Neutrino masses from beta decays after KamLAND and WMAP

F. R. Joaquim*
Centro de Física das Interacções Fundamentais (CFIF), Departamento de Física, Instituto Superior Técnico, Avenida Rovisco Pais, 1049-001 Lisboa, Portugal

(Received 12 May 2003; published 28 August 2003)

The first data released by the KamLAND Collaboration have confirmed the strong evidence in favor of the large mixing angle solution of the solar neutrino problem. Taking into account the ranges for the oscillation parameters allowed by the global analysis of the solar, CHOOZ and KamLAND data, we update the limits on the neutrinoless double beta decay effective neutrino mass parameter and analyze the impact of all the available data from neutrinoless double beta decay experiments on the neutrino mass bounds, in view of the latest WMAP results. For the normal neutrino mass spectrum the range (0.05–0.23) eV is obtained for the lightest neutrino mass if one takes into account the Heidelberg-Moscow evidence for neutrinoless double beta decay and the cosmological bound. It is also shown that under the same conditions the mass of the lightest neutrino may not be bounded from below if the spectrum is of the inverted type. Finally, we discuss how future experiments can improve the present bounds on the lightest neutrino mass set by the Troitsk, Mainz, and Wilkinson microwave anisotropy probe results.

DOI: 10.1103/PhysRevD.68.033019 PACS number(s): 14.60.Pq, 23.40.−s, 26.65.+t

I. INTRODUCTION

The first results reported by the Kamioka liquid scintillator antineutrino detector (KamLAND) experiment [1] have brought new light to the search of the true solution of the solar neutrino problem. Through large distance measurements, the KamLAND Collaboration has measured, for the first time, reactor $\overline{\nu_e}$ disappearance within the large mixing angle allowed region for the neutrino oscillation parameters. Moreover, the Super-Kamiokande (SK) atmospheric [2] and the KEK to Kamioka (K2K) accelerator [3] neutrino experiments strongly indicate that neutrino flavor oscillations in the $\nu_\mu \rightarrow \nu_\tau$ channel are the most natural explanation for ν_μ disappearance. These two evidences have changed our standard picture of particle physics in the sense that neutrino oscillations require neutrinos to be massive. In addition to this, neutrinos behave very differently from the other known elementary particles not only because they are much lighter but also due to the fact that their mixing pattern differs very much from the one observed in the quark sector.

In spite of all the great achievements of oscillation experiments, we are still far from a reasonable understanding of neutrino properties. Among all the unanswered questions in neutrino physics, the most fundamental one concerns the nature of neutrinos. In particular, it is of prior importance to know whether they are Dirac or Majorana particles [4]. While Dirac neutrinos require the existence of highly suppressed Yukawa couplings, which are difficult to accommodate on theoretical grounds, small Majorana neutrino masses naturally arise in the context of minimal extensions of the standard model where the seesaw mechanism [5] operates. Still, our present knowledge of the neutrino sector leaves too much room for speculation about the neutrino mass generation mechanism.

If neutrinos are massive, then the first question which immediately arises is concerned with the value of their absolute mass scale. The direct neutrino mass determination method relies on the detailed analysis of the end-point part of the beta decay spectrum of some nuclei [6]. At present, the most stringent experimental bound on the neutrino mass comes from the Mainz [7] and Troitsk [8] experiments which have set a maximum value for m_{ν_e} of 2.2 eV. The KATRIN experiment [9], which is planned to start taking data in 2007, will be able to improve the sensitivity to neutrino masses by approximately one order of magnitude.

The search for positive signs in the neutrinoless double beta $[(\beta\beta)_{0\nu}]$-decay mode of certain even-even nuclei seems to be, at present, the most reliable way to look for the Majorana nature of neutrinos. The phenomenological consequences of $(\beta\beta)_{0\nu}$ decays in the framework of neutrino oscillations have been widely studied in the literature [10,11]. The extraction of neutrino mass limits through $(\beta\beta)_{0\nu}$-decay measurements involves certain subtleties which are related to the fact that this method of probing on the absolute neutrino mass scale strongly depends on the results provided by neutrino oscillation experiments. Therefore, and in spite of not being sensitive to it, neutrino oscillation experiments turn out to be of great importance in the determination of the absolute neutrino mass scale.

All the outstanding developments in experimental neutrino physics were accompanied by an equally remarkable evolution of cosmological experiments. Recently, the results from the Wilkinson microwave anisotropy probe (WMAP) have brought new insights into the measurements of a large set of cosmological parameters with an incredible precision. When combined with the data from the two degree field (2dF) galactic redshift survey [12], the WMAP results severely constrain the neutrino masses. In particular, this simultaneous analysis leads to the following upper bound on the neutrino contribution to the Ω cosmological parameter, $\Omega_\nu h^2 < 0.0076$ where h is the Hubble constant [13]. This

*Electronic address: filipe@gtae3.ist.utl.pt

bound, together with the relation $\Sigma_i m_i = 91.5\,\Omega_\nu h^2$ (eV) [14] implies $\Sigma_i m_i < 0.70$ eV. Hence, in the framework of three light neutrino species, one obtains

$$m_i \lesssim 0.23 \text{ eV}, \tag{1}$$

for the mass of each neutrino, indicating that the KATRIN experiment may not have enough sensitivity to measure the electron neutrino mass. Although this is true for β-decay experiments it may not hold for $(\beta\beta)_{0\nu}$-decay searches which will be sensitive to even lower values of neutrino masses, in the future. This raises the interesting question on how competitive their results can be when compared with the cosmological ones.

