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PHYSICS LETTERS B
Physics Letters B 564 (2003) 175–182
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Particle physics implications of the WMAP neutrino mass bound

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Received 4 March 2003; received in revised form 14 April 2003; accepted 14 April 2003

Editor: P.V. Landshoff

Abstract

The recently published cosmological bound on the absolute neutrino masses obtained from the Wilkinson microwave anisotropy probe (WMAP) data has important consequences for neutrino experiments and models. Taken at face value, the new bound excludes the determination of the absolute neutrino mass in the KATRIN experiment and disfavors a neutrino oscillation interpretation of the LSND experiment. Combined with the KamLAND and Super-K data, the WMAP bound defines an accessible range for the neutrinoless double beta decay amplitude. The bound also impacts the Z-burst annihilation mechanism for resonant generation of extreme-energy cosmic rays on the cosmic neutrino background in two ways: it constrains the local over-density of neutrino dark matter which is not helpful, but it also limits the resonant energy to a favorable range. In R-parity violating SUSY models, neutrino masses are generated by trilinear and bilinear lepton number violating couplings. The WMAP result improves the constraints on these couplings by an order of magnitude.

PACS: 12.60.Jv; 14.60.Pq; 14.80.Ly; 23.40.Bw; 96.40.Tv; 98.80.Es

1. Introduction

With the recently published first data of the Wilkinson microwave anisotropy probe (WMAP) [1] on the cosmic microwave background (CMB) anisotropies the age of precision cosmology has arrived. A flat, vacuum-energy dominated cold dark matter (ΛCDM) universe seeded by nearly scale-invariant Gaussian primordial fluctuations appears to be firmly established as the standard cosmology. Moreover, when combined with additional CMB data-sets (CBI, ACBAR) [2] and observations of large scale structure from the 2dF galaxy redshift survey (2dFGRS) [3] to lift degeneracies, the WMAP data offers the potential of testing various extensions and sub-dominant components in the ΛCDM model, such as small non-flatness, quintessence, possible tensor-gravitational wave modes, and a massive cosmic neutrino background (CνB). Investigation of the latter has most important consequences for terrestrial physics experiments exploring the neutrino sector.

The power spectrum of early-Universe density perturbations is processed by gravitational instabilities.

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doi:10.1016/S0370-2693(03)00587-2

However, the free-streaming relativistic neutrinos suppress the growth of fluctuations on scales below the horizon (approximately the Hubble size $c/H(z)$) until they become non-relativistic at $z \sim m_j/3T_0 \sim 1000\,(m_j/\text{eV})$. When the amplitude of fluctuations is normalized to the WMAP data, the amplitude of fluctuations in the 2dFGRS places significant limits on the contribution of neutrinos to the energy density of the universe,

$$\Omega_\nu h^2 = \frac{\sum_i m_i}{93.5\ \text{eV}} < 0.0076 \quad (95\%\ \text{C.L.}), \tag{1}$$

which translates into

$$\sum_i m_i < 0.71\ \text{eV} \quad (95\%\ \text{C.L.}). \tag{2}$$

The new mass bound (2) impacts most directly four-neutrino mass models constructed to accommodate the LSND evidence for oscillation. Such models require the heaviest neutrino mass to be ~ 1 eV, and so at face value are disfavored by the new result [4,5]. However, there are several loopholes in the argument against an ~ 1 eV sterile neutrino. If there is only one isolated "heavy" sterile as in the $3+1$ model, then the WMAP/2dF data at face value allow the Δm^2_{LSND} region up to 0.5 eV^2, whereas relaxing the WMAP/2dF bound from 0.71 eV to 1 eV allows virtually the entire LSND region to co-exist. In a $2+2$ model, there are two heavy mass eigenstates, and the WMAP/2dF data at face value limit Δm^2_{LSND} to 0.1 eV^2. Still another possibility, not yet explored to the best of our knowledge, might be to model the heavier neutrinos as decaying to light flavors plus a light boson, with a lifetime much less than the age of the Universe at structure formation. In such a model, the decay products would be free-streaming particles with masses well below the WMAP bound. Relevant to this discussion is the limit from Big Bang nucleosynthesis (BBN) [6,7], that neutrinos beyond the three active could not have been in thermal equilibrium already at the BBN temperature ~MeV, long before the epoch of structure formation. So the more serious constraint for the sterile neutrino is the BBN limit. Overcoming this BBN limit automatically immunizes the sterile against the WMAP/2dF bound [8], since the depopulated states at BBN are not populated at a later time. One way to evade thermalization at the BBN epoch is via a tiny lepton asymmetry [9]. There are several other ways, conveniently summarized in [6]. In conclusion, MiniBooNE is still required to settle the fate of the sterile neutrino [10].

