Modulino as natural candidate for a sterile neutrino

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We discuss the possible generation of R-parity violating bilinear terms $\mu_i \bar{L}_i H_2$ in both cases of gravity and gauge mediated supersymmetry breaking. Some phenomenological aspects are reviewed. In particular for scenario where $\mu_i$ depend on the vacuum expectation values of some fields $S$, its fermionic partner $\tilde{S}$ plays the role of a sterile neutrino. This situation is quiet generic in the case of models arising from M-theory where $S$ is a modulus field. Observable effects are expected to be seen if the mass and mixing of $\tilde{S}$ with active neutrinos lie in an interesting region of parameters. This is naturally the case in gauge mediated supersymmetry breaking and some class of no-scale supergravity models where the gravitino mass is very small. For models with gravitino mass $m_{3/2} \sim$ TeV, we discuss the possibility that the modulino mass is of the order of $m_{3/2}/M_{Pl}$.

1 Introduction

Observations of the solar\(^1\), atmospheric\(^2\) anomalies and LSND events\(^3\) are experimental hints for non-zero neutrino mass and mixing. It is difficult to explain simultaneously all of these observations by masses and mixing of only the three known neutrinos. If the present data are confirmed\(^4\) one may need to introduce a very light fermion $\tilde{S}$, with a mass $m_s < 10$ eV\(^5\), called “sterile neutrino”. This fermion mixes with the active neutrinos leading to oscillation patterns that would explain the experimental data. Several models of the singlet fermions have been proposed recently\(^6\). In the present work we are interested in the investigation of the possibility that fermionic partners of moduli fields that appear in most of superstring and M-theory compactifications play the role of sterile neutrinos. We will argue that the desired neutrino-modulino mixing appears naturally in models where R-parity is broken through bilinear terms.

For the particle content of the Minimal Supersymmetric Standard Model the most general renormalizable superpotential contains an R-parity breaking part:

$$W_{nR} = \lambda_{ijk} L_i L_j E_k^c + \lambda'_{ij} L_i Q_j D_k^c + \lambda''_{ijk} D_i^c D_j^c U_k^c$$

\[ \mu_i L_i H_2 \]

(1)

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Here $m_{e_i}$, $m_{d_i}$, $m_{u_i}$ are the fermion masses while $v_{1,2}$ are the vacuum expectation values (v.e.v.) of the scalar components of the superfields $H_{1,2}$. While theories with R-parity conserving superpotential have been thoroughly investigated, the models where R-parity is broken have only recently been the focus of more attention. This is because the size of the couplings of the trilinear terms $\lambda_{ijk}$, $\lambda'_{ijk}$ and $\lambda''_{ijk}$ in $W_{nR}$ are strongly constrained by experiments. In particular they generate lepton and baryon number violation leading if combined to dangerous proton decay as well as other exotic low energy processes. The experimental bounds on these couplings makes more natural to assume that they vanish. This situation is somehow similar with the one encountered within the R-parity conserving part of the theory for the $\mu H_1 H_2$ bilinear term. The size of the mass parameter $\mu$ expected to be of the order of the Planck mass is suppressed and has to be of the order of the electroweak scale. The appearance of such scale at the tree level of the fundamental theory, where the only present scale is the Planck scale, needs a lot of fine-tuning and thus seems very unnatural. This is the so-called $\mu$-problem. The most satisfactory solution is to consider that the $\mu$-term is triggered by supersymmetry breaking. It is thus naturally of the same magnitude than the soft terms. In a similar way, if the superpotential $W_{nR}$ is generated at the supersymmetry breaking scale with a “reasonable” suppression of the different coefficients, it lies in the experimentally allowed region. Below we would like to consider the generation of the $\mu_i L_i H_2$ mixing terms and some of their implications for neutrino physics.

2 Origin of $\bar{L}H$ terms

For convenience we denote $H_1$ also by $L_0$ and we use the greek indices $\alpha = 0, 1, 2, 3$ while the latin indices take the value $i = 1, 2, 3$. The generic bilinear term allowed by the gauge interactions is then $\mu_\alpha L_\alpha H_2$. The existence and the magnitude of these terms depend crucially on the mechanism used to break supersymmetry and to communicate it to the observable sector. We give a quick overview below. More details will be given elsewhere.

