Leptogenesis in left-right symmetric theories

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Abstract

The masses and mixing of the light left–handed neutrinos can be related to those of the heavy right–handed neutrinos in left–right symmetric theories. Properties of the light neutrinos are measured in terrestrial experiments and the CP–violating decays of their heavy counterparts produce a baryon asymmetry via the well–known leptogenesis mechanism. The left–handed Higgs triplet, present in left–right symmetric theories, modifies the usual see–saw formula. It is possible to relate the lepton asymmetry to the light neutrino parameters when the triplet and the top quark through the usual see–saw mechanism give the dominant contribution to the neutrino mass matrix. We find that in this situation the small angle MSW and vacuum solutions produce reasonable asymmetry, whereas the large angle MSW case requires extreme fine–tuning of the three phases in the mixing matrix.

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1 Introduction

The explanation of the observed ratio of baryons to photons in the universe is one of the most challenging theoretical problems. In standard cosmology the ratio is explained as a disappearance of antimatter in the early universe as proposed by Sakharov [1]. The creation of a matter–antimatter asymmetry is, in many cases, suppressed by the conservation of the B-L quantum number. Fukugita and Yanagida observed [2] that a Majorana mass term provides an attractive possibility for the creation of a lepton asymmetry when heavy Majorana neutrinos decay at an epoch in which they are out of equilibrium. The effect is further enhanced by self–energy contributions which create relatively long–lived states [3]. The asymmetry is later converted into a baryon asymmetry via sphaleron processes [4]. This Majorana neutrinos come closer to explaining the observed ratio of baryons to photons of [5]

$$Y_B \simeq (0.1...1) \cdot 10^{-10}.$$
 (1)

The explanation of the baryon asymmetry seems to demand physics beyond the Standard Model [6]. In addition, the collected evidence for massive neutrinos [7] also demands physics beyond the Standard Model. The next logical step is to check if one can relate the data of the light left-handed neutrinos with the heavy right-handed ones and to obtain the correct order of magnitude for Y_B . Several recent papers dealt with this problem [8, 9, 10, 11, 12, 13, 14, 15, 16, 17], assuming specific structures for the mass matrices and symmetries of the theory. In this article we study a left-right symmetric model where in addition to the usual Higgs doublet there are left- and right-handed Higgs triplets. The breaking of the symmetry generates vacuum expectation values v_L and v_R which in turn generate neutrino mass matrices. For a natural choice of parameters, the left-handed Higgs triplet gives the main contribution to the neutrino mass matrix. Only the top quark contribution of the Dirac mass matrix entering through the see-saw mechanism is of comparable size. The important role played by the triplet Higgs was highlighted in [18]. In this case, the light left- and heavy right-handed neutrino sector are related naturally and no further assumptions are required. At the end we find that from the three solutions to the solar neutrino problem small angle MSW and vacuum oscillations generate a baryon asymmetry of the correct order of magnitude. The large mixing angle MSW solution yields a very high Y_B .

The paper is organized as follows: In Section 2 we review the conventional see—saw mechanism and its application to leptogenesis. In Section 3 we describe how the mechanism is modified in left—right symmetric theories. The experimental status of the left—handed neutrino mass matrix is included in Section 4, which is then used to calculate the right—handed mass matrix in Section 5. These results are collected together in Section 6 and figures for the asymmetry as function of the parameters are presented. In the last Section 7 we give our conclusions.

2 Conventional See–Saw mechanism

The conventional see—saw mechanism follows from the Lagrangian of the Standard Model enlarged by the addition of a singlet right—handed neutrino N'_{Ri} for each generation. The new part of the Lagrangian is

$$-\mathcal{L}_Y = \overline{l_{iL}} \frac{\Phi}{\langle \Phi \rangle} \tilde{m}_{Dij} N'_{Rj} + \frac{1}{2} \overline{N''_{Ri}} M_{Rij} N'_{Rj} + \text{ h.c.}$$
 (2)

with l_{iL} the leptonic doublet, $\langle \Phi \rangle$ the vacuum expectation value (vev) of the conventional Higgs doublet Φ , \tilde{m}_D a Dirac mass matrix operating in generation space and M_R is the symmetrical Majorana mass matrix for the right-handed neutrinos. We can go to the physical basis by diagonalizing M_R

$$U_R^* M_R U_R^{\dagger} = \text{diag}(M_1, M_2, M_3)$$
 (3)

and defining the physical states

$$N_R = U_R N_R'. (4)$$

In the new basis the Dirac mass matrix also changes to

$$m_D = \tilde{m}_D U_R. \tag{5}$$

Thus, the Dirac Yukawa couplings are also rotated by the matrix U_R .

