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Supergravity can reconcile dark matter with lepton number violating neutrino masses

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Supersymmetry offers a cold dark matter candidate, provided that lepton number is *not violated by an odd number of units*. On the other hand, lepton number violation by even (two) units gives us an attractive mechanism of neutrino mass generation. Here we offer an explanation of this, in a supergravity framework underlying a supersymmetric scenario, the essential feature being particles carrying lepton numbers, which interact only gravitationally with all other known particles. It is shown that one can have the right amount of $\Delta L = 2$ effect giving rise to neutrino masses, whereas the lifetime for $\Delta L = 1$ decays of the lightest supersymmetric particle can be prolonged beyond the present age of the Universe.

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I. INTRODUCTION

Speculations abound nowadays about a supersymmetric (SUSY) nature at the fundamental level, described by a theory invariant under boson-fermion transformations. One of the positive features of SUSY [1] is that in its minimal form it provides a candidate for cold dark matter in our Universe, since, in most models, the lightest SUSY particle (LSP) is stable, neutral, and weakly interacting only, and also lies in the mass range of the electroweak scale [2]. However, the stability of the LSP requires the SUSY theory to conserve R-parity, defined as $R = (-1)^{L+3B+2s}$, where L, B, and s stand for the lepton number, baryon number, and spin of a particle, respectively. Arguing in this line, lepton number is expected to be conserved in order to ensure that SUSY is the source of cold dark matter.

To be very precise, however, the conservation of R-parity means that L is not violated *by an odd number of units*. Thus one may conserve R-parity and retain stability of the LSP even if L is violated by an *even number of units*, which makes it possible to have $\Delta L = 2$ neutrino masses of the form $\bar{\nu}^c \nu$. Such mass terms form the seed for, say, the seesaw mechanism [3] which is a beautiful explanation of the smallness of neutrino masses vis-à-vis the masses of the charged leptons. So the smallness of neutrino masses and the dark matter in the Universe can be explained together in SUSY models if we assume lepton number is violated by two units. In contrast to this, with $\Delta L = 1$ terms one may explain the smallness of neutrino masses but additional sources for dark matter are required to be postulated. The question that one may ask now is: if $\Delta L = 2$ is allowed, is there any fundamental reason to believe that $\Delta L = 1$ terms either cannot occur or are very suppressed? In other words, can a cold dark matter candidate be reconciled with Majorana masses for neutrinos with the help of some fundamental principles? These questions form the central theme of the present work.

The superpotential of a lepton number conserving theory, including right-handed neutrino superfields N_i (required for the seesaw mechanism), is

$$W_{\rm MSSM} = Y_u^{ij} Q_i U_j^c H_2 + Y_d^{ij} Q_i D_j^c H_1 + Y_e^{ij} L_i E_j^c H_1 + Y_\nu^{ij} L_i N_j^c H_2 - \mu H_1 H_2, \qquad (1)$$

where the flavor indices, i, j run from 1 to 3 and SU(2) gauge indices have been suppressed. The Y's stand for various Yukawa couplings. μ is the Higgsino mass parameter. H_1, H_2 are the two SU(2) doublet Higgs superfields, with $H_1 = (H_1^0, H_1^-)$ and $H_2 = (H_2^+, H_2^0)$. Q, L are SU(2) doublet quarks and leptons, while U^c, D^c, E^c are SU(2) singlet up-quark, down-quark, and charged lepton superfields, respectively. If L is violated, then one can further add the terms [4]

$$W_{\not L} = \lambda_{ijk} L_i L_j E_k^c + \lambda'_{ijk} Q_i L_j D_k^c + \epsilon_i L_i H_2 + M_{ij} N_i^c N_j^c, \qquad (2)$$

where λ, λ' are some constants and ϵ, M are mass parameters. Here the first three terms violate L by one unit, and need to be forbidden for the stability of LSP. The last term, violating L by two units, gives Majorana masses for neutrinos. The aim, therefore, is to try to understand why the last term should be allowed but not the first three terms of Eq. (2).

There have been some explanations of the above claim in, say, supersymmetric grand unified theories (GUTs). For example, in a SUSY GUT model based on the gauge group $SU(5) \times SU(5)$, R-parity can automatically be conserved and seesaw masses for neutrinos can also be generated [5].

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There have been some SUSY models where R-parity arises from a continuous gauge symmetry [6]. Since R-parity is naturally conserved, the stability of LSP and Majorana nature of neutrinos can be understood in these models.

