Published by Institute of Physics Publishing for SISSA/ISAS

Received: *December 5, 2004*
Revised: *January 11, 2005*
Accepted: *February 9, 2005*

Effects of the neutrino B-term on the Higgs mass parameters and electroweak symmetry breaking

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Abstract: To embed the seesaw mechanism in the MSSM, two or three right-handed neutrino supermultiplets, N_i, have to be added to the model. In the presence of these new superfields, the soft supersymmetry breaking potential includes a lepton number violating mass term which is known as the neutrino B-term, $MB_\nu\tilde{N}\tilde{N}/2$. In this paper, we study the effects of B_ν on the Higgs mass parameters. Using the condition for the electroweak symmetry breaking, we derive an upper bound on B_ν which, for some range of parameters, is two orders of magnitude stronger than the previous bounds. We also propose a simple model in which it is natural to have large values of B_ν while the rest of the supersymmetry breaking terms are at the TeV scale or smaller.

Keywords: Higgs Physics, Supersymmetric Standard Model, Rare Decays.

Contents

1. Introduction

The Standard Model (SM) of the particle physics so far has had an enormous success. However, there are two cases that the SM fails to provide an explanation: the neutrino oscillation data and the stability of the electroweak scale.

There is a consensus among the physicists that to accommodate the solar, atmospheric and reactor data, neutrinos have to be massive. On the other hand, from the cosmological data as well as the studies of the end-point of the Tritium beta decay, we know that the masses of neutrinos are very tiny: $m_\nu/m_e < 10^{-6}$. One way to attribute a tiny but nonzero mass to neutrinos is the famous seesaw mechanism which adds heavy right-handed neutrinos to the SM.

To stabilize the electroweak symmetry breaking scale against radiative corrections, several models have been developed among which supersymmetry is arguably the most popular one. The simplest version of supersymmetry is the Minimal Supersymmetric Standard Model (MSSM).

To build a model that incorporates the both mechanisms, we can augment the MSSM to include three heavy right-handed neutrino supermultiplets, N_i. In the presence of the right-handed neutrinos the superpotential will be

$$W = Y_\ell^{ij}\epsilon_{\alpha\beta}H_d^\alpha l_{Ri}^c L_j^\beta + Y_\nu^{ij}\epsilon_{\alpha\beta}H_u^\alpha N_i L_j^\beta + \frac{1}{2}M_{ij}N_iN_j + (\mu H_u H_d + \text{H.c.})\,, \tag{1.1}$$

where L_j^β is the supermultiplet corresponding to the doublet $(\nu_{Lj},\ l_{Lj})$. Since the quark sector is irrelevant for our discussion, we have omitted the terms involving quarks in the above formula. Without loss of generality, we can rephase and rotate the fields to make the matrices Y_ℓ^{ij} and M_{ij} real and diagonal: $Y_\ell^{ij} = \text{diag}(Y_e, Y_\mu, Y_\tau)$ and $M^{ij} = \text{diag}(M_1, M_2, M_3)$. In this basis, Y_ν can have off-diagonal and complex elements. We will assume that M_1, M_2 and M_3 are much larger than m_{susy}.

For simplicity, in this paper we work in the framework of the mSUGRA for which at the GUT scale, the soft supersymmetry breaking terms are flavor blind; that is at the GUT scale

$$
\begin{aligned}
-\mathcal{L}_{\rm soft} = & \; m_0^2(\tilde{L}^\dagger_{L\alpha}\tilde{L}_{L\alpha} + \tilde{l}^\dagger_{R\alpha}\tilde{l}_{R\alpha} + \tilde{N}^\dagger_\alpha\tilde{N}_\alpha + H_d^\dagger H_d + H_u^\dagger H_u) + \frac{1}{2}m_{1/2}(\tilde{B}^\dagger\tilde{B} + \tilde{W^a}^\dagger\tilde{W^a}) + \\
& + (b_H^0 H_d H_u + h.c.) + A_\ell^{ij}\epsilon_{\alpha\beta}H_d^\alpha\tilde{l}^\dagger_{Ri}\tilde{L}^\beta_{Lj} + A_\nu^{ij}\epsilon_{\alpha\beta}H_u^\alpha\tilde{N}_i\tilde{L}^\beta_{Lj} + \\
& + \left(\frac{1}{2}B_\nu M_i\tilde{N}^i\tilde{N}^i + \text{h.c.}\right), \qquad (1.2)
\end{aligned}
$$

where $A_\ell = a_0 Y_\ell$ and $A_\nu = a_0 Y_\nu$. The generalization to a more general supersymmetry breaking scheme is straightforward. Due to the radiative corrections, at the electroweak scale the masses of different scalars; in particular, the masses of H_u and H_d ($m^2_{H_u}$ and $m^2_{H_d}$) will be different.