2.1 Giudice-Masiero mechanism

In the same way as for the $\mu H_1 H_2$ term Giudice-Masiero mechanism can be used for generating the R-parity breaking bilinear terms. The main idea is that the Kahler potential contains terms of the form:

$$\frac{\lambda_\alpha}{\mathcal{M}_{Pl}} \int d\theta^4 \bar{z} L_\alpha H_2 \quad \text{and} \quad \frac{\lambda'_\alpha}{\mathcal{M}_{Pl}} \int d\theta^4 \bar{z}^2 L_\alpha H_2$$

(2)
where $z$ is a field responsible of the breaking of supersymmetry $\langle F_z \rangle \sim m_{3/2} \mathcal{M}_{pl}$ so that after integration, one gets the right magnitude for both $\mu$ and $B\mu$ coefficients.

In M-theory, the Kahler potential generically contains terms $K_\alpha z^2 H^2$ where $K_\alpha$ are functions of the moduli $z$. Supersymmetry breaking by the auxiliary field of one of these moduli triggers $\mu$-terms of the form:

$$\mu_\alpha = (K_{HR} K_{LL})^{-1/2} (m_{3/2} K_\alpha - \bar{F}^z \partial_z K_{\alpha 2}).$$

(3)

The relative size of the $\mu_\alpha$s is governed by the dependence of $K_{\alpha 2}$ on $z$. As an example if $z$ is one of the $T$ moduli describing the size of a compactification, there is often in the large radius limit a scaling symmetry. Matter fields might have different charges (modular weights) under this symmetry and $T$ might appear with different powers, leading to different strengths of the $\mu_\alpha$s.

2.2 Higher-weight F-term

In the presence of a light singlet $N$ with a Yukawa coupling $\lambda_\alpha NL_\alpha H_2$ there is a new contribution to $\mu_s$. This is due to higher-weight F-terms and is not usually included in the standard two-derivatives supergravity. It leads to a contribution of the form:

$$\mu_\alpha = - (K_{HR} K_{LL})^{-1/2} \langle \lambda_\alpha \bar{F}^z K^{N\bar{N}} (f_1^1 f_2^2 + f_2^2 f_1^1) \rangle$$

(4)

where $K^{N\bar{N}}$ is the inverse metric (coefficient of the kinetic term) of the singlet, $f^{(1,2)}$ are two complex functions and $z$ is the modulus superfield whose auxiliary component vev breaks supersymmetry.

Note that this possibility was found to break $R$-parity in models where the two-derivative lagrangian seems to respect the symmetry. The same discussion than for the Giudice-Masiero contribution applies here concerning the relative strengths of the $\mu_\alpha$.

2.3 Superpotential induced $\mu$-terms

While the previous possibilities for generating $\mu$ terms lead to values at most of the order of $m_{3/2}$. In scenarios where the gravitino mass is very small this is not a satisfactory option. The only known solution in this case is to appeal to extra (elementary or composite) singlet(s) $N$ with a superpotential:

$$\lambda_\alpha NL_\alpha H_2 + W(N)$$

(5)

Models for $W(N)$ can be built where the inclusion of soft terms leads at the electroweak breaking scale to a vev is for $N$ generating the desired $\mu$
terms. The situation with the strength of the resulting $\mu_\alpha$ is analog with the problem of fermion masses: $\lambda_\alpha$ might remember that the $L_\alpha$ arises from different sectors of the theory (as twists in orbifolds) or carry different charges under a new (horizontal) symmetry.

One of the implications of the R-parity violating bilinear terms is that the neutrino-neutralino mixing leads after electroweak symmetry breaking to a vev for the sneutrino. This induces a neutrino mass:

$$m_\nu \approx 9 \frac{m_Z^2}{M_{\tilde{Z}}} \left( \frac{\mu_1}{\mu} \right)^2 \left[ \frac{h_B^2}{16\pi^2} \log \frac{M_X^2}{m_W^2} \right]^2 .$$

(6) 

where $M_X$ is the scale where the soft terms are universal (where supersymmetry is broken in the hidden sector), $h_B$ is the $b$ quark Yukawa coupling, $m_Z$, $m_W$, $M_{\tilde{Z}}$ are the $Z$, $W$ bosons and Zino masses. A rough estimation for small $\tan \beta \sim 1$ gives

$$m_\nu \sim 3 \cdot 10^{-8} \left( \frac{\mu_1}{\mu} \right)^2 \frac{m_Z^2}{M_{\tilde{Z}}} .$$

(7)

For $M_{\tilde{Z}} \sim 300$ GeV, this gives $m_\nu \sim \left( \frac{\mu_1}{\mu} \right)^2 10^4$eV. Models with $\left( \frac{\mu_1}{\mu} \right) \sim 10^{-2}$ might lead to masses in the eV range for the $\tau$-neutrino.

