Physics Letters B 695 (2011) 181–184

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Schizophrenic neutrinos and ν-less double beta decay

Rouzbeh Allahverdi [a], Bhaskar Dutta [b], Rabindra N. Mohapatra [c,*]

[a] Department of Physics and Astronomy, University of New Mexico, Albuquerque, NM 87131, USA
[b] Department of Physics and Astronomy, Texas A&M University, College Station, TX 77843-4242, USA
[c] Maryland Center for Fundamental Physics, Department of Physics, University of Maryland, College Park, MD 20742, USA

ARTICLE INFO

Article history:
Received 17 October 2010
Accepted 29 October 2010
Available online 6 November 2010
Editor: M. Cvetič

Keywords:
Neutrinos
Schizophrenic
Dirac–Majorana mixture
Double beta decay

ABSTRACT

We point out a new possibility for neutrinos where all neutrino flavors can be part Dirac and part Majorana. Our primary motivation comes from an attempt to use supersymmetric seesaw models to tie inflation, baryon asymmetry of the Universe and dark matter to the neutrino sector. The idea however could stand on its own, with or without supersymmetry. We present a realization of this possibility within an S_3 family symmetry for neutrino masses, where we obtain tri-bi-maximal mixing for neutrinos to the leading order. The model predicts that for the case of inverted hierarchy, the lower limit on the neutrino mass measured in neutrinoless double beta decay experiments is about a factor of two larger than the usual Majorana case.

1. Introduction

Experiments over the past two decades have conclusively established that neutrinos have mass. The true nature of the neutrino mass however is unknown since available observations based on flavor oscillations do not tell us whether it is its own anti-particle (Majorana type) or not (Dirac type). Unlike the quarks and charged leptons, both these possibilities are allowed for the neutrinos since they are electrically neutral. Numerous neutrinoless double beta decay experiments are under way or in preparation to settle this question.

An intermediate possibility that has been discussed in literature is known as the pseudo-Dirac case [1] where one includes a very tiny amount of the Majorana mass for each neutrino flavor which has dominantly Dirac type mass. The Majorana entry in this case must be very tiny ($\leqslant 10^{-10}$ eV) in order to be consistent with current solar neutrino observations [2]. In this note we point out a new class of possibilities where each neutrino flavor is a large admixture of both Dirac and Majorana masses under certain circumstances. We point out the experimental implications of this possibility as well as its possible theoretical origin.

While discussing the Dirac versus Majorana nature of neutrinos, it is usual to frame the discussion in terms of the neutrino flavor eigenstates that are emitted in beta decay. When the neutrinos travel in free space, they do so as a mass eigenstates, which are linear superpositions of the flavor eigenstates. This phenomenon is responsible for neutrino oscillation phenomena. In this note we point out that the possibility of one of the neutrino mass eigenstates having a Dirac mass at the tree level with the others having Majorana type mass, appears consistent with all current observations. Since in this case, each flavor eigenstate is a large admixture of both Dirac and Majorana masses, we call this "schizophrenic neutrino" alternative. We further notice that this model has distinct predictions for neutrinoless double beta decay searches compared to the case where the neutrinos are pure Majorana type.

On the theoretical side, the mass eigenstate having the Dirac mass must have a Dirac Yukawa coupling which extremely tiny ($\sim 10^{-12}$) whereas the other masses could arise from high mass scale physics as in seesaw models [3]. The Dirac type mass eigenstate would pair up with a right-handed (RH) neutrino (ν_s) to form the Dirac mass. A priori, we do not know which of the three mass eigenstates has the Dirac nature. In this Letter, we consider a specific model where we want to get tri-bi-maximal pattern [5] for the PMNS matrix. We therefore determine the eigenstates to be representations of an S_3 symmetry group. The model then picks the "middle" eigenstate ν_2 (the one that determines solar neutrino oscillations) as Dirac type whereas the other two are Majorana.