Besides the constraints imposed on neutrino masses, the precise cosmological measurements also place important bounds on the number of effective allowed neutrino species N_ν^{eff}. Furthermore, using the range of the baryon-to-photon ratio $\eta = 6.5^{+0.4}_{-0.3} \times 10^{-10}$ and the measurements of ^{4}He primordial abundances, it has been shown that $N_\nu^{eff} < 3.4$ [15]. This result turns out to be in serious conflict with the evidence for $\overline{\nu}_\mu \to \overline{\nu}_e$ oscillations reported by the liquid scintillator neutrino detector (LSND) Collaboration [16]. It is well known that the LSND results require the existence of at least four neutrino species which is incompatible with the bound $N_\nu^{eff} < 3.4$. Moreover, the neutrino oscillation data indicate that four neutrino oscillation scenarios are disfavored when compared with the three neutrino oscillation scheme [17]. The KARMEN experiment [18] already excludes part of the LSND parameter region but a conclusive statement about the validity of the LSND results can only be made by the MiniBooNE experiment [19].

In this paper we update the bounds on the effective Majorana neutrino mass parameter in view of the latest global analysis of all the solar, CHOOZ, and KamLAND data. This will be done in the framework of the normal and inverted neutrino mass schemes. Taking into account the WMAP constraint given in Eq. (1) and all the presently available neutrino oscillation and $(\beta\beta)_{0\nu}$-decay data we obtain the neutrino mass bounds for each type of neutrino mass spectrum. Finally, the impact of future $(\beta\beta)_{0\nu}$-decay projects on the determination of neutrino masses is discussed and the consequences of the first Heidelberg-Moscow evidence for $(\beta\beta)_{0\nu}$ decay, when considered simultaneously with the WMAP bound on m_i, are analyzed.

II. NEUTRINO OSCILLATION DATA: PRESENT STATUS

The presently available neutrino oscillation data, not including the results from LSND, can be accommodated in the framework of three mixed massive neutrinos. All the information about neutrino mixing is enclosed in the leptonic mixing matrix U which relates neutrino flavor and mass eigenstates in the following way

$$\nu_{L\alpha} = \sum_{i=1}^{3} U_{\alpha j} \nu_{Lj}, \quad \alpha = e, \mu, \tau, \quad j = 1,2,3. \tag{2}$$

The mixing matrix U is a 3×3 unitary matrix and its form depends on whether neutrinos are Dirac or Majorana particles. In the framework of three light Majorana neutrinos, the matrix U can be parametrized as

$$U = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & -c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13}e^{i\delta} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13}e^{i\delta} \end{pmatrix} \cdot \begin{pmatrix} 1 & 0 & 0 \\ 0 & e^{i\alpha} & 0 \\ 0 & 0 & e^{i\beta} \end{pmatrix}, \tag{3}$$

where $s_{ij} \equiv \sin\theta_{ij}$ and $c_{ij} \equiv \cos\theta_{ij}$. The phase δ is the leptonic Dirac CP-violating phase and α and β are Majorana phases [20]. The Dirac phase δ can induce CP-violating effects in neutrino oscillations, sizable enough to be measured by very long baseline neutrino oscillation experiments in the future [21]. On the contrary, oscillation experiments are blind to the physical effects associated to the Majorana phases α and β. Therefore, the determination of these phases is only viable in experiments sensitive to the Majorana nature of neutrinos like $(\beta\beta)_{0\nu}$-decay experiments. Nevertheless, this seems to be a difficult task to achieve since it requires not only the knowledge of all the neutrino mass and mixing parameters but also the understanding of the physics related with $(\beta\beta)_{0\nu}$-decays; namely, the uncertainties in the nuclear matrix elements involved in the calculation of the $(\beta\beta)_{0\nu}$-decay rates seem to be the major problem on the possible determination of the Majorana phases [22].

The neutrino oscillation experimental results provide us with information about the neutrino mixing angles θ_{ij} and mass squared differences $\Delta m^2_{ij} = m_i^2 - m_j^2$. The identification of these parameters with the experimentally measured ones depends on the neutrino mass ordering. In this paper we will always identify θ_{12}, θ_{23}, and θ_{13} with the "solar," "atmospheric," and the CHOOZ angles, respectively. There are two possible ways of ordering neutrino masses corresponding to the normal ($m_1 < m_2 < m_3$) and inverted ($m_3 < m_1 < m_2$) spectra. In both cases one can express two of the neutrino masses as a function of the remaining one and the Δm^2_{ij}'s. For the normal neutrino mass spectrum (NNMS)

$$m_2 = \sqrt{m_1^2 + \Delta m^2_{21}}, \quad m_3 = \sqrt{m_1^2 + \Delta m^2_{21} + |\Delta m^2_{32}|}, \tag{4}$$

and for the inverted spectrum (INMS)

$$m_2 = \sqrt{m_3^2 + |\Delta m^2_{32}|}, \quad m_1 = \sqrt{m_3^2 + |\Delta m^2_{32}| - \Delta m^2_{21}}. \tag{5}$$

TABLE I. Allowed ranges for the oscillation parameters Δm^2_{21} and $\sin^2\theta_{12}$ taken from the global analysis of all solar+CHOOZ and all solar+CHOOZ+KamLAND data performed in Ref. [24].

