Quasi Goldstone fermion as a sterile neutrino

Eung Jin Chun
International Center for Theoretical Physics, P.O. Box 586, 34100 Trieste, Italy

Anjan S. Joshipura
Theoretical Physics Group, Physical Research Laboratory, Navarangpura, Ahmedabad, 380 009, India

Alexei Yu. Smirnov
International Center for Theoretical Physics P.O. Box 586, 34100 Trieste, Italy
and Institute for Nuclear Research, Russian Academy of Sciences, 117312 Moscow, Russia
(Received 31 July 1995)

The existence of sterile neutrinos is hinted by the simultaneous presence of diverse neutrino anomalies. We suggest that the quasi Goldstone fermions (QGF's) arising in supersymmetric theory as a result of spontaneous breaking of global symmetry such as the Peccei-Quinn symmetry or the lepton number symmetry can play the role of the sterile neutrino. The smallness of the mass of QGF's ($m_S \sim 10^{-3}$–10 eV) can be related to the specific choice of superpotential or Kähler potential (e.g., no-scale kinetic terms for certain superfields). Mixing of QGF's with neutrinos implies R-parity violation. It can proceed via the coupling of QGF's with Higgs supermultiplets or directly with a lepton doublet. A model which accounts for the solar and atmospheric anomalies and dark matter is presented. [S0556-2821(96)00217-2]

PACS number(s): 14.60.St, 14.60.Pq, 26.65.+t, 96.60.Jw

I. INTRODUCTION

All the experimentally known fermions transform nontrivially 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 [2] neutrino fluxes, (b) the possible need for a significant hot component [3] in the dark matter of the Universe, and (c) some indication of $\overline{\nu_e}$-$\overline{\nu_\mu}$ oscillations in the laboratory [4]. These hints can be reconciled with each other if there exists a fourth very light ($\lesssim 1$ 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 the number of neutrino flavors coming both from the experiments at the CERN e^+e^- collider LEP as well as from primordial nucleosynthesis [5].

The existence of a very light sterile neutrino demands theoretical justification since, unlike the active neutrinos, the mass of the sterile state is not protected by the gauge symmetry of the SM and, hence, could be very large. Usually the 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 have been several attempts to construct models for sterile neutrinos which have their origin beyond the usual lepton structure [6–8]. In particular, in Ref. [6] we suggested the possibility that supersymmetry (SUSY) may be responsible for both the existence and the lightness of the sterile fermions.

One could consider three different ways in which supersymmetry can keep sterile states very light: (1) The combination of supersymmetry and the (continuous) R symmetry present in many supersymmetric models may not allow a mass term for the light sterile state; (2) the spontaneous breakdown of some other global symmetry in supersymmetric theory can lead to massless fermions which form the superpartners of the Goldstone bosons; (3) the spontaneous breakdown of the global supersymmetry itself would give rise to a massless fermion, the Goldstino.

Mechanism (1) and its phenomenological consequences were discussed in Ref. [6]. Mechanism (3) though appealing is not favored phenomenologically in view of the difficulties in building realistic models based on the spontaneously broken global SUSY. We discuss in this paper implications of the mechanism (2) concentrating for definiteness on the simplest case of a global $U(1)_G$.

The spontaneously broken global symmetries are required for reasons unrelated to the existence of light sterile states. The most interesting examples are the spontaneously broken lepton number symmetry [9] and the Peccei-Quinn (PQ) symmetry imposed [10] to solve the strong CP problem. PQ symmetry arises naturally in many supersymmetric models. Apart from solving the strong CP problem, this symmetry can also explain the smallness of the μ parameter [11,12]. Phenomenologically consistent breaking of these symmetries generally needs [13] Higgs fields which are singlets of $SU(3)\times SU(2)\times U(1)$. In the supersymmetric context this automatically generates a massless sterile fermion. While the existence of these quasi Goldstone fermions (QGF's) 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 $U(1)_L$ carries a nontrivial $U(1)_L$ charge and therefore it can directly couple to leptons if the charge is appropriate. In the case of the PQ symmetry $U(1)_{PQ}$, this superfield could couple to the Higgs supermultiplet. If theory contains small violation of R parity then this mixing with Higgs supermultiplet gets communicated to the neutrino sector. Thus the occurrence of the QGF can have implications for neutrino physics. We wish to discuss in this paper prospects for building realistic models based on this mechanism.

In the following section we elaborate upon the expected properties of the QGF's, especially their masses when SUSY is broken. Section III discusses various mechanisms of mixing of these fermions with the active neutrinos. An explicit model based on the scenario presented in Secs. II and III is given in Sec. IV and the last section presents our conclusions.

II. QUASI GOLDSTONE FERMIONS AND THEIR MASSES

In this section and subsequently, we will consider the general superpotential

$$W = W_{\rm MSSM} + W_S + W_{\rm mixing}, \tag{1}$$

where W is assumed to be invariant under some global symmetry $U(1)_G$. As we outlined in the Introduction, this symmetry may be identified with the PQ symmetry, lepton number symmetry, or combination thereof. The first term in Eq. (1) 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

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

where σ, σ' carry nontrivial G charges and f_G sets the scale of $U(1)_G$ breaking. The last term of Eq. (1) describes mixing of the singlet fields with the superfields of the MSSM.

