_3|^2 |S_1|^2 \right), \quad (15)$$

And the $D$-term scalar potential is

$$V_D = \frac{G^2}{8} \left( |H_u|^2 - |H_d|^2 \right)^2 + \frac{1}{2} g_{Z'}^2 \left( Q_S |S|^2 + Q_{H_d} |H_d|^2 + Q_{H_u} |H_u|^2 + \sum_{i=1}^3 Q_{S_i} |S_i|^2 \right)^2, \quad (16)$$

where $G^2 = g_1^2 + g_2^2$; $g_1$, $g_2$, and $g_{Z'}$ are the coupling constants for $U(1)$, $SU(2)_L$ and $U(1)'$, respectively; and $Q_\Phi$ is the $U(1)'$ charge of the field $\Phi$.\[9\]
In addition, we introduce the supersymmetry breaking soft terms

\[ V^H_{soft} = m_{H_d}^2 |H_d|^2 + m_{H_u}^2 |H_u|^2 + m_S^2 |S|^2 + \sum_{i=1}^3 m_{S_i}^2 |S_i|^2 - (A_h h S H_d H_u + A_\lambda \lambda S_1 S_2 S_3 + m_{SS_1}^2 S_1 S_1 + m_{SS_2}^2 S_2 S_2 + \text{H.C.}) . \]  

(17)

There is an almost $F$ and $D$ flat direction involving $S_i$, with the flatness lifted by a small Yukawa coupling $\lambda$. For a sufficiently small value of $\lambda$, the $Z'$ mass can be arbitrarily large. For example, if $h \sim 10\lambda$, one can generate the $Z - Z'$ mass hierarchy where the $Z'$ mass ($M_{Z'}$) is at the order of 1 TeV \cite{27}.

We shall consider the model where the sterile neutrinos are charged under the $U(1)'$ gauge symmetry. Because the sterile neutrinos must be decouple with the observable sector fields at Universe temperature about 1 GeV or at least before the QCD phase transition with chiral symmetry breaking, the $U(1)'$ gauge interaction must be decouple at the same temperature. For the $U(1)'$ interaction mediated by massless gauge boson, the interaction rate $\Gamma \sim n\sigma|v| \sim \alpha_{Z'}^2 T$ where $\alpha_{Z'} = g_{Z'}^2/(4\pi)$. And during radiation-dominated epoch, $H \sim T^2/M_{Pl}$. So, $\Gamma/H \sim \alpha_{Z'}^2 M_{Pl}/T$, and for $T \geq 10^{16}$ GeV such reaction is effectively decouple. For the $U(1)'$ interaction mediated by massive gauge boson, the interaction rate $\Gamma \sim G_{Z'}^2 T^5$ where $G_{Z'} = \alpha_{Z'}/M_{Z'}^2$. And then $\Gamma/H \sim G_{Z'}^2 M_{Pl} T^3$. Thus, if we require that the $U(1)'$ gauge interaction decouple at temperature about 1 GeV,

\[ T_{\text{decouple}} \sim \left( \frac{M_{Z'}}{100 \text{ GeV}} \right)^{4/3} \text{MeV} > 1 \text{ GeV} , \]  

(18)

we obtain that the $U(1)'$ gauge boson mass $M_{Z'}$ is larger than about 18 TeV. And if $T_{\text{decouple}} > 300$ MeV, we find that $M_{Z'} > 7.2$ TeV. These roughly estimations agree with the relevant detail calculations in Ref. \cite{29}.