	2ν Solar+CHOOZ			2ν Solar+CHOOZ+KamLAND	
	Δm^2_{21} ($\times 10^{-4}$ eV^2)	$\sin^2\theta_{12}$		Δm^2_{21} ($\times 10^{-4}$ eV^2)	$\sin^2\theta_{12}$
99% C.L.	(0.25−4.2)	(0.21−0.46)	LMA I	(0.52−1.0)	(0.23−0.46)
			LMA II	(1.2−2.1)	(0.23−0.39)
95% C.L.	(0.28−2.3)	(0.24−0.42)	LMA I	(0.57−0.92)	(0.24−0.4)
			LMA II	(1.4−1.8)	(0.27−0.33)
90% C.L.	(0.3−1.9)	(0.24−0.4)		(0.6−0.9)	(0.26−0.4)
Best fit	0.6	0.3		0.7	0.3

The SK and K2K neutrino data point towards the existence of neutrino oscillations in the $\nu_\mu \rightarrow \nu_\tau$ channel. The results of these two experiments constrain the parameters θ_{23} and $|\Delta m^2_{32}|$ which, considering the SK only and SK+K2K data, are found to lie in the ranges [23]

$$\text{SK(99\% C.L.)}: \quad 1.3\times 10^{-3}\,\text{eV}^2 \leqslant |\Delta m^2_{32}| \leqslant 5.0\times 10^{-3}\,\text{eV}^2, \quad \sin^2 2\,\theta_{23} > 0.85, \tag{6}$$

$$\text{SK+K2K(99\% C.L.)}: \quad 1.4\times 10^{-3}\,\text{eV}^2 \leqslant |\Delta m^2_{32}| \leqslant 3.8\times 10^{-3}\,\text{eV}^2, \quad \sin^2 2\,\theta_{23} > 0.85,$$

with the best-fit values

$$\text{SK}: |\Delta m^2_{32}| = 2.7\times 10^{-3}\,\text{eV}^2, \quad \sin^2 2\,\theta_{23} = 1,$$

$$\text{SK+K2K}: |\Delta m^2_{32}| = 2.6\times 10^{-3}\,\text{eV}^2, \quad \sin^2 2\,\theta_{23} = 1. \tag{7}$$

Although the impact of the K2K data is still not very significant, the most important fact to retain is that these results are compatible with the pre-K2K ones.

The presently available data from all solar and reactor neutrino experiments confirm that the solar neutrino problem can be explained through the existence of $\nu_e \rightarrow \nu_{\mu,\tau}$ oscillations. After the release of the first KamLAND results, several global analysis of all the solar, CHOOZ, and KamLAND data have been performed [24,25]. As an example we take the pre-KamLAND and post-KamLAND results obtained in [24] which are summarized in Table I. Besides selecting the LMA region as *the* solution to the solar neutrino problem, the KamLAND experiment significantly restricts the corresponding oscillation parameter space as can be seen from Table I. The absence of $\overline{\nu_e}$ disappearance reported by the CHOOZ [26] and Palo Verde [27] experiments imposes severe bounds on the θ_{13} angle. The global analysis performed in [24] shows that

$$\sin^2\theta_{13} \lesssim 0.05 \quad (99.73\% \text{ C.L.}), \tag{8}$$

being the best-fit value

$$(\sin^2\theta_{13})_{\text{BF}} \lesssim 0.01. \tag{9}$$

In spite of being insensitive to the absolute neutrino mass scale, neutrino oscillation experiments indicate that neutrinos oscillate with $\Delta m^2_{21} \ll |\Delta m^2_{32}|$. This allows the classification of the neutrino spectrum in hierarchical (HI), inverted hierarchical (IH), and quasidegenerate (QD). From Eqs. (4) and (5) one has[1]

$$\text{HI}: \quad m_1 \ll \Delta m^2_{21} \Rightarrow m_2 \simeq \sqrt{\Delta m^2_{21}}, \quad m_3 \simeq \sqrt{|\Delta m^2_{32}|},$$

$$\text{IH}: \quad m_3 \ll |\Delta m^2_{32}| \Rightarrow m_1 \simeq m_2 \simeq \sqrt{|\Delta m^2_{32}|},$$

$$\text{QD}: \quad m_1 \gg |\Delta m^2_{32}| \Rightarrow m_1 \simeq m_2 \simeq m_3. \tag{10}$$

Although, the presently available neutrino data do not discriminate between normal and inverted neutrino mass spectra, such a selection will be possible in future long baseline neutrino experiments; namely, the detailed study of earth matter effects in neutrino oscillations will allow for the determination of the sign of Δm^2_{32} and therefore give us a hint about the mass ordering of neutrino states [29].