In the supersymmetric limit the fermionic component of the Goldstone boson is massless. In the case (2) this Goldstone fermion is contained in

$$S = \frac{1}{\sqrt{2}}(\sigma - \sigma'). \tag{3}$$

The SUSY breakdown results in generation of the mass of the Goldstone fermion. In general, this mass can be as big as the SUSY breaking scale, $m_{\rm SUSY}$. Broken supersymmetry itself cannot automatically protect the masses of QGF in Eq. (3) much below $m_{\rm SUSY}$. However as discussed in Ref. [17] the mass resulting after the SUSY breaking is quite model dependent. It depends on the manner in which SUSY is broken and on the way this breaking is communicated to the singlet S. It also depends on the structure of superpotential and the scale f_G. Below we identify theories which can allow for very light QGF ($m_S < 1$ eV). As a case of special interest we will consider the mass of QGF and its mixing with the electron neutrino in the range

$$m_S \simeq (2-3)\times 10^{-3}\ \text{eV},$$

$$\sin\theta_{es} \simeq \tan\theta_{es} \simeq (2-6)\times 10^{-2}. \tag{4}$$

These values of parameters allow one to solve the solar neutrino problem through the resonance conversion $\nu_e \to S$ [14].

One could consider different mechanisms for the QGF mass generation. In models with spontaneously broken global SUSY the QGF generically acquires a mass of $O(m_{\rm SUSY}^2/f_G)$ [15]. But it can remain massless in spite of SUSY breaking (a) if SUSY is broken by a D term of the gauge field or (b) if the F terms that break SUSY do not carry any G charges. The latter is exemplified by a simple generalization of Eq. (2):

$$W_S = \lambda_1(\sigma\sigma' - f_1^2)y_1 + \lambda_2(\sigma\sigma' - f_2^2)y_2.$$

SUSY is broken in this example if $f_1^2 \neq f_2^2$. For a minimum with the F terms: $F_\sigma = F_{\sigma'} = 0$, the Goldstone fermion in Eq. (3) remains massless at the tree level in spite of the SUSY breakdown. As we noticed before this version has phenomenological problems and further on we will concentrate on possibilities related to supergravity.

The mass of the QGF in supergravity theory is typically of the order of gravitino mass $m_{3/2}$ ($= m_{\rm SUSY}$) [16–18]. For instance, the superpotential in Eq. (2) leads to $m_S \sim m_{3/2}$ when generic SUSY breaking soft terms are allowed [16]. However, the mass m_S can be much smaller for specific choices of (1) the superpotential and/or (2) soft SUSY breaking terms. Let us consider these possibilities in order.

(1) The superpotential

$$\lambda(\sigma\sigma' - X^2)y + \lambda'(X - f_G)^3$$

is shown [17] to generate the tree-level mass

$$m_S \sim \frac{m_{3/2}^2}{f_G} \tag{5}$$

as in the global case if the minimal kinetic terms of the fields are assumed. For the commonly accepted value of the PQ symmetry breaking scale, $f_G = f_{\rm PQ} = 10^{10} - 10^{12}$ GeV, one gets from Eq. (5) $m_S \sim (10 - 10^3)$ eV. On the other hand, the value of m_S in Eq. (4) desired for explanation of the solar neutrino deficit requires $f_G \sim 10^{16}$ GeV which can be related to the grand unification scale. To identify f_G with $f_{\rm PQ}$, one should overcome the cosmological bound $f_{\rm PQ} < 10^{12}$ GeV. The bound can be removed by axion mixing with some other Goldstone boson in their kinetic terms [19] or by a dilaton field driven to small values in an inflationary period [20]. In this case however, the axion cannot play the role of cold dark matter.

(2) Another possibility to get very light S is based on the idea of no-scale supergravity [21]. The Kähler potential and the superpotential can be arranged in such a way that supersymmetry breaking is communicated to the singlet S via a set of interactions. As a result, the mass of S appears in one, two, or even three loops.

Let us consider the Kähler potential

$$K = -3\ln(T + T^* - Z_a Z_a^*) + C_i C_i^*, \tag{6}$$

where T is the moduli field appearing in the underlying superstring theory, Z_a and C_i are the matter superfields which have the no-scale kinetic term (Z sector) and the minimal kinetic term (C sector), respectively. The corresponding scalar potential at the Planck scale reads

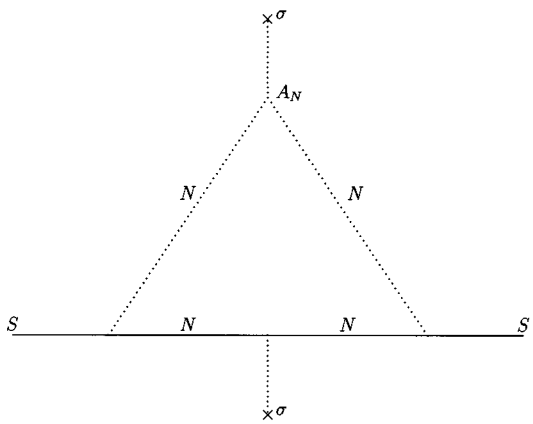

FIG. 1. One-loop diagram for the QGF mass. The solid lines are fermions and the dotted lines are bosons. A_N is the soft parameter of $NN\sigma$.

$$V=|W_i|^2+\{m_0C_iW_i+\text{ H.c.}\}+m_0^2|C_i|^2+|W_a|^2, \tag{7}$$

where $m_0=O(m_{3/2})$. The tree-level masses of the fermionic components of the fields Z_a are determined by the global supersymmetric results. Therefore, if the singlet fields triggering $U(1)_G$ breaking are in the Z sector, the QGF will be massless at the tree level [18]. The QGF will acquire a mass through the interactions with fields C_i having minimal kinetic terms, and consequently, the usual soft SUSY breaking terms. Moreover, S (or σ,σ') may not couple to C_i directly. It can interact with C_i via couplings with some other fields Z_a having no-scale kinetic terms. In this case S will get the mass in two or larger number of loops.

Let us consider realizations of this idea in the context of the seesaw mechanism, when σ,σ' couple with right-handed (RH) neutrinos N. Let us introduce the following terms in the superpotential:

$$W=\frac{m^D}{v_2}LNH_2+\frac{M}{f_G}NN\sigma+\lambda(\sigma\sigma'-f_G^2)y. \tag{8}$$

We will specify the generation structure of these terms in Sec. IV, when we describe a concrete model. The first term in Eq. (8) gives rise to the Dirac mass of the neutrino, whereas the second one gives the Majorana mass of the RH neutrino component. The scale $f_G\sim10^{10}-10^{12}$GeV generates $M\sim10^{10}-10^{11}$ GeV required by the hot dark matter (HDM) and atmospheric neutrinos.

