

## Correlating $\mu$ parameter and right-handed neutrino masses in $N = 1$ supergravity

This content has been downloaded from IOPscience. Please scroll down to see the full text.

JHEP03(2006)010

(<http://iopscience.iop.org/1126-6708/2006/03/010>)

View [the table of contents for this issue](#), or go to the [journal homepage](#) for more

### Download details:

IP Address: 193.140.249.2

This content was downloaded on 13/10/2016 at 13:56

Please note that [terms and conditions apply](#).

You may also be interested in:

[Low energy six-dimensional  \$N = 2\$  supersymmetric  \$SU\(6\)\$  models on  \$T^2\$  orbifolds](#)

Jing Jiang, Tianjun Li and Wei Liao

[Interplay of type I and type II seesaw contributions to neutrino mass](#)

Evgeny Kh. Akhmedov and Michele Frigerio

[Symplectic Symmetry of the Neutrino Mass For Many Neutrino Flavors](#)

Nurcan Öztürk

[Neutrino masses: evidences and implications](#)

J W F Valle

[Neutrino masses and mixing, lightest neutralino decays and a solution to the problem in supersymmetry](#)

Pradipta Ghosh and Sourov Roy

[Little supersymmetry and the supersymmetric little hierarchy problem](#)

Andreas Birkedal, Zackaria Chacko and Mary K. Gaillard

[Convolutated solutions in supergravity](#)

A M Ghezelbash

## Correlating $\mu$ parameter and right-handed neutrino masses in $N = 1$ supergravity

---

**Durmuş A. Demir**

*Department of Physics, Izmir Institute of Technology*

*TR35430, Izmir, Turkey*

*E-mail: demir@physics.iztech.edu.tr*

**Yasaman Farzan**

*Institute for Studies in Theoretical Physics and Mathematics (IPM)*

*P.O. Box 19395-5531, Tehran, Iran*

*E-mail: yasaman@theory.ipm.ac.ir*

**ABSTRACT:** The minimal supersymmetric standard model, when extended to embed the seesaw mechanism, obtains two dimensionful parameters in its superpotential: the  $\mu$  parameter and the right-handed neutrino mass  $M_N$ . These mass parameters, belonging to the supersymmetric sector of the theory, pose serious naturalness problems as their scales are left completely undetermined. In fact, for correct phenomenology,  $\mu$  must be stabilized at the electroweak scale while  $M_N$  lies at an intermediate scale. In this work we construct an explicit model of the hidden sector of  $N = 1$  supergravity for inducing both  $\mu$  and  $M_N$  at their right scales. The model we build utilizes lepton number conservation and continuous  $R$  invariance as two fundamental global symmetries to forbid bare  $\mu$  and  $M_N$  appearing in the superpotential, and induces them at phenomenologically desired scales via spontaneous breakdown of the global symmetries and the supergravity. We discuss briefly various phenomenological implications of the model.

**KEYWORDS:** Supersymmetry Breaking, Supergravity Models, Supersymmetry Phenomenology.

---

## Contents

<b>1. Introduction</b>	<b>1</b>
<b>2. The model</b>	<b>3</b>
<b>3. Phenomenological implications</b>	<b>7</b>
<b>4. Conclusion</b>	<b>12</b>

---

## 1. Introduction

Supergravity, once spontaneously broken at a scale  $M_{SUSY}$ , gives rise to a softly-broken globally supersymmetric theory at a scale  $M_{SUSY}^2/M_{Pl}$  which corresponds to the mass scale of the gravitino,  $m_{3/2}$  [2] (notice that  $M_{SUSY}$  differs from soft masses often denoted by  $m_{SUSY}$ ). If the gravitino mass is at the weak scale,  $m_{3/2} \sim (1-10)M_{EW}$ , or equivalently, if supergravity is spontaneously broken at an intermediate scale  $M_{SUSY}^2 \sim m_{3/2}M_{Pl}$ , the gauge hierarchy problem is solved: The fact that supersymmetry is broken only softly guarantees that the electroweak scale is radiatively stable, that is, the ratio  $M_{EW}/M_{Pl} \sim m_{3/2}/M_{Pl}$  is immunized against quantum fluctuations.

The most economic description of the observable sector is realized by the Minimal Supersymmetric Standard Model (MSSM) which essentially corresponds to a direct supersymmetrization of the SM spectrum. One of the crucial aspects of the entire mechanism is the breakdown of local supersymmetry at an intermediate scale  $M_{SUSY}^2 \sim m_{3/2}M_{Pl}$  which itself poses a new naturalness problem. The question of how such an intermediate scale has been formed can be answered only through a concrete modeling of the hidden sector. Just to give an idea, one can consider, for instance, dynamical supersymmetry breaking scenarios in which all energy scales in the infrared are generated from  $M_{Pl}$  via dimensional transmutation [3]. Right here one recalls that intermediate scales like  $M_{SUSY}$  are also necessitated by other phenomena not related to supersymmetry breaking. As an example, one can allude to the Peccei-Quinn mechanism [6] which is devised to solve the strong CP problem. This mechanism is based on the presence of an intermediate scale  $M_{PQ} \sim 10^{14}$  GeV.

Another example, on which we are going to concentrate in this paper, is the famous seesaw mechanism [5] which explains the tiny but nonzero neutrino masses. The seesaw mechanism is based on the existence of ultra heavy right-handed neutrinos,  $N_i$ , which are singlets of the  $SU(3) \times SU(2) \times U(1)$  symmetry of the standard model. Indeed, in spite of several alternative models [4], the seesaw mechanism is arguably the most popular way of explaining the tiny masses of neutrinos, partly because it provides an explanation for the baryon asymmetry of the universe through a mechanism called leptogenesis [7]. Successful leptogenesis requires masses of heavy right-handed neutrinos to be above  $\sim 10^9$  GeV [8].

Obviously, low-energy phenomena do not necessitate any correlation among the mass scales  $M_{SUSY}$ ,  $M_{PQ}$  and  $M_N$ . They show up as independent scales, needed to explain distinct phenomena. However, it would establish a strong case, besides superstrings, for the existence of a supersymmetric organizing principle operating at ultra high energies if they can be correlated within a specific model. Concerning this point, one here recalls [9] in which  $M_{PQ}$  and  $M_{SUSY}$  have been correlated by using a hidden sector composite axion. The subject matter of the present work will be essentially to relate  $M_{SUSY}$  and  $M_N$ , leaving aside  $M_{PQ}$ , within  $N = 1$  supergravity.

The superpotential of MSSM-RN is given by

$$\begin{aligned} \widehat{W}_{MSSM-RN} = & \mu \widehat{H}_u \cdot \widehat{H}_d + M_N \widehat{N}^c \widehat{N}^c \\ & + Y_u \widehat{Q} \cdot \widehat{H}_u \widehat{U}^c + Y_d \widehat{H}_d \cdot \widehat{Q} \widehat{D}^c + Y_e \widehat{H}_d \cdot \widehat{L} \widehat{E}^c + Y_\nu \widehat{L} \cdot \widehat{H}_u \widehat{N}^c \end{aligned} \quad (1.1)$$

where  $\widehat{N}^c$  stands for the anti right-handed neutrino supermultiplet [5]. This model classically preserves the baryon number while  $R$  parity and SM gauge group are exact symmetries of the model even at the quantum level.

