LIGHT SINGLET FERMIONS AND NEUTRINO PHYSICS

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\begin{abstract}

The existence of a light singlet fermion mixed with the electron neutrino is hinted by the simultaneous explanation of various neutrino anomalies. We show that supersymmetry can provide a natural framework for the existence and the desired properties of such a fermion. Quasi Goldstone fermions (QGF) of spontaneously broken global symmetries like the Peccei-Quinn symmetry or lepton number can mix properly with the neutrinos provided the presence of the $R$-parity breaking term $\epsilon L H_2$. The lightness of QGF can be a consequence of non-minimal Kähler potentials like that of no-scale supergravity. In order to keep $R$-parity, such a sterile component has to be placed in a new singlet superfield with no vacuum expectation value. In the context of the standard seesaw mechanism the lightness of such a singlet can be understood by imposing a $R$-symmetry.

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\end{abstract}
1 Introduction

All the experimentally known fermions transform non-trivially under the gauge group $SU(3) \times SU(2) \times U(1)$ of the standard model (SM). However there are experimental hints in the neutrino sector which suggest the existence of $SU(3) \times SU(2) \times U(1)$ - singlet fermions mixing appreciably with the known neutrinos. These hints come from (a) the deficits in the solar [1] and atmospheric [4] neutrino fluxes (b) possible need of significant hot component [3] in the dark matter of the universe and (c) some indication of $\bar{\nu}_e - \bar{\nu}_\mu$ oscillations in the laboratory [4]. These hints can be reconciled with each other if there exists a fourth very light ($< \mathcal{O}(\text{eV})$) neutrino mixed with some of the known neutrinos preferably with the electron one. The fourth neutrino is required to be sterile in view of the strong bounds on number of neutrino flavours coming both from the LEP experiment and from the primordial nucleosynthesis [3].

The existence of a very light sterile neutrino demands theoretical justification since unlike the active neutrinos, the mass of a sterile state is not protected by the gauge symmetry of the SM and hence could be very large. Usually a sterile neutrino is considered on the same footing as the active neutrinos and some ad hoc symmetry is introduced to keep this neutrino light. Recently there are several attempts to construct models for sterile neutrinos which have their origin beyond the usual lepton structure [6, 7, 8, 9, 10].

In this report, we discuss the role of supersymmetry (SUSY) in explaining both the existence and the lightness of a singlet fermion $S$ which can mix with the neutrinos. As a case of special interest we will concentrate on the mass of $S$ and its mixing with the electron neutrino in the range:

$$m_S \approx (2 - 3) \cdot 10^{-3} \text{eV}$$
$$\sin \theta_{es} \approx \tan \theta_{es} \approx (2 - 6) \cdot 10^{-2}.$$  \hspace{1cm} (1)

These values of parameters allow one to solve the solar neutrino problem through the resonance conversion $\nu_e \rightarrow S$ [11]. More discussions on simultaneous reconciliations of the diverse neutrino problems can be found in refs. [3, 4] on which this report is based.

2 Quasi Goldstone Fermion

The existence of SM-singlet fields is a common property in physics beyond the standard model. The most interesting examples are the Goldstone bosons of spontaneously broken
global symmetries required to solve the strong CP problem (the Peccei-Quinn symmetry) \[12\] and to explain the origin of neutrino masses (the lepton number symmetry) \[13\]. In the SUSY limit, a spontaneously broken global symmetry automatically generates a massless singlet (Goldstone) fermion being a superpartner of a Goldstone boson. However, SUSY breakdown results in generation of mass of a Goldstone fermion. While the existence of these quasi Goldstone fermions (QGF) is logically independent of neutrino physics, there are good reasons to expect that these fermions will couple to neutrinos. Indeed, in the case of lepton number symmetry the superfield which is mainly responsible for the breakdown of the lepton number symmetry carries nontrivial lepton number and therefore it can directly couple to leptons if the charge is appropriate. In the case of the PQ symmetry, this superfield could couple to the Higgs supermultiplet. If theory contains small violation of $R$-parity then this mixing with the Higgs gets communicated to the neutrino sector. Thus the occurrence of a QGF can have implications for neutrino physics. In the following subsections we elaborate upon the expected properties of the QGF: their masses arising after SUSY breaking and the mixing of these fermions with the electron neutrino.

