Neutrino - Modulino Mixing

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We suggest an existence of light singlet fermion, S, which interacts with observable matter only via Planck mass suppressed interaction: \( \sim m_{3/2}/M_p \), where \( m_{3/2} \) is the supergravity gravitino mass. If the mass of the singlet equals \( \sim m_{3/2}/M_p \), then \( \nu_e \to S \) resonance conversion solves the solar neutrino problem or leads to observable effects. The \( \nu S \)-mixing changes supernova neutrino fluxes and has an impact on the primordial nucleosynthesis. The singlet \( S \) can originate as the supersymmetric partner of the moduli fields in supergravity or low energy effective theory stemming from superstrings. The \( \nu S \)-mixing may be accompanied by observable R-parity breaking effects.

I. Introduction. Neutrinos played a key role in the construction and tests of the Standard model. It is believed that neutrino mass, if non-zero, implies physics beyond the Standard model. We argue that neutrino properties, being sensitive to the Planck scale suppressed interactions, may open a window to hidden world.

Some time ago it was marked that the Planck scale \( (M_p = 2.4 \cdot 10^{18} \text{ GeV}) \) suppressed interactions can be relevant for neutrino physics [4]. Namely, the coupling \( M_p^{-1} LLHH \), where \( L \) is the lepton doublet and \( H \) is the Higgs doublet of the Standard model, generates neutrino mass \( \left( \langle H \rangle^2 / M_p \sim 10^{-3} \text{eV} \right) \) which can lead to observable effects in solar and supernova neutrinos.

There is a number of statements that neutrinos may reveal novel very weak interactions with new particles. In last years this idea has taken rather concrete shape. Observations of the solar and atmospheric neutrinos, large scale structure of the Universe (the need of the hot component of the dark matter), LSND events etc., testify for non-zero neutrino mass and lepton mixing [5]. It is difficult to explain simultaneously all (or even some) of these observations by masses and mixing of only three known neutrinos. In this connection new very light (\( m < 10 \text{ eV} \)) neutral fermions \( S \) which mix with usual neutrinos are introduced [6]. The LEP bound on the number of neutrino species, implies that fermions \( S \) should be singlets of \( SU(2) \times U(1) \), i.e. "sterile" neutrinos. Several models of the singlet fermions have been proposed recently. The singlet can be a component of 27-plet in \( E'_6 \) models [7]. It may have a supersymmetric origin and its properties may be related to the \( R \)-symmetry [8]. It could be a Nambu-Goldstone fermion in the supersymmetric theory with spontaneously broken global symmetry like lepton number or Peccei-Quinn symmetry [9].

Another suggestion is that the singlet is a neutrino from a mirror world [10]. The mirror neutrinos mix with usual neutrinos via the Planck scale suppressed interactions: \( L^M LH^M H/M_p \), where \( L^M \) and \( H^M \) are the mirror lepton and Higgs doublets correspondingly.

In this letter we consider new possibilities related to superstring theories.

II. Observation. A majority of extensions of the Standard model contain singlets of the \( SU(2) \times U(1) \). We suggest that among these singlets there is at least one, \( S \), with the following properties:

(i) \( S \) has only the Planck mass suppressed, \( 1/M_p \), interactions with the observable matter. In the simplest version the only light scale relevant for singlets is the gravitino mass \( m_{3/2} \). Therefore a dimensionless coupling constants with observable sector could be as \( \lambda = \alpha m_{3/2}/M_p \), where \( \alpha = \mathcal{O}(1) \). The mixing of \( S \) with neutrinos involves the electroweak symmetry breaking, and the simplest appropriate effective operator is \( \chi LSH \).

This operator generates a mass term \( m_{\nu S} \nu S \) with

\[
m_{\nu S} = \eta \frac{m_{3/2} \langle H \rangle}{M_p}.
\]

Here \( \eta \) is the renormalization effect and \( \nu = \nu_e, \nu_\mu \) or \( \nu_\tau \).