In the $U(1)'$ model with a secluded $U(1)'$-breaking sector, by running our code, we have shown that we can indeed generate the 10 TeV scale mass for the $U(1)'$ gauge boson by choosing smaller $\lambda$ and suitable supersymmetry breaking soft parameters. Let us give an exemple. We choose the standard GUT value $g_{Z'} = \sqrt{5/3}g_1$ (It is $\sqrt{5/3}g_1$ that unifies with $g_2$ and $g_3$ in the simple GUT models.). With the input parameters in Tabel II, we obtain the VEVs for the Higgs fields after minimizing the Higgs potential numerically: $\langle H_u^0 \rangle = 123 \text{ GeV}, \langle H_d^0 \rangle = 123 \text{ GeV}, \langle S \rangle = 128 \text{ GeV}, \langle S_1 \rangle = 10846 \text{ GeV}, \langle S_2 \rangle = 10846 \text{ GeV}, \langle S_3 \rangle = 10846 \text{ GeV}$. And then, we get that the $U(1)'$ gauge boson mass is 17334 GeV.
TABLE II: The input parameters and Higgs VEVs in the $U(1)'$ model with a secluded $U(1)'$-breaking sector. For the mass parameters and the mass-squared parameters, the units are GeV and GeV$^2$, respectively.

|       | $h$  | $\lambda$ | $A_h$ | $A_\lambda$ | $m^2_{H_u}$ | $m^2_{H_d}$ |
|-------|------|-----------|-------|-------------|-------------|-------------|
|       | 0.9  | 0.01      | 219   | 219         | $-15.5^2$   | $-15.5^2$   |
| $m^2_S$ | $m^2_{S_1}$ | $m^2_{S_2}$ | $m^2_{S_3}$ | $m^2_{S_{S_1}}$ | $m^2_{S_{S_2}}$ | $-21.9^2$   | $19.3^2$     | $-10.9^2$    | $2.4^2$       | $2.4^2$       |
| $Q_{H_u}$ | $Q_{H_d}$ | $Q_S$ | $Q_{S_1}$ | $Q_{S_2}$ | $Q_{S_3}$ | $0.5$ | $0.5$ | -1 | 1 | 1 | -2 |
| $\langle H^0_u \rangle$ | $\langle H^0_d \rangle$ | $\langle S \rangle$ | $\langle S_1 \rangle$ | $\langle S_2 \rangle$ | $\langle S_3 \rangle$ | $123$ | $123$ | $128$ | $10846$ | $10846$ | $10846$ |

In addition, if $\phi$ is charged under the $U(1)'$ gauge symmetry, its supersymmetry breaking soft mass will be much larger than 100 keV without fine-tuning due to the low energy $U(1)'$ gauge interaction. Thus, unlike the intermediate-scale $U(1)'$ model, we assume that $\phi$ is neutral under the $U(1)'$.

To realize the MSSNM, we define a global $Z_3$ symmetry in our model

$$n_i \rightarrow n_i \ , \ S_i \rightarrow e^{-i2\pi/3}S_i \ , \ \Phi \rightarrow e^{i2\pi/3}\Phi \ ,$$

(19)

where $\Phi$ denotes all the other fields. One can easily check that the above superpotential and supersymmetry breaking soft terms satisfy this $Z_3$ global symmetry.

With the following $U(1)'$ charges for the relevant fields,

$$Q_{n_i} = -\frac{3}{2}Q_{S_1} \ , \ Q_{H_u} + Q_{L_i} = \frac{3}{2}Q_{S_1} \ ,$$

(20)

we can have the superpotential for neutrino sector

$$W_\nu = \lambda_{ij} L_i n_j H_u \frac{\phi}{M} + \frac{\kappa}{3} \phi^3 + \frac{1}{M'^2} \left(\bar{\lambda}_{ij}^1 S_1^3 + \bar{\lambda}_{ij}^2 S_2^3\right) n_i n_j \ ,$$

(21)

where the non-renormalizable operators can be obtained by integrating out the heavy fields at scales $M$ and $M'$. Here, $M \sim M' \sim 10^{8-9}$ GeV. Thus, after the $U(1)'$ gauge symmetry is broken, we obtain

$$m_{nij} = \frac{1}{M'^2} \left(\bar{\lambda}_{ij}^1 \langle S_1 \rangle^3 + \bar{\lambda}_{ij}^2 \langle S_2 \rangle^3\right) \ .$$