(i) Suppose that only σ,σ',y superfields belong to the Z sector, whereas all other superfields have minimal kinetic terms: $N,H_2,L\in C$. Then SUSY breaking induces the soft term

$$A_N\frac{M}{f_G}\tilde{N}\tilde{N}\sigma \tag{9}$$

which generates the mass of QGF in one loop (Fig. 1):

$$m_S\simeq\frac{1}{16\pi^2}\left(\frac{M}{f_G}\right)^2A_N. \tag{10}$$

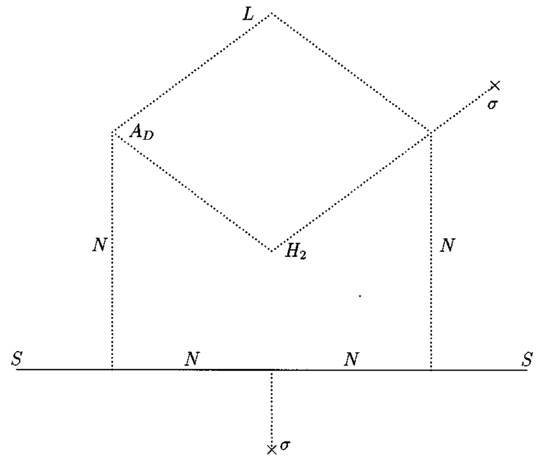

FIG. 2. Two-loop diagram for the QGF mass. A_D is the soft parameter of LNH_2.

This mechanism is similar to that of the axino mass generation by coupling of S with heavy quarks [18,22]. For $A_N=O(m_{3/2})$ and $(M/f_G)\sim10^{-3}$, m_S is in the keV range.

(ii) Let us suppose that not only σ,σ',y but also N have the no-scale kinetic terms. In this case $A_N=0$ at tree level, but nonzero A_N will be generated in one loop (see Fig. 2) by the soft breaking term related to the usual Yukawa interaction LNH_2: $A_Dm^D\tilde{L}\tilde{N}H_2$, and by the quartic coupling $\sigma\tilde{N}\tilde{L}^*H_2^*$ which follows from $|W_N|^2$ term of the supersymmetric scalar potential. As a result one has

$$A_N\sim\frac{1}{16\pi^2}\left(\frac{m^D}{v_2}\right)^2A_D. \tag{11}$$

Correspondingly, m_S appears in two loops (Fig. 2). Combining Eqs. (10) and (11) we get the estimation of m_S:

$$m_S\simeq\frac{1}{(16\pi^2)^2}\frac{A_DM^3}{v_2^2f_G^2}m_\nu. \tag{12}$$

Here $m_\nu=(m^D)^2/M$. For the HDM mass scale $m_\nu\simeq3$ eV, $A_D\simeq v_2\simeq100$ GeV and $f_G\simeq10^{12}$ GeV it follows from Eq. (12) that $m_S\simeq3\times10^{-3}$ eV can be achieved if the mass of the RH component is $M\simeq10^9$ GeV.

In this version of the model the left and right neutrino components have different kinetic terms which may look unnatural.

(iii) Finally we consider the case where all chiral superfields belong to the Z sector. This so-called strict no-scale model [23,24] has only one seed of SUSY breakdown (i.e., gaugino mass). In this case $A_D=0$ at the tree level and nonzero A_D is generated in one loop by gaugino exchange. Correspondingly, m_S appears in three loops (Fig. 3) and it can be estimated as

$$m_S\simeq\frac{\alpha_2}{(4\pi)^5}\frac{m_{1/2}M^3}{v_2^2f_G^2}m_\nu. \tag{13}$$

Here α_2 and $m_{1/2}$ are the SU(2) fine structure constant and gaugino mass, respectively. For $m_\nu\simeq3$ eV, $m_{1/2}\simeq v_2\simeq100$ GeV, and $f_G\simeq10^{12}$ GeV, one gets from Eq. (13) $m_S\simeq3\times10^{-3}$ eV with a value of $M\simeq10^{10}$ GeV.

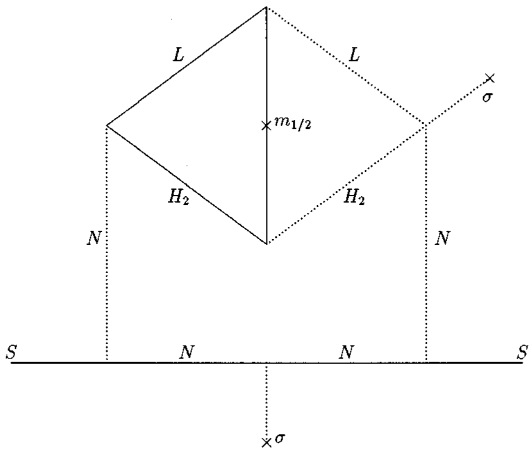

FIG. 3. Three-loop diagram for the QGF mass. The cross with $m_{1/2}$ denotes the gaugino mass insertion.

A contribution to the mass of the QGF can follow also from interactions, $W_{\rm mixing}$, which mix S with usual neutrinos (Sec. III).

III. NEUTRINO QGF MIXING

We now discuss possible ways which lead to mixing of the QGF with neutrinos. Such a mixing can occur only in the presence of either explicit or spontaneous violation of the R parity conventionally imposed in the MSSM [25]. Indeed, the Higgs field which breaks $U(1)_G$ may belong either to R even or odd superfield depending upon the nature of the $U(1)_G$. If it belongs to R even (i.e., Higgs-like) superfield then the corresponding QGF is R odd and its mixing with neutrinos implies the R violation. In contrast, if the QGF is R even, e.g., similar to the right-handed neutrino, then its scalar partner is R odd and the R symmetry gets broken together with the $U(1)_G$ symmetry. The first alternative is realized when the $U(1)_G$ is identified with the PQ symmetry. On the other hand, the lepton number symmetry containing a right-handed neutrinolike superfield would provide an example of the second alternative. We discuss these cases in turn.