Interference of tree level with one–loop vertex and self–energy diagrams leads to a lepton asymmetry in the decays of the lightest Majorana, $N_1 \to \Phi l^c$ and $N_1 \to \Phi^{\dagger} l$ [3]:

$$\varepsilon = \frac{\Gamma(N_1 \to \Phi \, l^c) - \Gamma(N_1 \to \Phi^{\dagger} \, l)}{\Gamma(N_1 \to \Phi \, l^c) + \Gamma(N_1 \to \Phi^{\dagger} \, l)} = \frac{1}{8 \pi \, v^2} \frac{1}{(m_D^{\dagger} m_D)_{11}} \sum_{j=2,3} \text{Im}(m_D^{\dagger} m_D)_{1j}^2 f(M_j^2 / M_1^2) \quad (6)$$

where $v \simeq 174 \text{ GeV}$ is the weak scale and the function f is defined as

$$f(x) = \sqrt{x} \left(1 + \frac{1}{1-x} - (1+x) \ln \left(\frac{1+x}{x} \right) \right) \simeq -\frac{3}{2\sqrt{x}}.$$

The approximation holds for $x \gg 1$. There can be a resonant enhancement of the asymmetry in case of the degenerate Majorana neutrinos. Obviously, the magnitude of the asymmetry is of great interest since it introduces a B-L violation in the theory.

As already mentioned, the interaction in Eq. (2) leads to the famous see—saw prediction for the light neutrino mass matrix [19]

$$m_{\nu} = -\tilde{m}_D M_R^{-1} \tilde{m}_D^T = -m_D \operatorname{diag}(M_1^{-1}, M_2^{-1}, M_3^{-1}) m_D^T.$$
 (7)

Note that \tilde{m}_D in Eq. (2) can always be written as $\tilde{m}_D = V_L \operatorname{diag}(m_{1D}, m_{2D}, m_{3D}) V_R^{\dagger}$. It then follows from Eqs. (5,6) that the asymmetry ε depends upon the right-handed mixing

matrices V_R and U_R rather than the experimentally accessible left-handed ones. This has lead to the conviction that the lepton asymmetry is independent of the low energy parameters [20]. However, the theoretical input of the left-right symmetry allows us to relate the right-handed mixing to the left-handed one and connects the baryon asymmetry to the parameters of the left-handed neutrinos.

If m_{ν} is given by Eq. (7) then knowing the neutrino masses and mixing angles from oscillation experiments does not help in determining m_D because the right hand side in Eq. (7) is quadratic in m_D . Given a specific model for \tilde{m}_D and/or M_R , one can always invert Eq. (7) to obtain the asymmetry ε as was done in e.g. [8, 9, 10, 11, 12, 13, 14, 15, 16, 17]. The left-right symmetry provides instead a more natural framework to obtain ε . In this case the unitary matrices diagonalizing m_{ν} and M_R are related. Furthermore, in an interesting situation the oscillation experiments provide us with m_{ν} , which is used for the derivation of M_R , as described in the next section.

3 Left-right symmetric models

The minimal left-right symmetric model [21] implementing the see-saw mechanism requires three Higgs fields, namely: a bi-doublet and a right-handed as well as a left-handed triplet¹. The presence of the latter is necessary in order to maintain the left-right symmetry. Both the triplets acquire vevs v_L and v_R , respectively, at the minimum of the potential. Each of them generates a Majorana mass term for left- and right-handed neutrinos:

$$m_L = f v_L$$
 and $M_R = f v_R$ (8)

with f being the coupling matrix in generation space. The conventional see—saw formula (7) is then modified to [21, 23]

$$m_{\nu} = m_L - \tilde{m}_D M_R^{-1} \tilde{m}_D^T.$$
 (9)

Frequently, the first term is neglected which however might not be justified, as we will argue below. In fact, whenever m_L is the dominant contribution to m_{ν} we will have

$$f = \frac{1}{v_L} U_L^T \operatorname{diag}(m_1, m_2, m_3) U_L.$$
 (10)

Here U_L is the matrix diagonalizing the neutrino mass matrix m_{ν} :

$$U_L^* m_\nu U_L^{\dagger} = \text{diag}(m_1, m_2, m_3) \tag{11}$$

and m_i are the light neutrino masses. We must however be careful not to ignore the second term in Eq. (9) in cases when it is important. Later on, we argue that when one identifies the Dirac mass matrix with the up quark mass matrix then only the top quark gives a

¹There are various models with triplets of Higgs fields [22].

sizable contribution.