3 Appearance of a sterile-active neutrino mixing

Superstring compactifications usually provide us (among their massless states) with massless fields singlets under the standard model gauge symmetry. These singlets can be divided into two classes according to the way they interact with the observable matter:

(i) The moduli fields: they couple to the light matter fields only through non-renormalizable interactions suppressed by power of $M_P$. Among these fields are the dilaton $S$, $T_i$ moduli, $U_i$ moduli, the continuous Wilson lines, the blowing-up modes of orbifolds. Moduli masses are induced by SUSY breaking.

(ii) The non-moduli fields: they can have renormalizable interactions with standard matter. String compactifications often lead to one anomalous $U(1)_A$ among several gauge factors $U(1)'s$ and to a number of chiral supermultiplets charged under them. To satisfy the $D$ and $F$ flatness conditions, some of these fields get large VEV’s. The resulting symmetry breaking generates a mass matrix which may have small or vanishing eigenvalues. In addition to the Higgs doublets a singlet $S$ could remain light.

Because of the dependence of the different couplings on $S$, one can write:
\[ \mu_\alpha = m_\alpha \frac{\langle S \rangle}{M_{Pl}} \]  
(8)

this leads to the effective mixing:

\[ \frac{\mu_\alpha \langle H_2 \rangle}{\langle S \rangle} \tilde{S} L_\alpha \]  
(9)

between the active neutrinos and the fermionic partner \( \tilde{S} \) of \( S \).

The magnitude of this mixing mass:

\[ m_{S\nu} = \frac{\mu_\alpha \langle H_2 \rangle}{\langle S \rangle} \]  
(10)

depends on the value of \( \langle S \rangle \).

There suppose for example that \( \mu_\alpha \sim m_W \) and consider the three natural regions to consider for the value of \( \langle S \rangle \):

1- For most of the string or M-theory moduli fields who have a geometrical interpretation as the size of most of the compactification scales:

\[ M_{GUT} \leq \langle S \rangle \leq M_{Pl} \rightarrow 10^{-5} \text{eV} \leq m_{S\nu} \leq 10^{-3} \text{eV} \]  
(11)

2- For the modulus giving the size of the segment in M-theory on \( S^1/Z_2 \):

\[ 10^{12} \text{GeV} \leq \langle S \rangle \leq 10^{15} \text{GeV} \rightarrow 10^{-2} \text{eV} \leq m_{S\nu} \leq 10 \text{eV} \]  
(12)

3- May be more exotic possibility for moduli, but other singlets leads to

\[ \langle S \rangle \sim m_{3/2} \rightarrow m_{S\nu} \sim M_W \]  
(13)

that is not in the range of our interest.

### 4 Singlet fermion mass

The supergravity mass matrix formula for the relevant fermions from chiral supermultiplets has the form:

\[ M^{\alpha\beta} = m_{3/2} (G^{\alpha\beta} - G^{\alpha\beta\dot{\gamma}} \tilde{G}_\dot{\gamma} + \frac{1}{3} \tilde{G}^{\alpha}_\dot{\gamma} \bar{G}^{\beta}_\dot{\gamma}) \]  
(14)

where \( G = K + \ln |w|^2 \), \( K \) is the Kähler potential and \( w \) is the superpotential, \( G^{\alpha} \equiv \partial G/\partial \phi^\alpha \), \( \tilde{G}^{\dot{\gamma}} \equiv \partial \tilde{G}/\partial \bar{\phi}^{\dot{\gamma}} \) etc.. The physical mass is obtained by dressing (14) with corresponding wave function renormalization factors. These factors are typically of order one and we will omit their effects here. The gravitino mass is given by \( m_{3/2} = (e^{K/2} w) \).
4.1 Slim gravitinos

The sterile neutrino mass is of the order of the gravitino mass. The present lower bound on the latter is around \(10^{-5}\) eV. Such low masses can be obtained for \(p \sim 2\) in a class of no-scale supergravity models where:

\[
K = -3\ln(z + \bar{z}) \tag{15}
\]

and

\[
f_{ab} = \delta_{ab} e^{iA z^q} \tag{16}
\]

where \(q\) is an integer lead to a relation between the gravitino mass and the gaugino mass \(M_{1/2}\):

\[
m_{3/2} \sim M_P (\frac{M_{1/2}}{M_P})^p \tag{17}
\]

with \(p = 1/(1 - 2/3q)\). The main idea behind these models is that gaugino and gravitino masses are given depend on two different functions \(K\) and \(f_{ab}\) and can be arranged to get a large hierarchy between the two scales.