(22)
Furthermore, we can construct the model where the sterile neutrino masses are also generated during late time phase transition. For example, with the following $U(1)'$ charges for the relevant fields,

$$Q_{n_i} = -\frac{1}{2}Q_{S_1}, \quad Q_{H_u} + Q_{L_i} = \frac{1}{2}Q_{S_1},$$

we can have the superpotential for neutrino sector

$$W_\nu = \lambda_{ij} L_i n_j H_u \frac{\phi}{M} + \frac{\kappa}{3}\phi^3 + \frac{\phi}{M'} (\tilde{\lambda}_{ij}^1 S_1 + \tilde{\lambda}_{ij}^2 S_2) n_i n_j.$$  \hspace{1cm} (24)

The discussions for this model are quite similar to those in Section II, so, we will not give them here.

V. DFSZ AND KSVZ AXION MODELS

As we know, the strong CP problem is solved elegantly by the Peccei–Quinn (PQ) mechanism \[^{30}\], in which a global axial symmetry $U(1)_{PQ}$ is introduced and broken spontaneously at some high energy scale. The original Weinberg–Wilczek axion \[^{31}\] is excluded by experiment, in particular by the non-observation of the rare decay $K \to \pi + a$ \[^{32}\]. And there are two viable “invisible” axion models in which the experimental bounds can be evaded: (1) the DFSZ axion model, in which a SM singlet and one pair of Higgs doublets are introduced, and the SM fermions and Higgs fields are charged under $U(1)_{PQ}$ symmetry \[^{18}\]; (2) the KSVZ axion model, which introduces a SM singlet and a pair of extra vector-like quarks that carry $U(1)_{PQ}$ charges while the SM fermions and Higgs fields are neutral under $U(1)_{PQ}$ symmetry \[^{19}\]. In addition, from laboratory, astrophysics, and cosmology constraints, the $U(1)_{PQ}$ symmetry breaking scale $f_a$ is limited to the range $10^{10}$ GeV $\leq f_a \leq 10^{12}$ GeV \[^{32}\].

The quantum gravitational effects, associated with black holes, worm holes, etc., are believed to violate all the global symmetries, while they respect all the gauge symmetries \[^{33}\]. These effects may destabilize the axion solutions to the strong CP problem due to the violation of the global Peccei–Quinn symmetry. However, after a gauge symmetry is spontaneously broken, there may exist a remnant discrete gauge symmetry which will not be violated by quantum gravity \[^{34}\]. Thus, we can avoid the destabilization problem associated with quantum gravity by introducing an additional approximate global symmetry arising from the broken gauge symmetry.
In string model buildings, there generically exists at least one anomalous $U(1)_A$ gauge symmetry with its anomalies cancelled by the Green–Schwarz mechanism \[35\]. The anomalous $U(1)_A$ gauge symmetry is broken near the string scale when some scalar fields, which are charged under $U(1)_A$, obtain VEVs and cancel the Fayet–Iliopoulos term of $U(1)_A$. Then the D-flatness for $U(1)_A$ is preserved and the supersymmetry is unbroken \[36\]. Usually, there is an unbroken discrete $Z_N$ subgroup of the $U(1)_A$ gauge symmetry, which is protected against quantum gravitational violation. We shall consider this $Z_N$ discrete symmetry as an additional global symmetry to forbid the dangerous non-renormalizable operators which can destabilize the axion solutions to the strong CP problem \[37, 38\].

For the gauge symmetry $\prod_i G_i \times U(1)_A$, the Green–Schwarz anomaly cancellation conditions from an effective theory point of view are \[39, 40\]

$$\frac{A_i}{k_i} = \frac{A_{\text{gravity}}}{12} = \delta_{GS} ,$$

(25)

where the $A_i$ are anomaly coefficients associated with $G_i^2 \times U(1)_A$, $k_i$ is the level of the corresponding Kac–Moody algebra, and $\delta_{GS}$ is a constant which is not specified by low-energy theory alone. For a non-Abelian group, $k_i$ is a positive integer, while for the $U(1)$ gauge symmetry, $k_i$ need not be an integer. All the other anomaly coefficients such as $G_i G_j G_k$ and $[U(1)_A]^2 \times G_i$ should vanish.