The superpotential (1.1) involves two dimensionful parameters  $\mu$  and  $M_N$ . These mass parameters are both nested in the superpotential, and thus, they bear no relation whatsoever to the supersymmetry breaking sector of the theory. Therefore, they pose a serious naturalness problem in that MSSM offers no mechanism, dynamical or otherwise, to enable one to know characteristic scales of  $\mu$  and  $M_N$ . In fact, present neutrino data already require  $M_N \sim \langle H_u \rangle^2 Y_\nu^2 / m_\nu \gtrsim 10^{14} (Y_\nu)^2$  GeV. On the other hand, LEP lower bound on chargino mass [10] leads ‘in a rather model independent way’ to a lower bound on the  $\mu$  parameter:  $\mu > 110$  GeV [11]. On the other hand, to have a successful electroweak symmetry breaking  $\mu$  has to be stabilized at the weak scale [12]. Consequently, the MSSM must be extended to provide a dynamical understanding of how and why  $\mu$  and  $M_N$  are stabilized to their phenomenologically favored scales. This is the naturalness problem we discuss in this work. Actually, part of the problem *i.e.* stabilization of the  $\mu$  parameter at the TeV scale has already been discussed in various contexts and various solutions have been devised [12]. What is left over is to understand the mechanism which stabilizes  $M_N$  at an intermediate scale of  $\sim 10^{14}$  GeV. In what follows we will attack on both naturalness problems by constructing a hidden sector model which leads to a dynamical determination of  $\mu$  and  $M_N$  upon breakdown of certain global symmetries [13] and supergravity [14]. It will turn out to be a hybrid model in that we will utilize both Kähler potential and superpotential of the hidden sector fields such that induction of the  $\mu$  parameter proceeds in a way similar to the Giudice-Masiero mechanism [14].

In section 2, we will attempt to formulate a model of hidden sector dynamics which induces  $\mu$  and  $M_N$  *simultaneously*, as a result of spontaneous local supersymmetry breaking. As will be seen, this can be accomplished by including logarithmic terms in the Kähler potential (which might be motivated by string theory [15]). In section 3, we will provide a brief discussion of certain phenomenological implications of this model. In section 4, we summarize our conclusions.

## 2. The model

Any attempt at answering the question put forward in the introduction should first provide a way of forbidding bare  $\mu$  and  $M_N$  appearing in the superpotential. The fact that  $\widehat{W}_{MSSM-RN}$  possesses additional continuous symmetries in the absence of  $\mu$  and  $M_N$  [13] implies that imposing global symmetries can protect  $\widehat{W}_{MSSM-RN}$  against the bare  $\mu$  and  $M_N$  parameters. In particular, imposing a global continuous  $R$  invariance and restoring lepton number conservation forbid the bare  $\mu$  and  $M_N$  parameters, respectively. This can be achieved by introducing new fields appropriately charged under these symmetries. Realization of this setup requires at least two hidden sector chiral superfields  $\widehat{z}_1, \widehat{z}_2$  with charges  $R(\widehat{z}_1) = 0, R(\widehat{z}_2) = 2, L(\widehat{z}_1) = 2$  and  $L(\widehat{z}_2) = -2$ . The  $R$  charge of each lepton and quark superfield equals 1 and  $R(\widehat{H}_u) = R(\widehat{H}_d) = 0$  so that the superpotential (1.1) acquires two units of  $R$  charge. The complete superpotential can be decomposed as

$$\widehat{W} = \widehat{W}_{obs} + \widehat{W}_{hid} + \widehat{W}_{obs-hid} \tag{2.1}$$

where

$$\begin{aligned} \widehat{W}_{obs} &= Y_u \widehat{Q} \cdot \widehat{H}_u \widehat{U}^c + Y_d \widehat{H}_d \cdot \widehat{Q} \widehat{D}^c + Y_e \widehat{H}_d \cdot \widehat{L} \widehat{E}^c + Y_\nu \widehat{L} \cdot \widehat{H}_u \widehat{N}^c \\ \widehat{W}_{hid} &= M_{hid} \widehat{z}_2 \widehat{z}_1 \\ \widehat{W}_{obs-hid} &= \lambda \widehat{z}_1 \widehat{N}^c \widehat{N}^c / 2 \end{aligned} \tag{2.2}$$

where, like  $M_N$ ,  $\lambda$  is a matrix in the space of right-handed neutrino flavors. This superpotential respects a global  $U(1)_{R-sym} \otimes U(1)_{Lepton}$  invariance in addition to the baryon number conservation and the MSSM gauge symmetries.

The supergravity Lagrangian is based on the Kähler potential [2]

$$\mathcal{G}(\widehat{\phi}, \widehat{\phi}^\dagger) = K(\widehat{\phi}, \widehat{\phi}^\dagger) + M_{Pl}^2 \ln \left| \frac{\widehat{W}(\widehat{\phi})}{M_{Pl}^3} \right|^2 \tag{2.3}$$

where  $\widehat{\phi} = (\widehat{z}_i; \widehat{H}_u, \widehat{H}_d, \dots, \widehat{N}^c)$  collectively denotes the chiral superfields in the hidden and observable sectors of the theory. The kinetic terms of the superfields are collected in  $K(\widehat{\phi}, \widehat{\phi}^\dagger)$ . The part of the scalar potential induced by  $F$ -terms is given by

$$V_F(\phi) = e^{K/M_{Pl}^2} \left\{ g^{ab^*} D_a W D_{b^*} W^* - 3 \frac{|W|^2}{M_{Pl}^2} \right\} \tag{2.4}$$

where

$$D_a W = \frac{\partial W}{\partial \phi^a} + \frac{W}{M_{Pl}^2} \frac{\partial K}{\partial \phi^a} \tag{2.5}$$

and  $g_{ab^*}$  is the Kähler metric:

$$g_{ab^*} = \frac{\partial^2 K}{\partial \phi^a \partial \phi^{*b}}. \tag{2.6}$$

In addition to  $V_F$ , there are contributions from  $D$ -terms as well. However,  $D$ -terms do not contribute to supersymmetry breaking since, in our model, none of the hidden sector superfields exhibits a gauge invariance.

Similarly to  $\widehat{W}$ , we decompose kinetic terms of the superfields as

$$\begin{aligned}
 K_{obs} &= \widehat{H}_u^\dagger \widehat{H}_u + \widehat{H}_d^\dagger \widehat{H}_d + \dots + \widehat{N}^{c\dagger} \widehat{N}^c \\
 K_{hid} &= \widehat{z}_1^\dagger \widehat{z}_1 + \widehat{z}_2^\dagger \widehat{z}_2 + C_1 M_{Pl}^2 \log \frac{\widehat{z}_1^\dagger \widehat{z}_1}{M_{Pl}^2} + C_2 M_{Pl}^2 \log \frac{\widehat{z}_2^\dagger \widehat{z}_2}{M_{Pl}^2} \\
 K_{obs-hid} &= \frac{1}{M_{Pl}^2} \left( \lambda_1 \widehat{z}_1^\dagger \widehat{z}_1 + \lambda_2 \widehat{z}_2^\dagger \widehat{z}_2 \right) \widehat{H}_u \cdot \widehat{H}_d + \text{h.c.}
 \end{aligned}
 \tag{2.7}$$

where  $K_{obs-hid}$  is similar to the operator used in the Giudice-Masiero mechanism [14] which solves the naturalness problem associated with the  $\mu$  parameter<sup>1</sup>. The logarithmic terms in  $K_{hid}$ , which might be inspired from strings [15], do not change the Kähler metric and are included to achieve a sensible vacuum configuration in the hidden sector<sup>2</sup>. The dimensionless couplings  $C_1$  and  $C_2$  are determined from the minimization of (2.4) and demanding zero (or very small) cosmological constant.