Also in the case of gauge mediated scenario, the gravitino mass is in a range of few eV to the the GeV scale. The modulinos might naturally remain as light as 0.1 to 10 eV.

These two cases provide examples where the modulinos (or other singlets) might arise in models from M-theory with masses and mixing of the desired sterile neutrino.

4.2 Fat gravitinos

The situation is more complicate in the case where \(m_{3/2}\) is of the order of TeV or heavier. A mass of order \(m_{3/2}\) comes from interactions of the form:

\[
\frac{\lambda}{M_P} \int d^2 \theta z SS \tag{18}
\]

with \(z\) the “goldstino superfield”.

A sufficiently light singlet needs \(\lambda\) to be very small. A particular example where this situation would exists is that of \(S\) is identified with a twisted state in orbifold compactifications. If \(z\) is an untwisted field then it was found in that \(\lambda\) vanishes. If in contrast \(z\) and \(S\) are both twisted with a \(Z_2\) or \(Z_4\) twist, then one could consider the reasonable Kahler potential expansion:

\[
K_2 = K_{s\bar{s}s} + \frac{z^2}{M_{Pl}^2} (ss + h.c.) + \ldots \tag{19}
\]

6
and the superpotential vanishes. Then according to (20) a desired mass of $S$ can be obtained for $\langle Z \rangle = \Lambda_{hid}$ and $\langle \mathcal{G}_i \rangle = 0$, or for $\langle Z \rangle = m_{3/2}$ and $\langle \mathcal{G}_i \rangle \approx 1$.

Diagonalizing mass matrix (14), and writing $\mathcal{G}$ in terms of mass eigenstates we get a necessary condition for the mass of the singlet $S$ to be of the order $m_{3/2}/M_P$:

$$\langle G^{ss} - G^{ss\dagger} \gamma \rangle \sim \frac{m_{3/2}}{M_P}$$

while typically it is of order $O(1)$. Allowing for arbitrary dependence of the superpotential on the moduli, it is possible to find functions $\mathcal{G}$ (that is, $K$ and $w$) which satisfy (20).

As an example consider the case of moduli with large ($\sim M_P$) VEV’s with Kahler potential of the form:

$$K = p \ln(\bar{s} + s) + ...$$

where $s \equiv S/M_P$; $p$ is an integer (typically $p = -1, -2, -3$). If a theory is invariant under shifts $S \to S + i$, then the full superpotential would have an expansion in powers of $e^{-2\pi s}$: $w \sim e^{-2\pi s} \sum a_n e^{-2\pi ns}$.

For $a \sim p/2\pi (\langle \bar{s} + s \rangle)$, it is easy to find solutions of (20) with $s \sim 1$ with $\mathcal{G}_3 = 0$ or $\mathcal{G}^\gamma \sim 1$ and $\langle Z \rangle \sim m_{3/2}$. For other values of $a$ (20) implies large coefficients $a_n \sim e^{2\pi n}$. Such large $a_n$ are not excluded. For instance, the modular forms like the $j$-function which can appear in theories with an $SL(2, \mathbb{Z})$ symmetry have such large coefficients.

Moduli fields usually describe some geometrical patterns of the compactification they might be subject to some discrete symmetries. The scalar and fermion singlet components masses depend on the supersymmetry breaking mechanism. The main point is that modulino mass is as induced by (18) depends on the interactions between $S$ and the goldstino direction. Recent developments in M-theory might help understanding the physics of moduli and supersymmetry breaking and thus shade a light on the question of their mass.

When might also try to build models where the smallness of $m_S$ can originate from mixing of $S$ with fields, $\phi$, getting a Planck scale mass, if the $S\phi$-mixing is the order $m_{3/2}$. The latter scale can appear in the same way as the $\mu$-term appears. It can be protected by additional $U(1)'$ gauge symmetry broken at $m_{3/2}$, if $S$ is charged under $U(1)'$, whereas $\phi$ is a singlet of this group. Then for the mass of $S$ we have the usual see-saw formula: $m_S = m_{3/2}^2/M_P$.

Another possibility is when the superfield $S$ charged under $U(1)'$ gets a VEV of the order $m_{3/2}$. This VEV will lead to mixing of the fermion $S$ and gaugino associated with $U(1)'$. If this gaugino has the Majorana mass of the order $M_P$, then again the see-saw mechanism leads to the desirable mass of $S$. 