III. $(\beta\beta)_{0\nu}$ DECAYS AND NEUTRINO MASS SPECTRA

The combined analysis of $(\beta\beta)_{0\nu}$ decay and neutrino oscillation experimental results may be of crucial importance on the clarification of some aspects related with massive neutrinos. In particular, the observation of these processes may not only reveal the Majorana character of neutrinos but also help in the determination of the absolute neutrino mass scale and spectra. If these decays occur due to the exchange of virtual massive Majorana neutrinos, their probability amplitudes are proportional to the so-called effective Majorana mass parameter

$$m_{ee} = |(\mathcal{M}_\nu)_{11}| = \left|\sum_{i=1}^{3} m_i U^2_{ei}\right|, \tag{11}$$

where $\mathcal{M}_\nu$ is the Majorana neutrino mass matrix, m_i is the mass of the neutrino mass eigenstate ν_i, and the U_{ei} are the elements of the first row of the leptonic mixing matrix. The

[1]We will not consider here the cases where $m_1 \simeq m_2 < m_3$ and $m_3 < m_1 \simeq m_2$ which correspond to the partial and partial inverted mass hierarchy spectra, respectively. The reader is addressed to Ref. [10] for a complete analysis on the subject.

most significant bounds on the value of the effective neutrino mass parameter come from the Heidelberg-Moscow [30] and IGEX [31] ^{76}Ge experiments. Taking into account the uncertainties in the nuclear matrix elements involved in the determination of the $(\beta\beta)_{0\nu}$-decay amplitudes, one has

$$m_{ee}\leqslant(0.35-1.24)\ \text{eV}\quad(\text{Heidelberg-Moscow}),\tag{12}$$

$$m_{ee}\leqslant(0.33-1.35)\ \text{eV}\quad(\text{IGEX}).\tag{13}$$

The reanalysis of the Heidelberg-Moscow data performed in [32] has been interpreted as an evidence of $(\beta\beta)_{0\nu}$ decay of ^{76}Ge. The deduced range for m_{ee} was

$$m_{ee}=(0.11-0.56)\ \text{eV}\quad 95\%\ \text{C.L.},\tag{14}$$

which is modified to

$$m_{ee}=(0.05-0.84)\ \text{eV}\quad 95\%\ \text{C.L.},\tag{15}$$

if a $\pm 50\%$ uncertainty of the nuclear matrix elements is considered. The interval

$$m_{ee}=(0.4-1.3)\ \text{eV}\tag{16}$$

has been obtained in Ref. [33] using a different set of nuclear matrix elements. After their publication, the results presented in Ref. [32] were criticized[2] by some authors [35]. At the same time, some phenomenological implications of this claimed evidence were explored [36]. In any case, future experiments will have enough sensitivity to clarify this situation, and if confirmed, we will surely have the first evidence in favor of the Majorana nature of massive neutrinos.

The definition of the effective Majorana neutrino mass parameter m_{ee} in terms of the physical parameters θ_{ij}, m_i, δ, α, and β, is easily obtained from Eq. (11) and the relation

$$\mathcal{M}_\nu=U^*\text{diag}(m_1,m_2,m_3)U^\dagger,\tag{17}$$

which comes from the diagonalization of the 3×3 Majorana mass matrix. Parametrizing the mixing matrix U as done in Eq. (3) one gets

$$\begin{aligned}m_{ee}&=|m_1U_{e1}^2+m_2U_{e2}^2+m_3U_{e3}^2|\\&=|m_1c_{12}^2c_{13}^2+m_2c_{13}^2s_{12}^2e^{-2i\alpha}+m_3s_{13}^2e^{-2i\beta}|.\end{aligned}\tag{18}$$

This expression can be further written in terms of the lightest neutrino mass m_1 (m_3) for the NNMS (INMS) and the Δm^2's, taking into account the mass definitions given in Eqs. (4) and (5)

$$\begin{aligned}m_{ee}^{\text{NS}}=&|m_1c_{12}^2c_{13}^2+\sqrt{m_1^2+\Delta m_{21}^2}c_{13}^2s_{12}^2e^{-2i\alpha}\\&+\sqrt{m_1^2+\Delta m_{21}^2+|\Delta m_{32}^2|}s_{13}^2e^{-2i\beta}|,\end{aligned}\tag{19}$$

$$\begin{aligned}m_{ee}^{\text{IS}}=&|\sqrt{m_3^2+|\Delta m_{32}^2|-\Delta m_{21}^2}c_{12}^2c_{13}^2\\&+\sqrt{m_3^2+|\Delta m_{32}^2|}c_{13}^2s_{12}^2e^{-2i\alpha}+m_3s_{13}^2e^{-2i\beta}|.\end{aligned}\tag{20}$$