2.1 masses of QGF

The supersymmetric standard model with some global symmetry $U(1)_G$ can be characterized by the following superpotential:

$$ W = W_{\text{MSSM}} + W_S + W_{\text{mixing}} ,$$

where $W$ is assumed to be invariant under $U(1)_G$. As we outlined in the above, this symmetry may be identified with the PQ symmetry, lepton number symmetry or combination thereof. The first term in eq. (2) refers to the superpotential of the minimal supersymmetric standard model (MSSM). The second term contains $SU(3) \times SU(2) \times U(1)$ singlet superfields which are responsible for the breakdown of $U(1)_G$. The minimal choice for $W_S$ is

$$ W_S = \lambda (\sigma \sigma' - f_G^2) y ,$$

where $\sigma, \sigma'$ carry non trivial $G$-charges and $f_G$ sets the scale of $U(1)_G$ breaking. The last term of eq. (2) describes mixing of the singlet fields with the superfields of the MSSM.

In the case (3) the Goldstone fermion is contained in $S \sim \sigma - \sigma'$ and is massless in the SUSY limit. Broken SUSY itself cannot automatically protect the mass of a QGF. It depends
on the structure of the superpotential $W_S$ \cite{14} and on the pattern of soft-terms \cite{15}. It also depends on the way this breaking is communicated to the singlet $S$ and the scale $f_G$ \cite{7}. The most natural framework for light QGF is no-scale supergravity \cite{16}. No-scale models contain only one kind of soft-terms, namely, gaugino masses. Therefore, the soft SUSY-breaking terms corresponding to $W_S$ in eq. (2) are absent at tree-level and thus QGF remains massless. However, the radiative mass can be triggered by the $SU(3) \times SU(2) \times U(1)$ gaugino masses through a set of interactions. A realistic example can be found in the context of the seesaw mechanism. The vacuum expectation value (VEV) of the field $\sigma$ (or $\sigma'$) may give rise to large masses of right-handed (RH) neutrinos $N$ as in the following superpotential invariant under $U(1)_G$:

$$W = \frac{m^D}{\langle H_2 \rangle} L N H_2 + \frac{M}{f_G} N N \sigma ,$$

where we have omitted the generation indices. The generation structure of the superpotential (4) will depend on the $U(1)_G$-charge assignment to the fields \cite{4}. This $U(1)_G$ symmetry is not necessarily the lepton number symmetry as we will discuss in subsection 2.2. The first term in eq. (4) gives rise to the Dirac masses of the neutrinos, whereas the second one gives the Majorana masses of RH neutrino components. The scale $f_G \sim 10^{10} - 10^{12}$ GeV generates $M \sim 10^{10} - 10^{11}$ GeV required by the hot dark matter and atmospheric neutrinos. If the soft-term $A_N N N \sigma$ with $A_N \sim m_{3/2}$ is present, there appears one-loop mass of the QGF proportional to $A_N$ \cite{17}. But in no-scale models $A_N = 0$ at tree-level and the QGF mass is indeed generated in three loops as shown in Figure 1. This three-loop mass can be estimated

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{Figure1.png}
\caption{Three-loop diagram for the QGF mass. The cross with $m_{1/2}$ denotes gaugino mass insertion.}
\end{figure}
as
\[ m_S \simeq \frac{\alpha_2}{(4\pi)^5} \frac{m_\nu M^3}{v^2 f_G} m_{1/2} . \] (5)

Here \( \alpha_2 \) and \( m_{1/2} \) are the \( SU(2) \) fine structure constant and gaugino mass respectively. For \( m_\nu \simeq 3 \text{ eV}, m_{1/2} \simeq v_2 \simeq 100 \text{ GeV}, \) and \( f_G \simeq 10^{12} \text{ GeV}, \) one gets \( m_S \simeq 3 \cdot 10^{-3} \text{ eV} \) with a value of \( M \simeq 10^{10} \text{ GeV}. \)

A contribution to the mass of the QGF can follow also from interactions, \( W_{\text{mixing}} \), which mix \( S \) with usual neutrinos (subsection 2.2).