From here on we focus on the consequences of the new WMAP bound for three-neutrino models. It was previously noted [11] that there are potentially four independent approaches for measuring the absolute neutrino mass. These are large-scale structure studies measuring the total mass in the CνB (as reported by WMAP), the Z-burst method measuring individual masses in the CνB, and the terrestrial measurements of the tritium end point spectrum and neutrinoless double beta decay rate. Of course, the results of these approaches are correlated in the sense that the experiments all attempt to determine the same neutrino masses. We will examine the impact of the new WMAP bound on the future of the other three approaches.

Neutrino oscillation studies have established three important facts of relevance here. The first is that the two mass-squared differences are small compared to the WMAP limit. Thus, when the WMAP limit is saturated, the three neutrinos are nearly degenerate in mass, and we have

$$m_i < 0.24\ \text{eV} \quad (95\%\ \text{C.L.}) \tag{3}$$

for each neutrino mass. The second important fact is that oscillation studies provide a *lower* bound on the heaviest neutrino mass, given by the minimum $\sqrt{\Delta m^2_{\text{atm}}} \sim 0.03$ eV. Thus, we may write

$$0.03\ \text{eV} \leqslant m_3 \leqslant 0.24\ \text{eV} \quad (95\%\ \text{C.L.}), \tag{4}$$

which shows the remarkable fact that knowledge of the heaviest neutrino mass (which we shall always denote by m_3) is now known to an order of magnitude! A plot of the total neutrino mass versus m_3 is shown in Fig. 1. The relation is linear, $\sum_i m_i = 3m_3$, except very near the smallest allowed m_3, $\sim \sqrt{\Delta m^2_{\text{atm}}}$. The third important fact is that the three angles parameterizing the unitary flavor-mass mixing-matrix, $U_{\alpha i}$, are well known. The one CP-violating Dirac phase and two CP-violating Majorana phases are not known. The angles and phases will be important when we look at neutrinoless double beta decay.

Absolute neutrino mass bounds also constrain all entries in the neutrino mass matrix in flavor space due to unitarity. This results in bounds on couplings in theories with lepton number violation [12]. As an

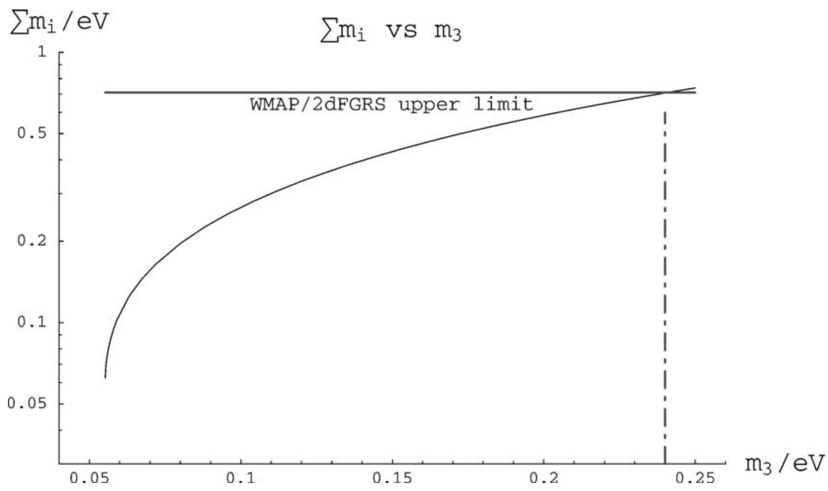

Fig. 1. Implications of the WMAP neutrino mass bound for the mass of the heaviest neutrino m_3. Here we take the best-fit value for $\Delta m^2_{\rm atm} = 3 \times 10^{-3}$ eV^2; $\Delta m^2_{\rm sun}$ is too small to be relevant.

example, we derive bounds on parameters of the R-parity violating ($\not{R}_P$) SUSY model, improving them by one order of magnitude over the existing values.