Here we take an alternative approach and seek an explanation of such a phenomenon in a supergravity (SUGRA) framework which is by far the most popular paradigm of SUSY breaking [7]. Apart from being the local extension of global SUSY, SUGRA can also have its root in some radically fundamental idea like superstring theory. In this framework, the SUSY-breaking soft terms have their origin in nonrenormalizable interactions of the observable fields with a hidden sector which is sterile under all known interactions except gravity. Also, nonrenormalizable interactions with the hidden sector have often been invoked to explain neutrino masses [8–10]. We utilize this framework to explain the Majorana nature of neutrinos and the dark matter content of the Universe. As described below, we postulate some symmetries applicable to hidden sector fields, suppress phenomenological effects of the unwanted $\Delta L = 1$ terms, and explain why lepton number is violated by two units in the neutrino sector. Unlike the models [5,6], where R-parity is naturally conserved, this model leads to R-parity violating terms when the global symmetries (lepton number and R-charge) are spontaneously broken below the Planck scale. We show that the effect of these R-parity violating terms is small and the lightest neutralino of minimal supersymmetric standard model (MSSM) can be a long-lived particle. The lifetime of this neutralino field can be larger than the age of the Universe and can serve as a candidate for dark matter. The two central features used by us are (a) the possibility of some hidden sector field carrying lepton number, and (b) the fact that the theory may have a nonminimal Kahler potential, thus accommodating certain desirable values of parameters in the observable sector.

In the proposed scenario, we make use of three hidden sector chiral superfields $S(L = 0)$, $S'(L = 0)$, and $X(L = 1)$, where L within the brackets indicates the lepton number of the field. As will be shown in detail later, the purpose of the fields S, S' is to give masses to the scalar and gaugino fields. Among the hidden sector fields, X carries lepton number $+1$, and we propose this field to establish the Majorana nature of neutrinos. While it may be unusual to attribute lepton number to a hidden sector field, such suggestions have been considered earlier in the context of leptogenesis [11]. It should also be remembered that a right-handed neutrino superfield itself is a gauge singlet carrying lepton number, and can pass off as a hidden sector field but for the Yukawa interactions. However, the lepton number assignments prevent X from entering into Yukawa couplings. To generate $\Delta L = 2$ terms for neutrinos, we assume that lepton number is conserved at the Planck scale. This assumption forbids the usual $N^c N^c$ term in the superpotential, but a nonrenormalizable term $\frac{XX}{M_P} N^c N^c$ is allowed. If X acquires a vacuum expectation value (vev) $\Delta L = 2$ terms can be generated for neutrinos, and to have right-handed neutrino mass at the TeV scale the vev of X should be at the intermediate scale (10^{10-11} GeV). The assumption of lepton number conservation may not be enough to postulate, since terms like XN^c are also allowed in the superpotential, and they generate unwanted $\Delta L = 1$ terms through the vev of X. To avoid these problems, we make use of the R-symmetry, assumed to be conserved at the Planck scale. The breakdown of these symmetries at somewhat lower scales is motivated by (a) the requirement of seesaw masses for neutrinos and (b) the need to achieve adequate suppression of $\Delta L = 1$ effects when the symmetries are broken. It may also be noted that such global symmetries have played a role in explaining neutrino masses in SUSY models, discussed earlier [8,9].

We describe our model in detail and essential features of it in Sec. II. We obtain the low-energy scalar potential of our model in Sec. III and argue that the sneutrinos cannot acquire nonzero vevs. Since we need a nonzero scalar vev for the field X which is a carrier of lepton number, there are $\Delta L = 1$ interaction terms in our model. This and some other consequences of our model are given in Sec. IV. Section V contains conclusions.

II. ESSENTIAL FEATURES OF THE MODEL

As is explained in the introduction, we assume the conservation of lepton number and R-symmetry at the Planck scale. We construct a Lagrangian which is invariant under these symmetries. Apart from the gauge kinetic terms, the Lagrangian of the model is

$$\mathcal{L} = \int K d^4\theta + \left(\int W d^2\theta + \text{H.c.} \right), \tag{3}$$

where W and K are the superpotential and the Kahler potential, respectively.

It is well known that in the low-energy limit of a spontaneously broken supergravity model, the cancellation of large contributions to the cosmological constant requires the presence of at least one scalar field (usually a singlet under the observable sector gauge group) with a vev of the order of Planck scale [1,12,13]. If the superpotentials for this (hidden sector) field and the observable sector fields are additive, then the cosmological constant is determined by the vevs of the hidden sector field(s) [13]. Given such vevs, and with appropriate choice of parameters in the hidden sector superpotential in such a scenario, supersymmetry can be broken at an intermediate scale with gravitino mass as low as $\sim$TeV. This in turn results in the generation of SUSY-breaking soft terms of the order of gravitino mass in the observable sector [14]. In our case, the field S plays this role in the hidden sector. Although the problem of cosmological constant is not the main focus of this paper, we need the scalar vev of S at the Planck scale to attain the gravitino and scalar masses at the TeV scale. We briefly comment on the cosmological constant in Sec. IV. As we shall see and is already mentioned in the introduction, the other hidden sector field, namely, X, on the other hand, has its vev at an intermediate scale, and being a carrier of lepton number, has an altogether different role to play in the observable sector phenomenology. Finally, the field S' should acquire zero vev in order for the gaugino fields to acquire masses at the TeV scale. We present below the explicit forms of the superpotential and the Kahler potential, through which we can attain consistent SUSY breaking.