While one would like to understand the small Dirac Yukawa coupling as a consequence of some high scale theory, it is comforting to know that it stable under radiative corrections due to chiral symmetry (or in this case under the symmetry $\nu_s \to -\nu_s$). There may be other motivations for the existence of such tiny Yukawa couplings. One such motivation in supersymmetric versions of such models comes from attempts to use the RH sneutrino to drive in-

* Corresponding author.
E-mail address: rmohapat@physics.umd.edu (R.N. Mohapatra).

0370-2693 © 2010 Elsevier B.V. Open access under CC BY license.
doi:10.1016/j.physletb.2010.11.006

flation [4]. In such a scenario, small Dirac coupling is essential to make the inflation predictions for density fluctuations consistent with observations. However the hypothesis schizophrenic neutrinos could be considered independently of this. As indicated earlier, a testable prediction of this model is that in the case of inverted hierarchy, the lower bound on $m_{\beta\beta}$ measured in neutrinoless double beta decay searches is roughly twice that of usual inverted hierarchy models in literature. This model will therefore be easier to rule out by the current generation of $\beta\beta_{0\nu}$ experiments if long base line oscillation searches indicate inverted neutrino mass ordering.

We hasten to point out that this kind of pattern for neutrino masses is not protected by a symmetry. As a result, when loop corrections are taken into account, tiny corrections of order $g^2 m_\tau^2/(32\sqrt{6}\pi^2 M_W^2) \sim 4\times 10^{-7}$ appear giving the Dirac eigenstate a pseudo-Dirac mass splitting of order 10^{-14} eV. These corrections have no impact on our prediction for $\beta\beta_{0\nu}$ decay.

2. Motivation from cosmology

In this section, we review the cosmological motivation for small neutrino Dirac coupling in a supersymmetric seesaw model. We consider a supersymmetric extension of MSSM based on the gauge group $SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ which requires that there be three RH neutrinos N ($\equiv \nu^c$ in SUSY language) to cancel anomalies (with the eventual possibility of SO(10) grand unification). A combination of superpartners $\phi \equiv (\tilde{N} + H_u + \tilde{L}/\sqrt{3})$ in this theory is a D-flat directions under the whole gauge symmetry, and is also F-flat when neutrino Yukawa couplings h are turned off. As shown in [4], this flat direction can act as the inflaton. The neutrino Yukawa couplings in combination with the soft mass and A-term for ϕ leads to a potential for the radial component of ϕ (denoted as φ) of the form [4]

$$V(\varphi) = \frac{m_\phi^2}{2}\varphi^2 + \frac{h^2}{12}\varphi^4 - \frac{Ah}{6\sqrt{3}}\varphi^3 \tag{1}$$

where $m^2 = (m_{\tilde{N}_0}^2 + m_{H_u}^2 + m_{\tilde{L}}^2)/3$ can lead to inflection point inflation [4] and the amplitude of observationally relevant density perturbations (as measured by COBE and WMAP) matches the observed value $\delta_H \sim 1.9\times 10^{-5}$ for weak scale masses $m_\phi \sim \mathcal{O}(100)$ GeV, provided that $h \sim 10^{-12}$ (for details, see [4]). The neutrino mass is intimately connected to h. For instance, if neutrinos are Dirac fermions, we have $m_\nu = h\langle H_u\rangle$, where $\langle H_u\rangle = (174\ \text{GeV})\sin\beta$ and $\tan\beta$ is the ratio of vacuum expectation value (VEV) of Higgs fields of the minimal supersymmetric standard model (MSSM). Then $h \sim 10^{-12}$ would give rise to $m_\nu \sim \mathcal{O}(0.1)$ eV, which is precisely in the range of interest for neutrino oscillations. We could take this as a hidden message from cosmology that at least one of the neutrinos can be dominantly of Dirac nature, and study its implications for neutrino masses and mixings.