2. Tritium beta decay

The mass to be inferred from β-decay is $m^2_{\nu_e} \equiv \sum_j |U_{ej}|^2 m_j^2$. The KATRIN project [13] plans to start taking data in 2007. The sensitivity aim after three years of measurement is 0.08 eV^2 at 1σ accuracy. This may be improved to 0.05–0.06 eV^2, when optimizing the data point distribution and resolution, which implies a final sensitivity of m_{ν_e} to be 0.4 eV at 3σ. Thus, the reach of this experiment includes only the nearly mass-degenerate neutrino case, for which unitarity allows one to write $m_{\nu_e} = m_3$.

Comparing the KATRIN reach to the WMAP limit in Eq. (3), one comes to the unfortunate conclusion that a positive signal is unlikely.

3. Neutrinoless double beta decay

The mass inferred in neutrinoless double-β decay is

$$m_{ee} = \left| \sum_i U_{ei}^2 m_i \right|. \tag{5}$$

Here one needs the neutrino mixing parameters explicitly. The most recent analysis of atmospheric neutrino data [14] yields

$$10^{-3}\ \mathrm{eV}^2 < \Delta m^2_{\rm atm} < 5 \times 10^{-3}\ \mathrm{eV}^2 \tag{6}$$

and

$$\sin^2 2\theta_{\rm atm} > 0.8. \tag{7}$$

On the other hand, a recent evaluation of solar neutrino data including the KamLAND reactor experiment [15] inferred

$$5 \times 10^{-5}\ \mathrm{eV}^2 < \Delta m^2_{\rm sun} < 1.1 \times 10^{-4}\ \mathrm{eV}^2, \tag{8}$$

and

$$0.3 < \tan^2 \theta_{\rm sun} < 0.8. \tag{9}$$

Thus, the LMA solar solution is confirmed. The neutrino mixing matrix is seen to be “bi-large”. It is also known that $|U_{e3}|^2 \approx 0$, which means that the third mixing angle is negligibly small [16].

The cases of degenerate, hierarchical, and inverse hierarchical neutrinos (see Fig. 2) must be considered separately (for a detailed discussion, see, e.g., [17]). The WMAP limit is sufficiently large that it impacts only the case of degenerate neutrinos.

- Degenerate neutrinos: $m_1 \simeq m_2 \simeq m_3$. With $|U_{e3}|^2 \approx 0$, one has a mass proportional to $|U_{e1}^2 +$

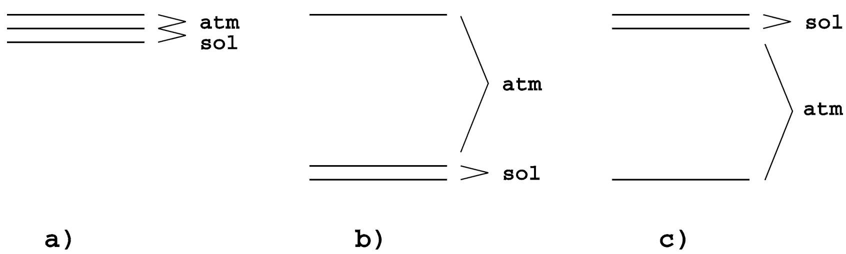

Fig. 2. Neutrino mass spectra for the three neutrino case: (a) degenerate, (b) hierarchical and (c) inverse hierarchical spectrum.

$U_{e2}^2| = |\cos^2\theta + e^{2i\delta}\sin^2\theta|$, which, upon extremizing the unknown phase, leads to

$$\cos 2\theta_{\text{sun}} m_3 < m_{ee} < m_3. \tag{10}$$

Inputting the new WMAP bound, and the solar angle, one gets

$$0.1\, m_3 < m_{ee} < 0.24\ \text{eV}; \tag{11}$$

- Hierarchical neutrinos: $m_1 \ll m_2 \ll m_3$ and $\Delta m^2_{\text{sun}} = \Delta m^2_{12}$. Here a lower limit is obtained by taking $|U_{e3}|^2 = 0$ and $m_1 = 0$. The result is