The scalar potential of the hidden sector fields (2.4) takes the form

$$\begin{aligned}
 V_F(z_1, z_2) &= M_{hid}^2 \exp \left[ \frac{|z_1|^2 + |z_2|^2}{M_{Pl}^2} \right] \left[ \frac{|z_1|^2}{M_{Pl}^2} \right]^{C_1} \left[ \frac{|z_2|^2}{M_{Pl}^2} \right]^{C_2} \\
 &\times \left[ |z_1|^2 \left( 1 + C_2 + \frac{|z_2|^2}{M_{Pl}^2} \right)^2 + |z_2|^2 \left( 1 + C_1 + \frac{|z_1|^2}{M_{Pl}^2} \right)^2 - 3 \frac{|z_1|^2 |z_2|^2}{M_{Pl}^2} \right]
 \end{aligned}
 \tag{2.8}$$

where  $z_{1,2}$  stand for the scalar components of  $\widehat{z}_{1,2}$ , respectively. (In what follows, we will denote their fermionic partners by  $\psi_{z_{1,2}}$ .) Clearly,  $V_F$  diverges as  $|z_1|, |z_2| \rightarrow \infty$ . Moreover, when  $C_1, C_2 < 0$  and  $C_1 + C_2 < -1$  potential is not minimized for vanishing  $z_1$  and  $z_2$ . For determining the vacuum configuration, we should solve

$$\frac{\partial V_F}{\partial |z_1|^2} = 0, \quad \frac{\partial V_F}{\partial |z_2|^2} = 0, \quad V_F(z_1, z_2) = 0
 \tag{2.9}$$

---

<sup>1</sup>In principle, the Kähler potential can include higher order terms such as

$$\Delta K_{obs-hid} = \sum_{m,n>1} \frac{\beta_{mn}}{(m!n!)^2} \left( \frac{\widehat{z}_1^\dagger \widehat{z}_1}{M_{Pl}^2} \right)^m \left( \frac{\widehat{z}_2^\dagger \widehat{z}_2}{M_{Pl}^2} \right)^n \widehat{H}_u \cdot \widehat{H}_d$$

where  $\beta_{mn}$  are dimensionless constants. In our analysis, we restrict ourselves to the minimal case and neglect such higher order effects noticing that they do not alter the scale of observable sector parameters *e.g.* the  $\mu$  parameter and Higgs bilinear term to be derived in this section.

<sup>2</sup>In general, quantum gravitational interactions do not respect global symmetries [16], and thus, Kähler potential above can receive corrections of the form

$$\Delta K_{hid} = \sum_{m,n,p,q>2} \frac{\alpha_{mnpq}}{M_{Pl}^{m+n+p+q-2}} \frac{(\widehat{z}_1^\dagger)^m \widehat{z}_1^n (\widehat{z}_2^\dagger)^p \widehat{z}_2^q}{m!n!p!q!}$$

where  $\alpha_{mnpq}$  are dimensionless constants. Recalling the fact that global symmetry breaking effects get strongly suppressed if gravity is modified near the Planckian scale [17], throughout this work we neglect such terms.

where the first two determine the extremum of the potential whereas the third is needed for nullifying the cosmological constant. These conditions lead to the constraints

$$1 + C_2 = \frac{|\langle z_2 \rangle|^2}{M_{Pl}^2}, \quad 1 + C_1 = \frac{|\langle z_1 \rangle|^2}{M_{Pl}^2}, \quad |\langle z_1 \rangle|^2 + |\langle z_2 \rangle|^2 = \frac{3}{4} M_{Pl}^2 \quad (2.10)$$

where  $F$  components of  $\widehat{z}_{1,2}$  also develop VEVs

$$\begin{aligned} \langle F_{z_1} \rangle &= 2M_{hid}^* \langle z_2^* \rangle \frac{|\langle z_1 \rangle|^2}{M_{Pl}^2} \exp \left[ \frac{K}{2M_{Pl}^2} \right] \sim M_{Pl} M_{hid} \\ \langle F_{z_2} \rangle &= 2M_{hid}^* \langle z_1^* \rangle \frac{|\langle z_2 \rangle|^2}{M_{Pl}^2} \exp \left[ \frac{K}{2M_{Pl}^2} \right] \sim M_{Pl} M_{hid}. \end{aligned} \quad (2.11)$$

The vacuum configuration is symmetric under simultaneous ( $C_1 \leftrightarrow C_2$ ) and ( $|\langle z_1^2 \rangle| \leftrightarrow |\langle z_2^2 \rangle|$ ) exchanges. Clearly,  $C_1$  and  $C_2$  have to add up to  $-5/4$ . The vanishing of the cosmological constant puts stringent constraints on the allowed ranges of  $C_1$  and  $C_2$ . Indeed, in order to have a nontrivial vacuum with  $\langle z_{1,2} \rangle, \langle F_{z_{1,2}} \rangle \neq 0$  and with vanishing energy, one needs  $-1 < C_2 < -1/4$ . (Notice that if we set  $V_F$  nonzero but equal to an exceedingly small value corresponding to the observed cosmological constant,  $C_1, C_2$  and  $\langle z_{1,2} \rangle$  get modified only slightly, leaving the overall argument similar to the case  $V_F = 0$ .) As an explicit example, let us consider the case  $(C_1, C_2) = (-5/8, -5/8)$  which gives rise to  $|\langle z_1^2 \rangle| = |\langle z_2^2 \rangle| = (3/8)M_{Pl}^2$  at which  $V(z_1, z_2) = 0$  as desired, and the matrix  $M_{i,j} = \partial^2 V_F(z_1, z_2) / \partial |z_i|^2 \partial |z_j|^2$  acquires positive eigenvalues (5.41, 5.41), guaranteeing thus the minimization of the potential. (Obviously, there is nothing special about these numerical values we assign to  $C_1$  and  $C_2$ ; they are picked up just for a fast analysis of the potential landscape.) Another example is  $(C_1, C_2) = (-75/76, -5/19)$  which yields  $|\langle z_1^2 \rangle| = 0.013M_{Pl}^2$  and  $|\langle z_2^2 \rangle| = 0.737M_{Pl}^2$  at which  $V(z_1, z_2) = 0$  as desired, and  $M_{ij}$  develops positive eigenvalues (242.71, 4.34). These case studies illustrate the behavior of the potential landscape as a function of  $C_1$  and  $C_2$ .

We can rewrite the  $F$ -terms in (2.11) as

$$\begin{aligned} \langle F_{z_1} \rangle &= 2e^{3/8} (1 + C_1)^{1+\frac{1}{2}C_1} (1 + C_2)^{\frac{1}{2}C_2} \times \langle z_2^* \rangle \times M_{hid}^* \\ \langle F_{z_2} \rangle &= 2e^{3/8} (1 + C_1)^{\frac{1}{2}C_1} (1 + C_2)^{1+\frac{1}{2}C_2} \times \langle z_1^* \rangle \times M_{hid}^* \end{aligned} \quad (2.12)$$

which are  $\mathcal{O}(M_{hid} M_{Pl})$ . Indeed, for  $(C_1, C_2) = (-5/8, -5/8)$  it turns out that  $|\langle F_{z_1} \rangle| = |\langle F_{z_2} \rangle| \approx 1.25 M_{hid} M_{Pl}$ , and for  $(C_1, C_2) = (-75/76, -5/19)$  it is that  $|\langle F_{z_1} \rangle| \approx 0.30 M_{hid} M_{Pl}$  and  $|\langle F_{z_2} \rangle| \approx 2.16 M_{hid} M_{Pl}$ . In the latter,  $|\langle F_{z_1} \rangle|$  is smaller than  $|\langle F_{z_2} \rangle|$  by an order of magnitude due to relative smallness of  $|\langle z_1^2 \rangle|$ . The gravitino mass obeys the relation

$$m_{3/2}^2 = M_{Pl}^2 e^{\mathcal{G}/M_{Pl}^2} = e^{3/4} (1 + C_1)^{1+C_1} (1 + C_2)^{1+C_2} M_{hid}^2 \quad (2.13)$$

which yields  $m_{3/2} \sim M_{hid}$ . Indeed, it gives  $m_{3/2} = M_{hid}$  for  $(C_1, C_2) = (-5/8, -5/8)$ , and  $m_{3/2} \approx 1.26 M_{hid}$  for  $(C_1, C_2) = (-75/76, -5/19)$ . Consequently, the mass parameter  $M_{hid}$  in  $\widehat{W}_{hid}$  corresponds to the gravitino mass. In other words,  $M_{hid} \simeq m_{3/2}$  in the superpotential and the fundamental scale of gravity  $M_{Pl}$  combine to break supersymmetry at the intermediate scale  $M_{\text{SUSY}}^2 \sim M_{hid} M_{Pl}$  as suggested by the sizes of the associated  $F$ -terms (2.12).