It is important to emphasize that the inflation model constrains the coupling of only one of the RH neutrinos whose superpartner is responsible for inflation. That RH neutrino could be the Dirac partner of one of the light neutrino combinations making it a Dirac neutrino. The other two RH neutrinos have unconstrained Yukawa couplings that take natural values ($\sim 10^{-5}$–0.1), and hence their mass must be heavy. Note however that the heavy RH neutrinos must not mix with the RH neutrino whose superpartner is part of the inflaton so as not to spoil the picture of inflation mentioned above. The simplest possibility for neutrino masses in this case would therefore appear to be that one linear combination of the flavor eigenstates is a Dirac fermion whereas the other two will be Majorana and get their mass via the seesaw mechanism. Below we suggest this as new picture for neutrino masses.

3. An S_3 model for schizophrenic neutrinos

One of the challenges in neutrino mass physics is to understand the observed near tri-bi-maximal mixing pattern among different flavors. Discrete symmetries have been discussed extensively as a way to address this issue [6] and the group S_3 is one of the symmetries that appears promising in this context and we use it in our discussion. The basic assumptions of our neutrino model can therefore be summarized as follows:

- The extended gauge group responsible for neutrino masses consists of a local B–L symmetry [7], which requires the existence of three RH neutrinos for anomaly cancellation.
- One of the RH neutrinos (whose superpartner is the inflaton field) couples to a linear combination of all neutrino flavors with a Yukawa coupling of order 10^{-12} so that it gets a Dirac mass without any need for seesaw, whereas the remaining orthogonal combinations get their masses from the conventional seesaw mechanism.

The first assumption is quite well motivated and has been widely discussed in literature. It also naturally incorporates $\tilde{N}$ along with H_u and $\tilde{L}$ into a single D-flat direction that can drive inflection point inflation. The second assumptions is motivated by the cosmological scenario discussed above.

As already mentioned, our neutrino model is based on the idea that only one of the neutrino flavor combinations corresponding to a mass eigenstate has a small Yukawa coupling to one RH neutrino whereas the other two combinations get their masses from the seesaw mechanism. If we take the tri-bi-maximal matrix as the leading order PMNS matrix, then one might start thinking of a discrete symmetry group which has one singlet and one doublet as part of its irreducible representations and the singlet one being the Dirac neutrino whereas the doublet combinations becoming Majorana. One such example used in literature is the S_3 group [8] which proves convenient for our discussion.

We assume the S_3 to permute the three families of leptons (L_e, L_μ, L_τ) among themselves. Of course, it is well known that this is a reducible representation of S_3 group but the following linear combinations of these fields transform as singlet and two dimensional representations of S_3:

$$\begin{aligned} L_2 &= \frac{1}{\sqrt{3}}(L_e + L_\mu + L_\tau) \quad \text{(Singlet)}, \\ L_1 &= \frac{1}{\sqrt{6}}(2L_e - L_\mu - L_\tau) \quad \text{(Doublet)}, \\ L_3 &= \frac{1}{\sqrt{2}}(L_\mu - L_\tau) \quad \text{(Doublet)}. \end{aligned} \tag{2}$$

We assume that muon type RH neutrino is the S_3 singlet whereas (N_e, N_τ) form a doublet. The masses of these doublet RH neutrinos can be chosen different by appropriate choice of symmetry breaking (see below). The effective lepton Yukawa coupling after integrating out N_e and N_τ can then be written as

$$\mathcal{L}_\nu = hL_2 H_u N_\mu + \frac{h_1^2}{M_{N_e}}(L_1 H_u)^2 + \frac{h_3^2}{M_{N_\tau}}(L_3 H_u)^2 + \text{h.c.} \tag{3}$$

After the electroweak symmetry breaking, this gives rise to one Dirac neutrino corresponding to the mass eigenstate ν_2 and two Majorana eigenstates ν_1, ν_3 and clearly leads to tri-bi-maximal form for the PMNS matrix provided the charged lepton mass matrix is diagonal.