The forms of the superpotential and the Kahler potential follow from the conservation of R-symmetry and lepton number, which we have assumed to be valid at the Planck scale. The R-charges of the fields are assigned as: $R(S) = 2$, $R(S') = 0$, $R(X) = 1$, $R(Q_i) = R(L_i) = R(U_i^c) = R(N_i^c) = 0$, $R(D_i^c) = R(E_i^c) = 2$, $R(H_1) = 0$, $R(H_2) = 2$. With this, the superpotential has the form

$$W = W_h + W_{\rm MSSM} + \frac{XX}{2M_P} N_i^c a_{ij} N_j^c, \tag{4}$$

where M_P is the Planck scale and a_{ij} are $\mathcal{O}(1)$ constants, with $a_{ij} = a_{ji}$. $W_{\rm MSSM}$ is defined in Eq. (1). W_h is the hidden sector part of the superpotential which takes the following leading terms:

$$W_h = \Lambda^2 S + \Lambda' S' S + q S'^2 S, \tag{5}$$

where $\Lambda \sim 10^{10-11}$ GeV, $\Lambda' \sim$ TeV, and q is a $\mathcal{O}(1)$ constant. For phenomenological consistency of our model we need the Λ' to be at the TeV scale. But in this work we are not justifying it. There is another TeV scale parameter in the superpotential, which is the μ-term of $W_{\rm MSSM}$. Here we are not addressing the origin of this term. In the past there were attempts to understand why the μ-parameter is at the TeV scale [8,15]. We may follow these approaches and address the origins of Λ' and μ terms, but that is not the main focus of this paper. The last term of Eq. (4), which arises from the conservation of lepton number, is especially noteworthy; such a nonrenormalizable term can obviously lead to $\Delta L = 2$ neutrino masses once the scalar component of X acquires vev. The role of R-symmetry in the superpotential is that it forbids the following lepton number conserving terms, such as: $\frac{S}{M_P} Q U^c H_2$, $\frac{S}{M_P} Q D^c H_1$, $\frac{S}{M_P} L E^c H_1$, $\frac{S}{M_P} L N^c H_2$. Since, as explained before, S acquires a scalar vev of the order of Planck scale, these terms generate unacceptably high masses for quarks and leptons, expect for the top quark. However, terms of the form: $\frac{S'}{M_P} Q U^c H_2$, $\frac{S'}{M_P} L E^c H_1$, etc. are allowed both by lepton number and R-symmetry conservations, but since S' acquires zero vev they do not contribute to masses of fermions. R-symmetry also forbids the term $S H_1 H_2$, which generates a high value for the μ parameter. R-symmetry also forbids the term $X N^c$, which generates a $\Delta L = 1$ term in the low-energy regime through the scalar vev of X. The role of lepton number conservation in the superpotential is that it forbids the following terms: $L L E^c$, $Q L D^c$, and $L H_2$, which are R-symmetric invariant, and these are the terms that should be avoided for the stability of the LSP. It also forbids terms such as XX or $X X N^c$, which are in principle allowed by R-symmetry. Terms in higher powers of S as well as X in W are also absent via R-symmetry as well as the assumption of lepton number conservation at the Planck scale.

Next, we suggest a specific form of the Kahler potential. In general, the Kahler manifold is a real function of the fields Y and $Y^\dagger$, where $Y = S, S', X$ in our case. Since X carries lepton number, here one has a Kahler potential where X enters only in the form $X^\dagger X$ (for the conservation of lepton number) and considers a Kahler potential of the form

$$K = K_0(S, S', XX^\dagger) + \sum_i \Phi_i^\dagger \Phi_i. \tag{6}$$

Here, $\Phi_i = Q_i, U_i^c, D_i^c, L_i, E_i^c, N_i^c, H_1, H_2$. K_0 is some function of hidden sector fields, which should be chosen so that the Kahler potential is invariant under both the R-symmetry and lepton number. In general, this function depends nontrivially on the hidden sector fields, and thus the Kahler potential has nonminimal character [16]. The Kahler potential can contain terms involving both hidden and observable fields, such as

$$\frac{S^\dagger S}{M_P^2} \Phi_i^\dagger \Phi_i, \qquad \frac{S'}{M_P} \Phi_i^\dagger \Phi_i, \qquad \frac{X^\dagger X}{M_P^2} \Phi_i^\dagger \Phi_i. \tag{7}$$