$$m_{ee} > \sqrt{\Delta m^2_{\text{sun}}}\, \sin^2\theta_{\text{sun}} = 2 \times 10^{-3}\ \text{eV}, \tag{12}$$

and $m_{ee} \ll m_3 \sim \sqrt{\Delta m^2_{\text{atm}}} \lesssim 0.07$ eV;
- Inverse hierarchical neutrinos: $m_1 \ll m_2 \simeq m_3$ and $\Delta m^2_{\text{sun}} = \Delta m^2_{23}$. The situation is analogous to the degenerate case, but with the scale of m_3 fixed by the atmospheric neutrino evidence, rather than the WMAP result. One gets

$$\cos 2\theta_{\text{sun}} \sqrt{\Delta m^2_{\text{atm}}} < m_{ee} < \sqrt{\Delta m^2_{\text{atm}}}, \tag{13}$$

i.e.,

$$3 \times 10^{-3}\ \text{eV} < m_{ee} < 0.07\ \text{eV}. \tag{14}$$

In summary, neglecting unnatural cancellations due to a conspiracy of δ, m_1 and mixing angles, the predicted range of m_{ee} is given by

$$2 \times 10^{-3}\ \text{eV} < m_{ee} < 0.24\ \text{eV}. \tag{15}$$

Fortunately, the whole region can be covered by the most ambitious double beta decay proposals [18] (for an overview of the experimental status see [19]). The lower limit is not impacted by the WMAP result, whereas the upper limit comes directly from the WMAP data. The central value of the recent discovery claim of the Heidelberg–Moscow experiment [20], $m_{ee} = 0.39^{+0.45}_{-0.34}$ eV, exceeds the WMAP bound, but the reported lower range does not (this fact has been pointed out already in Ref. [4]). We point out, though, that double beta decay mechanisms other than the standard neutrino mass mechanism are not affected by this bound. A particular interesting possibility to accommodate the Heidelberg–Moscow result involves singlet neutrinos propagating in large extra dimensions in which case a mechanism decorrelating the neutrino mass eigenstates from the double beta decay amplitude is operative [21]. Exchange of superpartners in R-parity violating SUSY, leptoquarks, or right-handed W bosons constitute other possibilities to account for a sizable neutrinoless double beta decay signal (for a review see [22]).

4. The Z-burst model for EECR's

The Z-burst mechanism [23] generates extreme-energy cosmic rays (EECRs) by resonant annihilation of a EECR neutrino on the CνB neutrinos. The resonant energy is

$$E^R = \frac{4 \times 10^{21}\ \text{eV}}{(m_\nu/\text{eV})}. \tag{16}$$

The decay products of Z-bursts include on average two nucleons and, from ten neutral pions, twenty photons. The decay multiplicity is $N \sim 30$. The nucleons lose $f \sim 20\%$ of their energy for each $\lambda \sim 6$ Mpc traveled in the CMB, so the average energy of a secondary nucleon arriving at Earth from distance D is

$$E_P \sim \frac{10^{21}\ \text{eV} \times (0.8)^{D/6\,\text{Mpc}}}{(m_\nu/0.1\ \text{eV})}. \tag{17}$$

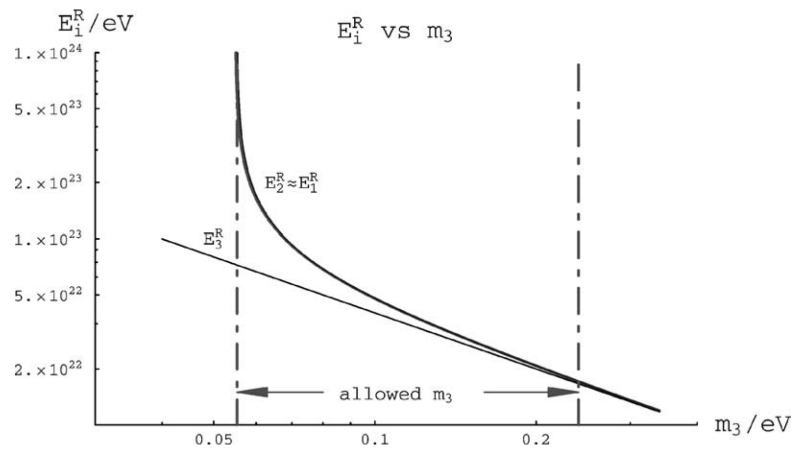

Fig. 3. Resonant energies for different neutrino mass eigenstates in the Z-burst model as a function of the largest neutrino eigenmass m_3. The Δm^2's used here are the same as in Fig. 1.