(ii) The mass of singlet, \( m_s \), is induced when supersymmetry is broken. We suggest that \( m_s \) is absent at the level \( m_{3/2} \) and appears as

\[
m_s = \beta \frac{m_{3/2}^2}{M_p},
\]

where \( \beta = \mathcal{O}(1) \). It turns out that for supergravity value \( m_{3/2} \sim 1 \text{ TeV} \) masses (1) and (2) are in the range needed for a solution of the solar neutrino (\( \nu_S \)-) problem.

III. The singlet with properties (1,2) has a rich phenomenology. Taking \( m_{3/2} = (0.1 - 3) \text{ TeV} \) and \( \alpha, \beta \) in the interval 0.5 - 2 we find from (1,2): \( m_{\nu S} = (0.2 - 10) \cdot 10^{-4} \text{ eV} \) and \( m_s = (4 \cdot 10^{-6} - 4 \cdot 10^{-8}) \text{ eV} \). Manifestations of \( S \) depend on mixing angle \( \theta \) with neutrino: \( \tan 2\theta = 2m_{\nu S}(m_S - m_\nu)^{-1}, \) where \( m_\nu \) is the neutrino mass. From this we find a relation between mass \( m_{\nu S} \), and the oscillation parameters, \( \Delta m^2 \approx m_S^2 - m_\nu^2, \sin^2 2\theta \):

\[
\Delta m^2 \approx 4m_{\nu S}^2 \frac{\cos 2\theta}{\sin^2 2\theta}.
\]

According to (1) a spread of possible values \( m_{\nu S} \) fixes region (band) of the oscillation parameters \( \Delta m^2, \sin^2 2\theta \) which can follow from \( \nu S \)- mixing (fig. 1).

For \( m_{3/2} \sim (1 - 2) \text{ TeV} \), and \( \alpha, \beta \sim 1 - 2 \) we get from (1,2) values \( \Delta m^2 \approx m_S^2 \) and \( \sin^2 2\theta \) in the range of small mixing solution of the \( \nu_S \)-problem via the resonance conversion \( \nu_e \to S \) in the Sun (fig. 1). Notice that mixing angle relevant for solar neutrinos is
determined by the ratio of the electroweak scale and the gravitino mass: \( \theta \sim \mu_S/m_S \sim (H)/m_{3/2} \). Forthcoming experiments, and in particular SNO [8], will be able to establish whether this conversion takes place or not.

From (3) and (4) we find that the sensitivity of the neutrino mass spectrum has a hierarchy \( m_3 > m_2 \gg m_1 \) with \( m_2 \sim (2 - 4) \cdot 10^{-3} \) eV in the range of solution of the \( \nu_S \)-problem via \( \nu_e \rightarrow \nu_\mu \) conversion. A presence of \( \nu_S \)-mixing will modify this solution in the following way:

(i). Final neutrino flux contains not only the electron and muon components but also the S-component. Moreover, the content (relative values of different fluxes) depends on neutrino energy. For example, if \( m_S > m_2 \), we find [4] that flavor composition of the final flux can change with increase of neutrino energy as \( (\nu_e) \rightarrow (\nu_e, \nu_\mu) \rightarrow (\nu_\mu, S) \to (\nu_e, \nu_\mu, S) \).

Future detection of the neutral current interactions, and measurements of the ratio of neutral to charged current events, \((NC/CC)\), in different parts of the energy spectrum will allow to check the presence of \( S \)-flux.

(ii). A dependence of the \( \nu_\mu \)-suppression factor on energy (“suppression pit”) is modified. One may expect an appearance of second pit or narrow dip in the non-adiabatic or adiabatic edges of the two neutrino suppression pit [4]. This can be revealed in measurements of energy spectra of the boron- or \( pp \)-neutrinos.

For \( m_S < m_1 < m_{\nu_S} \) the \( \nu_S \)-mixing is large, so that vacuum oscillations \( \nu_e \leftrightarrow S \) on the way from the Sun to the Earth become important. If \( \Delta m^2 \gg 10^{-10} \) eV\(^2\), the \( \nu_S \)-mixing gives additional 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 \( m_S < m_{\nu_S} \sim 10^{-5} \) eV, the \( \nu_e \leftrightarrow S \) oscillations alone can explain the \( \nu_\mu \)-data.