Furthermore, since $\Delta m_{21}^2\ll|\Delta m_{32}^2|$, the following approximations hold:

$$\begin{aligned}m_{ee}^{\text{NS}}\simeq m_1\Bigg|&c_{12}^2c_{13}^2+\sqrt{1+\frac{\Delta m_{21}^2}{m_1^2}}c_{13}^2s_{12}^2e^{-2i\alpha}\\&+\sqrt{1+\frac{|\Delta m_{32}^2|}{m_1^2}}s_{13}^2e^{-2i\beta}\Bigg|,\end{aligned}\tag{21}$$

$$\begin{aligned}m_{ee}^{\text{IS}}\simeq m_3\Bigg|&\sqrt{1+\frac{|\Delta m_{32}^2|}{m_3^2}}c_{12}^2c_{13}^2\\&+\sqrt{1+\frac{|\Delta m_{32}^2|}{m_3^2}}c_{13}^2s_{12}^2e^{-2i\alpha}+s_{13}^2e^{-2i\beta}\Bigg|.\end{aligned}\tag{22}$$

The above equations are valid for any neutrino mass spectrum since the lightest neutrino mass m_1 (or m_3) is not constrained. With the help of Eqs. (10) one can obtain approximate expressions of the effective neutrino mass parameter m_{ee} depending on the type of neutrino mass spectrum. From Eqs. (19) and (20) and the definitions given in Eq. (10) one has

$$m_{ee}^{\text{HI}}\simeq|\sqrt{\Delta m_{21}^2}s_{12}^2c_{13}^2+\sqrt{|\Delta m_{32}^2|}s_{13}^2e^{2i(\alpha-\beta)}|,\tag{23}$$

and

$$m_{ee}^{\text{IH}}\simeq\sqrt{|\Delta m_{32}^2|}c_{13}^2|c_{12}^2+s_{12}^2e^{-2i\alpha}|,\tag{24}$$

for the HI and IH neutrino mass spectra, respectively. In the case of three quasidegenerate neutrinos with a common mass approximately equal to m, Eqs. (19) and (20) reduce to

$$m_{ee}^{\text{QD}}\simeq m|c_{13}^2(c_{12}^2+s_{12}^2e^{-2i\alpha})+s_{13}^2e^{-2i\beta}|.\tag{25}$$

The allowed ranges for the effective neutrino mass parameter in each case depend not only on the values of the Δm_{ij}^2's and mixing angles but also on the Majorana phases α and β. In fact, depending on whether CP is conserved or violated the value of m_{ee} can drastically change [28]. Here we will only be concerned with the absolute bounds on m_{ee} which can be easily obtained from Eqs. (23)–(25) with the appropriate choice of the phases α and β, leading to

$$\begin{aligned}(m_{ee}^{\text{HI}})_{\text{low}}&\simeq|\sqrt{\Delta m_{21}^2}s_{12}^2c_{13}^2-\sqrt{|\Delta m_{32}^2|}s_{13}^2|,\\(m_{ee}^{\text{HI}})_{\text{up}}&\simeq\sqrt{\Delta m_{21}^2}s_{12}^2c_{13}^2+\sqrt{|\Delta m_{32}^2|}s_{13}^2,\end{aligned}\tag{26}$$

$$\begin{aligned}(m_{ee}^{\text{IH}})_{\text{low}}&\simeq\sqrt{|\Delta m_{32}^2|}(1-s_{13}^2)(1-2s_{12}^2),\\(m_{ee}^{\text{IH}})_{\text{up}}&\simeq\sqrt{|\Delta m_{32}^2|}(1-s_{13}^2),\end{aligned}\tag{27}$$

[2] For a complete discussion on proofs and disproofs the reader is addressed to Ref. [34], and references therein.

TABLE II. Allowed ranges for the effective Majorana mass parameter m_{ee} for each type of neutrino mass spectrum. The values presented here were determined taking into account the SK+K2K data summarized in Eqs. (6)–(7) and the solar+CHOOZ+KamLAND results given in Table I and in Eqs. (8) and (9).

		$m_{ee}^{\rm HI}$ ($\times 10^{-3}$ eV)	$m_{ee}^{\rm IH}$ ($\times 10^{-2}$ eV)	$m_{ee}^{\rm QD}/m$
99% C.L.	LMA I	$\lesssim 7.5$	$(0.3-6.2)$	$(0.026-1)$
	LMA II	$\lesssim 8.5$	$(0.8-6.2)$	$(0.16-1)$
95% C.L.	LMA I	$\lesssim 6.5$	$(0.7-6.2)$	$(0.14-1)$
	LMA II	$\lesssim 7.3$	$(1.2-6.2)$	$(0.27-1)$
90% C.L.		$\lesssim 6.7$	$(0.7-6.2)$	$(0.14-1)$
Best fit		$(2.0-3.0)$	$(2.0-5.0)$	$(0.39-1)$

$$(m_{ee}^{\rm QD})_{\rm low} \simeq m|c_{13}^2(1-2s_{12}^2)-s_{13}^2|, \quad (m_{ee}^{\rm QD})_{\rm up} \simeq m. \tag{28}$$