(1) *PQ symmetry*. The supersymmetric theories with Peccei-Quinn symmetry may contain a term

$$\lambda H_1 H_2 \sigma, \tag{14}$$

with σ being a superfield transforming nontrivially under the PQ symmetry. If the axionic superfield, S, predominantly consists of the field σ, the vacuum expectation value (VEV) $\langle\sigma\rangle \sim f_{\rm PQ}$ would be large $\sim 10^{10}-10^{12}$ GeV. Since this VEV generates the parameter $\mu=\lambda\langle\sigma\rangle$ of the MSSM through the interaction (14), one would need to fine-tune λ in order to understand the smallness of μ. The coupling of axionic supermultiplet S to Higgs superfield is then given by

$$\frac{1}{\sqrt{2}}\frac{\mu}{f_{\rm PQ}} H_1 H_2 S. \tag{15}$$

The smallness of μ can be understood if σ couples to the Higgs field through a nonrenormalizable term [11]

$$\lambda H_1 H_2 \frac{\sigma^2}{M_P}, \tag{16}$$

where M_P is the Planck scale mass. In this case, $\mu=\lambda\langle\sigma\rangle^2/M_P$ is naturally about the weak scale. Since $f_{\rm PQ}\simeq\langle\sigma\rangle$, the axionic coupling following from Eq. (16) can be written as

$$\sqrt{2}\frac{\mu}{f_{\rm PQ}} H_1 H_2 S. \tag{17}$$

Alternatively, the σ may acquire a small VEV $\sim m_{3/2}$ and the scale of the PQ symmetry may be set by some other field which would predominantly contain the axionic multiplet [12]. The μ parameter is naturally of the order $m_{3/2}$ in this case. As long as the field σ transforms nontrivially under PQ symmetry, it will contain a small admixture $\sim\langle\sigma\rangle/f_{\rm PQ}$ of the axionic field S. The interaction in Eq. (14) results in the coupling

$$c_\mu \frac{\mu}{f_{\rm PQ}} H_1 H_2 S, \tag{18}$$

c_μ being of order 1.

It follows from Eqs. (15), (17), and (18) that the axionic coupling to the Higgs superfield is insensitive to mechanism of implementation of the PQ symmetry. We can, therefore, consider the generic effective term

$$c_\mu \frac{\mu}{f_{\rm PQ}} H_1 H_2 S + \mu H_1 H_2 + \epsilon L_e H_2. \tag{19}$$

Here we also have included the explicit R violating coupling $L_e H_2$ which will induce $\nu_e - S$ mixing. The effective terms (19) can be generated starting from $U(1)_G$ invariant superpotential as the result of $U(1)_G$ symmetry breaking. By this one can explain the smallness of the parameters μ, ϵ for certain G charges of the fields. Before constructing the model let us consider the phenomenology of Eq. (19).

The superpotential (19) leads to the following mass matrix in the basis (ν_e, S, h_1, h_2):

$$\begin{pmatrix} 0 & 0 & 0 & \epsilon \\ 0 & m_S^0 & c_\mu \mu v \sin\beta/f_{\rm PQ} & c_\mu \mu v \cos\beta/f_{\rm PQ} \\ 0 & c_\mu \mu v \sin\beta/f_{\rm PQ} & 0 & \mu \\ \epsilon & c_\mu \mu v \cos\beta/f_{\rm PQ} & \mu & 0 \end{pmatrix}, \tag{20}$$

where $v\equiv\sqrt{v_1^2+v_2^2}$ is the weak scale, $\tan\beta\equiv v_2/v_1$ and $v_{1,2}$ are the VEV's of $H_{1,2}$. In the matrix (20) we have included also the direct axino mass m_S^0 that can be generated by the mechanisms of Sec. II. We have neglected the contribution from the interactions with the gauginos in Eq. (20). In general gauginos mix with Higgsino through $v_{1,2}$. This mixing will not change the qualitative results which follow from Eq. (20). Moreover, the mixing can be small if the gaugino mass is chosen much larger than the μ parameter. Gauginos will also mix with neutrinos through the VEV of sneutrino field which may arise due to the presence of the ϵ coupling in Eq. (19) and soft SUSY breaking terms. This mixing generates [26] neutrino mass of order $g^2\langle\tilde{\nu}_e\rangle^2/m_{1/2}$ [g is the SU(2) coupling constant]. For $m_{1/2}>100$ GeV and $\langle\tilde{\nu}_e\rangle<10$ keV, this contribution is much smaller than $m_S^0\sim 3\times 10^{-3}$ eV which can result from the radiative corrections.[1] From the condition $m(\nu_e)<m_s^0$ implied by resonance conversion of solar neutrinos one gets (unless the masses of S and ν_e are strongly degenerate) the upper bound on the sneutrino VEV: $\langle\tilde{\nu}_e\rangle<0.1$ MeV. This in turn restricts ϵ which is of the order $\langle\tilde{\nu}_e\rangle$, although precise relation between them depends on SUSY breaking parameters. As a conservative bound one can use $\epsilon<0.1-1$ MeV.