If $|B_\nu|$ is large ($B_\nu \gg m_{\rm susy}$), several interesting phenomena can occur. For example, as discussed in [1], a large B_ν can give rise to $\tilde{\nu}-\bar{\tilde{\nu}}$ oscillation in the linear colliders. Moreover, large B_ν can give a radiative correction to the masses of left-handed neutrinos [1]. The neutrino B-term also gives a correction to the mass matrix of left-handed sleptons which, in principle, can be Lepton Flavor Violating (LFV), inducing LFV rare decays [2]. Finally, the neutrino B-term gives a correction to the A-term of the charged leptons, A_ℓ, proportional to B_ν. As it is well-known, an imaginary A_ℓ can induce Electric Dipole Moment (EDM) for the charged leptons so this way imaginary B_ν induces EDMs for the charged leptons.

The strong upper bounds on the neutrino mass and branching ratios of LFV rare decays and the EDM of the electron give bounds on different combinations of B_ν, Y_ν and M. However, as we shall discuss, these bounds are not enough to restrict the value of B_ν, independently of the values of other parameters.

In this paper, we will show that B_ν can also give radiative corrections to the Higgs mass parameters. The corrections to $m^2_{H_u}$ and b_H turn out to be finite and of order of $B_\nu m_{\rm susy}Y_\nu^2/(16\pi^2)$. Assuming that there is no significant cancelation between the tree level values of b_H and $m^2_{H_u}$ and the loop corrections, the condition for the electroweak symmetry breaking puts a strong bound on the neutrino B-term: $B_\nu Y_\nu^2/(4\pi)^2 < m_{\rm susy}/\tan\beta$.

In section 2, we suggest a simple model which makes a large neutrino B-term naturally possible. In section 3, we study the effects of the neutrino B-term on the Higgs mass parameters and derive a bound on B_ν. In section 4, we review the effects of B_ν on different observables and discuss how the bound we find in section 3 complements our knowledge of this sector. Finally, in section 5, we summarize our conclusions.

2. Theoretical framework for a large B_ν

In this section, we first review what the "natural" size of the neutrino B-term in the context of the mSUGRA is. Then, we suggest a model in which larger values of B_ν are acceptable.

In the context of the mSUGRA, the soft supersymmetry breaking terms originate from the interaction of a chiral superfield S with the super-potential:

$$\int d^2\theta S(\theta)W(\theta)\,. \qquad (2.1)$$

The scalar and F-components of S develop vacuum expectation values $\langle S\rangle = 1 + F_S\theta^2$ and

$\langle F_S \rangle$ determines the scale of the soft supersymmetry breaking terms. Within this model we expect $B_\nu \sim a_0 \sim m_{\rm susy}$. Remember that we have parameterized the neutrino B-term as $MB_\nu \tilde{N}\tilde{N}/2$ so, in this model, we expect $\sqrt{B_\nu M} \gg m_{\rm susy}$.

In this paper, we are more interested in large values of B_ν ($B_\nu Y_\nu^2/16\pi^2 \sim m_{\rm susy}$) and we propose a model that can provide us with large values of B_ν.