Among the above terms, the first two at most contribute to the SUSY-breaking soft terms for the scalar fields and they do not generate undesired $\Delta L = 1$ terms. So we do not consider them in the Kahler potential as they do not affect the main conclusions of our work. Moreover, the soft terms for scalar fields can be obtained from the Kahler potential that we have chosen in Eq. (6), and it will be shown in Sec. III. However, the last term of Eq. (7) can give rise to $\Delta L = 1$ terms through the scalar vev of X. But this term is suppressed by two powers of Planck mass. In Sec. IV we argue that in our model $\Delta L = 1$ terms dominantly come due to the last term of Eq. (4). It is clear that this term in the superpotential is suppressed by one power of Planck mass, so it gives dominant effects compared to the last one of Eq. (7). To present our ideas in a simple fashion we omit this possible term in the Kahler potential. The Kahler potential can also contain the term $\frac{S^\dagger}{M_P} H_1 H_2$ which is consistent with the lepton number and R-symmetry conservations. This term effectively generates the μ-term if the auxiliary vev of S is nonzero. It will be shown below that the auxiliary vev of S is at the intermediate scale and so the effective μ parameter is at the TeV scale. Such a μ-term is already there in the superpotential and here our main motivation is not on the explanation of the origin of the μ-term. So without loss of generality the term $\frac{S^\dagger}{M_P}H_1H_2$ can be excluded from the Kahler potential. Notice that the assumption of conservation of R-symmetry at the Planck scale enables us to avoid the terms of the form $\frac{X^\dagger}{M_P^2}LLE^c$, $\frac{X^\dagger}{M_P^2}QLD^c$, and $\frac{X^\dagger}{M_P}LH_2$ which, via a vev of X, violate lepton number by one unit.

To completely specify our model, we need to fix the gauge kinetic function which determines the interactions of gauge and gaugino fields. We do not study them in detail since they do not affect our conclusions. But to be phenomenologically consistent, gaugino fields should have masses which are determined by the form of gauge kinetic function. In our model its form is

$$F_{ab} = \delta_{ab}\left(\frac{1}{g_a^2} + \frac{1}{M_P} f_a S' + \cdots\right), \tag{8}$$

where g_a are the three gauge couplings of the standard model gauge group, f_a are $\mathcal{O}(1)$ constants, and the indices a, b run over 1, 2, 3. The dots in Eq. (8) are higher order terms which can be neglected. The second term of Eq. (8) gives masses to gauginos. We will see below that the auxiliary vev of S' is at the intermediate scale and hence the gaugino fields have masses of the order of $\frac{\langle F_{S'}\rangle}{M_P} \sim$ TeV.

So far, everything included in W as well K conserve lepton number. Now, the very form of W tells us that $\langle F_S\rangle = \langle\frac{\partial W}{\partial S}\rangle$ is on the order of Λ^2 and $\langle F_{S'}\rangle = \langle\frac{\partial W}{\partial S'}\rangle = \Lambda'\langle S\rangle + 2q\langle S'\rangle\langle S\rangle$ is at the intermediate scale if $\langle S\rangle \sim M_P$ and $\langle S'\rangle = 0$. We will make the scalar vev of S' to be zero in order to get its auxiliary vev at the intermediate scale, and this we require to make sure that the gauginos have masses at the TeV scale which is explained in the previous paragraph. While $\langle F_X\rangle = \langle\frac{\partial W}{\partial X}\rangle = 0$ if the right-handed sneutrinos $\tilde{N}_i$ have no vev, something that we need to establish in order to eliminate the possibility of $\Delta L = 1$ terms.

III. SCALAR POTENTIAL

Let us now consider the scalar potential of this theory and place our claims about the vevs of S, S', and X on firmer ground. The reasons for $\langle S'\rangle = 0$ and $\langle X\rangle \sim \Lambda$ have already been explained. We will show below that in order for the gravitino mass to be at the TeV scale we need $\langle S\rangle \sim M_P$. These are the demands that we are making on the vevs of hidden sector fields in order to be consistent with the phenomenological masses of the supersymmetric fields. One crucial thing to be shown is that the sneutrino fields should not acquire nonzero vevs. We show below that these demands can be satisfied by minimizing the scalar potential and choosing appropriate values of the parameters of the model. The contribution to the scalar potential from the superpotential and the Kahler potential is given by [1]

$$V = M_P^4 e^G[M_P^2 G_M K^{M\bar{N}} G_{\bar{N}} - 3], \tag{9}$$

where

$$G = \frac{K}{M_P^2} + \ln\left|\frac{W}{M_P^3}\right|^2, \tag{10}$$

$G_M = \frac{\partial G}{\partial \phi_M}$ and $G_{\bar{N}} = \frac{\partial G}{\partial \phi_N^*}$, ϕ being a chiral superfield. The matrix $K^{M\bar{N}}$ is the inverse of $\frac{\partial^2 K}{\partial \phi_N^* \partial \phi_M}$.