Block diagonalization of the matrix (20) leads to the following effective mass matrix for the neutrino and the axino, (ν_e,S):

$$\begin{pmatrix} 0 & -c_\mu \epsilon v \sin\beta/f_{PQ} \\ -c_\mu \epsilon v \sin\beta/f_{PQ} & m_S^0-c_\mu^2\mu v^2\sin 2\beta/f_{PQ}^2 \end{pmatrix}. \tag{21}$$

If $m_S^0=0$ in Eq. (21), the QGF mass, $m_S=(2-3)\times 10^{-3}$ eV can be obtained for the marginally allowed value of the PQ scale:

$$f_{PQ}\approx v\sqrt{\frac{\mu\sin 2\beta}{m_S}}\lesssim 4\times 10^9 \text{ GeV}. \tag{22}$$

In this case, however, axions cannot provide the cold dark matter of the Universe. Note that the lightest supersymmetric particles cannot be cold dark matter either because of their instability due to the R-parity violation or due to their decay into the lighter axino. For $f_{PQ}>10^{10}$ GeV the QGF mass generated via μ 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. The axion can however serve as cold dark matter provided $f_{PQ}\sim 10^{12}$ GeV. In this case, the seesaw contribution to m_S is very small and one needs a nonvanishing mass m_S^0.

If m_S^0 is the dominant contribution to the mass of S, $m_S\simeq m_S^0$ one obtains, from Eq. (21) for the ν_e-S mixing,

$$\tan\theta_{es}\sim\frac{c_\mu \epsilon v\sin\beta}{m_S^0 f_{PQ}}. \tag{23}$$

Then the desired value, $\tan\theta_{es}\sim(2-6)\times 10^{-2}$ (4), can be obtained if the R-parity breaking parameter ϵ equals

$$\epsilon=\frac{m_S^0 f_{PQ}\tan\theta_{es}}{c_\mu v\sin\beta}\approx(2-6)\times 10^{-16}\frac{f_{PQ}}{\sin\beta}. \tag{24}$$

For $f_{PQ}\sim 10^{12}$ GeV one has $\epsilon\sim 0.1$ MeV. In general, the appropriate range of ϵ is $(10^{-3}-10)$MeV. It can be generated as a radiative correction: $\epsilon\sim h^2 m_{3/2}/16\pi^2$. Alternatively, ϵ may arise through the coupling of the product L_eH_2 to some fields carrying a nonzero lepton number. In this case the required smallness of ϵ may be understood in analogy with that of the μ parameter.

(2) *Lepton number symmetry.* Let us identify $U(1)_G$ with the lepton number symmetry. Unlike in the previous case, it is possible now to couple the QGF directly to neutrino through the term

$$hL_eH_2\sigma \tag{25}$$

or through a nonrenormalizable term analogous to Eq. (16). Equation (25) is similar to Eq. (14) but now the scalar component of σ is R odd and its VEV breaks R parity. Electroweak symmetry breaking $v_2\neq 0$ leads through the term (25) to the direct coupling between QGF's and neutrinos. Note that σ is similar to the RH neutrino components. Just as the interaction in Eq. (14) generates the μ, the interaction (25) generates the parameter ϵ. Thus it is possible to correlate their origin of ϵ to the breaking of the lepton number symmetry. The smallness of ϵ may be due to (i) fine-tuning of h or (ii) smallness of the VEV of σ or due to (iii) occurrence of the nonrenormalizable coupling analogous to that in Eq. (16). All these possibilities lead to the following effective coupling of ν to QGF:

$$c_\epsilon\frac{\epsilon}{f_L}L_eH_2S+\epsilon L_eH_2, \tag{26}$$

where f_L denotes the scale associated with the spontaneous breaking of the lepton number symmetry and c_ϵ is a parameter of order unity. The mass matrix generated by Eq. (26) is

$$\begin{pmatrix} 0 & c_\epsilon\epsilon v\sin\beta/f_L \\ c_\epsilon\epsilon v\sin\beta/f_L & m_S^0 \end{pmatrix}. \tag{27}$$

and the desired ν_e-S mixing can be obtained for $\epsilon\simeq 0.1$ MeV and $f_L\sim 3\times 10^{11}$ GeV.

Let us give an example of a model which leads to the mixing term of Eq. (26). Consider the $U(1)_L$ charge assignments $(1,-1,-3)$ for the fields (σ,σ',L), respectively. All other fields are taken neutral. The relevant part for the $U(1)_G$ invariant superpotential is given as

$$W=\lambda(\sigma\sigma'-f_L^2)y+\frac{\delta_\epsilon}{M_P^2}L_eH_2\sigma^3, \tag{28}$$

where the first term breaks the lepton symmetry and generates a majoron supermultiplet of Eq. (3). The second term in Eq. (28) generates the effective interaction displayed in Eq. (26) with $c_\epsilon=3/\sqrt{2}$ and $\epsilon\sim(\delta_\epsilon/M_P^2)f_L^3$. Thus the specific choice for the lepton charges allows one to correlate ϵ to the scale f_L. In particular, for $\delta_\epsilon\sim 0.3$ and $f_L\simeq 3\times 10^{11}$ GeV, one has $\epsilon\sim 0.1$ MeV.

(3) *PQ as the lepton number symmetry.* If both Higgs bosons and leptons transform nontrivially under the $U(1)_G$ symmetry then the latter can play a dual role of the PQ symmetry and the lepton number symmetry as in Ref. [27]. In this case one can correlate the origin of ϵ and μ to the same symmetry breaking scale f_{PQ}. The neutrino coupling to QGF is given by the combination of Eqs. (19) and (26):

---

[1]If there is also the coupling $L_eN_eH_2$ then N_e gets a VEV: $\langle N_e\rangle\sim A\langle\nu_e\rangle\langle H_2\rangle/M^2$, where A is the soft symmetry breaking parameter and M is the mass of N. For $M=10^{12}$ GeV one gets $\langle N\rangle\sim 10^{-15}$ eV which is negligibly small. Note that in the model of Sec. IV the coupling $L_eN_eH_2$ itself is strongly suppressed.

$$\mu H_1 H_2 + \epsilon L_e H_2 + c_\mu \frac{\mu}{f_{\rm PQ}} H_1 H_2 S + c_\epsilon \frac{\epsilon}{f_{\rm PQ}} L_e H_2 S. \quad (29)$$