At the minimum of the potential, the left– and right–handed triplets assume their vevs and produce masses for the gauge bosons. Then, in general [21] the following relation holds:

$$v_L v_R \simeq \gamma M_W^2, \tag{12}$$

where the constant γ is a model dependent parameter of $\mathcal{O}(1)$. Substituting the results from Eqs. (8) and (12) into (9) yields

$$m_{\nu} = v_L \left(f - \tilde{m}_D \frac{f^{-1}}{\gamma M_W^2} \tilde{m}_D^T \right). \tag{13}$$

This result exhibits the strength of the left-right symmetric theory. The oscillation experiments allow one to estimate several matrix elements of f through Eqs. (9) and (10). Once we identify \tilde{m}_D with the up quark matrix and decide that the contribution from the top quark alone is important, we can determine f, whose diagonalization gives U_R , which in turn gives m_D and therefore the lepton asymmetry via Eq. (6).

In the next section we will shortly discuss the current status of the neutrino mass matrix and will then take up the task of calculating f, estimate the magnitude of $v_{L,R}$ and determine the baryon asymmetry within the situation described before.

4 Current status of m_{ν}

The experimental data on neutrino oscillations can be used to derive the neutrino mass matrix [24]. The mixing matrix U_L may be parameterized as

$$U_{L} = U_{\text{CKM}} \operatorname{diag}(1, e^{i\alpha}, e^{i(\beta+\delta)})$$

$$= \begin{pmatrix} c_{1}c_{3} & s_{1}c_{3} & s_{3}e^{-i\delta} \\ -s_{1}c_{2} - c_{1}s_{2}s_{3}e^{i\delta} & c_{1}c_{2} - s_{1}s_{2}s_{3}e^{i\delta} & s_{2}c_{3} \\ s_{1}s_{2} - c_{1}c_{2}s_{3}e^{i\delta} & -c_{1}s_{2} - s_{1}c_{2}s_{3}e^{i\delta} & c_{2}c_{3} \end{pmatrix} \operatorname{diag}(1, e^{i\alpha}, e^{i(\beta+\delta)}),$$

$$(14)$$

where $c_i = \cos \theta_i$ and $s_i = \sin \theta_i$. The "CKM-phase" δ may be probed in oscillation experiments, as long as the LMA solution is the solution to solar oscillations [25]. The other two "Majorana phases" α and β can be investigated in neutrinoless double beta decay [26, 27]. The choice of the parameterization in Eq. (14) reflects this fact since the ee element of the mass matrix $\sum_i U_{Lei}^2 m_i$ is only depending on the phases α and β . In a hierarchical scheme, to which we will limit ourselves, there is no constraint on the phases from neutrinoless double beta decay [27]. Thus, we can choose them arbitrarily. The mass eigenstates are given as

$$m_3 \simeq \sqrt{\Delta m_A^2 + m_2^2}$$
 $m_2 \simeq \sqrt{\Delta m_\odot^2 + m_1^2} \gg m_1.$
(15)

The values of θ_2 and Δm_A^2 are known to a good precision, corresponding to maximal mixing $\theta_2 \simeq \pi/4$ and $\Delta m_A^2 \simeq 3 \times 10^{-3} \text{ eV}^2$. Regarding θ_1 and Δm_{\odot}^2 three distinct areas in the parameter space are allowed, small (large) mixing, denoted SMA (LMA) and quasi-vacuum oscillations (QVO):

SMA:
$$\tan^2 \theta_1 \simeq 10^{-4} \dots 10^{-3}$$
, $\Delta m_{\odot}^2 \simeq 10^{-6} \dots 10^{-5} \,\mathrm{eV^2}$
LMA: $\tan^2 \theta_1 \simeq 0.1 \dots 4$, $\Delta m_{\odot}^2 \simeq 10^{-5} \dots 10^{-3} \,\mathrm{eV^2}$. (16)
QVO: $\tan^2 \theta_1 \simeq 0.2 \dots 4$, $\Delta m_{\odot}^2 \simeq 10^{-10} \dots 10^{-7} \,\mathrm{eV^2}$

For the last angle θ_3 there exists only a limit of about $\sin^2 \theta_3 \lesssim 0.08$. For a recent three–flavor fit to all available data see [28].