The photons have shorter absorption lengths, except above 10^{21} eV, and so are not expected to contribute much. For a neutrino mass in the range of Eq. (4), the mechanism is optimized: a larger mass would move Z-burst secondaries down below the GZK energy $\sim$ few $\times$ 10^{19} eV where the "background" of normal EECR events is huge, whereas a smaller mass would move the resonant energy beyond the reach of any realistic neutrino flux. The Z-burst resonant energies as a function of the heaviest neutrino mass m_3 are shown in Fig. 3. Note that over most of the allowed m_3 range, all three neutrinos contribute to annihilation with a resonant energy within a factor of two of each other.

In the simplest approximation, the spectrum of arriving nucleons is

$$\frac{dN}{dE} \sim \frac{1}{D^2} \times \frac{dN}{dD} \times \frac{dD}{dE} \propto E^{-1} \tag{18}$$

from sources uniformly distributed out to

$$D_{\mathrm{GZK}} \sim \lambda \frac{\ln\left(\frac{N E_{\mathrm{GZK}}}{E^R}\right)}{\ln(1-f)}, \tag{19}$$

with a pileup at E_{GZK} resulting from all primaries originating beyond this distance. The $1/E$ spectrum extends from E_{GZK} out to the maximum nucleon energy $\sim E^R/30 \sim 10^{21}(\frac{0.1\ \mathrm{eV}}{m_\nu})$ eV. More realistic simulations including energy-loss processes, cosmic expansion, and boosted Z-boson fragmentation functions give a softer spectrum, but a characteristic feature of the Z-burst mechanism remains that the super-GZK spectrum is considerably harder than the sub-GZK spectrum having power-law index -2.7.

What is not known is whether nature has provided the large neutrino flux at E^R to allow an appreciable event rate in future EECR detectors. It is conceivable, although unlikely, that the flux is so large that present EECR events are initiated by Z-bursts. A recent analysis [24] of this possibility gave a best fit with $m_\nu = 0.26^{+0.20}_{-0.14}$ eV, nicely consistent with the WMAP bound. Another analysis [25] fits the EECR spectrum down to the ankle with Z-burst generated events and a neutrino mass of 0.07 eV, again in accord with the WMAP bound. The flux requirements for the Z-burst mechanism can be ameliorated if there is an over-density of relic neutrinos, as would happen if (i) there was a significant chemical potential, or (ii) neutrinos were massive enough to cluster in "local" structures such as the galactic supercluster. Large chemical potentials have been ruled out recently [26], and this exclusion is confirmed by the WMAP data. Local clustering has been studied [27], with the conclusion being that a significant over density on the supercluster scale requires a neutrino mass in excess of 0.15 eV. Such a mass is marginally allowed by the new WMAP/2dF limit.

5. WMAP neutrino mass bound on $\not{R}_P$ SUSY

Supersymmetry without R-parity [28] provides an elegant mechanism for generating neutrino (Majorana) masses and mixings. In these models, there are mainly two sources of neutrino mass generation. In one scenario, products of trilinear λ and/or λ' couplings generate a complete neutrino mass matrix through one-loop self-energy graphs [29,30]. In the other scenario, the bilinear R-parity-violating terms induce sneutrino vacuum expectation values (VEVs) allowing neutrinos to mix with the neutralinos.

The L-violating part of the $\not{R}_P$ superpotential can be written as

$$\mathcal{W}_{RPV} = \frac{1}{2}\lambda_{ijk} L_i L_j E_k^c + \lambda'_{ijk} L_i Q_j D_k^c + \mu_i L_i H_u, \tag{20}$$

where i, j and k are quark and lepton generation indices. In Eq. (20), L_i and Q_i denote SU(2)-doublet lepton and quark superfields, E_i^c and D_i^c are SU(2)-singlet charged lepton and down-quark superfields, and H_u is the Higgs superfield responsible for the mass generation of the up-type quarks, respectively. There are 9 λ-type (due to an antisymmetry in the first two generation indices), 27 λ'-type and 3 μ_i couplings. Stringent upper limits exist on all these couplings from different experiments [31,32].