Let us assume that besides S which couples to the lepton number conserving part of the superpotential, there is a spurion field, X, that carries lepton number equal to two. Then we can write the following term in the superpotential

$$\int d^2\theta \lambda_i X N_i N_i\,. \tag{2.2}$$

However, terms such as $\int d^2\theta X H_u H_d$ are forbidden by lepton number conservation. Moreover terms such as $\int X^\dagger X \Phi^\dagger \Phi d^4\theta$ in the Kähler potential are suppressed by powers of M_{pl}^{-1}. Let us assume that the self-interaction of the hidden sector is such that both the scalar- and F-components of X develop nonzero vacuum expectation values. The vacuum expectation values of the components of X break the lepton number symmetry of the model. The vacuum expectation value of the scalar component of X, $\langle \tilde{X} \rangle$, corresponds to the Majorana mass term of the right-handed neutrinos while the vacuum expectation value of the F-component, $\langle F_X \rangle$, gives the neutrino B-term. With our parameterization of the neutrino B-term,

$$B_\nu = \frac{\langle F_X \rangle}{\langle \tilde{X} \rangle}\,. \tag{2.3}$$

Both $\langle F_X \rangle$ and $\langle \tilde{X} \rangle$ can be large, giving rise to large right-handed neutrino masses and B_ν while other supersymmetry breaking terms, being given by $\langle F_S \rangle$, are at TeV scale or smaller. Notice that in this model, in the basis that the mass matrix of the right-handed neutrinos is real diagonal, the neutrino B-term is also diagonal so the parameterization that we are using for the neutrino B-term is appropriate.

Notice that this model is very similar to the singlet Majoron model [3], with the difference that here $\langle F_X \rangle$ is also nonzero.

If the X field is very heavy, it has to decouple from the theory [4]. In particular although, the term in (2.2) gives rise to a quartic right-handed sneutrino self-interaction term ($|\sum_i \lambda_i \tilde{N}_i \tilde{N}_i|^2$), based on the decoupling theorem we expect that there has to be a term (for example, coming from the mass term of S) canceling this effect. Even in the absence of such a cancelation, as far as $|B_\nu| < |M|$, right-handed sneutrinos do not develop non-zero vacuum expectation values and the scheme of the electroweak symmetry breaking will be similar to the MSSM. In the next section, we will take an agnostic approach and will not care about the origin of B_ν.

3. Effects of the neutrino B-term on the Higgs mass parameters

In this section, we study the effects of a large B_ν on the Higgs mass parameters and derive bounds on its value from the fulfillment of the electroweak symmetry breaking condition. Then we compare this bound with the limit derived by considering the effect of B_ν on the neutrino mass.

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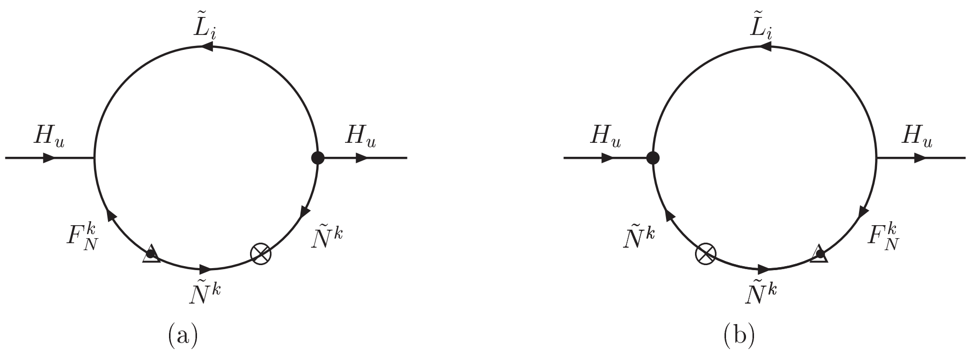


Figure 1: Diagrams contributing to $m^2_{H_u}$. F_N^k represents the auxiliary field associated with the right-handed neutrino, N_k. The A_ν vertices are marked with black circles. The neutrino B-term and M insertions are indicated by $\otimes$ and Δ, respectively.

As we shall see, the effects of the neutrino B-term on the Higgs mass parameters are finite so they cannot be derived by integrating the renormalization group equations. However, since the momentum propagating in the loop is of order of M_i, to calculate the effect precisely, we have to insert the values of the relevant parameters at the energy scale of the right-handed neutrinos, M_i, instead of those at the GUT scale. The running of the A-terms, Yukawa coupling and the right-handed (s)neutrino masses have been discussed in detail in the literature [5]. The running of the neutrino B-term is given by

$$16\pi^2\mu\frac{d\ B_\nu M}{d\ \mu} = 2A_\nu Y_\nu^\dagger M + 2MY_\nu^* A_\nu^T - B_\nu M Y_\nu^\dagger Y_\nu - Y_\nu^T Y_\nu^* B_\nu M\,. \tag{3.1}$$