Note that we have identified the neutrino mixing matrix in Eq. (14) with the matrix U_L diagonalizing the neutrino mass matrix Eq. (11). This assumes implicitly that the charged lepton mixing is small. We shall work with this assumption in what follows.

5 Determination of f and the baryon asymmetry

As mentioned before, we argue that only the top quark gives a sizable contribution to the conventional see–saw formula $\tilde{m}_D M_R^{-1} \tilde{m}_D^T$. Identifying \tilde{m}_D with the up quark mass matrix and neglecting mixing among up quarks, the relative magnitude of both terms contributing to m_{ν} can be written as

$$\frac{|\tilde{m}_D M_R^{-1} \tilde{m}_D^T|}{|m_L|} \simeq \frac{m_q^2 / v_R}{v_L} \simeq \frac{m_q^2}{\gamma M_W^2},\tag{17}$$

where we only used Eq. (12) and assumed that the matrix elements of f and f^{-1} are of the same order of magnitude. One sees immediately that only the top quark mass makes the ratio in Eq. (17) of order one. In practically all models [29] the heaviest mass is the (33) entry of the mass matrix, which means that only the (33) element of m_{ν} has a contribution from the term $\tilde{m}_D M_R^{-1} \tilde{m}_D^T$. The matrix \tilde{m}_D may therefore be taken as

$$\tilde{m}_D \simeq \operatorname{diag}(0, 0, m_t). \tag{18}$$

There might be a common factor of order 1 for the complete matrix, but in light of the factor γ in Eq. (12) and the uncertainty in the oscillation parameters we can safely work with this form of \tilde{m}_D . Later on we will comment on the dependence of the results on this factor.

It is helpful to repeat the argument with typical numbers. The maximal scale of m_{ν} is $\sqrt{\Delta m_A^2} \simeq 0.1$ eV. Then the relations $v_L v_R = \gamma M_W^2$ and $m_{\nu} \simeq f v_L \simeq 0.1$ eV are compatible for $v_R \simeq 10^{15}$ GeV and thus $v_L \simeq 0.1$ eV as long as $f \simeq 0.1...1$. The scale

of $v_L f$ is again matched by the factor m_t^2/v_R . This means that v_R is close to the grand unification scale and v_L is of order of the neutrino masses, which one expects since m_L is the dominating contribution to m_{ν} .

We can now proceed to calculate the contribution to the Yukawa coupling matrix f in this situation. Since only the (33) element of m_{ν} has a contribution from the see–saw term we have

$$f_{ij} = \frac{(m_{\nu})_{ij}}{v_L}$$
 for all i, j except for $i = j = 3$. (19)

For the last term we adopt

$$f_{33} = \frac{(m_{\nu})_{33} + s}{v_L},\tag{20}$$

where the parameter s denotes the contribution arising from the see–saw term. The parameter is consistently determined by using Eqs. (12,13)

$$s = \left(\frac{\tilde{m}_D f^{-1} \tilde{m}_D^T}{v_R}\right)_{33} = \frac{m_t^2}{v_R} \frac{F_{33}}{\det f} = \frac{m_t^2}{v_R} \frac{F_{33}}{\tilde{F} + F_{33} f_{33}}$$
(21)

where $F_{33} = f_{11} f_{22} - f_{12}^2$ and $\tilde{F} = 2 f_{12} f_{13} f_{23} - f_{13}^2 f_{22} - f_{23}^2 f_{11}$. Using Eq. (20) we can solve for s and find

$$s \simeq \pm \frac{\sqrt{\gamma} M_W m_t}{v_R}.$$
 (22)

As expected, s is of the order of 0.1 to 0.01 eV.

With the matrix f now determined completely, we diagonalize it and calculate the baryon asymmetry in the following way. From ε the baryon asymmetry Y_B is obtained by

$$Y_B = c \kappa \frac{\varepsilon}{q^*},\tag{23}$$

where $g^* \simeq 110$ is the effective number of massless degrees of freedom at $T = M_1$. The factor c indicates the fraction of the lepton asymmetry converted into a baryon asymmetry via the sphaleron processes [4]. It depends on the group structure of the theory and is of order one. For three lepton families and one (also two) Higgs doublets it is approximately equal to -0.55. Finally, κ is a dilution factor due to lepton–number violating wash–out processes. It can be obtained by integrating the Boltzmann equations and depends strongly on