This effective superpotential generates the following mass matrix for ν_e and S which is the combination of Eq. (21) and Eq. (27):

$$\begin{pmatrix} 0 & (c_\epsilon - c_\mu)\epsilon v \sin\beta/f_{\rm PQ} \\ (c_\epsilon - c_\mu)\epsilon v \sin\beta/f_{\rm PQ} & m_S^0 - c_\mu^2 \mu v^2 \sin 2\beta/f_{\rm PQ}^2 \end{pmatrix}. \quad (30)$$

According to Eq. (30) the $\nu - S$ mixing angle $\theta_{\nu s}$ is determined by

$$\tan\theta_{\nu s} \sim \frac{(c_\mu - c_\epsilon)\epsilon v \sin\beta}{m_S^0 f_{\rm PQ} - c_\mu^2 \mu v^2 \sin 2\beta/f_{\rm PQ}}. \quad (31)$$

Let us give an example of underlying $\mathrm{U}(1)_G$ invariant superpotential which generates neutrino-QGF mixing. The G-charge prescription $(-1,-1,\ 1,-1,-2)$ for $(H_1,\ H_2,\ \sigma,\ \sigma',\ L_e)$ permits the following $\mathrm{U}(1)_G$ invariant superpotential:

$$W = \lambda(\sigma\sigma' - f_{\rm PQ}^2)y + \frac{\delta_\mu}{M_P} H_1 H_2 \sigma^2 + \frac{\delta_\epsilon}{M_P^2} L_e H_2 \sigma^3. \quad (32)$$

It gives the terms displayed in Eq. (29) with $c_\epsilon = 3/\sqrt{2}, c_\mu = \sqrt{2}$.

IV. MODEL

Let us put together the basic ingredients discussed in Secs. II and III into a model which simultaneously explains the solar, atmospheric, and dark matter problems. In principle the sterile state, such as the axino, could mix with any of the neutrinos but the possibility of the ν_e-S mixing solving the solar neutrino problem seems most preferred phenomenologically. The required range of the ν_e-S mixing and S mass is given in Eq. (4). The alternative possibility of large ν_μ-S mixing accounting for the atmospheric neutrino deficit can conflict with the cosmological bound coming from the nucleosynthesis. This conflict can be avoided [28] if the lepton asymmetry is much larger than the baryon asymmetry. It is argued [29] that for a suitable range of parameters the asymmetry can be enhanced by $\nu_\tau - \nu_s$, $\overline{\nu_\tau} - \overline{\nu_s}$ oscillations. In what follows we will concentrate on a safer possibility with sterile neutrino mixing only with the ν_e.

We consider here a specific model in which $\mathrm{U}(1)_G$ is identified with $\mathrm{U}(1)_{\rm PQ}$. The latter is chosen to act nontrivially on leptons also and is also required to be generation dependent. This is needed in order to suppress the mixing of S with $\nu_{\mu,\tau}$ and to get pseudo-Dirac structure for the ν_μ-ν_τ system.[2] Specifically, the model contains the following fields with their $\mathrm{U}(1)_{\rm PQ}$ assignments as given below:

H_1	H_2	σ	σ'	L_e	L_μ	L_τ	N_e	N_μ	N_τ	y
-1	-1	1	-1	-2	$-1/2$	3/2	0	3/2	$-1/2$	0.

[2]One can introduce for this an additional horizontal symmetry, suggesting that $\mathrm{U}(1)_G$ is generation blind.

This choice gives rise to the desired phenomenological results as we now discuss. [One can get more symmetric or regular charge prescription by introducing more singlet fields or a horizontal symmetry in addition to $\mathrm{U}(1)_G$.]

Our superpotential consists of two parts. The part containing the electron family is based on Eq. (32) and is given explicitly as

$$W_e = \lambda(\sigma\sigma' - f_{\rm PQ}^2)y + \frac{\delta_\mu}{M_P} H_1 H_2 \sigma^2 + \frac{\delta_\epsilon}{M_P^2} L_e H_2 \sigma^3 + M_e N_e N_e + h N_e N_e \sigma\sigma'/M_P. \quad (33)$$

The part of superpotential involving the μ and τ family is given by the superpotential

$$W_{\mu\tau} = \sum_{\alpha=\mu,\tau} \frac{m_\alpha^D}{v_2} L_\alpha N_\alpha H_2 + \frac{M_\tau}{f_{\rm PQ}} N_\tau N_\tau \sigma + \frac{M_{\mu\tau}}{f_{\rm PQ}} N_\mu N_\tau \sigma'. \quad (34)$$

The model is assumed to have the strict no-scale kinetic term

$$K = -3\ln(T + T^* - Z_a Z_a^*), \quad (35)$$

where Z_a denotes all the other chiral super fields. Neglecting the D-term contribution, the scalar potential takes the supersymmetric form $V \sim |W_a|^2$ [21] which immediately leads us to the vacuum expectation values of the fields:

$$\langle\sigma\rangle \simeq \langle\sigma'\rangle \simeq f_{\rm PQ}, \quad \langle y\rangle = 0. \quad (36)$$

The rest of the fields have zero VEV in the supersymmetric limit.

The second relation in Eq. (36) is crucial to ensure the masslessness of the quasi Goldstone fermion S at tree level. Being a singlet under $\mathrm{U}(1)_G$ the field y can receive tadpole divergence generic in supergravity theories [30]. Even if no term in Eqs. (33)–(35) is responsible for such a divergence, one may write nonrenormalizable terms like $y^2\sigma\sigma'/M_P$ in the superpotential which may potentially induce a large vacuum expectation value $\langle y\rangle \sim m_{3/2}$ due to two-loop tadpole [30]. This kind of nonrenormalizable terms can be forbidden by a symmetry. Inspecting the superpotentials (33) and (34) one finds that $\mathrm{U}(1)_R$ symmetry under which the fields carry the charges:

H_1	H_2	σ	σ'	y	L_e	L_μ	L_τ	N_e	N_μ	N_τ
1	1	0	0	2	1	0	0	1	1	1

can play such a role.