The expression $(m_{ee}^{\rm QD})_{\rm up}$ together with the WMAP bound on neutrino masses given in Eq. (1) implies $m_{ee} \lesssim 0.23$ eV which can be interpreted has the *cosmological* bound on the effective Majorana mass parameter. Cancellations in m_{ee} can in principle occur in the HI and QD cases if the following conditions are fulfilled:

$$\text{HI:} s_{13}^2 = \frac{\sqrt{\Delta m_{21}^2} s_{12}^2}{\sqrt{|\Delta m_{32}^2|} + \sqrt{\Delta m_{21}^2} s_{12}^2} \simeq (0.03-0.15), \tag{29}$$

$$\text{QD:} s_{13}^2 \simeq \frac{\cos 2\theta_{12}}{1+\cos 2\theta_{12}} \simeq (0.07-0.35). \tag{30}$$

Comparing these ranges with the bound given in Eq. (8) one immediately concludes that cancellations cannot occur if neutrinos are almost degenerate. Still, the condition $m_{ee}=0$ is compatible with a hierarchical neutrino mass spectrum. From Eq. (27) one can see that this can be accomplished in the IH case if $s_{12}^2=0.5$, which is already excluded by the data at 99% C.L. The allowed ranges for the Majorana mass parameter are shown in Table II for each type of neutrino mass spectrum and the dependence of m_{ee} on the lightest neutrino mass for the NNMS (INMS) is shown in Figs. 1(a,b) where the shaded regions indicate the possible values for m_{ee} in each case. From Fig. 1 and Table II it becomes evident that the intervals determined for m_{ee} may overlap when considering the different types of neutrino mass spectra. This fact raises the interesting question which are the conditions that must be fulfilled in order to get a clear separation of the different types of neutrino mass schemes. The imposition of such constraints can be translated into

$$(m_{ee}^{\rm HI})_{\rm up} < (m_{ee}^{\rm IH})_{\rm low}, \quad (m_{ee}^{\rm IH})_{\rm up} < (m_{ee}^{\rm QD})_{\rm low},$$

$$(m_{ee}^{\rm HI})_{\rm up} < (m_{ee}^{\rm QD})_{\rm low}. \tag{31}$$

Making use of Eqs. (26)–(28) one can show that this is equivalent to

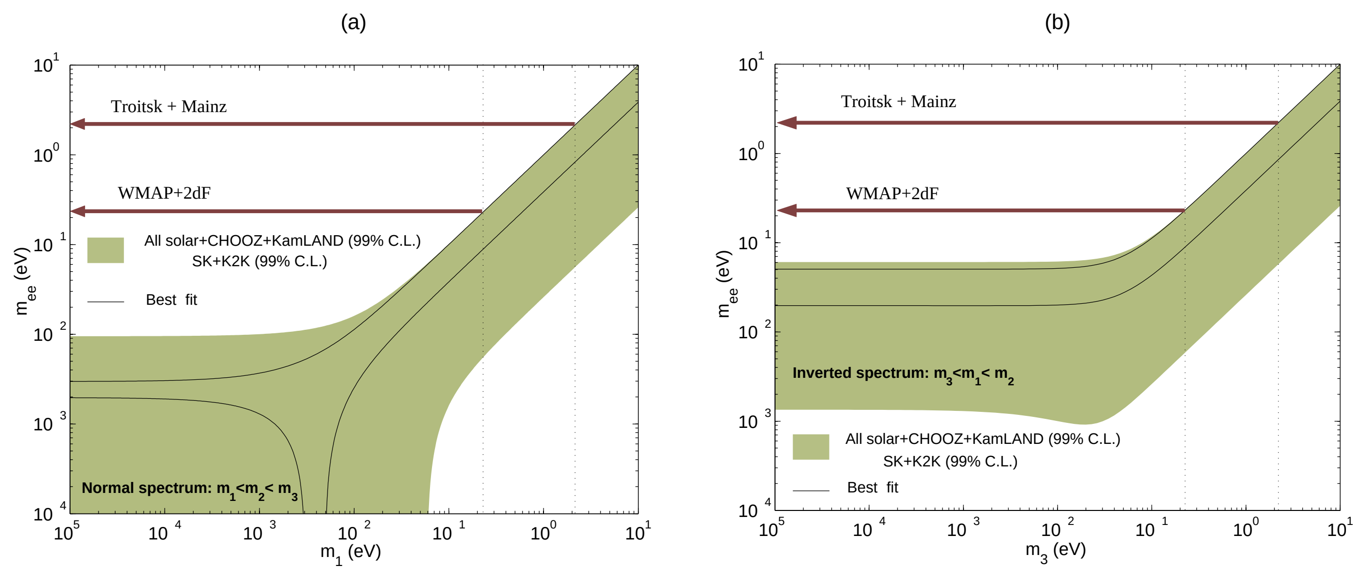

FIG. 1. Dependence of the effective Majorana neutrino mass parameter m_{ee} on the lightest neutrino mass for the (a) NNMS and (b) INMS. The shaded regions were obtained considering the SK+K2K results and the solar+CHOOZ+KamLAND at 99% C.L. given in Eqs. (6) and Table I, respectively, and the bound (8) on s_{13}^2. The solid lines correspond to the case where the best-fit values given in Eqs. (7), (9) and Table I are taken into account.