Let us consider possible consequences of the \( \nu_S \)-mixing for the supernova neutrinos. Using density distribution \( \rho \propto R^{-4} \) below the envelope of star \((R \text{ is the distance from the center})\) we get from the adiabaticity condition the sensitivity region

\[ \Delta m^2 \cdot \sin^2 2\theta > A \cdot 10^{-8} \text{ eV}^2. \]  

Here \( A \sim O(1) \) depends on a model of star. As follows from fig.1, the \( \nu_S \)-mixing can lead to appreciable transitions for \( \Delta m^2 < 10^{-11} \) eV\(^2\). This inequality corresponds via the resonance condition to densities \( \rho < 10^8 \) g/cm\(^3\). Therefore \( \nu_S \)-mixing does not influence both dynamics of collapse \((\rho > 10^8 \) g/cm\(^3\)) and the supernova nucleosynthesis \((\rho > 10^6 \) g/cm\(^3\)) [11] which occur in the central regions of star. The \( \nu_S \)-mixing can, however, induce a resonance conversion in external regions of star thus strongly modifying properties of neutrino fluxes which can be detected on the Earth. If neutrinos have the mass hierarchy: \( m_3 = 1 - 10 \) eV, \( m_1 \ll m_2 = 10^{-3} - 10^{-1} \) eV and \( m_S < m_1 \), the resonance conversion \( \nu_e \to S \) will lead to partial or complete disappearance of the \( \nu_e \)-signal. The observation of the \( \nu_e \) signal from SN 87A allows one to put a bound on \( \nu_e \to S \) transition [11]. Furthermore, if the adiabaticity condition is fulfilled in \( \nu_\mu \)-resonance, the transitions \( \nu_e \to \nu_\tau \) and \( \nu_\mu \to S \) lead also to disappearance of the \( \nu_e \)-flux.

The \( \nu \to S \) oscillations in the Early Universe generate S components which increases the expansion rate and therefore influences the primordial nucleosynthesis [12]. As follows from fig.1, the \( \nu_S \)-mixing is important for \( \Delta m^2 < 10^{-11} \) eV\(^2\). This mixing can produce large leptonic asymmetry of the Universe even for larger values of \( \Delta m^2 \) which correspond to \( \sin^2 2\theta > 10^{-8} \) [12]. Note that \( m_S^2 - m_{\nu_e}^2 > 0 \) implied by the \( \nu_S \)-problem corresponds to weaker bound from PNS (line a in fig.1) [4].
IV. Origin of $S$. There are several possibilities to get a singlet with desired properties \cite{1} and \cite{2}. The latter imply its origin in a hidden sector of the supergravity theory. Superstring compactifications lead to existence of massless singlets which can be divided into two classes.

Moduli fields (dilaton, $T_i$ and $U_i$ moduli, continuous Wilson lines, blowing-up modes of orbifolds) couple to the observable matter fields only through non-renormalizable interactions suppressed by power of $M_P$. At string perturbative level moduli have flat potential: their VEV’s are unfixed and masses are zero. It is believed that non-perturbative effects fix VEV’s of moduli ($\sim M_P$ for those having geometrical interpretation) and generate masses. If the same non-perturbative phenomena are also responsible for SUSY breaking, one expects that masses of modulinos are at most $O(m_{3/2})$ \cite{14},

Non-moduli singlet fields can have renormalizable interactions with observable matter. The compactifications typically lead to several additional $U(1)'$ gauge factors (one of which can be anomalous) and to a number of chiral supermultiplets singlets of standard symmetry group but charged under $U(1)'$ factors. Some of singlets acquire large ($\sim M_P$) VEV’s thus breaking $U(1)'$ factors. The mass matrix of the chiral fields generated as the result of this breaking may have small eigenvalues.

An interesting possibility is one that the non-anomalous $U(1)'$ can be broken at low scale: $O(m_{3/2})$ \cite{15}, so that mass and mixing of $S$ (charged under this $U(1)'$) are protected by this symmetry.