The part $W_{\mu\tau}$ of the superpotential leads to the mass matrix in the $(\nu_\mu, \nu_\tau, N_\mu, N_\tau)$ basis:

$$\mathcal{M} = \begin{pmatrix} 0 & 0 & m_\mu^D & 0 \\ 0 & 0 & 0 & m_\tau^D \\ m_\mu^D & 0 & 0 & M_{\mu\tau} \\ 0 & m_\tau^D & M_{\mu\tau} & M_\tau \end{pmatrix}. \quad (37)$$

The above mass matrix gives rise to pseudo-Dirac neutrino with a common mass

$$m_{\rm DM} \sim \frac{m_\mu^D m_\tau^D}{M_{\mu\tau}}. \tag{38}$$

This mass can be in the eV range as required for the solution of the dark matter problem by taking the values $m_\mu^D \sim 0.1$ GeV, $m_\tau^D \sim 50$ GeV, and $M_{\mu\tau} \sim 10^9$ GeV. The mass splitting is given by

$$\frac{\Delta m^2}{m_{\rm DM}^2} \simeq 2\left(\frac{m_\mu^D}{m_\tau^D}\right)\left(\frac{M_\tau}{M_{\mu\tau}}\right). \tag{39}$$

Taking $(M_\tau/M_{\mu\tau}) \sim 1$, one reproduces both mixing and Δm^2 required to explain the atmospheric anomaly.

The superpotential W_e generates the mixing of the ν_e with QGF contained in S of Eq. (3). The mixing angle (31) following from the superpotential (32) can fall in the required range (4) if $\epsilon \sim 1$ MeV and $\langle\sigma\rangle = \langle\sigma'\rangle = f_{\rm PQ} \sim 10^{12}$ GeV.

The mass of the QGF depends on the SUSY breaking terms. In no-scale models, this mass can arise in two or three loops via the mechanism described in Sec. II. For the specific choice made in Eq. (35), the gauginos provide the only seed of SUSY breaking. Then the coupling of the QGF to the RH components $N_{\mu,\tau}$ displayed in Eq. (34) induce the mass at three loops (Fig. 3). The dominant contribution comes from N_τ in the loop due to the hierarchy $m_\mu^D \ll m_\tau^D$ and is approximately given [see Eq. (13)] by

$$m_S \simeq \frac{\alpha_2}{(4\pi)^5}\left(\frac{M_\tau}{f_{\rm PQ}}\right)^2\left(\frac{m_\tau^D}{v_2}\right)^2 m_{1/2}. \tag{40}$$

With the values of the parameter chosen earlier namely, $M_\tau \sim 10^9$ GeV, $m_\tau^D \sim 50$ GeV, and $f_{\rm PQ} \sim 10^{12}$ GeV and for $m_{1/2} \sim 200$ GeV, $v_2 \sim 100$ GeV, one gets $m_S \sim 10^{-3}$ eV. Thus W_e together with the kinetic terms specified in Eq. (35) lead to the Mikheyev-Smirnov-Wolfenstein (MSW) solution of the solar neutrino puzzle.

Notice that the specific $U(1)_G$ symmetry of the model does not allow Yukawa coupling $L_e H_2 N_e$ for the electron neutrino, in contrast with muon and τ neutrinos. The Dirac mass term of electron neutrino can be generated by a higher-order nonrenormalizable term: $h L_e N_e H_2 \sigma^3/M_P^3$, and therefore, $m_e^D \sim m_e (f_{\rm PQ}/M_P)^3$ is negligibly small. At the same time the charge prescription, $G(N_e)=0$, permits the mass terms for N_e in Eq. (33) which will produce $M_e \sim 10^6 - 10^{18}$ GeV.

The model presented above does not contain any mixing between ν_e and $\nu_{\mu,\tau}$. Such mixing can be induced, for example, by adding a new Higgs field which could generate a Dirac mass term $m_{e\tau}\nu_e N_\tau$. This gives rise to the ν_e-ν_μ mixing angle $\theta_{e\mu} \sim m_{e\tau}/m_\mu$ which can be in the range of sensitivity of KARMEN and the Liquid Scintillation Neutrino Detector [4] for $m_{e\tau} \sim 30$ MeV, $m_\mu \sim$ GeV [6].

V. CONCLUSIONS

Simultaneous presence of different neutrino anomalies points to the existence of a sterile neutrino. We have considered the possibility that the sterile neutrino is a quasi Goldstone fermion appearing in supersymmetric theory as a result of the 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, or the horizontal symmetry.

The mass of QGF generated by SUSY breaking can be as small as 10^{-3} eV so that the $\nu_e \to S$ resonance conversion solves the solar neutrino problem. The smallness of m_S can be attributed in supergravity theory either to special forms of the superpotential and the $U(1)_G$ breaking scale $f_G \gtrsim 10^{16}$ GeV or to no-scale kinetic terms for certain superfields. In the last case, m_S is generated in two or three loops.

The mixing of QGF with the neutrinos implies spontaneous or explicit violation of the 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).

The $U(1)_G$-symmetry being generation dependent can simultaneously explain the dominance of QGF coupling with electron neutrino and pseudo Dirac structure of the ν_μ-ν_τ system needed to explain the atmospheric neutrino problem and HDM.

The PQ breaking scale $f_{\rm PQ} \sim 10^{10} - 10^{12}$ GeV determines several features of the model presented here. It provides simultaneous explanation of the parameters ϵ and μ and thus leads to small R-parity violation required in order to solve the solar neutrino problem in our approach. It also provides the intermediate scale for the RH neutrino masses which is required in order to solve the dark matter and the atmospheric neutrino problem. Finally, 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. Thus the basic scenario presented here is able to correlate variety of phenomena.