We first consider the effects of the λ' interactions. The relevant part of the Lagrangian can be written as

$$-\mathcal{L}_{\lambda'} = \lambda'_{ijk}\big[\bar{d}_k P_L \nu_i \tilde{d}_{jL} + \bar{\nu}_i^c P_L d_j \tilde{d}_{kR}^*\big] + \text{h.c.} \tag{21}$$

P_L is the left-helicity projector. Majorana mass terms for the left-handed neutrinos, given by $\mathcal{L}_M = -\frac{1}{2} \times m_{\nu_{ii'}} \bar{\nu}_{Li} \nu_{Ri'}^c + \text{h.c.}$, are generated at one loop. Fig. 4 show the corresponding diagrams. The masses so induced are given by

$$m_{\nu_{ii'}} \simeq \frac{N_c \lambda'_{ijk} \lambda'_{i'kj}}{16\pi^2} m_{d_j} m_{d_k} \times \left[\frac{f\big(m_{d_j}^2/m_{\tilde{d}_k}^2\big)}{m_{\tilde{d}_k}} + \frac{f\big(m_{d_k}^2/m_{\tilde{d}_j}^2\big)}{m_{\tilde{d}_j}}\right], \tag{22}$$

where $f(x) = (x\ln x - x + 1)/(x-1)^2$. Here, m_{d_i} is the down quark mass of the ith generation inside the loop, $m_{\tilde{d}_i}$ is some kind of an average of $\tilde{d}_{Li}$ and $\tilde{d}_{Ri}$ squark masses, and $N_c = 3$ is the color factor. In deriving Eq. (22), we assumed that the left-right squark mixing terms in the soft part of the Lagrangian are diagonal in their physical basis and are proportional to the corresponding quark masses, i.e., $\Delta m_{\text{LR}}^2(i) = m_{d_i} m_{\tilde{d}_i}$. The small effect of quark mixing is neglected in order not to complicate the discussion unnecessarily.

With λ-type interactions, one obtains exactly similar results as in Eqs. (21) and (22). The quarks and squarks in these equations will be replaced by the leptons and sleptons of the corresponding generations. The color factor $N_c = 3$ would be replaced by 1. We do not explicitly write them down.

For numerical purpose, we have assumed the mass of whatever scalar is relevant to be 100 GeV throughout, to be consistent with common practice and, in particular, to compare with the old bounds. While for sleptons this sounds a reasonable approximation, for squarks the present lower limit, even in $\not{R}_P$ scenarios, is around 250 GeV [33]. In any case, for different squark masses one can easily derive the appropriate bounds by straightforward scaling. It should be noted that the product couplings under consideration contribute to charged lepton masses as well, but with the present limits those contributions are too small to be of any relevance. The resulting bounds are

$$\begin{aligned}
&\lambda'_{i33}\lambda'_{i'33} < 3.6 \times 10^{-8}, && \lambda'_{i32}\lambda'_{i'23} < 8.9 \times 10^{-7},\\
&\lambda'_{i22}\lambda'_{i'22} < 2.2 \times 10^{-5}, && \lambda_{i33}\lambda_{i33} < 6.3 \times 10^{-7},\\
&\lambda_{i32}\lambda_{i23} < 1.1 \times 10^{-5}, && \lambda_{i22}\lambda_{i22} < 1.7 \times 10^{-4}.
\end{aligned} \tag{23}$$

There is one combination which receives a more stringent limit from μe conversion in nuclei [34], namely $\lambda'_{122}\lambda'_{222} < 3.3 \times 10^{-7}$. The chirality flips in Fig. 4 explain why with heavier fermions inside the loop the bounds are tighter. For this reason, we have presented the bounds only for $j, k = 2, 3$.