The vevs that we get after minimizing the scalar potential determine the SUSY breaking of the theory. SUSY breaking requires that, expressed in terms of these, the vev of

$$\mathcal{F}_\phi = \frac{\partial W}{\partial \phi} + \frac{W}{M_P^2}\frac{\partial K}{\partial \phi} \tag{11}$$

should be nonzero for some hidden sector field(s) ϕ [17]. In our case, we have $\langle \mathcal{F}_S\rangle = \Lambda^2 + \frac{\Lambda^2\langle S\rangle}{M_P^2}\langle\frac{\partial K}{\partial S}\rangle$, after putting $\langle S'\rangle = \langle\tilde{N}_i^c\rangle = 0$. If one has $\langle S\rangle$ at the Planck scale, together with Λ at an intermediate scale, one can not only have a nonzero $\langle \mathcal{F}_S\rangle$ but also ensure $\langle \mathcal{F}_S\rangle$ of an order which is required by a phenomenologically consistent SUSY spectrum. One of the consequences of SUSY breaking in supergravity is that the gravitino field acquires a nonzero mass, which is given below

$$m_{3/2}^2 \sim M_P^2 e^{\langle G\rangle} = e^{\langle K\rangle/M_P^2}\frac{\langle W\rangle^2}{M_P^4} = e^{\langle K\rangle/M_P^2}\frac{(\Lambda^2\langle S\rangle)^2}{M_P^4}. \tag{12}$$

The mass of the gravitino can be of the order of TeV provided if $\langle S\rangle \sim M_P$.

Substituting the forms of K and W in Eqs. (9) and (10), one obtains the form of the scalar potential [18] as

$$V = V_0 + V_1 \tag{13}$$

with

$$\begin{aligned} V_0 = e^{K/M_P^2}\Big\{ & K_0^{S\bar{S}}|\Lambda^2 + \Lambda' S' + qS'^2|^2 + K_0^{S\bar{S}}\frac{\partial K_0}{\partial S}\Big(\frac{\partial K_0}{\partial S}\Big)^* \frac{W_h W_h^*}{M_P^4} + K_0^{S'\bar{S}'}|\Lambda' S + 2qS'S|^2 - \frac{3}{M_P^2}W_h W_h^* \\ & + K_0^{S'\bar{S}'}\frac{\partial K_0}{\partial S'}\Big(\frac{\partial K_0}{\partial S'}\Big)^*\frac{W_h W_h^*}{M_P^4} + \Big(K_0^{S\bar{S}}(\Lambda^2 + \Lambda' S' + qS'^2)^*\frac{\partial K_0}{\partial S}\frac{W_h}{M_P^2} + K_0^{S\bar{S}'}(\Lambda^2 + \Lambda' S' + qS'^2)(\Lambda' S + 2qS'S)^* \\ & + K_0^{S\bar{S}'}(\Lambda' S + 2qS'S)^*\frac{\partial K_0}{\partial S}\frac{W_h}{M_P^2} + K_0^{S\bar{S}'}(\Lambda^2 + \Lambda' S' + qS'^2)\Big(\frac{\partial K_0}{\partial S'}\Big)^*\frac{W_h^*}{M_P^2} + K_0^{S\bar{S}'}\frac{\partial K_0}{\partial S}\Big(\frac{\partial K_0}{\partial S'}\Big)^*\frac{W_h W_h^*}{M_P^4} \\ & + K_0^{S'\bar{S}'}(\Lambda' S + 2qS'S)\Big(\frac{\partial K_0}{\partial S'}\Big)^*\frac{W_h^*}{M_P^2} + \text{H.c.}\Big) + K_0^{X\bar{X}}\frac{\partial K_0}{\partial X}\Big(\frac{\partial K_0}{\partial X}\Big)^*\frac{W_h W_h^*}{M_P^4} \\ & + \Big(K_0^{S\bar{X}}(\Lambda^2 + \Lambda' S' + qS'^2)\Big(\frac{\partial K_0}{\partial X}\Big)^*\frac{W_h^*}{M_P^2} + K_0^{S\bar{X}}\frac{\partial K_0}{\partial S}\Big(\frac{\partial K_0}{\partial X}\Big)^*\frac{W_h W_h^*}{M_P^4} + K_0^{S'\bar{X}}(\Lambda' S + 2qS'S)\Big(\frac{\partial K_0}{\partial X}\Big)^*\frac{W_h^*}{M_P^2} \\ & + K_0^{S'\bar{X}}\frac{\partial K_0}{\partial S'}\Big(\frac{\partial K_0}{\partial X}\Big)^*\frac{W_h W_h^*}{M_P^4} + \text{H.c.}\Big)\Big\} \end{aligned} \tag{14}$$