The spontaneous breakdown of local supersymmetry induces soft breakdown of global supersymmetry in the observable sector (presumably with additional hidden sector fields different from  $\widehat{z}_i$  which facilitate induction of  $\mu$  and  $M_N$  in the present model) such that each of the scalar fields acquires a mass-squared  $\sim m_{3/2}^2$  and each Yukawa interaction in  $\widehat{W}_{obs}$  gives rise to a triscalar coupling  $\sim m_{3/2}$  [2].

Concerning the parameters pertaining to the right-handed neutrino sector, from eq. (2.2), one finds that the right-handed neutrinos acquire a mass term in the superpotential

$$\frac{M_N}{2} \widehat{N}^c \widehat{N}^c \quad \text{where} \quad M_N = \lambda' \langle z_1 \rangle \tag{2.14}$$

and

$$\lambda' \equiv e^{\frac{K}{2M_{Pl}^2}} \lambda = e^{3/8} (1 + C_1)^{C_1/2} (1 + C_2)^{C_2/2} \lambda \tag{2.15}$$

in accord with the fact that, in supergravity framework, observable Yukawa couplings are related to the bare ones in the superpotential via Kähler dressing. Moreover, scalar right-handed neutrinos acquire the bilinear coupling

$$\frac{B_N}{2} \widetilde{N}^c \widetilde{N}^c \quad \text{where} \quad B_N = \lambda' \langle F_{z_1} \rangle. \tag{2.16}$$

In terms of  $\lambda'$ ,  $M_N \approx 0.61 \lambda' M_{Pl}$  and  $B_N \approx 1.25 \lambda' M_{hid} M_{Pl}$  for  $(C_1, C_2) = (-5/8, -5/8)$ , and  $M_N \approx 0.11 \lambda' M_{Pl}$  and  $B_N \approx 0.30 \lambda' M_{hid} M_{Pl}$  for  $(C_1, C_2) = (-75/76, -5/19)$ . The model, in general, predicts

$$B_N M_N^{-1} = 2m_{3/2} \tag{2.17}$$

from which it follows that  $B_N$  falls in a range that is too large for “soft leptogenesis” to be effective [18]. On the other hand,  $B_N$  is too small for inducing significant radiative effects [19]. These numerical estimates are intended for consistency check of the model.

The seesaw-induced neutrino masses are given by [5]

$$m_\nu = Y'^T_\nu M_N^{-1} Y'_\nu \langle H_u \rangle^2. \tag{2.18}$$

where again  $Y'_\nu = e^{\frac{K}{2M_{Pl}^2}} Y_\nu$ . One finds that in order to have  $m_\nu \simeq 0.05$  eV, for  $|Y'_\nu| \lesssim 1$ , the right-handed neutrino masses must satisfy  $M_N \lesssim 10^{15}$  GeV. This implies that  $(0.1 - 1) \lambda' \lesssim 10^{-3}$  or  $\lambda' \lesssim 10^{-2} (0.1 - 1)$ . For the neutrino Yukawa matrix  $Y'_\nu \sim \mathcal{O}(1)$ ,  $M_N$  reaches  $10^{15}$  GeV level and hence  $\lambda'$  takes its maximal value.

With the minimal form of the Kähler potential given in (2.7), the effective  $\mu$  parameter of the MSSM is generated to be

$$\begin{aligned} \mu &= \frac{1}{M_{Pl}^2} (\lambda_1 \langle z_1 \rangle \langle F_{z_1}^* \rangle + \lambda_2 \langle z_2 \rangle \langle F_{z_2}^* \rangle) \\ &= 2 \frac{M_{hid} \langle z_1 \rangle \langle z_2 \rangle}{M_{Pl}^4} e^{K/2M_{Pl}^2} (\lambda_1 |\langle z_1 \rangle|^2 + \lambda_2 |\langle z_2 \rangle|^2) \simeq M_{hid} \end{aligned} \tag{2.19}$$

which lies at the desired scale. This solution for naturalness of the scale of the  $\mu$  parameter<sup>3</sup> is similar to the one proposed in [14]. The corresponding soft supersymmetry breaking Higgs bilinear mass-squared parameter reads as

$$\begin{aligned}
 B_H &= \frac{1}{M_{Pl}^2} \left( \lambda_1 |\langle F_{z_1} \rangle|^2 + \lambda_2 |\langle F_{z_2} \rangle|^2 \right) \\
 &= 4 \frac{|M_{hid} \langle z_1 \rangle \langle z_2 \rangle|^2}{M_{Pl}^6} e^{K/M_{Pl}^2} (\lambda_1 |\langle z_1 \rangle|^2 + \lambda_2 |\langle z_2 \rangle|^2) \simeq M_{hid}^2 \quad (2.20)
 \end{aligned}$$

which is again at the right scale for keeping the MSSM Higgs sector sufficiently light.

The main implication of the model at hand is that  $\mu$  parameter of the MSSM and right-handed neutrino masses are correlated with each other. It might be instructive to illustrate this correlation explicitly, and this is done in Fig. 1 by plotting  $\mu$  and  $m_{3/2}$  as functions of  $M_N$  for given values of  $\lambda'$ ,  $\lambda_1$  and  $\lambda_2$ . We take  $\lambda' \sim 10^{-4}$  so that for any value of  $M_N$ ,  $\langle z_1 \rangle = (\lambda')^{-1} M_N$  and in turn  $\langle z_2 \rangle$  (see eq. (2.10)) can be calculated. Then, as suggested by the figure, for  $M_{hid} = 0.5$  TeV the  $\mu$  parameter exhibits a strong dependence on  $M_N$  depending on the values of  $\lambda_1$  and  $\lambda_2$ . Indeed, as can be confirmed by using (2.19), the  $\mu$  parameter remains around a TeV for  $\lambda_1 = \lambda_2 = 1$  whereas it swings between the unphysical value zero and the desired value TeV when either  $\lambda_1$  or  $\lambda_2$  vanishes. The reason is that  $\mu$  vanishes at zero  $\langle z_1 \rangle$ . In conclusion, the model offers a manifest correlation between  $\mu$  and  $M_N$ , and  $\mu$  gets properly stabilized to lie at a TeV when both  $\lambda_1$  and  $\lambda_2$  are nonvanishing.

For a detailed analysis of constraints on the model discussed in this section, it is necessary to confront it with the available laboratory, astrophysical and cosmological experimental data. In the next section we will provide a brief tour of the implications of the present model for a number of observables.