$$K \equiv \frac{\Gamma_1}{H(T=M_1)} = \frac{(m_D^{\dagger} m_D)_{11} M_1}{8 \pi v^2} \frac{M_{\text{Pl}}}{1.66 \sqrt{g^*} M_1^2}, \tag{24}$$

where Γ_1 is the width of the lightest Majorana neutrino and $H(T=M_1)$ the Hubble constant at the temperature of the decay. $M_{\rm Pl}$ is the Planck mass. A convenient parame-

terization is [30]

$$-\kappa \simeq \begin{cases} \sqrt{0.1 \, K} \, \exp(-4/3 \, (0.1 \, K)^{0.25}) & \text{for} \quad K \gtrsim 10^6 \\ \frac{0.3}{K \, (\ln K)^{0.6}} & \text{for} \quad 10 \lesssim K \lesssim 10^6 \\ \frac{1}{2\sqrt{K^2 + 9}} & \text{for} \quad 0 \lesssim K \lesssim 10 \end{cases}$$
 (25)

Typically, values for κ lie in the range of 10^{-3} to 0.1.

In addition to the decay of the right handed neutrino, the out-of-equilibrium decay of the Higgs triplet has also been considered as possible mechanism for generating lepton asymmetry [31]. This needs CP violation in the Higgs sector and hence an enriched Higgs sector to implement it. For example, the models in [31, 32, 33] need the presence of two left-handed Higgs triplets as opposed to one triplet considered here. If an asymmetry is produced with several higgs triplets, their subsequent decays it will tend to be erased as long as the triplets are heavier than the lightest right-handed neutrino. This mass pattern happens to be the natural possibility in the present scenario. The mass of the Higgs triplet is given by λv_R where λ is a typical quartic coupling of the Higgs potential. In contrast, the mass of the lightest right-handed neutrino is given within our approximation [18] by $M_1 \sim \alpha v_R$, where $\alpha \equiv \frac{m_2 s_1^2}{m_3}$ lies in the range $(10^{-3} - 10^{-6})$, depending upon the chosen solution for the solar neutrino problem. Hence, the triplet will be heavier than the lightest right-handed neutrino as long as the quartic coupling λ is $\mathcal{O}(1)$. The lepton asymmetry created through triplet will be washed out in this case, according to the usual damping of any preexisting asymmetry. For the above reasons the asymmetry originating from the triplet decay does not contribute to the following numerical analysis.

6 Results

The main variables are the parameters Δm_{\odot}^2 and $\tan^2\theta_1$ which specify the solar solution, as given in Eq. (16). Below, we analyze the dependence of Y_B on these parameters. It is found that the value of γ and the sign of s do not play a decisive role. Also, the value of Δm_A^2 (varied within $(3\pm5)\cdot 10^{-3}~{\rm eV}^2$) has little influence on Y_B . The same is true for changing $\tan^2\theta_2$. The asymmetry decreases (increases) with decreasing (increasing) top quark mass, though not much. For the SMA case the dependence on the phase α is not as strong as on the other two phases, whereas it is equally strong for the LMA and QVO case. The conclusions we draw now will be only changed if all these parameters conspire and take rather extreme values within their allowed ranges.

We work now with positive s from Eq. (22) and apply maximal atmospheric mixing with $\Delta m_A^2 = 3 \cdot 10^{-3} \text{ eV}^2$. The parameter γ is fixed to one and the top quark mass at 175 GeV. We shall work with $v_R = 10^{15}$ GeV from which v_L and s are obtained via Eqs. (12)

and (22). We find that K from Eq. (24) is always below 10 and thus κ lies between 0.17 and 0.05. There is practically no dependence on the lightest mass eigenstate m_1 . Fig. 1 shows the behavior of Y_B as a function of Δm_{\odot}^2 for different $\sin^2 \theta_3$. One sees that for lower masses the asymmetry decreases. In Fig. 2 we display the dependence on $\tan^2 \theta_1$. The baryon asymmetry decreases with decreasing $\tan^2 \theta_1$. This dependence is stronger than the one on Δm_{\odot}^2 . In both plots it is seen that Y_B approximately increases with increasing $\sin^2 \theta_3$.