and

$$\begin{aligned} V_1 = e^{K/M_P^2}\Big\{ & \Big(\frac{\partial W_{\text{MSSM}}}{\partial \Phi_i}\Big)^*\frac{\partial W_{\text{MSSM}}}{\partial \Phi_i} + m_0^2(S, S')\Phi_i^*\Phi_i + \Big(A_1(S, S')\frac{\partial W_{\text{MSSM}}}{\partial \Phi_i}\Phi_i + A_2(S, S')W_{\text{MSSM}} + B_N(S, S', X)\tilde{N}_i^c a_{ij}\tilde{N}_j^c \\ & + \frac{XX}{M_P}\Big(\frac{\partial W_{\text{MSSM}}}{\partial \tilde{N}_i^c}\Big)^* a_{ij}\tilde{N}_j^c + \frac{XXX^*X^*}{2M_P^2}a_{ij}a_{ik}^*\tilde{N}_j^c\tilde{N}_k^{c^*} + \text{H.c.}\Big)\Big\}, \end{aligned} \tag{15}$$

where

$$\begin{aligned} m_0^2(S, S') &= \frac{W_h W_h^*}{M_P^4}, \\ A_1(S, S') &= \frac{W_h^*}{M_P^2}, \\ A_2(S, S') &= K_0^{S\bar{S}}\frac{(\Lambda^2 + \Lambda' S' + qS'^2)^*}{M_P^2}\frac{\partial K_0}{\partial S} + K_0^{S\bar{S}}\frac{\partial K_0}{\partial S}\Big(\frac{\partial K_0}{\partial S'}\Big)^*\frac{W_h^*}{M_P^4} + K_0^{S\bar{S}'}\frac{(\Lambda' S + 2qS'S)^*}{M_P^2}\frac{\partial K_0}{\partial S} \\ &\quad + K_0^{S'\bar{S}}\frac{(\Lambda' S + 2qS'S)^*}{M_P^2}\frac{\partial K_0}{\partial S'} + K_0^{S\bar{S}'}\frac{W_h^*}{M_P^4}\frac{\partial K_0}{\partial S}\Big(\frac{\partial K_0}{\partial S'}\Big)^* + K_0^{S'\bar{S}'}\frac{(\Lambda' S + 2qS'S)^*}{M_P^2}\frac{\partial K_0}{\partial S'} \\ &\quad + K_0^{S'\bar{S}'}\frac{W_h^*}{M_P^4}\frac{\partial K_0}{\partial S'}\Big(\frac{\partial K_0}{\partial S'}\Big)^* - 3\frac{W_h^*}{M_P^2}, \\ B_N(S, S', X) &= \frac{XX}{2M_P}A_2(S, S') + \frac{X}{M_P}K_0^{X\bar{S}}(\Lambda^2 + \Lambda' S' + qS'^2)^* + \frac{X}{M_P}K_0^{X\bar{S}}\Big(\frac{\partial K_0}{\partial S}\Big)^*\frac{W_h^*}{M_P^2} + \frac{X}{M_P}K_0^{X\bar{S}'}(\Lambda' S + 2qS'S)^* \\ &\quad + \frac{X}{M_P}K_0^{X\bar{X}}\Big(\frac{\partial K_0}{\partial X}\Big)^*\frac{W_h^*}{M_P^2} + \frac{XX}{M_P}\frac{W_h^*}{M_P^2}. \end{aligned} \tag{16}$$

It should be noted that the chiral superfields and their scalar components have been denoted by the same set of symbols here. The contribution to scalar potential from the gauge kinetic function is neglected here, since the scalar vev of S' is made to be zero which is shown below. By minimizing the scalar potential we have to achieve the vevs of hidden sector fields as $\langle S\rangle \sim M_P$, $\langle X\rangle \sim \Lambda$, and $\langle S'\rangle = 0$. We can achieve this by choosing an appropriate form of the function $K_0(S, S', X)$. There are six independent double derivatives and three independent single derivatives of K_0 in the scalar potential. We can fix them in such a way that the desired vevs of hidden sector fields arise after minimizing the potential. In deriving V_1, we have used $\langle \Phi_i\rangle \ll \langle X\rangle \ll \langle S\rangle$, and terms suppressed by higher powers of M_P have been neglected. Substituting the vevs of the hidden sector fields in V_1, we get the low-energy scalar potential. The first term in Eq. (15) is the F-term contribution to the scalar potential of MSSM. Remaining terms in Eq. (15) are SUSY-breaking soft masses. The soft masses, which are given in Eq. (16), are determined by the high scale parameters of our model. It can be noticed that these soft masses are at the TeV scale after plugging the vevs of hidden sector fields in their respective formulas. So our proposed model is consistent with SUSY breaking as it produces a viable low-energy scalar potential. For simplicity, we list the vevs of the scalar and auxiliary components

TABLE I. Vev's of the scalar and auxiliary components of the additional chiral superfields.