Assuming $|A_\nu| \ll |B_\nu|$, the first term is negligible. The running of the mass matrix of the right-handed neutrinos is

$$16\pi^2\mu\frac{d\ M}{d\ \mu} = -MY_\nu^\dagger Y_\nu - Y_\nu^T Y_\nu^* M\,. \tag{3.2}$$

Comparing eqs. (3.1) and (3.2), we observe that, up to a correction of $\mathcal{O}(A_\nu/B_\nu)$, the neutrino B-term remains proportional to the mass matrix of the right-handed neutrinos.

Diagrams shown in figure 1 give a correction to $m^2_{H_u}$ which is equal to

$$-i\Delta m^2_{H_u} = 2\sum_k \int \frac{M_k^2\,\mathrm{Re}[B_\nu \sum_i (Y_\nu)_{ki}(A_\nu^*)_{ki}]}{k^2(k^2-M_k^2)^2}\frac{d^4k}{(2\pi)^4} = -i2\sum_{k,i}\mathrm{Re}\,\frac{[B_\nu\,\mathrm{Tr}(Y_\nu A_\nu^\dagger)]}{16\pi^2}\,. \tag{3.3}$$

Notice that this result is independent of M_k. That is because in the diagram (1) there are two vertices proportional to M_k (the B-term, $MB_\nu\tilde{N}\tilde{N}$, and $\langle F_N N\rangle \propto M$) which cancel the factors of M in the denominator. Had we defined the B-term as $B_\nu\tilde{N}\tilde{N}$ (with B_ν being of dimension of two), the result would have been inversely proportional to M.

At one loop level, there is no correction proportional to $|B_\nu|^2$ because the corresponding diagrams (shown in figure 2) cancel each other. The cancelation seems to be accidental because there is no symmetry forbidding such a contribution. At the two-loop level, there is only one diagram proportional to $[\mathrm{Trace}[Y_\nu^\dagger Y_\nu]]^2$. This diagram (which we will call

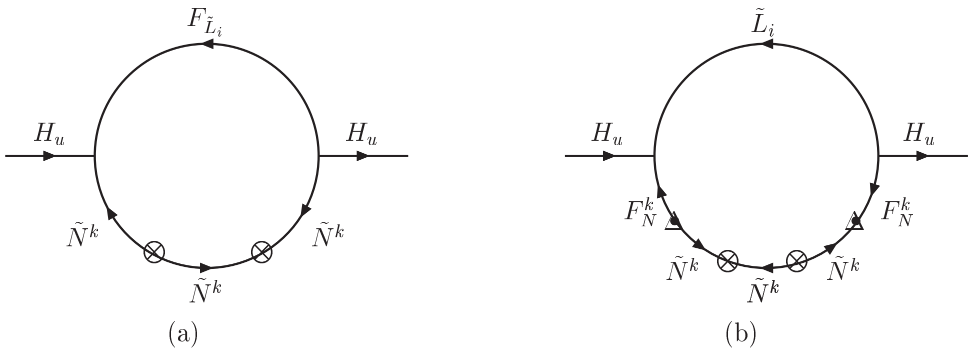


Figure 2: Diagrams proportional to $|B_\nu|^2$ contributing to $m^2_{H_u}$. F^k_N and $F_{\tilde{L}_i}$ represent the auxiliary fields associated with N_k and $\tilde{L}_i$, respectively. The neutrino B-term and M insertions are indicated by $\otimes$ and Δ, respectively. These diagrams cancel each other.