The effective Lagrangian in (3) could for instance arise in an $SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ theory supplemented by a global discrete symmetry $S_3 \times D$, (where D is a product of extra Z_n symmetries) if we take an S_3 doublet Higgs fields (Δ_1^c, Δ_2^c) and a singlet field Δ_0^c (all with B–L charge -2 and I_{3R} charge $+1$) with VEVs $\langle\Delta_1\rangle = 0$ and others with non-zero VEV. This will generate different Majorana masses M_{N_e} and M_{N_τ} for the S_3 doublet RH neutrinos.

In this model, inflation occurs along the flat direction corresponding to the first superpotential term in Eq. (3). The coupling between N_μ and (N_e, N_τ) can be forbidden by e.g., a Z_8 symmetry contained in D under which $N_\mu \to -iN_\mu$ and a gauge singlet field X with $X \to e^{i\pi/4}X$ with all other fields invariant. The Dirac coupling of N_μ can be obtained from a higher dimensional coupling $(\lambda L_2 H_u N_\mu X^2/M_{Pl}^2)$, where $\langle X\rangle \sim 10^{12}$ GeV or so. At the inflection point VEV ($\sim 10^{12}$ GeV [4]), this interaction then yields the effective Dirac coupling of N_μ in Eq. (3). An additional RH neutrino mixing term $(N_\mu N_{e,\tau}\Delta_{1,2}^c X^2/M_{Pl}^2)$ is allowed by the Z_8 symmetry, but has negligible contribution to masses and mixings and can be ignored. We need to add the fields $(\overline{\Delta}_1^c, \overline{\Delta}_2^c)$ and a singlet field $\overline{\Delta}_0^c$ to preserve supersymmetry below the B–L scale as well as to cancel anomalies.

Turning to the charged lepton mass matrix, neutrino oscillation data require that it be nearly diagonal. We employ the technique used in the second reference in [8]. We add three gauge singlet superfields $(\sigma_e, \sigma_\mu, \sigma_\tau)$ and three extra Z_n symmetries, i.e. $Z_{n,e} \times Z_{n,\mu} \times Z_{n,\tau}$, with RH lepton fields e^c, μ^c, τ^c transforming as $\omega_{e,\mu,\tau}^p$ and singlet fields transforming as $\omega_{e,\mu,\tau}^{-p}$. Both sets of three fields also transform under S_3 exactly like the lepton doublet fields. We can write down the corresponding Yukawa superpotential as

$$\mathcal{W}_{l,Y} = \frac{1}{M} h_e H_d \left(L_e \sigma_e e^c + L_\mu \sigma_\mu \mu^c + L_\tau \sigma_\tau \tau^c\right). \qquad (4)$$

There can be another term where the L_e, L_μ, L_τ are permuted among themselves. This will contribute to the off-diagonal elements of the charged lepton mass matrix after symmetry breaking. We set this coupling to zero. Now by adjusting the VEVs of the singlet fields, we can get diagonal mass matrix for the charged leptons. On the other hand, if the small contributions to the mass matrix coming from the permuted terms are kept, there will corrections to the tri-bi-maximal form e.g. it will lead to non-zero θ_{13}.

4. Implications for neutrinoless double beta decay

This neutrino mass model has an interesting implication for neutrinoless double beta decay. Recall that in the conventional all Majorana neutrino case, the light neutrino contribution to $\beta\beta_{0\nu}$ decay takes the form $m_{\beta\beta} = \sum_i U_{ei}^2 m_i$ $(i = 1, 2, 3)$, where U_{ei} are entries of the PMNS matrix. In the inverted hierarchy scenario, this leads to the following lower bound for the conventional three Majorana neutrino case [9]

$$|m_{\beta\beta}| \simeq \left|\left(\cos^2\theta_\odot + e^{i\alpha}\sin^2\theta_\odot\right)m_{\rm atm}\right| \geqslant \frac{m_{\rm atm}}{3} \approx 17\ \text{meV} \quad \text{(Conventional)}. \qquad (5)$$