Next we consider the bilinear μ_i terms. Such terms lead only to one massive eigenstate as a result of tree level mixing between neutrinos and neutralinos. The induced neutrino mass [35] is given by $m \sim \mu_i^2/\mu$. Assuming the Higgsino mixing parameter $\mu = 100$ GeV, one obtains

$$\mu_i/\mu < 1.5 \times 10^{-6}. \tag{24}$$

The bounds in Eqs. (23) and (24) obtained using the recent WMAP bound are more stringent than the existing ones by one order of magnitude, precisely to the extent that the WMAP data have improved the absolute neutrino mass bound.

We make a note in passing that even our improved bounds on trilinear couplings do not invalidate the

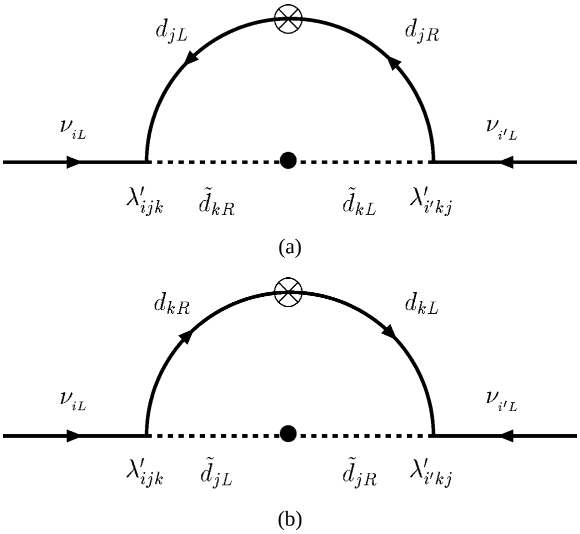

Fig. 4. The λ'-induced one loop diagrams contributing to Majorana masses for the neutrinos. The λ-induced diagrams are analogue with sleptons propagating in the loop.

$\not{R}_P$ SUSY search strategies proposed by the authors of [36]. Their suggestion is that at the Tevatron collider the production and decay of sparticles would occur in R-parity conserving modes except that the neutralino LSP would decay via $\not{R}_P$ channel into multi-b and missing energy final states constituting the signal.

6. Conclusions

We have discussed implications of the WMAP neutrino bound on future neutrino mass studies, including Tritium beta decay, neutrinoless double beta-decay, and the Z-burst mechanism for EECRs. We have shown that the Tritium beta decay project KATRIN is unlikely to measure an absolute neutrino mass, and that the WMAP bound in combination with the neutrino oscillation data defines a predicted range for the double beta-decay observable m_{ee}, which is accessible in the most ambitious proposed experiments. The WMAP bound also impacts the Z-burst mechanism for cosmic rays above the GZK cutoff. It constrains local over-densities, but it also limits the resonant energy to a favorable range.

Turning to model building, WMAP constrains theories with $\Delta L = 1$ lepton number violation, since in these theories Majorana neutrino masses are generated radiatively. Taking $\not{R}_P$ SUSY as our example, we have derived the upper limits on many individual and product couplings of the λ- and λ'-types, and also the bilinear μ_i terms, from their contribution to neutrino masses. Using the recent WMAP bound the limits have been improved by an order of magnitude. Finally, we remark that the new WMAP bound on neutrino mass coincides nicely with the one arising from the requirement of successful baryogenesis in the context of the neutrino see-saw model [37].

Note added

The new mass upper-limit expressed in Eqs. (2)–(4) depends on priors. In particular, the bound depends on the inclusion of Lyα data to estimate the power spectrum of intervening hydrogen clouds. Given the complexity of the Lyα analysis, some have questioned the reliability of this data set. Without the Lyα prior, the WMAP mass limit is relaxed to $\mathcal{O}(1)$ eV [38]. Accordingly, the upper dot-dash lines in Figs. 1 and 3, and the upper bounds in Eqs. (11), (23) and (24), would also relax by $\sim 40\%$.

Acknowledgements

We thank Nicole Bell and Bill Louis for helpful remarks. G.B. acknowledges hospitality of LPT, Univ. de Paris XI, Orsay, where a part of the work has been done. This research has been supported by (i) US Department of Energy grant number DE-FG05-85ER40226, (ii) the Bundesministerium für Bildung und Forschung (BMBF, Bonn Germany) under contract number 05HT1WWA2 (for H.P.), and (iii) in part, the Department of Science and Technology, India, project number SP/S2/K-10/2001 (for G.B.).

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