### 2.2 Neutrino-QGF mixing

We now discuss how the QGF can mix with neutrinos. Such a mixing implies the violation of \( R \)-parity conventionally imposed in the MSSM \[18\]. This is simply because that the leptons being ordinary matter fields are \( R \)-even and the QGF being a fermionic partner of a Goldstone boson is \( R \)-odd. The violation of \( R \)-parity may destabilize the lightest supersymmetric particle (LSP) which is usually considered as the cold dark matter (CDM) of the Universe. For this reason, we consider the PQ symmetry as a good candidate for \( U(1)_G \) since the coherent oscillation of the axion can provide the CDM for \( f_{PQ} \sim 10^{12} \text{ GeV} \[19\]. Therefore, the PQ mechanism required for a resolution of the strong CP problem can supply both the CDM and the sterile neutrino.

The best way to implement the PQ symmetry in the MSSM is to extend the Higgs mass term in such a way that the smallness of the Higgs mass parameter \( \mu \) can be naturally obtained. For instance, let us consider the non-renormalizable term \[20\]
\[ \lambda H_1 H_2 \sigma^2 M_P , \] (6)

where \( M_P \) is the Planck mass \[.\] Here the VEV of \( \sigma, \langle \sigma \rangle \sim f_{PQ} \), spontaneously breaks the PQ symmetry. In this case, \( \mu = \lambda \frac{\langle \sigma \rangle^2}{M_P} \) can be about the weak scale. When the axion superfield \( S \) is predominantly consists of \( \sigma \), the PQ symmetry breaking yields the Higgs mass term and the coupling of \( S \) to the Higgs superfields
\[ W_{\text{mixing}} = c_\mu \frac{\mu}{f_{PQ}} H_1 H_2 S + \mu H_1 H_2 \] (7)

\[1\text{One can also introduce the renormalizable term to generate } \mu \simeq m_{3/2} \[21\].\]
with $c_\mu$ being $O(1)$. In order to have the mixing of $S$ with neutrinos, one needs the lepton number violating term $\epsilon LH_2$. It is remarkable to notice that the PQ scale is in the right range for the RH neutrino masses. The PQ symmetry can indeed play a role of the lepton number symmetry if both the Higgs and leptons transform non-trivially under the PQ symmetry as in ref. [22]. In this case one can correlate the origin of $\epsilon$ and $\mu$ to the same symmetry breaking scale $f_{PQ}$. The neutrino and Higgs coupling to QGF is then given by

$$W_{\text{mixing}} = \mu H_1 H_2 + \epsilon L_e H_2 + \frac{c_\mu}{f_{PQ}} H_1 H_2 S + \frac{\epsilon}{f_{PQ}} L_e H_2 S,$$

where $L_e$ is the electron doublet. If the PQ symmetry is the standard one unrelated to the lepton sector, the parameter $\epsilon$ vanishes. On the other hand, the global $U(1)$ symmetry becomes the usual lepton number symmetry when $c_\mu = 0$ and the bare $\mu$-term is introduced.

An example of models which leads to the mixing terms of eq. (8) can be obtained by the PQ-charge prescription $(−1,−1, 1,−1,−2)$ for $(H_1, H_2, \sigma, \sigma', L_e)$. It permits the following $U(1)_{PQ}$ invariant superpotential:

$$W = \lambda (\sigma \sigma' - f_{PQ}^2) y + \frac{\delta_\mu}{M_P} H_1 H_2 \sigma^2 + \frac{\delta_\epsilon}{M_P^2} L_e H_2 \sigma^3,$$

which gives the terms displayed in eq. (8) with $c_\mu = \frac{3}{\sqrt{2}}$, $c_\epsilon = \sqrt{2}$.