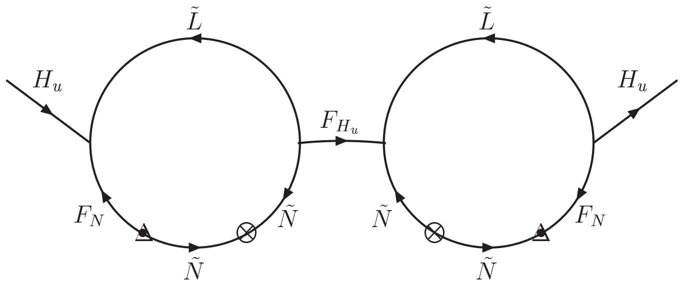


Figure 3: The "eyeglasses-diagram." The neutrino B-term and M insertions are indicated by $\otimes$ and Δ, respectively. This diagram gives a correction to $m^2_{H_u}$ equal to $|B_\nu|^2 \left[\mathrm{Tr}\{Y^\dagger_\nu Y_\nu\}/16\pi^2\right]^2$. Diagrams with a different topology have a different dependence on Y_ν and cannot cancel the effect of the "glasses-diagram." This demonstrates that the cancelation at one-loop level (see figure 2) is completely accidental.

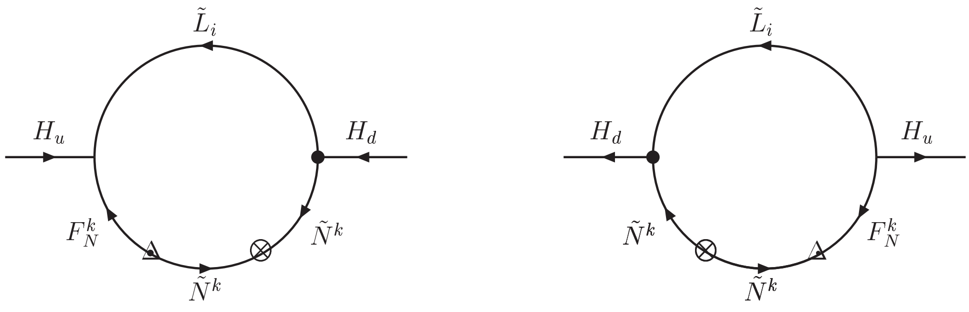
the "eyeglasses" diagram) is shown in figure 3. Diagrams with a different topology are proportional to other combinations of the Yukawa couplings (Y_{ki} and Y^*_{ki}) and cannot cancel the effect of the eyeglasses diagram. So, at the two-loop level, there is a non-zero effect proportional to $|B_\nu|^2$. Performing the full two-loop calculation is beyond the scope of this paper. The two-loop effect becomes significant only if $B_\nu Y^2_\nu/(16\pi^2) > m_{\rm susy}$ but as we shall see, such a possibility is quite unlikely.

The presence of a large neutrino B-term also induces non-negligible corrections to b_H as it is shown in Fig 4. The correction is finite and equal to

$$-i\Delta b_H = -B_\nu \sum_k \int \frac{M^2_k \,\mathrm{Tr}\left[(Y_\nu)_{ki}(Y^*_\nu)_{ki}\right]\mu}{k^2(k^2-M^2_k)^2}\frac{d^4k}{(2\pi)^4} = \frac{iB_\nu\mu\,\mathrm{Tr}[Y_\nu Y^\dagger_\nu]}{(4\pi)^2}\,. \tag{3.4}$$

By dimensional analysis we can show that any correction due to B_ν to the quartic Higgs interaction is suppressed by B_ν/M which is negligible. The contribution to the cubic Higgs term is also zero. So the potential of H^0_u and H^0_d is (see *e.g.*, [7])

$$V = (|\mu|^2 + m^2_{H_u} + \Delta m^2_{H_u})|H^0_u|^2 + (|\mu|^2 + m^2_{H_d})|H^0_d|^2 +$$



(a) (b)

Figure 4: Diagrams contributing to the Higgs B-term. F_N^k represents the auxiliary field associated with the right-handed neutrino, N_k. The two vertices marked with black circles are given by $\mu Y_\nu^\dagger$ and $\mu^* Y_\nu$. The neutrino B-term and M insertions are indicated by $\otimes$ and Δ, respectively.

$$+[(b_H + \Delta b_H) H_u^0 H_d^0 + \text{H.c.}] + \frac{g^2 + g'^2}{8}(|H_u^0|^2 - |H_d^0|^2)^2 \,. \tag{3.5}$$