In our model, however, the second neutrino mass eigenstate is a Dirac type state and therefore has no contribution to $\beta\beta_{0\nu}$ decay. This leads to the following lower bound for inverted case:

$$|m_{\beta\beta}| \simeq \cos^2\theta_\odot m_{\rm atm} \geqslant \frac{2m_{\rm atm}}{3} \approx 34\ \text{meV} \quad \text{(Dual)}, \qquad (6)$$

which is roughly twice the value of the conventional case (5). This makes it easier to rule out our model in the current generation of neutrinoless double beta decay searches, provided we have independent evidence, e.g. from long base line neutrino experiments for inverted mass hierarchy.

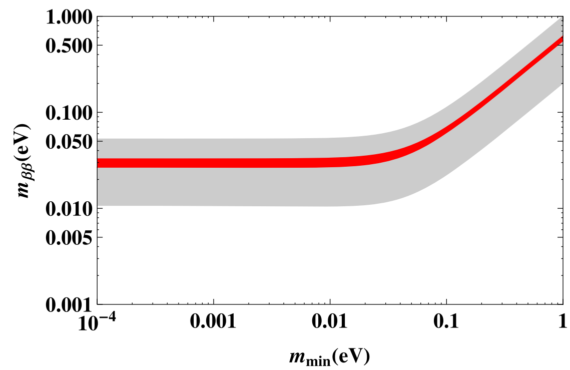

Fig. 1. We plot $|m_{\beta\beta}|$ as a function of the lightest neutrino mass for the case of inverted hierarchy. The dark (red in the web version) band shows the prediction from the two Majorana and one Dirac neutrino scenario and the gray shaded region shows the conventional three Majorana neutrino case.

In the normal hierarchy scenario, the corresponding formula becomes $m_{\beta\beta} \simeq (m_1\cos^2\theta_\odot + e^{i\alpha'}\sin^2\theta_{13}m_{\rm atm})$, which is different from the conventional three Majorana case. The precise value in this case depends on the unknown Majorana mass of ν_1 as well as the value of θ_{13}.

In Fig. 1, we plot $|m_{\beta\beta}|$ as a function of the lightest neutrino mass $\min(m_j) = m_{min}$ (which sets the absolute neutrino mass scale in the case of inverted hierarchy). The dark (red in the web version) band shows the prediction of our scenario and the gray shaded region shows the usual three Majorana neutrino scenario in the case of inverted hierarchy. The masses and mixing angles used for the figure are as follows [10]: $\Delta m_\odot^2 = (7.59 \pm 0.20_{-0.69}^{+0.61}) \times 10^{-5}$ eV^2, $\Delta m_{\rm atm}^2 = (2.46 \pm 0.12 \pm 0.37) \times 10^{-3}$ eV^2, $\theta_{13} < 12.5°$, $\theta_\odot = 34.4° \pm 1_{-2.9}^{+3.2}$ and $\theta_{\rm atm} = 42.8°{}_{-2.9-7.3}^{+4.7+10.7}$. We can see that in the case of inverted hierarchy (corresponding to $m_{min} < 0.05$ eV) the lower limit on $|m_{\beta\beta}|$ measured in neutrinoless double beta decay experiments is about a factor of two larger than the conventional case.

5. Comments

We now make some comments on the model described here.