Field	Scalar vev	Auxiliary vev
S	$\sim M_P$	$\sim\Lambda^2$
S'	0	$\sim\Lambda^2$
X	$\sim\Lambda$	0
N^c	0	0

TABLE II. The different parameters of low-energy SUSY and their sources.

Parameter	Source	Order of magnitude
m_0^2	$m_0^2(S, S')$ in V_1	TeV^2
A	$A_1(S, S')$, $A_2(S, S')$ in V_1	TeV
B_μ	$\mu A_{1(2)}(S, S')$ in V_1	TeV^2
$m_{1/2}$	$\frac{F_{S'}}{M_P}$ from gauge kinetic terms	TeV

of the additional superfields that we require in our model, in Table I. In Table II we present a list of sources for soft terms in the SUSY Lagrangian, completely determined by the superpotential and scalar potential, and a clear demonstration of how they are governed by the vevs of the hidden sector fields.

While our main conclusions depend critically on various vevs in the hidden sector, it may be argued that their values can be altered through gravitational effects. However, these effects should in general be suppressed by $\mathcal{O}(\Lambda/M_P)$ and therefore can be ignored in the preliminary proposal.

As mentioned above, we need to ascertain that neither the left- nor the right-chiral sneutrinos develop any vev. To ensure this, one has to fulfill the minimization conditions [19] for the low-energy scalar potential V_1 (i.e. vanishing of the first derivatives, positivity of the eigenvalues of the second derivatives, etc.) for $\langle\tilde{\nu}_i\rangle$ and $\langle\tilde{N}_i\rangle$, simultaneously. We have checked that such solutions can be guaranteed for appropriate values of the parameters in $A_{1(2)}$, B_N, B_μ, and a_{ij} as well as the vev of X and S. Ensuring this is relatively easy, since the right-chiral sneutrinos do not occur in quartic terms (except those suppressed by M_P^4), and can develop vev only through terms linear in the vevs of the left-chiral sneutrinos. Thus it is enough to make the latter zero through an appropriate choice of parameters.

IV. CONSEQUENCES

One of the consequences of our model is that we get $\Delta L = 2$ mass terms for neutrinos. After giving vev to X the last term of Eq. (4) gives right-handed Majorana neutrino mass of the form $M_R \sim \frac{\langle X\rangle^2}{M_P} \sim$ TeV. If the neutrino Yukawa couplings: $Y_\nu \sim 10^{-7}$, then the Dirac mass for neutrinos turns out to be $m_D \sim Y_\nu v_2 \sim 10^{-4}$ GeV, where $\langle H_2^0\rangle = v_2$. It may be legitimate to take the neutrino Yukawa couplings of $\mathcal{O}(10^{-7})$, since this is the same as that of the electron Yukawa coupling and we do not understand why the electron mass is that small. If we put the Yukawa couplings of the electron and neutrinos on the same footing, we can explain the smallness of neutrino masses. In our model, since the right-handed neutrino mass is much heavier than the Dirac mass for neutrinos, the seesaw mass formula for light neutrinos is $m_\nu = -m_D^2/M_R \sim 0.1$ eV. This is the right magnitude for the neutrino mass which has been estimated from the neutrino oscillation experiments. So we can explain consistently the Majorana nature of the neutrino and its smallness of mass in our model.

Another consequence of our model is that the fermionic state belonging to the chiral field X, which we denote as ψ_X, becomes massless at the tree level. This statement follows from the last term of Eq. (4), since the right-chiral sneutrinos are made to have zero vevs. Our model requires a nonzero vev for X and the last term of Eq. (4) generates an effective term of the form

$$\frac{\langle X\rangle}{M_P} X N^c N^c \tag{17}$$

in the superpotential. Through this $\Delta L = 1$ term and the neutrino Yukawa couplings, ψ_X mixes with the neutralino states through one-loop diagrams, if the neutral H_2^0 state acquires vev. This loop diagram gives very small contribution and the mass eigenvalue of ψ_X is less than the mass of any supersymmetric particle. As a result of this, through the same one-loop diagram the lightest neutralino of MSSM can decay to ψ_X and a neutral Higgs boson. We have found that the decay width of this process is approximately given by

$$\Gamma \sim \frac{1}{8\pi}\left(\frac{g}{16\pi^2} Y_\nu^2 \frac{\langle X\rangle}{M_P}\right)^2 m_{\chi^0}, \tag{18}$$

where g is the SU(2) gauge coupling strength and m_{χ^0} is the mass of the lightest neutralino. For typical values of parameters in the above equation, the lifetime of the lightest neutralino is $\tau = \frac{1}{\Gamma} \sim 2 \times 10^{11}$ years. This value is 1 order of magnitude greater than the age of the Universe. Such a long lifetime of neutralino in our model is due to the fact that the effective $\Delta L = 1$ term in the superpotential, Eq. (17), is suppressed by a factor of $\frac{\langle X\rangle}{M_P} \sim 10^{-7}$ and the neutrino Yukawa couplings give further suppression in the loop induced decay.