7
5 Experimental signature of a sterile neutrino

A manifestation of $\tilde{S}$ depends on its mixing angle $\theta$ with active neutrinos:

$$\tan 2\theta = \frac{2m_{\nu S}}{m_S - m_\nu},$$  \hspace{1cm} (22)

where $m_\nu$ is the neutrino mass and oscillation parameters, $\Delta m^2 \equiv m_S^2 - m_\nu^2$.

Consider first scenarios like the no-scale models where $m_3/2$ can be of the order of $10^{-3}$ eV, or suppressed modulino masses as for $\sim m_3^2/2M_{Pl}$ in case of a heavy gravitino. One gets values $\Delta m^2 \approx m_S^2$ and $\sin^2 2\theta$ in the range of small mixing solution of the $\nu_\odot$-problem via the resonance conversion $\nu_e \to \tilde{S}$ in the Sun.

Let us consider the mixing of $\tilde{S}$ with the electron neutrino $m_\tilde{S} > m_1$ ($m_1$ is the mass of the dominant component of $\nu_\mu$). Forthcoming experiments, and in particular SNO, will be able to establish whether conversion of solar neutrinos to singlet state takes place.

Sensitivity of the $\nu_\odot$ data to the neutrino parameters is determined by the adiabaticity condition for lowest detectable energies ($E \sim 0.2$ MeV):

$$\Delta m^2 > \frac{4 \cdot 10^{-10} \text{ eV}^2}{\sin^2 2\theta}$$  \hspace{1cm} (23)

We find that $\nu_\odot$ experiments are sensitive to $m_{\nu S} > 10^{-5}$ eV.

If on the contrary the mass and mixing of $\tilde{S}$ are outside the region of solutions of the $\nu_\odot$-problem and some other mechanism is responsible for the deficit of the $\nu_\odot$-fluxes, then the $\nu\tilde{S}$-mixing gives corrections to the main solution.

Let us for example assume for that the neutrino mass spectrum has a hierarchy with $m_3 \sim (2 - 4) \cdot 10^{-3}$ eV in the range of solution of the $\nu_\odot$-problem via $\nu_e \to \nu_\mu$ conversion, $m_3 \gg m_2$ and $m_1 \ll m_2$. There are two generic consequences of the $\nu\tilde{S}$-mixing:

(i). Final neutrino flux contains not only the electron and muon components but also the $\tilde{S}$- component. Moreover, the content (relative values of different fluxes) depends on neutrino energy. Future measurements of the neutral current interactions, and in particular, the ratio of neutral to charged current events, $(NC/CC)$, in different parts of the energy spectrum will allow to check the presence of $\tilde{S}$-flux.

(ii). A dependence of the $\nu_e$-suppression factor on energy (so called “suppression pit”) is modified. In particular, one may expect an appearance of second pit or the narrow dip in the non-adiabatic or adiabatic edges of the two
This can be revealed in measurements of energy spectra of the boron- or pp-neutrinos.

The system of three states, $\tilde{S}$, $\nu_e$, $\nu_\mu$, relevant for the problem, has in general three resonances. The interplay of transitions in these resonances leads to a variety of possible effects which depend on the adiabaticity conditions in different resonances and on the mass of $\tilde{S}$.

If $m_S > m_2$, then system has two resonances $\nu_e - \tilde{S}$ and $\nu_e - \nu_\mu$. Analyzing level crossing scheme we find that flavor composition of the final flux can change with increase of neutrino energy in the following way: $(\nu_e) \rightarrow (\nu_e, \nu_\mu \text{ or } \nu_\mu) \rightarrow (\nu_e, \nu_\mu, \tilde{S})$ (here dominant components are indicated only).

If $m_S < m_2$ the system has three resonances and a change of the flavor composition with increase of neutrino energy can be as follows: $(\nu_e) \rightarrow (\nu_e, \tilde{S} \text{ or } \nu_e) \rightarrow (\tilde{S}, \nu_\mu \text{ or } \nu_\mu) \rightarrow (\nu_e, \nu_\mu, \tilde{S}) \rightarrow (\nu_e, \nu_\mu, \tilde{S})$.