Requiring $m_Z^2 = (g^2 + g'^2)(\langle H_u \rangle^2 + \langle H_d \rangle^2)/2$ and $\partial V/\partial H_u^0 = \partial V/\partial H_d^0 = 0$, we find

$$|\mu|^2 + m_{H_d}^2 = |b_H + \Delta b_H| \tan\beta - \left(\frac{m_Z^2}{2}\right)\cos 2\beta \tag{3.6}$$

and

$$|\mu|^2 + m_{H_u}^2 + \Delta m_{H_u}^2 = |b_H + \Delta b_H| \cot\beta + \left(\frac{m_Z^2}{2}\right)\cos 2\beta \tag{3.7}$$

where $\tan\beta = \langle H_u \rangle / \langle H_d \rangle$. Assuming $|\mu|^2 \sim m_{H_u}^2 \sim m_{\text{susy}}^2$, (3.6) gives

$$\left| b_H - B_\nu \mu \frac{\text{Tr}[Y_\nu Y_\nu^\dagger]}{16\pi^2} \right| \sim \frac{m_{\text{susy}}^2}{\tan\beta} \,. \tag{3.8}$$

From the LEP data [8], we know that $\tan\beta > 2$ and the data favors large values of $\tan\beta$ ($\tan\beta > 10$). Based on naturalness condition, it seems quite unlikely that b_H and Δb_H cancel each other, so we expect that

$$\frac{B_\nu Y_\nu^2}{(16\pi^2)} < \frac{m_{\text{susy}}}{\tan\beta} \,. \tag{3.9}$$

Notice that if $Y_\nu \ll 1$, B_ν can be still several orders of magnitude larger than m_{susy}.

4. Bounds on B_ν from different observables

In this section, we review the effects of B_ν on different observables and discuss how the bound we have found in the previous section completes our knowledge of this sector.

As it is discussed in [1], B_ν gives a correction to the neutrino mass equal to[1]

$$-\frac{g^2}{32\pi^2 \cos^2\theta_W} \frac{2B_\nu}{m_{\tilde{\nu}}} Y_\nu^T \frac{\langle H_u \rangle^2}{M} Y_\nu \sum_j f(y_j) |Z_{jZ}|^2 \,, \tag{4.1}$$

[1]Note that, within this model, there is another contribution to m_ν; namely, the standard seesaw effect given by $Y_\nu^T (\langle H_u \rangle^2 / M) Y_\nu$. In general, the neutrino mass is the sum of the two contributions and, a priori, we do not know which one is dominant.

where

$$f(y_j) = \frac{\sqrt{y_j}[y_j - 1 - \log(y_j)]}{(1-y_j)^2},$$

$y_j \equiv m_{\tilde{\nu}}^2/m_{\tilde{\chi}_j^0}^2$ and $Z_{jZ} \equiv Z_{j2}\cos\theta_W - Z_{j1}\sin\theta_W$ is the neutralino mixing matrix element that projects out the $\tilde{W}^0$ eigenstate from the jth neutralino. Using (4.1), authors of ref. [1] have concluded that

$$B_\nu < 10^3 m_{\rm susy}\,. \tag{4.2}$$

For $Y_\nu \sim 1$ (which is a value suggested by the unification models), the bound that we found in the previous section [see eq. (3.9)] is two orders of magnitude stronger. Even if $b_H \sim m_{\rm susy}$,[2] the bound from the electroweak symmetry breaking [$B_\nu \lesssim 16\pi^2 m_{\rm susy}/({\rm Tr}\{Y_\nu Y_\nu^\dagger\})$] is more restrictive. At first sight, it may seem that unlike the bound from the electroweak symmetry breaking, the bound in eq. (4.2) is independent of Y_ν but the fact is that to derive eq. (4.2), it has been assumed that the B_ν-induced contribution to m_ν is subdominant and the main contribution to m_ν is given by the standard tree-level seesaw formula: $m_\nu \simeq Y_\nu^T(\langle H_u\rangle^2/M)Y_\nu$. However, in principle, $Y_\nu^T(\langle H_u\rangle^2/M)Y_\nu$ can be much smaller than m_ν and the dominant contribution to the neutrino mass can be the B_ν-induced one-loop effect given in eq. (4.1). In this case, we can write

$$\sum_k (Y_\nu)_{ki}(Y_\nu)_{kj}({\rm GeV}/M_k)B_\nu/m_{\rm susy} \sim 10^{-12}(m_{ij}/0.1~{\rm eV})\sin^{-2}\beta\,. \tag{4.3}$$