In our models, the gauge symmetry, which we are interested in, is $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_A$. So, the relevant Green–Schwarz anomaly cancellation conditions are

$$\frac{A_3}{k_3} = \frac{A_2}{k_2} = \delta_{GS} ,$$

(26)

where $A_3$ and $A_2$ are the $[SU(3)_C]^2 \times U(1)_A$ and $[SU(2)_L]^2 \times U(1)_A$ anomaly coefficients. In this paper, we do not consider the anomalies involving the $U(1)_Y$, because its associated Kac–Moody level $k_1$ is not an integer in general and this condition is not very useful from an effective low energy theory point of view \[41\]. Similarly, the $[U(1)_A]^3$ anomaly can be cancelled by the Green–Schwarz mechanism, but this condition also has an arbitrariness from the normalization of $U(1)_A$. And the $[U(1)_A]^2 \times U(1)_Y$ anomaly does not give any useful low energy constraint.
A. The Supersymmetric DFSZ Axion Model

We introduce two SM singlet fields $S$ and $S'$ to break the PQ symmetry, and two SM singlet fields $X$ and $\overline{X}$ to generate the dimension-5 operators for the active neutrino masses after they are integrated out. The superpotential is

$$W_{\text{tot}} = W_o + W_\nu ,$$

where

$$W_o = y_{ij}^u Q_i u_j^c H_u + y_{ij}^d Q_i d_j^c H_d + y_{ij}^e L_i e_j^c H_d + h \frac{S^2}{M_{Pl}} H_d H_u + y_S \frac{(SS')^2}{M_{Pl}} ,$$

$$W_\nu = y_i L_i H_u X + y'_j n_j \phi \overline{X} + y_X S X X + \frac{\kappa}{3} \phi^3 + \overline{\lambda}_{ij} \frac{S^3}{M_{Pl}^2} n_i n_j .$$

Similar to the discussions in the intermediate-scale $U(1)'$ model, we assume that the sum of the supersymmetry breaking soft masses for $S$ and $S'$ is negative, i.e., $m_S^2 + m_{S'}^2 < 0$. Then, the $S$ and $S'$ fields can acquire intermediate-scale VEVs around $10^{10}$ GeV, which gives us the PQ symmetry breaking scale $f_a$. And the $\mu$ term is given by

$$\mu = h \frac{(S)^2}{M_{Pl}} \sim 10^2 \text{ GeV} .$$

To forbid the other renormalizable operators and the dangerous non-renormalizable operators in the superpotential which are allowed by the gauge symmetry and can destabilize the axion solutions to the strong CP problem, we introduce a $Z_{102}$ discrete symmetry arising from the breaking of an anomalous $U(1)_A$ gauge symmetry. Under the $U(1)_{\text{PQ}}$ symmetry and $Z_{102}$ discrete symmetry, the charges for the particles in this model are given in Table [III]. Because the anomaly coefficients $A_3$ and $A_2$ are equal to 60 and 40, respectively, the anomalies can be cancelled by Green–Schwarz mechanism if $k_3 = 3$ and $k_2 = 2$, i.e., $A_3/k_3 = A_2/k_2$.

In addition, using mathematica code, we have shown that up to dimension-5 operators in the superpotential (dimension-6 operators in the Lagrangian), the terms in Eqs. (28) and (29) are the only operators which are allowed by the gauge symmetry and $Z_{102}$ discrete symmetry. Moreover, this $Z_{102}$ symmetry forbids to high orders the dangerous terms of the following type in the superpotential

$$\frac{S^m (S')^{n-m}}{M_{Pl}^{n-3}} ,$$

$$14$$
which can potentially destabilize the axion solutions \cite{37}. However, as pointed out in Ref. \cite{38}, the axion solutions to the strong CP problem may be destabilized by the non-renormalizable terms in the Kähler potential. Therefore, how to stabilize the axion solutions from the non-renormalizable terms in the Kähler potential is still an interesting question which deserves further study.