$$s_{12}^2 < \frac{1-\tan^2\theta_{13}}{2+\sqrt{\Delta m_{21}^2/|\Delta m_{32}^2|}},$$

$$s_{12}^2 < \frac{1}{2}(1-r_{32}-\tan^2\theta_{13}),$$

$$s_{12}^2 < \frac{1-\tan^2\theta_{13}(1+r_{32})}{2+r_{21}}, \quad (32)$$

respectively, with $r_{32}\equiv\sqrt{|\Delta m_{32}^2|}/m$ and $r_{21}\equiv\sqrt{\Delta m_{21}^2}/m$. Choosing $m \gtrsim 0.2$ eV [10] and $m \lesssim 2.2$ eV (Troitsk + Mainz) one has, considering the 99% C.L. bounds on the oscillation parameters, $\sqrt{\Delta m_{21}^2/|\Delta m_{32}^2|}\simeq(0.12-0.39)$, $r_{32}\simeq(0.02-0.31)$, and $r_{21}\simeq(0.3-7)\times10^{-2}$. Together with the inequalities (32) this leads to

$$s_{12}^2 \lesssim (0.39-0.47), \quad s_{12}^2 \lesssim (0.32-0.49), \quad s_{12}^2 \lesssim (0.45-0.50). \quad (33)$$

Taking into account that the WMAP result implies $m \simeq 0.23$ eV, the first and third relations above remain practically unchanged whereas the second one is modified to $s_{12}^2 \lesssim (0.34-0.39)$. This discussion shows that discriminating between the HI and IH or HI and QD leads to less restrictive conditions than the one needed for the discrimination between the IH and QD neutrino mass spectra. The possibility of determining the type of neutrino mass spectrum has been studied in detail in Ref. [37] where the uncertainties in the measured value of m_{ee} due to the imprecise knowledge of the nuclear matrix elements were taken into account. In particular, it has been shown that depending on the uncertainty factor, the above ranges of s_{12}^2 get modified, being even more restrictive.

At this point one may wonder if the results in favor of $(\beta\beta)_{0\nu}$ decay reported by the Heidelberg-Moscow Collaboration may help us in the determination of the neutrino mass spectrum type. Comparing the results plotted in Fig. 1 with the ranges (14)–(16) it becomes obvious that the HI spectrum is ruled out. Requiring $(m_{ee}^{\rm IH})_{\rm up} \lesssim m_{ee}^{\rm min} \simeq 0.05$, one can show that the IH spectrum is excluded if

$$|\Delta m_{32}^2| \lesssim \left[\frac{(m_{ee})_{\rm min}}{1-s_{13}^2}\right]^2 \simeq 2.8\times10^{-3}\ {\rm eV}^2. \quad (34)$$

Hence, considering the SK+K2K ranges for $|\Delta m_{32}^2|$ given in Eq. (7), one concludes that there is still a small window allowed for the inverted hierarchical neutrino mass spectrum. On the other hand, the bounds presented in Eqs. (14) and (16) select the QD spectrum as the only possible scenario for neutrino masses since in these cases $(m_{ee}^{\rm IH})_{\rm up} < m_{ee}^{\rm min}$.

IV. NEUTRINO MASS BOUNDS FROM $(\beta\beta)_{0\nu}$ DECAYS

In this section we analyze the impact of all the available data from $(\beta\beta)_{0\nu}$-decay experiments on the neutrino mass spectrum, taking into account the present neutrino oscillation results and the cosmological bound on neutrino masses.

Let us suppose that m_{ee} is found in the range $m_{ee}^{\rm min} \leqslant m_{ee} \leqslant m_{ee}^{\rm max}$. This, together with Eq. (19) leads to

$$m_{ee}^{\rm max} \geqslant m_{ee}^{\rm NS} \geqslant c_{13}^2 \left| m_1(1-2s_{12}^2) - s_{13}^2\sqrt{m_1^2+|\Delta m_{32}^2|}\right|, \quad (35)$$

and

$$m_{ee}^{\rm min} \leqslant m_{ee}^{\rm NS} \leqslant m_1 + \left(\sqrt{m_1^2+|\Delta m_{32}^2|}-m_1\right)s_{13}^2, \quad (36)$$

for the NNMS and $m_{ee}^{\rm max} \geqslant \sqrt{\Delta m_{21}^2}s_{12}^2$. The above inequalities allow one to find approximate expressions for the upper and lower bounds of the lightest neutrino mass m_1, which we will denote by $m_{\rm NS}^{\rm up}$ and $m_{\rm IS}^{\rm low}$, respectively. We obtain in this case

$$m_{\rm NS}^{\rm up} \simeq \frac{m_{ee}^{\rm max}\cos 2\theta_{12}c_{13}^2 + s_{13}^2\sqrt{(m_{ee}^{\rm max})^2+|\Delta m_{32}^2|(\cos^2 2\theta_{12}c_{13}^4 - s_{13}^4)}}{\cos^2 2\theta_{12}c_{13}^4 - s_{13}^4}, \quad (37)$$

$$m_{\rm NS}^{\rm low} \simeq \frac{m_{ee}^{\rm min}c_{13}^2 - s_{13}^2\sqrt{(m_{ee}^{\rm min})^2+|\Delta m_{32}^2|(1-2s_{13}^2)}}{1-2s_{13}^2}. \quad (38)$$