In addition, the last term of Eq. (4) generates some scalar interactions in the low-energy scalar potential, which are the last three terms of Eq. (15). They can generate $\Delta L = 1$ terms through the vev of X, which have the following schematic forms:

$$\frac{\langle X\rangle}{M_P}AX\tilde{N}^c\tilde{N}^c, \qquad \frac{\langle X\rangle}{M_P}\frac{(XX)^*}{M_P}X\tilde{N}^c\tilde{N}^{c^*}, \qquad \frac{\langle X\rangle}{M_P}X(LH_2)^*\tilde{N}^c. \tag{19}$$

Here, $A \sim$ TeV. All the above three terms have a suppression factor of $\frac{\langle X\rangle}{M_P} \sim 10^{-7}$. We have found that they do not induce any tree-level decay of the neutralino state. The loop induced decays due to them will have an additional suppression of neutrino Yukawa couplings. Thus, in our model the last term of Eq. (4) is only the source for $\Delta L = 1$ terms. All of them are suppressed by Planck mass and the lepton number violation by one unit is confined only in the neutrino sector. Because of these reasons, the lightest neutralino of MSSM can have a lifetime exceeding the age of the Universe, and it can still provide a candidate for the dark matter content. Now it can be easily justified why the last one of Eq. (7) is neglected in the Kahler potential: since that term is suppressed by two powers of Planck mass, it gives subleading contributions to the $\Delta L = 1$ terms and to the neutralino decay.

The low-energy scalar potential is also a consequence of our model. The SUSY-breaking soft parameters in the scalar sector come out in the phenomenologically expected range, as listed in Table II. An intermediate scale vev of the auxiliary component of S' justifies gaugino masses in the same scale as well, through S' participating in the gauge kinetic function. The parameter B_μ, which is the coefficient of the bilinear term H_1H_2 in the scalar potential, is at the TeV scale provided the μ parameter lies around the TeV scale. While we have not justified the value of μ in this work, an explanation of the twin parameters μ and Λ' in the proposed scenario, both belonging to the superpotential and still around the electroweak scale, is hoped to come from a deeper understanding of the "μ-problem."

It may be retread that the suggested orders of magnitude of the scalar and auxiliary components vevs of the fields X, S, and S' are all consistent with observable sector SUSY-breaking parameters, being all in the TeV range. It is intimately related to the fact that the same set of choices yields a gravitino mass on the same order.

Another interesting possibility which is kept alive by such choice concerns the cosmological constant which, as is well known, needs to be fine-tuned to a miniscule value in SUSY scenarios. The dominant contribution to this constant in our model comes from the part V_0 of the scalar potential. After giving vevs to hidden sector fields, the first four lines of V_0 give a contribution of the order of Λ^4 and the last two lines give contribution of the order of Λ^2 TeV^2. Although we are not claiming to solve the cosmological constant problem, the choice of fields and orders of their vevs make it possible to envision the mutual cancellation of the dominant terms contributing to it, by proper adjustment of the dimensionless parameters occurring in the Kahler potential.

V. CONCLUSIONS

To conclude, we have suggested a supersymmetric scenario where a field carrying lepton number but otherwise immune to standard model interactions can generate $\Delta L = 2$ neutrino mass terms and also keeps the lightest neutralino of MSSM long-lived particle. In this scenario we have proposed three kinds of hidden sector fields: S, S', and X, where only X carries a lepton number. The purposes served by these fields are (i) the generation of $\Delta L = 2$ mass terms for neutrinos, (ii) the elongation of the lifetime of the LSP, decaying through $\Delta L = 1$ interactions, beyond the present age of the Universe, and (iii) the occurrence of SUSY-breaking parameters around the TeV scale, thus yielding a phenomenologically viable SUSY spectrum. The lepton number carried by X as well as the R-charge assignments of various fields ensure this, both the charges being broken at energies below the Planck scale, along a line frequently taken in SUSY models of neutrino masses. It is also demonstrated that the scenario proposed here can accommodate a cancellation of the leading terms contributing to the value of the cosmological constant. This shows the potency of supergravity theories in reconciling seesaw masses for neutrinos with the observed cold dark matter of the Universe, and underscores the importance of attempts to derive scenarios such as the aforesaid one from more fundamental principles.

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