In general Yukawa couplings are complex numbers and as a result different terms in the summation can cancel each other, allowing very large values of B_ν. Let us assume that there is not any significant cancelation between the different terms in eq. (4.3). In this case, if $M < 10^{14}$ GeV$/\tan\beta$, the bound from the neutrino mass on B_ν will be stronger than (3.8); otherwise, the bound from the electroweak symmetry breaking will be more restrictive. In general, these two bounds are complementary because, a priori, we do not know the values of M_k. Moreover the different terms contributing to eq. (4.3) can cancel each other; that is while, in the case of the bound from the electroweak symmetry breaking, the combination of the Yukawa couplings that appears is $\sum_{ki}|(Y_\nu)_{ki}|^2$ and therefore such a cancelation is not possible.

As it is discussed in ref. [2], B_ν also gives a contribution to the mass matrix of the left-handed sleptons:

$$\Delta m_{\tilde{L}}^2 = -2\,{\rm Re}(B_\nu^* a_0)\frac{Y_\nu^\dagger Y_\nu}{16\pi^2}\,. \tag{4.4}$$

The off-diagonal elements of this matrix are LFV and, in principle, can induce LFV rare decays. As it is well-known there are strong bounds on the branching ratio of the rare LFV decays [9, 10]. There are two possibilities to satisfy these bounds: (i) the off-diagonal elements of the matrix in eq. (4.4) are much lower than $m_{\rm susy}$; (ii) $B_\nu Y_\nu^2/16\pi^2 \gg m_{\rm susy}$ and therefore $\tilde{L}$ becomes so heavy that the relevant diagrams become suppressed. The bound we found in the previous section rules out the second possibility. Then, the bounds

[2]Note that, for $\tan\beta > 10$, if $b_H \sim m_{\rm susy}$, there has to be a fine-tuned cancelation between the two terms in eq. (3.8) which seems unnatural.

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Neutrino mass	$\lvert B_\nu \rvert < 10^3 m_{\rm susy}$ [1]
$\mathrm{Br}(\mu \to e\gamma)$	$\mathrm{Re}(a_0^* B_\nu)(Y_\nu^\dagger Y_\nu)_{\mu e}/(16\pi^2) < 10^{-4} m_{\rm susy}^2$ [2]
$\mathrm{Br}(\tau \to e\gamma)$	$\mathrm{Re}(a_0^* B_\nu)(Y_\nu^\dagger Y_\nu)_{\tau e}/(16\pi^2) < 0.1 m_{\rm susy}^2$ [2]
$\mathrm{Br}(\tau \to \mu\gamma)$	$\mathrm{Re}(a_0^* B_\nu)(Y_\nu^\dagger Y_\nu)_{\tau \mu}/(16\pi^2) < 0.01 m_{\rm susy}^2$ [2]
$d_e < 1.4 \times 10^{-27}$ e cm	$\mathrm{Im}(B_\nu)(Y_\nu^\dagger Y_\nu)_{ee}/(16\pi^2) < 0.1 m_{\rm susy}$ [2]
Electroweak symmetry breaking	$\lvert \mu B_\nu \rvert \,\mathrm{Tr}[Y_\nu Y_\nu^\dagger]/(16\pi^2) < m_{\rm susy}^2/\tan\beta$

Table 1: A summary of different bounds on B_ν. For simplicity, we have assumed that all the supersymmetry breaking parameters except B_ν are given by $m_{\rm susy} \sim 200\,\mathrm{GeV}$. The analytical formula in the text provide the exact dependence on the different susy breaking parameters.