If future solar neutrino experiments establish that the ν_e-S conversion is the cause of the solar neutrino deficit then one might be seeing indirect evidence for the PQ-like symmetry or for that matter of SUSY itself.

ACKNOWLEDGMENTS

A.S.J. wants to thank ICTP for its hospitality during his visit.

[1] GALLEX Collaboration, P. Anselmann *et al.*, Phys. Lett. B **327**, 377 (1994); SAGE Collaboration, J. N. Abdurashitov *et al.*, *ibid.* **328**, 234 (1994); Homestake Collaboration, B. T. Cleveland *et al.*, in *Neutrino '94*, Proceedings of the International Conference on Neutrino Physics and Astrophysics, Eilat, Israel, edited by A. Dar, G. Eilam, and M. Gronou [Nucl. Phys. **B38** (Proc. Suppl.), 47 (1995)]; Kamiokande Collaboration, Y. Suzuki *et al.*, *ibid.* p. 55.

[2] Kamiokande Collaboration, H. S. Hirata *et al.*, Phys. Lett. B **205**, 416 (1988); **280**, 146 (1992); Kamiokande Collaboration, Y. Fukuda, *et al.*, *ibid.* **335**, 237 (1994).

[3] J. R. Primack, J. Holtzman, A. Klypin, and D. O. Caldwell, Phys. Rev. Lett. **74**, 2160 (1995), and references therein.

[4] LSND Collaboration, C. Anthanassopoulos *et al.*, Phys. Rev. Lett. **75**, 2650 (1995).
[5] K. Enqvist, J. Maalampi, and V. B. Semikoz, Nucl. Phys. **B456**, 339 (1995); for previous works see, X. Shi, D. N. Schramm, and B. D. Fields, Phys. Rev. D **48**, 2563 (1993), and references therein.
[6] E. J. Chun, A. S. Joshipura, and A. Yu. Smirnov, Phys. Lett. B **357**, 608 (1995).
[7] R. Foot and R. R. Volkas, Phys. Rev. D **52**, 6595 (1995); Z. G. Berezhiani and R. N. Mohapatra, *ibid.* **52**, 6607 (1995).
[8] E. Ma, Report No. UCRHEP-T149, hep-ph/9507348 (unpublished).
[9] Y. Chikashige, R. N. Mohapatra, and R. D. Peccei, Phys. Lett. **98B**, 265 (1981).
[10] R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. **38**, 1440 (1977).
[11] J. E. Kim and H. P. Nilles, Phys. Lett. **138B**, 150 (1984); E. J. Chun, J. E. Kim, and H. P. Nilles, Nucl. Phys. **B370**, 105 (1992); J. E. Kim and H. P. Nilles, Mod. Phys. Lett. A **9**, 3575 (1995).
[12] E. J. Chun, Phys. Lett. B **348**, 111 (1995).
[13] M. Dine, W. Fischler, and M. Srednicki, Phys. Lett. **104B**, 199 (1981).
[14] S. P. Mikheyev and A. Yu. Smirnov, Sov. J. Nucl. Phys. **42**, 913 (1985); Sov. Phys. JETP **64**, 4 (1986); L. Wolfenstein, Phys. Rev. D **17**, 2369 (1978).
[15] K. Tamvakis and D. Wyler, Phys. Lett. **112B**, 451 (1982); J. F. Nieves, Phys. Rev. D **33**, 1762 (1986).
[16] E. J. Chun and A. Lukas, Phys. Lett. B **357**, 43 (1995).
[17] E. J. Chun, J. E. Kim, and H. P. Nilles, Phys. Lett. B **287**, 123 (1992).
[18] T. Goto and M. Yamaguchi, Phys. Lett. B **276**, 103 (1992).
[19] K. S. Babu, S. M. Barr, and D. Seckel, Phys. Lett. B **336**, 213 (1994).
[20] G. Dvali, Report No. IFUP-TH 21/95, hep-ph/9505253 (unpublished).
[21] For a review, see A. B. Lahanas and D. V. Nanopoulos, Phys. Rep. **145**, 1 (1987).
[22] P. Moxhay and K. Yamamoto, Phys. Lett. **151B**, 363 (1985).
[23] R. Arnowitt and P. Nath, Phys. Rev. D **46**, 3981 (1992); P. Nath and R. Arnowitt, Phys. Lett. B **287**, 89 (1992).
[24] J. L. Lopez, D. V. Nanopoulos, and A. Zichichi, Phys. Rev. D **49**, 343 (1994); J. L. Lopez, D. V. Nanopoulos, G. Park, X. Wang, and A. Zichichi, Phys. Rev. D, **50**, 2164 (1994).
[25] For a recent review, see L. Hall, Mod. Phys. Lett. A **7**, 467 (1990).
[26] L. Hall and M. Suzuki, Nucl. Phys. **B231**, 419 (1984); A. S. Joshipura and M. Nowakowski, Phys. Rev. D **51**, 2421 (1995).
[27] R. Mohapatra and G. Senjanović, Z. Phys. C **17**, 53 (1983); P. Langacker, R. D. Peccei, and T. Yanagida, Mod. Phys. Lett. A **1**, 541 (1986); A. S. Joshipura, Z. Phys. C **38**, 479 (1988).
[28] R. Foot and R. R. Volkas, Phys. Rev. Lett. **75**, 4350 (1995).
[29] R. Foot, M. J. Thomson, and R. R. Volkas, Phys. Rev. D **53**, 5349 (1996).
[30] J. Bagger, E. Poppitz, and L. Randall, Nucl. Phys. **B455**, 54 (1995); J. Bagger presented at the International Workshop on Supersymmetry and Unification of Fundamental Interactions (SUSY95), Palaiseau, France, 1995 (unpublished), Report No. hep-ph 9511213 (unpublished).