V. Singlet fermion mass. The supergravity mass matrix formula for the fermions from (singlet) chiral supermultiplets has the form:

$$M^{\alpha\beta} = m_{3/2} N \left( G^{\alpha\beta} - G^{\alpha\bar{\gamma}} G_{\bar{\gamma}} + \frac{1}{3} G^2 G^{\bar{\gamma}} \right),$$

where $G \equiv K + \ln |w|^2$, $K$ is the Kähler potential and $w$ is the superpotential, $G^{\alpha} \equiv \partial G/\partial \phi^\alpha$, $G^{\bar{\gamma}} \equiv \partial G/\partial \bar{\phi}^\gamma$ etc., $N$ is the wave function renormalization factors (typically of the order one), $m_{3/2} = O(K^{1/2} w)$.

The conditions for modulino to be very light take a simple form if the Kähler function $G$ is written in terms of mass eigenstates. The singlet should be in the superfield which does not break supersymmetry, that is,

$$\langle G^S \rangle = 0$$

(7) (otherwise it will be eaten by the gravitino through the superHiggs mechanism). The condition \cite{1} ensures the minimum of the potential: $V^S = 0$.

Using \cite{1} we can write a necessary condition for the mass of the singlet $S$ to be of the order $m_{3/2}/M_P$:

$$\langle G^{SS} - G^{SS\bar{\gamma}} G_{\bar{\gamma}} \rangle \sim m_{3/2}/M_P.$$  

If $S$ does not mix with fields which break SUSY, that is, $\langle G^{SS} \rangle = 0$ for all $\langle G_{\bar{\gamma}} \rangle \neq 0$, then \cite{1} is reduced to $\langle G^{SS} \rangle \approx 0$, while usually one expects $\langle G^{SS} \rangle \sim O(1)$ \cite{14}.

Let us consider how the conditions \cite{13} could be implemented for some simple Kähler potentials which are known to arise from string compactifications.

If $\Phi$ is one of the moduli describing geometry of the compactified space of orbifolds or Calabi-Yau (like the $T$ moduli) then $\langle \Phi \rangle \sim 1$ and in the large volume approximation the Kähler potential has the form:

$$K = p \ln(\Phi + \Phi - z^\gamma z^\gamma) + K_{ho}.$$  

Here $z^\gamma$ represent Wilson line moduli and matter fields. $K_{ho}$ stands for all higher order corrections and unknown non-perturbative contributions. The fields are in units of Planck mass; $p$ is an integer (typically $p = -1, -2, -3$).

We find $DetK_{\gamma\gamma} = 0$, and in the case of one field $z$ the state with zero eigenvalue equals $S = \cos \alpha \Phi + \sin \alpha z$, where $\tan \alpha = -1/(z)$. If $G \approx K$, then fermion $S$ satisfying condition \cite{8} will have zero mass. A finite contribution to $m_S$ can follow from $K_{ho}$ or/and non-perturbative part of the superpotential related to SUSY breaking.

If VEV of $z$ is small or zero the conditions \cite{13} for $S \approx \Phi$ can be satisfied by cancellation of contributions from the Kähler potential and superpotential. The cancellation can be easily realized for polynomial superpotentials. In particular, for $w = a(S - \langle S \rangle)$ the condition \cite{9} is fulfilled automatically. It is believed, however, that the whole theory is invariant under the shift $S \rightarrow S + i \alpha \gamma$. In this case the superpotential has a general form $w = e^{-2\alpha S} \sum_n a_n \exp(-\pi n(S - \langle S \rangle))$. If $a \approx p/2\pi(S + \tilde{\gamma})$ the series converges very quickly and we can write the superpotential explicitly as:

$$w \approx A e^{2\alpha S} \left[ 1 + \frac{p}{4\pi^2} e^{-\pi n(S - \langle S \rangle)} \right].$$

Here $A \sim m_{3/2} M_P^2$. For other values of $a$ the coefficients in the expansion are large.

The mass of the singlet $m_S$ can be generated by second term in \cite{8}. For the Kähler potential \cite{9} we get $m_S = -2pm_{3/2}(z)\langle G_{\bar{\gamma}} \rangle/(S + \langle S \rangle)\gamma$ and a correct order of magnitude is achieved for $\langle z \rangle = m_{3/2}$ and $\langle G_{\bar{\gamma}} \rangle = 1$.