We analyze next the three distinct solutions to the solar neutrino problem in detail. Fig. 3 shows the SMA case for different values of the parameters. All four combinations yield Y_B in the right magnitude and seem to prefer a $\sin^2\theta_3$ lower than about 10^{-3} . Fig. 4 shows again that the LMA case results in a very high asymmetry. Here, fine—tuning of the parameters, specifically the CP violating phases is required to get a Y_B within its experimental limits. It is also seen that $\tan^2\theta_1 > 1$ gives a smaller asymmetry than $\tan^2\theta_1 < 1$. The QVO case, displayed in Fig. 5, might also produce an acceptable asymmetry. Note the different choice of the phases in this plot and Fig. 1. We note that the latest SuperKamiokande data seems to favor the LMA solution [34], using however a two-flavor analysis. For a more definitive conclusion additional data has to be waited for.

All our results are based on identifying \tilde{m}_D with the up quark mass matrix and retaining only the top quark contribution. The importance of the ordinary see—saw contribution will be less in any other models in which the largest scale of \tilde{m}_D is set by a fermion mass other than the top quark, i.e. the bottom quark mass. In the extreme case of completely neglecting the ordinary see—saw contribution, one will have $U_R = U_L$ and $M_{Ri} = \frac{v_R}{v_L} m_i$. In this case, the lepton asymmetry will be completely controlled by the left—handed neutrino masses and mixings as well as the ratio $\frac{v_R}{v_L}$. Thus, unlike in the present case, the results will be sensitive to the value of m_1 which has to be less than or similar to the solar scale. We have checked that the required asymmetry can be generated in this case with a proper choice of m_1 .

It is instructive to see the dependence of some other models on the solar solution. Models based on SO(10) were used e.g. in [14] where it was found that only the QVO solution gives acceptable baryon asymmetry. A slightly different analysis in [15] finds that also the SMA case gives an acceptable asymmetry. This solution has also been favored in the models presented in [11, 12, 17], which all use quite different symmetries. The LMA solution, which we disfavor, has been shown to be the only solution producing an acceptable Y_B in [16], using SU(5) inspired mass matrices. We stress again that the main difference to the present paper lies in the fact that the left–handed Higgs triplet plays a dominant role in producing the light neutrino mass matrix. Once one solution for the solar oscillation is established, more definite statements about the symmetry relations can be made, which is an important by–product of the analysis of relations between leptogenesis and neutrino oscillations.

7 Conclusion

Using very general properties of left–right symmetric theories we connected the light left–handed neutrino sector as measured in neutrino oscillations with heavy right–handed neutrinos, whose decay is responsible for the baryon asymmetry of the universe. Identifying the Dirac mass matrix with the up quark mass matrix we found that only the top quark has a significant contribution to the neutrino mass matrix. The main contribution to m_{ν} comes from the left–handed triplet, which is neglected in most papers dealing with this subject. In our scenario, the SMA and QVO case yield in reasonable asymmetry, whereas the LMA solution produces an asymmetry which is too high.

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References

- [1] A. D. Sakharov, JETP Lett. 5, 24 (1967).
- [2] M. Fukugita, T. Yanagida, Phys. Lett. B 174, 45 (1986).
- [3] M. Flanz, E. A. Paschos, and U. Sarkar, Phys. Lett. B 345, 248 (1995); E. Roulet, L. Covi, and F. Vissani, Phys. Lett. B 384, 169 (1996); M. Flanz et al., Phys. Lett. B 389, 693 (1996); A. Pilaftsis, Phys. Rev. D 56, 5431 (1997); W. Buchmüller and M. Plümacher, Phys. Lett. B 431, 354 (1998).
- [4] V. A. Kuzmin, V. A. Rubakov, and M. E. Shaposhnikov, Phys. Lett. **B 155**, 36 (1985).
- [5] K. A. Olive, G. Steigman, and T. P. Walker, Phys. Rep. **333**, 389 (2000).
- [6] G. R. Farrar, M. E. Shaposhnikov, Phys. Rev. **D** 50, 774 (1994).
- [7] K. Zuber, Phys. Rep. 305, 295 (1998); S. M. Bilenky, C. Giunti, and W. Grimus, Prog. Nucl. Part. Phys. 43, 1 (1999); J. Ellis, Talk given at the 19th International Conference on Neutrino Physics and Astrophysics Neutrino 2000, Sudbury, Ontario, Canada, 16-21 Jun 2000, Nucl. Phys. B (Proc. Suppl.) 91, 503 (2000); P. Langacker, Summary talk for Europhysics Neutrino Oscillation Workshop (NOW 2000), Conca Specchiulla, Otranto, Lecce, Italy, 9-16 Sep 2000, hep-ph/0101244.