### 3. Phenomenological implications

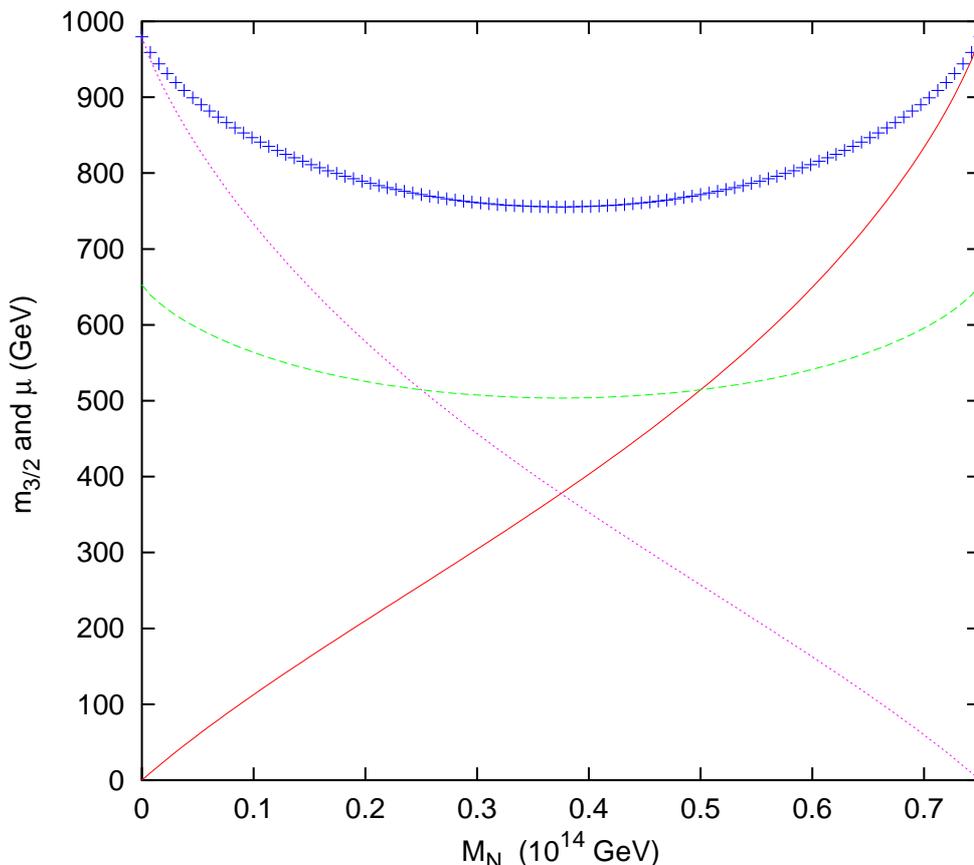
In this section we briefly discuss some phenomenologically interesting aspects of the model. The following remarks are in order:

- First of all, nonzero  $\langle z_1 \rangle$  and  $\langle z_2 \rangle$  lead to a spontaneous breakdown of the lepton number conservation and  $R$  invariance in the hidden sector. The spontaneous lepton number breaking releases a massless Goldstone boson, the Majoron  $J$ :

$$z_1 = (|\langle z_1 \rangle| + \phi_1) \times \exp i \left( \text{Arg} [\langle z_1 \rangle] + \frac{J}{|\langle z_1 \rangle|} \right) \quad (3.1)$$

---

<sup>3</sup>It might be instructive to contrast the model advocated here with low-scale MSSM extensions (by which we mean (i) the NMSSM which extends the MSSM with a chiral singlet superfield by forbidding a bare  $\mu$  in the superpotential with a  $Z_3$  symmetry, and (ii) the  $U(1)'$  models which extend the MSSM by both a chiral singlet and an additional Abelian invariance) which also solve the  $\mu$  problem in a dynamical fashion. The NMSSM suffers from the cosmological domain wall problem, and  $U(1)'$  models spoil gauge coupling unification unless one introduces a number of exotic or family non-universal charge assignments. The present model is devoid of such problems as it utilizes global continuous symmetries operating at high scale.



**Figure 1:** Variations of  $\mu$  and  $m_{3/2}$  with  $M_N$  for different values of the model parameters. We have set  $\lambda' = \frac{10^{14} \text{ GeV}}{M_{Pl}}$  and  $M_{hid} = 500 \text{ GeV}$ . The solid and dotted lines show  $\mu$  for  $(\lambda_1 = 1, \lambda_2 = 0)$  and  $(\lambda_1 = 0, \lambda_2 = 1)$ , respectively. The curve depicted with crosses shows  $\mu$  for  $(\lambda_1 = \lambda_2 = 1)$ . The dashed line stands for  $m_{3/2}$ .

where the real scalar fields  $\phi_1$  and  $J$  denote fluctuations about the vacuum state. In the flavor basis, the Majoron couples to right-handed neutrinos via  $\lambda J N^T C N$ ; however, it does not couple to the left-handed active neutrinos. This means that in the mass basis, the light active neutrinos couple to the Majoron with an exceedingly small strength  $\sim \lambda Y_\nu^2 \langle H_u \rangle^2 / M_N^2$ . This coupling is too small to have any significance [20].

In this model, due to the spontaneous breakdown of  $R$ -symmetry there is an additional Goldstone boson,  $J'$ ,

$$z_2 = (|\langle z_2 \rangle| + \phi_2) \times \exp i \left( \text{Arg}[\langle z_2 \rangle] + \frac{J'}{|\langle z_2 \rangle|} \right) \quad (3.2)$$

where we have used a parametrization similar to (3.1). The coupling of  $J'$  to the SM particles is determined by the Kähler potential and is suppressed by  $m_{SUSY}/M_{Pl} \sim 10^{-15} - 10^{-16}$  which is again too small to have any phenomenological consequences.

In the supersymmetric Majoron model infamous smajoron problem arises. In the following, we contrast the present model with supersymmetric Majoron model and compare status of smajoron problem in two scenarios. In the singlet Majoron model [21] there exists a new superfield  $\widehat{S}$  which carries 2 units of lepton number and couples to the right-handed neutrinos via  $W_1 = \widehat{S}\widehat{N}^c\widehat{N}^c$ . In similarity with the model proposed in this work, the right-handed neutrinos acquire masses through the VEV of  $\widehat{S}$ . However, in this singlet Majoron model the mechanism responsible for nonvanishing VEVs is different from the one in our model: To develop VEVs additional superfields  $\widehat{S}'$  and  $\widehat{\Lambda}$ , with respective lepton numbers equal to -2 and 0, are introduced such that they interact through the superpotential  $W_2 = \widehat{\Lambda}(\widehat{S}\widehat{S}' - M^2)$ . It can be shown that  $\langle \widehat{\Lambda} \rangle = 0$  while  $\langle \widehat{S} \rangle, \langle \widehat{S}' \rangle \neq 0$ . Since  $\langle \widehat{\Lambda} \rangle = 0$ , one out of three linear combinations of the fermionic components of  $\widehat{\Lambda}$ ,  $\widehat{S}$  and  $\widehat{S}'$  is massless, and can be interpreted as Goldstino. However, in the present model, the mass matrix of  $\psi_{z_1}$  and  $\psi_{z_2}$  (the fermionic components of  $\widehat{z}_1$  and  $\widehat{z}_2$ ) is equal to

$$e^{\frac{K}{2M_{Pl}^2}} \begin{pmatrix} 0 & M_{hid} \\ M_{hid} & 0 \end{pmatrix} \quad (3.3)$$

and does not possess any vanishing eigenvalue. Namely, massless Majoron does not have any fermionic counterpart. This difference between the two models is not surprising at all since in the singlet Majoron model, supersymmetry is preserved [ $\langle F_S \rangle = \langle F_{S'} \rangle = \langle F_\Lambda \rangle = 0$ ] and a massless bosonic particle has to have a fermionic partner whereas in our model supersymmetry is broken ( $\langle F_{z_1} \rangle, \langle F_{z_2} \rangle \neq 0$ ).