If the field $S$ has a small ($\ll M_P$) or vanishing VEV, the Kähler potential can be expanded as:

$$K = K^{SS} SS + \frac{zz}{M_P^2} (SS + h.c.) + K_{ho}.$$  

and the superpotential can be $a priori$ an arbitrary holomorphic function of $S$. The Kähler potential \cite{11} mixes $S$ with the field $z$. If $G^S = 0$ and $\langle w^{SS} \rangle = 0$ (which is easy to satisfy), the mass of $S$ can be written as $m_S = m_{3/2}(2zz - 2\bar{z}^2 G^2 - S^2)$ . Now there are different ways to get a desired value of $m_S$: (i) $\langle z \rangle = \Lambda_{hid} \sim (m_{3/2} M_P)^{1/2}$, $\langle G_{\bar{\gamma}} \rangle = 0$, $\langle S \rangle < \Lambda_{hid}$. (ii) $\langle G_{\bar{\gamma}} \rangle \sim 1$, $\langle z \rangle = m_{3/2}$. (iii) $\langle S \rangle = \Lambda_{hid}$ (without mixing with $z$ field).

The mass $m_S$ can originate from mixing of $S$ with fields $\Phi$ having a Planck scale mass, provided the $S\Phi$-mixing is the order $m_{3/2}$. The latter scale can appear from the Kähler potential 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. In this case for the mass of $S$ we have usual see-saw formula: $m_S = m_{3/2}'/M_P$.

Another possibility is when the superfield $S$ charged under $U(1)'$ gets a VEV of the order $m_{3/2}$. This VEV leads 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 results in the desirable mass of $S$.

VI. Mixing. The interaction $\bar{L}H_2S$ with the coupling constant $\lambda \sim m_{3/2}/M_P$ can be generated either through non-renormalizable interactions in the superpotential or from the K"ahler potential in a way similar to appearance of the $\mu$-term for the Higgs doublets \[18\], \[19\]. Let us consider the following coupling:

$$K = -\frac{1}{M_P} P(S)LH_2 + h.c...,$$

(12)

where $P$ is some function of moduli $S$. For $P(S) = S$ we find immediately the desired mixing mass: $m_{\nu_S} \sim m_{3/2}K^{SL} \sim m_{3/2}(H_2)/M_P$. In general,

$$m_{\nu_S} = (K_{S3}K_{LL})^{-1/2} (H_2(e^{\bar{\mu}} + m_{3/2}P^S - F^2\partial_{\bar{\nu}}P)).$$

(13)

Here $\bar{\mu}_i$ is the effective mass induced by non-perturbative effects or higher derivative terms \[19\], and $F^2$ appears if $P = P(S, z, \bar{z})$. $F^2$ and $\bar{\mu}_i$ are functions of $S$.

A generic feature of the neutrino-modulino mixing is the R-parity violation by dimension three operators. Indeed, the same term of the K"ahler potential (12) generates the lepton violating coupling $m_{LH}LH_2$. The mass $m_{LH}$ has a general expression \[12\]. Performing its expansion in series of $S/M_P$ we can write:

$$m_{3/2} \left[ \alpha' + \beta' \frac{\langle S \rangle}{M_P} \right] LH_2.$$  

(14)

The terms in (14) can be separately small, and $m_{LH}$ can range from zero to $m_{3/2}$, although for some moduli (like $T$) one expects $m_{LH} \sim O(m_{3/2})$.

Explicit R-parity violation by dimension two operators has interesting phenomenological consequences \[20\]. It generates lepton number violating Yukawa couplings $\lambda$ and $\lambda'$. Even in the case of universal soft symmetry breaking terms at high scales it leads due to renormalization group effect to nonzero VEV for sneutrino $\langle L \rangle \neq 0$. This in turn results in mixing of neutrino and neutralinos and generation of masses for light neutrinos \[20\].

The present scenario is based on gravity mediated SUSY breaking. Its signature is the neutrino transitions into singlet state and the R-parity violating effects.

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