- [8] J. Ellis, S. Lola, and D. V. Nanopoulos, Phys. Lett. **B** 452, 87 (1999).
- [9] G. Lazarides, N. D. Vlachos, Phys. Lett. **B 459**, 482 (1999).
- [10] M. S. Berger and B. Brahmachari, Phys. Rev. **D** 60, 073009 (1999); M. S. Berger, Phys. Rev. **D** 62, 013007 (2000); W. Buchmüller and M. Plümacher, hep-ph/0007176.
- [11] K. Kang, S. K. Kang, and U. Sarkar, Phys. Lett. B 486, 391 (2000).
- [12] H. Goldberg, Phys. Lett. **B** 474, 389 (2000).
- [13] R. Jeannerot, S. Khalil, and G. Lazarides, Phys. Lett. **B 506**, 344 (2001).
- [14] E. Nezri, J. Orloff, hep-ph/0004227.
- [15] D. Falcone, F. Tramontano, Phys. Rev. **D** 63, 073007 (2001).
- [16] D. Falcone, F. Tramontano, hep-ph/0101151.
- [17] H. B. Nielsen, Y. Takanishi, hep-ph/0101307.
- [18] A. S. Joshipura, E. A. Paschos, hep-ph/9906498.
- [19] M. Gell-Mann and P. Ramond, and R. Slansky, in Supergravity, P. van Nieuwenhuizen & D. Z. Freedman (eds.), North Holland Publ. Co., 1979 p 315; T. Yanagida, Proc. of the Workshop on Unified Theories and the Baryon Number of the Universe, edited by O. Sawada and A. Sugamoto, KEK, Japan 1979; R. N. Mohapatra, G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980).
- [20] M. A. Luty, Phys. Rev. D 45, 455 (1992); W. Buchmüller and M. Plümacher, Phys. Lett. B 389, 73 (1996).
- [21] R. N. Mohapatra, G. Senjanovic, Phys. Rev. D 23, 165 (1981); R. N. Mohapatra and P. B. Pal, Massive neutrinos in physics and astrophysics, Singapore, Singapore: World Scientific (1991) 318 p. (World Scientific lecture notes in physics, 41).
- [22] G. Lazarides, Q. Shafi, and C. Wetterich, Nucl. Phys. B 181, 287 (1981); R. Holman,
 G. Lazarides and Q. Shafi, Phys. Rev. D 27, 995 (1983); G. B. Gelmini, M. Roncadelli,
 Phys. Lett. B 99, 411 (1981); C. Wetterich, Nucl. Phys. B 187, 343 (1981).
- [23] D. Caldwell and R. N. Mohapatra, Phys. Rev. D 48, 3259 (1993); A. S. Joshipura,
 Z. Phys. C 64, 31 (1994); B. Brahmachari and R. N. Mohapatra, Phys. Rev. D 58, 015001 (1998).
- [24] S. M. Bilenky, C. Giunti, Phys. Lett. B 444, 379 (1998); W. Rodejohann, Phys. Rev. D 62, 013011 (2000).
- [25] M. Koike, J. Sato, Phys. Rev. D 61, 073012 (2000); V. Barger et al., Phys. Rev. D 62, 073002 (2000).

- [26] T. Fukuyama *et al.*, hep-ph/0012357.
- [27] W. Rodejohann, Nucl. Phys. **B** 597, 110 (2001).
- [28] M. C. Gonzalez-Garcia et al., Phys. Rev. **D** 63, 033005 (2001).
- [29] See e.g. B. P. Desai, A. R. Vaucher, hep-ph/0007233.
- [30] E. W. Kolb, M. S. Turner, *The early universe*, Redwood City, USA: Addison-Wesley (1990), (Frontiers in physics, 69); A. Pilaftsis, Int. J. Mod. Phys. A 14, 1811 (1999);
 E. A. Paschos, M. Flanz, Phys. Rev. D 58, 113009 (1998).
- [31] E. Ma, U. Sarkar, Phys. Rev. Lett. 80, 5716 (1998); U. Sarkar, Phys. Rev. D 59, 031301 (1999), E. Ma, U. Sarkar, and T. Hambye, hep-ph/0011192.
- [32] G. Lazarides, Phys. Lett. **B** 452, 227 (1999).
- [33] E. J. Chun and S. K. Kang, hep-ph/0001296.
- [34] S. Fukuda *et al.*, hep-ex/0103033.