- Notice that spontaneous breakdowns of lepton number conservation and  $R$  invariance in the hidden sector reflect themselves as explicit breaking in the observable sector. Indeed, as is clear from  $\widehat{W}_{obs-hid}$ ,  $\langle z_1 \rangle \neq 0$  leads to explicit lepton number breaking via the induced right-handed neutrino mass. On the other hand,  $\langle F_{z_{1,2}} \rangle \neq 0$  lead to spontaneous breakdown of supergravity whereby inducing explicit soft-breaking terms [13]. These soft terms lead to explicit breaking of  $R$  invariance down to its  $Z_2$  subgroup, the  $R$ -parity, which is an exact discrete symmetry of the observable sector. Indeed,  $\langle F_{z_1} \rangle$ , for instance, induces a neutrino B-term which explicitly breaks the  $R$  invariance but conserves the  $R$  parity. As a result, the model accommodates a natural candidate for cold dark matter which is the famous lightest supersymmetric particle.
- In the present model,  $z_i$  and  $\psi_{z_i}$  ( $i = 1, 2$ ) acquire masses of order of  $m_{3/2}$ . On the other hand, their couplings to MSSM spectrum are rather weak. Therefore, their decay rates are expected to be suppressed. For instance, their decay into Higgs fields occur with a rate

$$\Gamma(z_i \rightarrow HH) \sim \frac{M_{hid}}{4\pi} \left( \frac{M_{hid}}{M_{Pl}} \right)^2 \quad (3.4)$$

which is suppressed by the small ratio  $\left( \frac{M_{hid}}{M_{Pl}} \right)^2 \sim 10^{-32}$ . The smallness of  $\widehat{z}_i$  decay rates into SM species is rather generic.

At first sight, it may seem that this low but nonzero decay rate may be problematic for the big bang nucleosynthesis, especially in the face of the fact that at temperatures higher than  $M_N$  these particles can be produced via  $NN \rightarrow zz$  at a rate not suppressed by  $M_{Pl}$ . Fortunately, a simple estimate shows that initial abundance of these particles cannot take too high values:

$$\begin{aligned} \frac{m_z n_z}{n_\gamma} &\sim m_z \Gamma(N \rightarrow zz) t |_{T=M_N} \sim \left( \frac{|\lambda'|^4}{4\pi^3} T \right) \left( 0.2 \frac{M_{Pl}}{\text{MeV}^2} \left( \frac{\text{MeV}}{T} \right)^2 \right) m_z \\ &\sim \left( \frac{m_z}{1000 \text{ GeV}} \right) \left( \frac{\lambda'}{10^{-4}} \right)^3 \times 10^{-12} \text{ GeV}. \end{aligned} \quad (3.5)$$

According to Fig. 3 of [25], the late decay of such particles cannot destroy the results of big bang nucleosynthesis.

- It might be instructive to analyze under what conditions one can make  $\langle z_1 \rangle$  and  $\langle z_2 \rangle$  hierarchically split so that  $M_N \sim 10^{15} \text{ GeV}$  arises with  $\lambda' \sim \mathcal{O}(1)$ . This can be achieved with a sufficiently small  $1 + C_1$ . Setting  $C_1 = -1 + \epsilon^2$  with  $\epsilon^2 \ll 1$ , eq. (2.10) implies

$$\begin{aligned} \langle z_1 \rangle &\sim \epsilon M_{Pl}, \quad \langle F_{z_1} \rangle \sim \epsilon M_{Pl} M_{hid} \\ \langle z_2 \rangle &\sim M_{Pl}, \quad \langle F_{z_2} \rangle \sim M_{Pl} M_{hid} \end{aligned} \quad (3.6)$$

which make it manifest that, for a sufficiently small  $\epsilon$ ,  $\langle z_1 \rangle$  and  $\langle F_{z_1} \rangle$  can be substantially smaller than, respectively,  $\langle z_2 \rangle$  and  $\langle F_{z_2} \rangle$ . In fact, the VEVs in (3.6) suggest that

$$M_N = \lambda' \langle z_1 \rangle \sim \epsilon \lambda' M_{Pl}, \quad B_N = \lambda' \langle F_{z_1} \rangle \sim \lambda' \epsilon M_{Pl} M_{hid} \quad (3.7)$$

both of which involve  $\epsilon M_{Pl}$  rather than  $M_{Pl}$  itself. This  $\epsilon$  dependence, however, is not present in the Higgs sector

$$\mu \simeq \lambda_2 \frac{\langle z_2 \rangle \langle F_{z_2} \rangle}{M_{Pl}^2} \sim M_{hid}, \quad B_H \simeq \lambda_2 \frac{|\langle F_{z_2} \rangle|^2}{M_{Pl}^2} \sim M_{hid}^2 \quad (3.8)$$

and  $m_{3/2} \sim M_{hid}$  still holds. Finally, one notes that  $z_i$  and their fermionic counterparts weigh now  $M_{hid}/\epsilon$  rather than  $M_{hid}$ . Taking  $\epsilon \sim 10^{-3}$ , the existing neutrino data can be explained with  $\lambda' \sim 1$ .

- Let us analyze the Higgs sector in more detail. Note that, without loss of generality, we can absorb the phase of  $M_{hid}$  by rephasing  $\hat{z}_2$ . Therefore, hereon we are going to assume that  $M_{hid}$  is real. First consider the phase of the  $\mu$  parameter. From (2.20) and (2.19) it follows that the relative phase between  $\mu$  and  $B_H$  is determined solely by the total phase of  $\langle z_1 \rangle \langle z_2 \rangle$ . More explicitly,

$$\text{Arg} \left[ \frac{\mu}{B_H} \right] = \text{Arg} [\langle z_1 \rangle] + \text{Arg} [\langle z_2 \rangle] \quad (3.9)$$

as follows from (3.1) and (3.2). This is a physical basis-independent phase that cannot be rotated away by redefinition of the relative phase between  $\widehat{H}_u$  and  $\widehat{H}_d$ . In the context of both constrained [24] and unconstrained [26] MSSM, the present bounds on electric dipole moments imply that  $\text{Arg}[\langle z_1 \rangle \langle z_2 \rangle]$  cannot exceed a few percent. Therefore, the VEVs  $\langle z_1 \rangle$  and  $\langle z_2 \rangle$  must be nearly back-to-back in the vacuum of the theory.

The structure of the model entails certain correlations among certain parameters which would bear no correlation in the MSSM, constrained or otherwise. An interesting case concerns the phase of the  $\mu$  parameter. To see this, consider the Dirac coupling of sneutrinos in the soft-breaking sector:

$$\mathcal{L} \ni a_0 Y'_\nu \widetilde{L} \cdot H_u \widetilde{N}^c + \text{h.c.} \tag{3.10}$$

which respects both the lepton number and  $R$ -parity. Here  $a_0$  is the universal trilinear coupling, as appropriate for the constrained MSSM. It can be shown that at two-loop level this operator induces a bilinear interaction between  $z_1$  and  $z_2$  as follows

$$\mathcal{L} \ni c_z^* M_{hid} z_1 z_2 + \text{h.c.} \tag{3.11}$$

where

$$c_z^* = (\text{an } \mathcal{O}(1) \text{ real number}) \times a_0 \frac{\lambda'^T Y'_\nu{}^* Y'_\nu{}^T \lambda'^*}{(16\pi^2)^2} \tag{3.12}$$

which is a small perturbation in size. However, an interaction of the form (3.11) leads to direct alignment of  $\langle z_1 \rangle \langle z_2 \rangle$  towards  $c_z$  for scalar potential (in terms of  $H_u^0, H_d^0, \widetilde{\nu}_L^i, \widetilde{N}^c_i$ ) to possess a local minimum. Therefore, the phase of  $\mu$  (i.e.  $\text{Arg}[\langle z_1 \rangle \langle z_2 \rangle]$ ) relaxes to that of  $c_z$  and hence to the phase of  $a_0$  according to (3.12). (Note that the combination  $\lambda'^T Y'_\nu{}^* Y'_\nu{}^T \lambda'^*/(16\pi^2)^2$  is real.) Consequently, the phase of  $\mu$  gets traded for that of the trilinear couplings, and as a by-product of this correlation, in future if a nonzero  $d_e$  is measured, within this model, we can extract the phase of  $\mu$  which corresponds to the phase of  $a_0$  and then, in principle, we can predict the values  $d_{Hg}$  and  $d_n$  and test these predictions in the laboratory experiments [27].