After the PQ symmetry breaking, we obtain the superpotential for neutrino sector

\[ W_\nu = \lambda_{ij} L_i n_j H_u \frac{\phi}{M_X} + \frac{\kappa}{3} \phi^3 + m_{nij} n_i n_j \ , \quad (32) \]

where

\[ \lambda_{ij} = y_i y_j', \quad M_X = y_X \langle S \rangle \ , \quad m_{nij} = \tilde{\lambda}_{ij} \frac{\langle S \rangle^3}{M_{Pl}^2} \ . \quad (33) \]

Thus, we obtain the superpotential in the MSSNM.

B. The Supersymmetric KSVZ Axion Model

In addition to those particles in above DFSZ axion model, we introduce one pair of vector-like fields $\Psi$ and $\Psi$ which belong to the 5 and $\overline{5}$ representations, respectively, in the $SU(5)$ language. The superpotential is

\[ W_{tot} = W_o + W_\nu + y_\psi S' \Psi \overline{\Psi} \ , \quad (34) \]

where $W_o$ and $W_\nu$ are given in Eqs. \(28\) and \(29\), respectively.

Similar to the DFSZ axion model, to forbid the other renormalizable operators and the dangerous non-renormalizable operators in the superpotential which are allowed by the gauge symmetry and can destabilize the axion solutions to the strong CP problem, we introduce a $Z_{102}$ discrete symmetry arising from the breaking of an anomalous $U(1)_A$ gauge symmetry.
TABLE IV: Under the $U(1)_{PQ}$ symmetry and $Z_{102}$ discrete symmetry, the charges for the particles in the supersymmetric KSVZ axion model.

| $PQ$ | $Z_{102}$ |
|------|-----------|
|      | 0 0 0 0 0 0 0 0 0 0 -1 0 0 0 1 |
|      | 23 85 11 73 63 21 96 68 20 31 35 47 34 71 |

Under the $U(1)_{PQ}$ symmetry and $Z_{102}$ discrete symmetry, the charges for the particles in this model are given in Table IV. Note that the non-renormalizable term $y_S(SS')^2/M_{Pl}$ in the superpotential, which violates the $U(1)_{PQ}$ symmetry, can be generated from quantum gravity interaction.

Because of the additional fields $\Psi$ and $\bar{\Psi}$, we obtain that the anomaly coefficients $A_3$ and $A_2$ are equal to 80 and 60, respectively. And the anomalies can be cancelled by Green–Schwarz mechanism by choosing $k_3 = 4$ and $k_2 = 3$.

The rest discussions are similar to those in above subsection, so, we will not present them here.

VI. CONCLUSIONS AND DISCUSSIONS

We proposed the minimal supersymmetric sterile neutrino model with late time phase transition. The masses for the sterile neutrinos are about 1 eV, and the active neutrino masses and the mixings among the active and sterile neutrinos are generated during late time phase transition. We can also forbid the dangerous operators by introducing $Z_3 \times Z_2$ discrete symmetry. In the MSSNM, the current neutrino data from all the experiments include the LSND can be explained simultaneously, and one can automatically evade the constraints on sterile neutrinos from the BBN, large scale structure surveys and WMAP. However, how to produce the 1 eV masses for the sterile neutrinos is an interesting question. Also, the supersymmetry breaking is mediated by gauge interactions, then, the $\mu$ problem is still a severe problem in the MSSNM. To realize the MSSNM naturally, we considered the supersymmetric intermediate-scale $U(1)'$ model, the supersymmetric low energy $U(1)'$ model with a secluded $U(1)'$-breaking sector, and the supersymmetric DFSZ and KSVZ axion models. In these models, the 1 eV masses for the sterile neutrinos can be obtained.
via the high-dimensional operators by integrating out the heavy fields, and the dimension-5 operators for the active neutrino masses and the mixings among the active and sterile neutrinos can also be generated by integrating out the heavy fields. Moreover, the $\mu$ problem can be solved elegantly. Furthermore, for the low energy $U(1)'$ model with a secluded $U(1)'$-breaking sector, we briefly gave a scenario where the sterile neutrino masses are also generated during late time phase transition.

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