on the branching ratio of the rare LFV decays [$\mathrm{Br}(\mu \to e\gamma) < 1.2 \times 10^{-11}$ [9], $\mathrm{Br}(\tau \to e\gamma) < 2.7 \times 10^{-6}$ [9], $\mathrm{Br}(\tau \to \mu\gamma) < 3.1 \times 10^{-7}$ [10]] give strong bounds on the off-diagonal elements of the matrix

$$\frac{\mathrm{Re}(a_0 B_\nu^*)}{(4\pi)^2 m_{\rm susy}^2}(Y_\nu^\dagger Y_\nu)_{\mu e} < \text{few} \times 10^{-4} \qquad \frac{\mathrm{Re}(a_0 B_\nu^*)}{(4\pi)^2 m_{\rm susy}^2}(Y_\nu^\dagger Y_\nu)_{\tau e} < 0.1 \tag{4.5}$$

and

$$\frac{\mathrm{Re}(a_0 B_\nu^*)}{(4\pi)^2 m_{\rm susy}^2}(Y_\nu^\dagger Y_\nu)_{\tau \mu} < \text{few} \times 10^{-2}\,. \tag{4.6}$$

Here, we have implicitly assumed that different possible LFV effects do not cancel each other. Note that these bounds are stronger than the bound we found in the previous section. However, these bounds apply only to the off-diagonal elements of the matrix $Y_\nu^\dagger Y_\nu$. In principle, the off-diagonal elements of $Y_\nu^\dagger Y_\nu$ can be much smaller than its diagonal elements. On the other hand, the dependence of b_H on B_ν is through $\sum_{ki} \lvert (Y_\nu)_{ki} \rvert^2$ which is larger than the maximum $\lvert (Y_\nu)_{ki} \rvert^2$. Thus, bounds from the LFV rare decay discussed in [2] and from the condition of electroweak symmetry breaking discussed in this paper are complementary.

The phase of B_ν is a source of CP-violation and it can induce electric dipole moment for charged leptons. At one-loop level, the effect is suppressed by inverse powers of M [11]. However, at the two-loop level there is not such a suppression: The neutrino B-term gives a correction to the A-term of the charged leptons, A_ℓ, independent of M [2]. Then, the imaginary part of A_ℓ [which, in the case $\mathrm{Im}(a_0) = 0$ is given by $\mathrm{Im}(B_\nu)$] gives a correction to the EDMs of charged leptons [2] while the real part gives a contribution to the MDM of charged leptons.[3] As in the case of LFV rare decays, we can argue that for $B_\nu Y_\nu^2/(16\pi)^2 \gg m_{\rm susy}$, these effects are suppressed because the mass of $\tilde{L}$ becomes too large. The possibility of such large $\lvert B_\nu \rvert$ is ruled out by the bound that we discussed in the previous section so the bounds on the EDMs of charged leptons can be translated into bounds on $\mathrm{Im}(B_\nu)$ without this ambiguity. The present bound on the electric dipole moment of the electron ($d_e < 1.4 \times 10^{-27}$ e cm [9]) implies $\mathrm{Im}(B_\nu) \sum_i \lvert (Y_\nu)_{ie} \rvert^2/(16\pi^2) < m_{\rm susy}$ which is again complementary to the bound we found in the previous section.

All these bounds are summarized in table 1. Notice that in the near future the bounds on d_e and the branching ratios of the LFV rare decays will be dramatically improved.

[3] I would like to thank Gilad Perez for drawing my attention to this point.

5. Summary

In this paper, we have first proposed a simple model in which $B_\nu \gg m_{\rm susy}$, introducing a new singlet chiral superfield, X. The model is very similar to the singlet Majoron model with the difference that the F-component of X also develops a VEV.

We have then shown that B_ν gives corrections to the Higgs mass parameters b_H and $m^2_{H_u}$. We have discussed that, to satisfy the condition for the electroweak symmetry breaking, $|b_H - B_\nu \mu \,{\rm Tr}[Y_\nu Y_\nu^\dagger]/(16\pi^2)|$ has to be of order of $m^2_{\rm susy}/\tan\beta$ and, as a result, $B_\nu \sum_{ki} |(Y_\nu)_{ki}|^2$ cannot be much larger than $16\pi^2 m_{\rm susy}/\tan\beta$ [see eq. (3.9)]. We then have discussed how this bound complements our knowledge of this sector, arguing that without this piece of information one could evade all the bounds on B_ν that previously had been discussed in the literature [1, 2].

Acknowledgments

I would like to thank B. Bajc, A. Brignole, A. Masiero, A. Rossi and G. Senjanovic for useful discussion. I am also very grateful to M.E. Peskin for his useful remarks. The early stage of this work was supported by DoE under contract DE-AC03-76SF00515.

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