Having discussed the CP-odd phase of the  $\mu$  parameter, we now analyze the role of the  $B_H$  parameter in some depth. In general,  $\mu$  parameter can be determined by measuring masses and mixing angles of the charginos and neutralinos. On the other hand by studying the mass and decay modes of the CP-odd Higgs boson,  $A_0$ , we can derive the value of  $B_H$  [28]. Therefore, the ratio  $|B_H/\mu|$  can be measured in a rather model-independent way. It is straightforward to show that

$$2M_{hid} < |B_H/\mu| = 2m_{3/2} < 2.6M_{hid} \tag{3.13}$$

in the present model. Since interactions of gravitino are suppressed by  $M_{Pl}^{-1}$ , it cannot, in practice, be detected at colliders. However, gravitino mass affects cosmological observations, opening a window for testing this relation. Indeed, from the relation

$$|\mu|^2 + m_{H_d}^2 = B_H \tan \beta - \frac{m_Z^2}{2} \cos 2\beta \tag{3.14}$$

we expect  $|B_H/\mu| \sim |\mu|/\tan\beta$  which means that at large  $\tan\beta$ , there is a “little hierarchy” among the parameters of the model. Within our model this implies that gravitino, rather than the lightest neutralino, with mass  $\sim |\mu|/\tan\beta$ , might be the lightest supersymmetric particle and hence a candidate for cold dark matter (see, for instance, [29] for gravitino dark matter in constrained MSSM).

The discussions above provide a brief summary of the implications of the model constructed in section 2. For a proper description of the phenomenology of this model, it is necessary to perform a detailed analysis of various quantities of phenomenological interest.

#### 4. Conclusion

In this work we have constructed a hidden sector model, within  $N = 1$  supergravity, for generating, upon supergravity breakdown, the  $\mu$  parameter of the MSSM and the right-handed neutrino mass  $M_N$  at their right scales. The model utilizes global lepton number conservation and  $R$  invariance to forbid bare  $\mu$  and  $M_N$  parameters appearing in the superpotential. Moreover, the model employs a non-minimal Kähler potential exhibiting logarithmic dependencies on the hidden sector fields. The origin of these non-minimal contributions are left unexplained, yet string compactification has been an inspiring source. We have determined parameter regions where  $M_N$  and  $\mu$  come out to lie at their right scales, and found that the VEVs of the hidden sector fields can exhibit a hierarchical splitting so as to reduce unnatural tunings of the superpotential parameters. As footnoted in the text, the model at hand neither suffers from domain wall problem, nor exhibits any tension with gauge coupling unification as encountered, respectively, in NMSSM and  $U(1)'$  models.

We have confronted the model put forward with a number of observables, and identified distinctive features and ways of evading the existing bounds from various sources. The model predicts  $|B_H/\mu| = 2m_{3/2}$  which means, for a large part of parameter space, the gravitino is the lightest supersymmetric particle and thus a candidate for dark matter. Consequently, the mechanism advocated in this work possesses potential implications for various observables, a global analysis of which is yet to be performed.

#### Acknowledgments

Authors would like to thank Lotfi Boubekur, Lisa Everett and Goran Senjanovič for fruitful discussions. We are also grateful to M.M. Sheikh-Jabbari for careful reading of the manuscript and useful remarks. D.A.D. is grateful to Institute for Studies in Theoretical Physics and Mathematics (IPM) for its generous hospitality while this work was started, and International Centre for Theoretical Physics (ICTP), Turkish Academy of Sciences (through GEBIP grant), and Scientific and Technical Research Council of Turkey (through project 104T503) for their financial supports. Y.F. would like to thank ICTP for kind hospitality where part of this work was done.

## References

- [1] L. Susskind, *Dynamics of spontaneous symmetry breaking in the Weinberg-Salam theory*, *Phys. Rev. D* **20** (1979) 2619.
- [2] See the review volume: H.P. Nilles, *Supersymmetry, supergravity and particle physics*, *Phys. Rept.* **110** (1984) 1.
- [3] E. Witten, *Dynamical breaking of supersymmetry*, *Nucl. Phys. B* **188** (1981) 513;  
I. Affleck, M. Dine and N. Seiberg, *Calculable nonperturbative supersymmetry breaking*, *Phys. Rev. Lett.* **52** (1984) 1677.
- [4] A.Y. Smirnov, *Alternatives to the seesaw mechanism*, [hep-ph/0411194](http://arxiv.org/abs/hep-ph/0411194).
- [5] M. Gell-Mann, P. Ramond and R. Slansky, *Complex spinors and unified theories*, Print-80-0576, CERN;  
see also: P. Ramond, *The family group in grand unified theories*, [hep-ph/9809459](http://arxiv.org/abs/hep-ph/9809459);  
T. Yanagida, *Horizontal gauge symmetry and masses of neutrinos*, in *Proceedings of the Workshop on the Baryon Number of the Universe and Unified Theories, Tsukuba, Japan, 13-14 Feb 1979*;  
S.L. Glashow, *The future of elementary particle physics*, HUTP-79-A059 based on lectures given at Cargese Summer Inst., Cargese, France, Jul 9-29, 1979;  
R.N. Mohapatra and G. Senjanovic, *Neutrino mass and spontaneous parity nonconservation*, *Phys. Rev. Lett.* **44** (1980) 912.
- [6] R.D. Peccei and H.R. Quinn, *CP conservation in the presence of instantons*, *Phys. Rev. Lett.* **38** (1977) 1440;  
M. Dine, W. Fischler and M. Srednicki, *A simple solution to the strong CP problem with a harmless axion*, *Phys. Lett. B* **104** (1981) 199;  
A.R. Zhitnitsky, *On possible suppression of the axion hadron interactions*, *Sov. J. Nucl. Phys.* **31** (1980) 260;  
J.E. Kim, *Weak interaction singlet and strong CP invariance*, *Phys. Rev. Lett.* **43** (1979) 103;  
M.A. Shifman, A.I. Vainshtein and V.I. Zakharov, *QCD and resonance physics. Sum rules*, *Nucl. Phys. B* **147** (1979) 385.
- [7] M. Fukugita and T. Yanagida, *Baryogenesis without grand unification*, *Phys. Lett. B* **174** (1986) 45.
- [8] W. Buchmuller, P. Di Bari and M. Plumacher, *Leptogenesis for pedestrians*, *Ann. Phys. (NY)* **315** (2005) 305 [[hep-ph/0401240](http://arxiv.org/abs/hep-ph/0401240)];  
G.F. Giudice, A. Notari, M. Raidal, A. Riotto and A. Strumia, *Towards a complete theory of thermal leptogenesis in the SM and MSSM*, *Nucl. Phys. B* **685** (2004) 89 [[hep-ph/0310123](http://arxiv.org/abs/hep-ph/0310123)];  
see however, Y. Farzan and J.W.F. Valle, *R-parity violation assisted thermal leptogenesis in the seesaw mechanism*, *Phys. Rev. Lett.* **96** (2006) 011601 [[hep-ph/0509280](http://arxiv.org/abs/hep-ph/0509280)].
- [9] E.J. Chun, J.E. Kim and H.P. Nilles, *A natural solution of the  $\mu$  problem with a composite axion in the hidden sector*, *Nucl. Phys. B* **370** (1992) 105.
- [10] The LEP SUSY working group, <http://lepsus.web.cern.ch/lepsusy/>.
- [11] G. Belanger, F. Boudjema, A. Cottrant, A. Pukhov and S. Rosier-Lees, *Lower limit on the neutralino mass in the general MSSM*, *JHEP* **03** (2004) 012 [[hep-ph/0310037](http://arxiv.org/abs/hep-ph/0310037)].