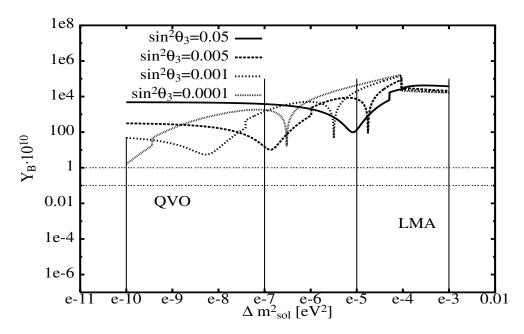


Figure 1: Behavior of the baryon asymmetry as a function of Δm_{\odot}^2 for different $\sin^2 \theta_3$. For this plot we fixed $\Delta m_A^2 = 3 \cdot 10^{-3} \text{ eV}^2$, $\theta_1 = \theta_2 = \pi/4$, $\alpha = \pi/3$, $\beta = \pi/5$ and $\delta = \pi/6$.

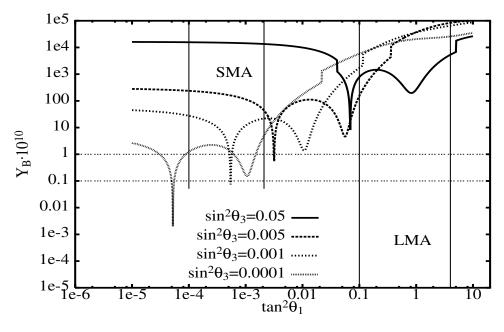


Figure 2: Behavior of the baryon asymmetry for $\Delta m_{\odot}^2 = 10^{-5} \text{ eV}^2$ as a function of $\tan^2 \theta_1$ for different $\sin^2 \theta_3$. The phases are $\alpha = \pi/3$, $\beta = \pi/4$, $\delta = \pi/6$ and the other parameters are as in the previous plot.

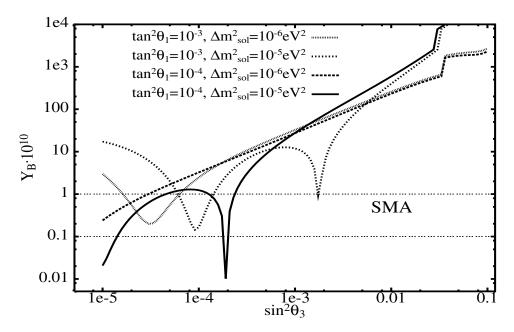


Figure 3: Behavior of the baryon asymmetry as a function of $\sin^2 \theta_3$ for different Δm_{\odot}^2 and $\tan^2 \theta_1$ for the case of the SMA solution. The other parameters are as in the previous plot.

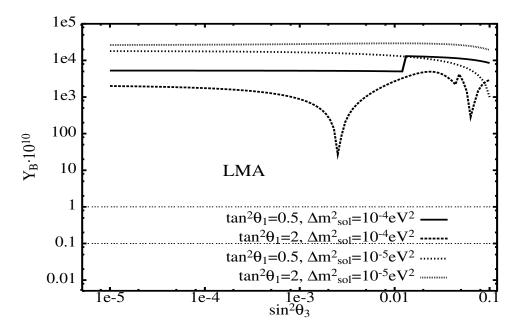


Figure 4: Behavior of the baryon asymmetry as a function of $\sin^2 \theta_3$ for different Δm_{\odot}^2 and $\tan^2 \theta_1$ for the case of the LMA solution. For this plot we fixed the atmospheric parameters as before and $\alpha = \pi/5$, $\beta = \pi/6$ and $\delta = \pi/3$.

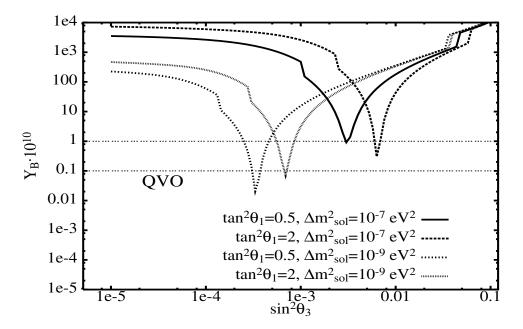


Figure 5: Behavior of the baryon asymmetry as a function of $\sin^2 \theta_3$ for different Δm_{\odot}^2 and $\tan^2 \theta_1$ for the case of the QVO solution. For this plot we fixed the atmospheric parameters as before and $\alpha = \pi/5$, $\beta = \pi/4$ and $\delta = \pi/6$.