(i) Since the Dirac nature of the second neutrino mass eigenstate is not protected by any symmetry, radiative corrections will induce Majorana component to its mass. The self-energy corrections to ν_i masses due to $W^+\ell^-$ intermediate states will lead to kinetic mixings between the different mass eigenstates that depends on the charged lepton masses: $\epsilon_{ij} \sim (\sum_k U_{ik}U_{kj}^* g^2 m_{\ell_k}^2 \times 32\sqrt{6}\pi^2 M_W^2)$. This mixing is of order 4×10^{-7}. When the kinetic energy term in the Lagrangian is diagonalized, this leads to mixing terms (for the normal hierarchy case), e.g. $m_\delta \nu_3 \nu_2 + \cdots$, where $m_\delta \sim m_{\nu_3}\epsilon_{23} + \cdots$. This introduces a Majorana mass term $\delta m_2 \nu_2 \nu_2$ with $\delta m_2 \sim 10^{-14}$ eV. It, being very small, keeps the Dirac nature of ν_2 to very high precision. This is also consistent with a bound $\sim 10^{-9}$ eV on this parameter from solar neutrino observations [2]. The same result holds for the inverted hierarchy case with ν_1 and ν_3 interchanged. In the SUSY version, quantum corrections that mix the slepton states introduce a Majorana component for the Dirac neutrino. At one loop this effect results in $\delta m_{ij}^2 \sim [(Y_\ell^\dagger Y_\ell)_{ij}/16\pi^2]m_0^2\ln(M^2/M_Z^2)$, which is of the same order as that mentioned before.

(ii) We wish to emphasize that our scenario is different from the usual pseudo-Dirac scenario [1] discussed in the literature. Our light neutrino mass matrix is a 4×4 matrix such that one of its eigenstates forms a Dirac pair with the sterile neutrino and other two eigenstates are Majorana. The Dirac eigenstate gives rise to a pseudo-Dirac pair only at the one-loop level.

(iii) The masses of the two heavy RH neutrinos depend on the scale at which B–L is broken, and can be as low as $\mathcal{O}(1)$ TeV. The decay of heavy Majorana neutrinos, and their SUSY partners, can generate baryon asymmetry of the Universe. If M_{N_e}, M_{N_τ} are of order TeV, resonant leptogenesis will be a relevant solution. However, in the S_3 symmetric model, this does not work since it will require the first and the third neutrino masses be almost equal. The oscillation data will be hard to fit with this pattern. However, soft leptogenesis [11] can work well in the model for a wide range of Majorana masses.

(iv) In this model, either the MSSM neutralino or the superpartner of the RH component of the Dirac neutrino can play the role of dark matter. The latter is naturally the lightest of the RH sneutrinos since its mass receives contribution from SUSY breaking alone. If the B–L is broken around TeV, it can obtain the correct relic density via thermal freeze out [4]. This also makes the corresponding Z' accessible at the LHC. On the other hand, for a high scale B–L the usual MSSM neutralino is a good dark matter candidate. The role of B–L in this case is to provide the R-parity symmetry naturally.

(v) The impact of the RH neutrino, which is responsible for the Dirac mass, on Big Bang Nucleosynthesis also depends on the scale at which B–L is broken. For example, for $M_{Z'} \sim 10$ TeV, the RH neutrinos decouple at $T_D \sim 1$ GeV, while for $M_{Z'} \sim 1$ TeV we have $T_D \sim 100$ MeV. In the latter case this amounts to $N_\nu^{\mathrm{eff}} \simeq 4$, whereas $N_\nu^{\mathrm{eff}} \simeq 3.1$ in the former case. For a high scale B–L the RH neutrinos decouple much earlier, and hence $N_\nu^{\mathrm{eff}} \approx 3$.

In summary, motivated by cosmology, we have pointed out a new picture for neutrino masses with the novel feature that one of the mass eigenstates is a Dirac fermion (at the tree-level) whereas the other two are Majorana type. We presented an S_3 realization of this idea that leads to tri-bi-maximal mixing for leptons in the leading order. This model can be ruled out by the current generation of neutrinoless double beta decay searches if inverted mass hierarchy is indicated by long base line neutrino oscillation experiments and neutrinoless double beta decay searches give $|m_{\beta\beta}| \lesssim 34$ meV.