- [12] P. Fayet, *Spontaneously broken supersymmetric theories of weak, electromagnetic and strong interactions*, *Phys. Lett.* **B 69** (1977) 489;  
 J.E. Kim and H.P. Nilles, *The  $\mu$  problem and the strong CP problem*, *Phys. Lett.* **B 138** (1984) 150;  
 D. Suematsu and Y. Yamagishi, *Radiative symmetry breaking in a supersymmetric model with an extra  $U(1)$* , *Int. J. Mod. Phys.* **A 10** (1995) 4521 [[hep-ph/9411239](#)];  
 M. Cvetič and P. Langacker, *New gauge bosons from string models*, *Mod. Phys. Lett.* **A 11** (1996) 1247 [[hep-ph/9602424](#)];  
 V. Jain and R. Shrock,  *$U(1)_a$  models of fermion masses without a  $\mu$  problem*, [hep-ph/9507238](#);  
 Y. Nir, *Gauge unification, Yukawa hierarchy and the  $\mu$  problem*, *Phys. Lett.* **B 354** (1995) 107 [[hep-ph/9504312](#)].
- [13] S. Dimopoulos and S.D. Thomas, *Dynamical relaxation of the supersymmetric CP-violating phases*, *Nucl. Phys.* **B 465** (1996) 23 [[hep-ph/9510220](#)];  
 D.A. Demir, *Dynamical relaxation of the CP phases in next-to-minimal supersymmetry*, *Phys. Rev.* **D 62** (2000) 075003 [[hep-ph/9911435](#)].
- [14] G.F. Giudice and A. Masiero, *A natural solution to the  $\mu$  problem in supergravity theories*, *Phys. Lett.* **B 206** (1988) 480.
- [15] L.J. Dixon, V. Kaplunovsky and J. Louis, *On effective field theories describing (2,2) vacua of the heterotic string*, *Nucl. Phys.* **B 329** (1990) 27;  
 J.P. Derendinger, S. Ferrara, C. Kounnas and F. Zwirner, *On loop corrections to string effective field theories: field dependent gauge couplings and sigma model anomalies*, *Nucl. Phys.* **B 372** (1992) 145.
- [16] M. Kamionkowski and J. March-Russell, *Planck scale physics and the Peccei-Quinn mechanism*, *Phys. Lett.* **B 282** (1992) 137 [[hep-th/9202003](#)];  
 R. Kallosh, A.D. Linde, D.A. Linde and L. Susskind, *Gravity and global symmetries*, *Phys. Rev.* **D 52** (1995) 912 [[hep-th/9502069](#)];  
 R. Holman et al., *Solutions to the strong CP problem in a world with gravity*, *Phys. Lett.* **B 282** (1992) 132 [[hep-ph/9203206](#)];  
 S.M. Barr and D. Seckel, *Planck scale corrections to axion models*, *Phys. Rev.* **D 46** (1992) 539.
- [17] R. Kallosh, A.D. Linde, D.A. Linde and L. Susskind, *Gravity and global symmetries*, *Phys. Rev.* **D 52** (1995) 912 [[hep-th/9502069](#)].
- [18] Y. Grossman, T. Kashti, Y. Nir and E. Roulet, *Leptogenesis from supersymmetry breaking*, *Phys. Rev. Lett.* **91** (2003) 251801 [[hep-ph/0307081](#)];  
 G. D'Ambrosio, G.F. Giudice and M. Raidal, *Soft leptogenesis*, *Phys. Lett.* **B 575** (2003) 75 [[hep-ph/0308031](#)].
- [19] Y. Farzan, *Effects of the neutrino  $B$ -term on the Higgs mass parameters and electroweak symmetry breaking*, *JHEP* **02** (2005) 025 [[hep-ph/0411358](#)]; *Effects of the neutrino  $B$ -term on slepton mixing and electric dipole moments*, *Phys. Rev.* **D 69** (2004) 073009 [[hep-ph/0310055](#)]; *On effects of large neutrino  $B$ -term on low energy physics*, [hep-ph/0505004](#).
- [20] G.F. Giudice, A. Masiero, M. Pietroni and A. Riotto, *The supersymmetric singlet majoron*, *Nucl. Phys.* **B 396** (1993) 243 [[hep-ph/9209296](#)];

- Y. Farzan, *Bounds on the coupling of the majoron to light neutrinos from supernova cooling*, *Phys. Rev. D* **67** (2003) 073015 [[hep-ph/0211375](#)].
- [21] R.N. Mohapatra and X. Zhang, *Cosmological constraint on the scale of the supersymmetric singlet majoron*, *Phys. Rev. D* **49** (1994) 1163 [[hep-ph/9307231](#)].
- [22] M. Bolz, W. Buchmuller and M. Plumacher, *Baryon asymmetry and dark matter*, *Phys. Lett. B* **443** (1998) 209 [[hep-ph/9809381](#)];  
M. Bolz, A. Brandenburg and W. Buchmuller, *Thermal production of gravitinos*, *Nucl. Phys. B* **606** (2001) 518 [[hep-ph/0012052](#)].
- [23] W. Buchmuller, R.D. Peccei and T. Yanagida, *Leptogenesis as the origin of matter*, [hep-ph/0502169](#).
- [24] T. Ibrahim and P. Nath, *The neutron and the lepton EDMs in MSSM, large CP-violating phases and the cancellation mechanism*, *Phys. Rev. D* **58** (1998) 111301 [[hep-ph/9807501](#)];  
D.A. Demir, O. Lebedev, K.A. Olive, M. Pospelov and A. Ritz, *Electric dipole moments in the MSSM at large  $\tan\beta$* , *Nucl. Phys. B* **680** (2004) 339 [[hep-ph/0311314](#)].
- [25] J.R. Ellis, G.B. Gelmini, J.L. Lopez, D.V. Nanopoulos and S. Sarkar, *Astrophysical constraints on massive unstable neutral relic particles*, *Nucl. Phys. B* **373** (1992) 399;  
see, however R. Allahverdi and A. Mazumdar, *Supersymmetric thermalization and quasi-thermal universe: consequences for gravitinos and leptogenesis*, [hep-ph/0512227](#);  
R. Allahverdi and A. Mazumdar, *Quasi-thermal universe: from cosmology to colliders*, [hep-ph/0505050](#).
- [26] M. Brhlik, G.J. Good and G.L. Kane, *Electric dipole moments do not require the CP-violating phases of supersymmetry to be small*, *Phys. Rev. D* **59** (1999) 115004 [[hep-ph/9810457](#)];  
S. Abel, S. Khalil and O. Lebedev, *Edm constraints in supersymmetric theories*, *Nucl. Phys. B* **606** (2001) 151 [[hep-ph/0103320](#)].
- [27] D.A. Demir and Y. Farzan, *Can measurements of electric dipole moments determine the seesaw parameters?*, *JHEP* **10** (2005) 068 [[hep-ph/0508236](#)].
- [28] J.F. Gunion, H.E. Haber, G.L. Kane and S. Dawson, *The Higgs hunter's guide*, SCIPP-89/13, 1989
- [29] J.R. Ellis, K.A. Olive, Y. Santoso and V.C. Spanos, *Gravitino dark matter in the CMSSM*, *Phys. Lett. B* **588** (2004) 7 [[hep-ph/0312262](#)].