For $m_S < m_1 < m_{\nu_S}$ the $\nu_e \tilde{S}$ mixing is large, so that vacuum oscillations $\nu_e \leftrightarrow \tilde{S}$ on the way from the Sun to the Earth become important. If $\Delta m^2 \gg 10^{-1} \text{ eV}^2$, there is an energy independent suppression of the $\nu_e$-flux by factor $1 - 0.5 \sin^2 2\theta_{eS}$ for the energies outside $\nu_e - \nu_\mu$ suppression pit. For smaller values of $\Delta m^2$ one expects non trivial interplay of the vacuum oscillations and resonance conversion. If $m_S < m_{eS} \sim 10^{-5} \text{ eV}$, then $\nu_e \leftrightarrow S$ vacuum oscillations alone can explain the $\nu_\odot$ data.

Let us now consider other possible consequences of the $\nu \tilde{S}$-mixing. Models of supernovas predict power dependence of density $\rho \propto R^{-3}$ below the envelope, in contrast with exponential dependence for the Sun. Therefore the dependence of the adiabaticity condition on the oscillation parameters differs from (23), and consequently, for the border of the sensitivity region we get

$$\Delta m^2 > A \frac{10^{-8} \text{ eV}^2}{\sin^2 2\theta}. \quad (24)$$

Here $A \sim \mathcal{O}(1)$ depends on the model of star. As follows from fig.1, the $\nu \tilde{S}$-mixing can lead to appreciable transitions for $\Delta m^2 < 10^{-1} \text{ eV}^2$. This inequality corresponds via the resonance condition to densities $\rho < 10^5 \text{ g/cm}^3$. (For larger $\Delta m^2$ and larger densities the mixing mass $m_{\nu_S}$ can not satisfy the adiabaticity condition.) Therefore $\nu \tilde{S}$-mixing does not influence dynamics of collapse. For this the resonance transition should take place at $\rho > 10^{11} \text{ g/cm}^3$. Also this mixing has no impact on the supernova nucleosynthesis which occurs in central regions where $\rho > 10^6 \text{ g/cm}^3$. The $\nu \tilde{S}$-mixing can, however, lead to a resonance conversion in external regions of star thus strongly modifying properties of neutrino fluxes which can be detected on the Earth. For instance, if the neutrinos have the mass hierarchy: $m_3 = 1 - 10 \text{ eV}$, $m_1 \ll m_2 =$
$10^{-3} - 10^{-1}$ eV and $m_S < m_1$, then the resonance conversion $\bar{\nu}_e \to \bar{S}$ will lead to partial or complete disappearance of the $\bar{\nu}_e$-signal. The observation of the $\bar{\nu}_e$ signal from SN87A allows to put a bound on $\bar{\nu}_e \to \bar{S}$ transition. Furthermore, if the adiabaticity condition is fulfilled in $\nu_\mu S$-resonance, then conversions $\nu_e \to \nu_\tau$ and $\nu_\mu \to \bar{S}$ lead also to disappearance of $\nu_e$-flux. Notice that without $\nu_\mu S$-mixing one would expect $\nu_\mu \to \nu_e$ transition which results in a hard (corresponding to initial $\nu_\mu$) $\nu_e$-spectrum.

For $m_1 < m_S < m_2$ the resonance conversion $\nu_e \to \tilde{S}$ occurs in supernovas. A simultaneous transition $\nu_\mu \to \nu_e$ will lead to hard $\nu_e$-spectrum and disappearance of the $\nu_\mu$-flux.

A different possibility is that all neutrinos are massless, then $\nu \to \tilde{S}$ transition can solve the $\nu_\odot$-problem. For supernova neutrinos it can lead to disappearance of the $\nu_e$-flux and to conversion $\bar{\nu}_\mu \to \tilde{S}$.

The $\nu \to \tilde{S}$ oscillations in the Early Universe generate $\tilde{S}$ components which increases the expansion rate of the Universe and therefore influences the primordial nucleosynthesis. The $\nu S$-mixing is important for $\Delta m^2 < 10^{-1}$ eV$^2$.

Scenario like the gauge mediated supersymmetry breaking (and some class of no-scale models) may prefer a modulino mass (roughly of the order of the gravitino mass) in the range 0.1-10 eV. The natural mixing mass $m_{S\nu}$ lies between $10^{-3}m_S$ and $m_S$. Phenomenological implications of this scenario are under investigation.

6 Conclusion

In conclusion, M-theory (or string) compactification often leads to moduli superfields, singlet of standard model group, which couple with observable matter via the Planck mass suppressed constant. If some of them remain very light as suggested in this work then they will have a number of manifestations in neutrino physics. Forthcoming experiments will be able to check whether neutrinos have mixing with such singlets.

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