The $W_{\text{mixing}}$ in eq. (8) generates the following effective mass matrix for $\nu_e$ and $S$

$$
\begin{pmatrix}
0 & (c_\epsilon - c_\mu) \epsilon v \sin \beta / f_{PQ} \\
(c_\epsilon - c_\mu) \epsilon v \sin \beta / f_{PQ} & m_S^0 - c_\mu^2 \mu v^2 \sin 2\beta / f_{PQ}^2
\end{pmatrix},
$$

where we added the direct mass $m_S^0$ which can be generated by the mechanism of subsection 2.1. According to eq. (10) the $\nu_e - S$ mixing angle $\theta_{es}$ is determined by

$$\tan \theta_{es} \sim \frac{(c_\mu - c_\epsilon) \epsilon v \sin \beta}{m_S^0 f_{PQ} - c_\mu^2 \mu v^2 \sin 2\beta / f_{PQ}^2}.$$

For $f_{PQ} \simeq 10^{12}$ GeV, $m_S^0 \simeq 3 \cdot 10^{-3}$ eV is the dominant contribution to the mass of $S$. In this case one obtains from eq. (11) for the $\nu_e - S$ mixing

$$\tan \theta_{es} \sim \frac{\epsilon v \sin \beta}{m_S^0 f_{PQ}}.$$
Then the desired value, $\tan \theta_{es} \sim (2-6) \cdot 10^{-2}$ eV (1), can be obtained if the $R$-parity breaking parameter $\epsilon$ equals

$$\epsilon \sim \frac{m^0_S}{v \sin \beta} \approx (2-6) \cdot 10^{-16} \frac{f_{PQ}}{\sin \beta} .$$

(13)

For $f_{PQ} \sim 10^{12}$ GeV one has $\epsilon \sim 0.1$ MeV.

Let us remark the other possibilities for the QGF mass. If $m^0_S = 0$ in eq. (10), the QGF mass, $m_S = (2-3) \cdot 10^{-3}$ eV can be obtained for the marginally allowed value of the PQ scale:

$$f_{PQ} \approx v \sqrt{\frac{\mu \sin 2\beta}{m_S}} \lesssim 4 \cdot 10^9 \text{ GeV} .$$

(14)

For $f_{PQ} > 10^{10}$ GeV the QGF mass generated via $\mu$-term is too small for the MSW solution. For $f_{PQ} \sim 10^{11}$ GeV, $m_s \approx 10^{-5}$ eV is in the region of “just-so” solution of the solar neutrino problem. In these cases, however, axions cannot provide the CDM as we noted before.

3 A light singlet in the standard seesaw structure

In the previous case, the QGF mixes with the electron neutrino directly ($c \epsilon_c \neq 0$) or via its coupling to the Higgses ($c \mu \neq 0$). The small mass of the QGF was related to the multi-loop effect or the suppression by $1/f^{2}_{PQ}$ due to the Goldstone property. An important consequence was the $R$-parity violation leading to destabilization of the LSP.

In this section, we will suggest another scheme in which $R$-parity is preserved. For this, one should place the singlet $S$ in the superfield with zero VEV. This implies that the singlet has to be introduced from outside. Being a singlet $S$ can mix with neutrinos via its coupling to the right-handed neutrinos. In this case, the existence of $S$ cannot be explained but the smallness of its mass can be understood in terms of the seesaw mechanism. In order to implement a light singlet fermion in the standard seesaw structure, we will suggest to use $R$-symmetry which occurs in many SUSY theories. The (unbroken) $R$-parity is then embedded in the $R$-symmetry.

Let us first determine the parameters appearing in the phenomenological superpotential

$$W = \frac{m_e}{\langle H_2 \rangle} L_e N_e H_2 + \frac{M_e}{2} N_e N_e + m_{es} N_e S ,$$

(15)

where $N_e$ is the right-handed neutrino component. The Dirac mass $m_e$ and the mixing mass $m_{es}$ are much smaller than the Majorana mass $M_e$: $m_e, m_{es} \ll M_e$. The superpotential (15)
leads to the mass matrix in the basis \((S, \nu_e, N_e)\):

\[
M = \begin{pmatrix}
0 & 0 & m_{es} \\
0 & 0 & m_e \\
m_{es} & m_e & M_e
\end{pmatrix}.
\] (16)

The diagonalization of (16) is straightforward. One combination of \(\nu_e\) and \(S\) is massless and the orthogonal combination acquires a mass via the see-saw mechanism:

\[
m_1 \simeq -\frac{m_e^2 + m_{es}^2}{M_e}.
\] (17)