Acknowledgements

We wish to thank A. de Gouvea, S.T. Petcov and L. Wolfenstein for useful discussions and helpful comments. B.D. is supported by the DOE grant DE-FG02-95ER40917. R.N.M. is supported by the NSF under grant PHY-PHY0968854.

References

[1] L. Wolfenstein, Nucl. Phys. B 186 (1981) 147;
S.T. Petcov, Phys. Lett. B 110 (1982) 245;
M. Doi, M. Kenmoku, T. Kotani, H. Nishiura, E. Takasugi, Prog. Theor. Phys. 70 (1983) 1331.
[2] A. de Gouvea, W.C. Huang, J. Jenkins, Phys. Rev. D 80 (2009) 073007.
[3] P. Minkowski, Phys. Lett. B 67 (1977) 421;
T. Yanagida, in: Workshop on Unified Theories, KEK Report 79-18, 1979, p. 95;
M. Gell-Mann, P. Ramond, R. Slansky, Supergravity, North-Holland, Amsterdam, 1979, p. 315;
S.L. Glashow, 1979 Cargese Summer Institute on Quarks and Leptons, Plenum, New York, 1980, p. 687;
R.N. Mohapatra, G. Senjanovic, Phys. Rev. Lett. 44 (1980) 912.
[4] R. Allahverdi, A. Kusenko, A. Mazumdar, JCAP 0707 (2007) 018;
R. Allahverdi, B. Dutta, A. Mazumdar, Phys. Rev. Lett. 99 (2007) 261301.
[5] P.F. Harrison, D.H. Perkins, W.G. Scott, Phys. Lett. B 530 (2002) 167;
Z.Z. Xing, Phys. Lett. B 533 (2002) 85;
X.G. He, A. Zee, Phys. Lett. B 645 (2007) 427;
Similar mixing patterns were discussed also in: L. Wolfenstein, Phys. Rev. D 18 (1978) 958;
Y. Yamanaka, H. Sugawara, S. Pakvasa, Phys. Rev. D 25 (1982) 1895;
Y. Yamanaka, H. Sugawara, S. Pakvasa, Phys. Rev. D 29 (1984) 2135, Erratum.
[6] For reviews, see: H. Ishimori, T. Kobayashi, H. Ohki, H. Okada, Y. Shimizu, M. Tanimoto, arXiv:1003.3552 [hep-th];
G. Altarelli, F. Feruglio, arXiv:1002.0211 [hep-ph].
[7] R.E. Marshak, R.N. Mohapatra, Phys. Lett. B 91 (1980) 222.
[8] For an incomplete list of references that use S_3 symmetry to understand neutrino mixings, see: F. Caravaglios, S. Morisi, arXiv:hep-ph/0503234;
R.N. Mohapatra, S. Nasri, H.B. Yu, Phys. Lett. B 639 (2006) 318;
R. Jora, J. Schechter, M. Naeem Shahid, Phys. Rev. D 80 (2009) 093007;
M. Mitra, S. Choubey, Phys. Rev. D 78 (2008) 115014;
D. Meloni, S. Morisi, E. Peinado, arXiv:1005.3482 [hep-ph].
[9] F. Feruglio, A. Strumia, F. Vissani, Nucl. Phys. B 637 (2002) 345;
F. Feruglio, A. Strumia, F. Vissani, Nucl. Phys. B 659 (2003) 359, Addendum;
S.M. Bilenky, S. Pascoli, S.T. Petcov, Phys. Rev. D 64 (2001) 053010.
[10] M.C. Gonzalez-Garcia, M. Maltoni, J. Salvado, JHEP 1004 (2010) 056.
[11] G. D'Ambrosio, G.F. Giudice, M. Raidal, Phys. Lett. B 575 (2003) 75;
Y. Grossman, T. Kashti, Y. Nir, E. Roulet, Phys. Rev. Lett. 91 (2003) 251801;
Y. Grossman, T. Kashti, Y. Nir, E. Roulet, JHEP 0411 (2004) 080.