The mass of the heavy neutrino is \(\simeq M_e\). The \(\nu_e-S\) mixing angle is determined by

\[
\tan \theta_{es} = \frac{m_e}{m_{es}}.
\] (18)

Taking for \(m_e\) the typical Dirac mass of the first generation: \(m_e \sim (1-5)\) MeV, and suggesting that \(\nu_e \rightarrow S\) conversion explains the solar neutrino problem with \(m_1 = m_S\) as in (1), we find

\[
m_{es} = \frac{m_e}{\tan \theta_{es}} \simeq (0.02 - 0.3) \text{ GeV}.
\] (19)

According to (17) the RH mass scale is

\[
M_e \simeq \frac{m_{es}^2}{m_1} = \frac{m_e^2}{m_1 \tan^2 \theta_{es}} \simeq (10^8 - 3 \cdot 10^{10}) \text{ GeV}.
\] (20)

One has now to understand how the mixing mass (19) arises without introducing new mass scales. One also has to ensure that there is no direct coupling of \(S\) with \(L_e\), and the mass term \(SS\) is absent or negligibly small.

Our prescription is quite simple. Consider the superpotential

\[
W = \frac{m_e}{\langle H_2 \rangle} L_e N_e H_2 + f N_e N_e \sigma + f' N_e S y - \frac{\lambda}{2} (\sigma^2 - M^2)y.
\] (21)

whose structure is determined by the \(R\)-symmetry under which the fields \((L_e, N_e, S, y, \sigma, H_2)\) carry the \(R\)-charges \((1, 1, -1, 2, 0, 0)\). Note that the \(R\)-symmetry forbids the bare mass terms \(SS\) as well as the coupling \(SS\sigma\). The last term in eq. (21) can be replaced by \((\sigma \sigma' - M^2)y\) to implement the lepton number symmetry. In the global SUSY limit, \(\sigma\) gets non-zero VEV \(\langle \sigma \rangle \simeq M \sim 10^{11} \) GeV which generates the Majorana mass of \(N_e\): \(M_e = f \langle \sigma \rangle\). The point
is that $y$ develops a VEV as a consequence of SUSY breaking. Broken SUSY produces the following soft-breaking terms in the scalar potential:

$$V_{\text{soft}} = \{ A_L \frac{m_e}{(H_2)} L_e N_e H_2 + f A_Y N_e N_e \sigma + f' A_S N_e S y - \frac{\lambda}{2} (A_y \sigma^2 - B_y M^2) y + \text{h.c.} \} + \sum_i m_i^2 |z_i|^2 ,$$  \hspace{1cm} (22)

where $z_i$ denotes the fields appearing in the superpotential (21) and $A_L$, etc., are the soft-breaking parameters. Minimization of the potential shows the following: (1) The fields $L_e, N_e, S$ do not develop VEV and therefore $R$-parity is unbroken. (2) The field $y$ acquires non-zero VEV due to the soft-breaking terms. Consequently, the mixing mass for $S$ and $N_e$ appears:

$$m_{es} = \frac{f'}{2\lambda} (A_y - B_y)$$  \hspace{1cm} (23)

Since $m_{es} \gg m_1$, no strong tunning of $A_y - B_y$ is needed. For $A_y - B_y \sim O(m_{3/2})$, the desired value of $m_{es}$ (13) can be obtained by choosing $f'/\lambda \sim 10^{-3} - 10^{-2}$. However, more elegant possibility is that $A_y = B_y$ at tree level but a non-zero value for $A_y - B_y$ is generated due to radiative corrections through the differences in interactions of $\sigma$ and $y$. In this case one expects

$$m_{es} \sim \frac{\bar{\lambda}^2}{16\pi^2} m_{3/2} ,$$  \hspace{1cm} (24)

where $\bar{\lambda}$ represents a combination of the constants $\lambda, f$ and $f'$. As a consequence, the value $m_{es} \sim 0.1$ GeV does not require smallness of $\bar{\lambda}$ or $f'$.

The equality $A_y = B_y$ at tree level can be achieved by the introduction of non-minimal K"ahler potential allowing mixings between the observable and hidden sectors. Let us introduce the following K"ahler potential:

$$K = CC + CC(a\frac{Z}{M_{Pl}} + \tilde{a}\frac{\tilde{Z}}{M_{Pl}}) + ZZ ,$$  \hspace{1cm} (25)

where $C$ and $Z$ represent an observable and hidden sector field, respectively. Then usual assumption that the observable sector has no direct coupling to the hidden sector in superpotential, $W = W(C) + W(Z)$, leads to the universal soft-terms:

$$V_{\text{soft}} \sim m_{3/2} W(C) + \text{h.c.} ,$$  \hspace{1cm} (26)

provided $\tilde{a} = \langle W(Z) \rangle/(M_{Pl}\partial W/\partial Z + W(Z)\tilde{Z}/M_{Pl})$. Note also that the field $C$ does not acquire a soft-breaking mass. This mechanism can be generalized to arbitrary number of
observable sector fields. For our purpose $C \equiv \sigma, y$, i.e., we couple $\sigma$ and $y$ to the hidden sector field $Z$ with the above-mentioned choice for $a$.

4 Conclusions

Simultaneous presence of different neutrino anomalies points to the existence of a sterile neutrino. In particular, the resonance conversion of the electron neutrino into such a singlet fermion $S$ can explain the solar neutrino problem provided its mass and mixing are appropriate \[1\). Supersymmetry is shown to provide a framework within which the existence and the desired properties of such a light fermion follow naturally.

We have considered first a possibility that the sterile neutrino is a quasi Goldstone fermion appearing in supersymmetric theories as a result of spontaneous breaking of a global $U(1)_G$ symmetry. This global $U(1)_G$ symmetry can be identified with the PQ symmetry, the lepton number symmetry. The smallness of $m_S$ can be attributed in supergravity theory to no-scale kinetic terms for certain superfields. The mixing of QGF with the neutrinos implies spontaneous or explicit violation of $R$-parity. QGF can mix with neutrino via interaction with Higgs multiplets (in the case of PQ symmetry) or directly via coupling with the combination $LH_2$ (in the case of lepton number symmetry). In the case of the PQ symmetry, the PQ-scale $f_{PQ} \sim 10^{10} - 10^{12}$ GeV determines several features of the model presented here. It provides simultaneous explanation of the parameters $\epsilon$ and $\mu$ and thus leads to small $R$-parity violation ($\epsilon LH_2$ with $\epsilon \sim 0.1$ MeV) required in order to solve the solar neutrino problem in our approach. It also provides the intermediate scale for the right-handed neutrino masses which is required in order to solve the dark matter and the atmospheric neutrino problem. Furthermore, it controls the magnitude of the radiatively generated mass of the QGF and allows it to be in the range needed for the MSW solution of the solar neutrino problem. Finally, the CDM can consist of the axion if $f_{PQ} \sim 10^{12}$ GeV. Thus the basic scenario presented here is able to correlate variety of phenomena.

The conservation of $R$-parity requires for the fermion $S$ to be a component of singlet superfield which has no VEV. This allows to construct simple model \[21\) in which the properties (mass and mixing) of $S$ follow from the conservation of $R$-symmetry. The singlet field is mixed with RH neutrinos by the interaction with the field $y$ which can acquire VEV radiatively after soft SUSY breaking.
Let us finally comment on the other phenomenological consequences of the existence of such a sterile state $S$. An $U(1)_G$ symmetry being generation-dependent \cite{6, 7} can provide simultaneous explanations for the predominant coupling of $S$ to the first generation (thus satisfying the nucleosynthesis bound) and for the pseudo-Dirac structure of $\nu_\mu - \nu_\tau$ needed in solving the atmospheric neutrino and the hot dark matter problem. In this case, it appears nontrivial to accommodate the parameters of $\bar{\nu}_\mu \rightarrow \bar{\nu}_e$ oscillations in the region of sensitivity of LSND and KARMEN experiments. The simplest way is to introduce a slight violation of the $U(1)_G$ symmetry through which such parameters can be incorporated.

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