On the other hand, if $m_{ee}^{\rm min} \lesssim \sqrt{\Delta m_{21}^2}s_{12}^2$, the lower bound on m_1 does not exist. This can be understood taking into account that this condition is compatible with $m_1=0$. In particular, from Eq. (26) one has $(m_{ee}^{\rm HI})_{\rm low} \lesssim \sqrt{\Delta m_{21}^2}s_{12}^2$, which means that for $m_{ee}^{\rm min} \lesssim \sqrt{\Delta m_{21}^2}s_{12}^2 \simeq 0.0025$eV the value of m_1 is not bounded from below. The expressions for $m_{\rm NS}^{\rm up}$ and $m_{\rm NS}^{\rm low}$ can be further simplified if the terms proportional to $|\Delta m_{32}^2|$ are negligible or more specifically if

$$(m_{ee}^{\rm max})^2 \gg |\Delta m_{32}^2|(\cos^2 2\theta_{12}c_{13}^4 - s_{13}^4),$$

$$(m_{ee}^{\rm min})^2 \gg |\Delta m_{32}^2|(1-2s_{13}^2). \quad (39)$$

If this is the case, one gets

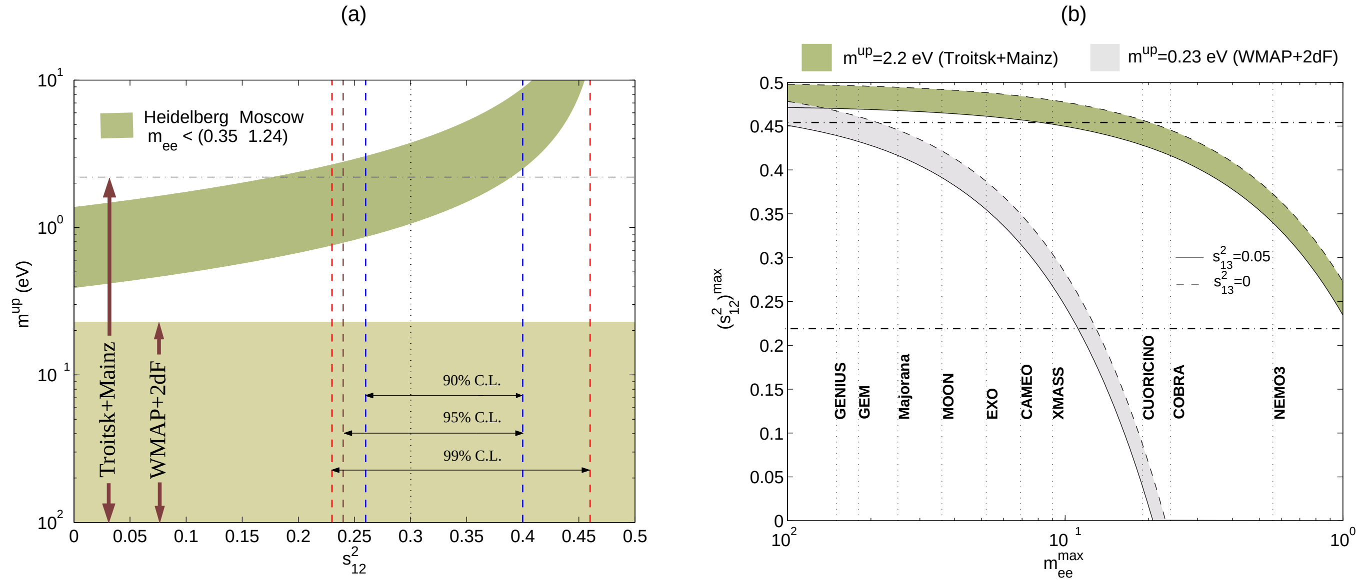

FIG. 2. (a) Allowed region for the upper bound on the lightest neutrino mass for both the NNMS and INMS when the Heidelberg-Moscow result $m_{ee} \lesssim (0.35-1.24)$ eV is taken into account. The horizontal dash-dotted line indicates the bound on $m^{\rm up}$ set by the Troitsk and Mainz experiments. The vertical dashed lines delimit the s_{12}^2 ranges given in Table I, and the vertical dotted line corresponds to the best-fit value $(s_{12}^2)_{\rm BF}=0.3$. The limit on $m^{\rm up}$ imposed by the WMAP+2dF results is also shown. (b) Values of $(s_{12}^2)_{\rm max}$ under which the bounds $m^{\rm up}=2.2$ eV (Trotsk+Mainz) and $m^{\rm up}=0.23$ eV (WMAP+2dF) can be improved for a given $m_{ee}^{\rm max}$. The horizontal dash-dotted lines delimit the presently allowed regions for s_{12}^2 at 99% C.L. and the vertical dotted lines correspond to the sensitivities of future $(\beta\beta)_{0\nu}$-decay experiments.

$$m_{\rm NS}^{\rm up} \simeq \frac{m_{ee}^{\rm max}}{\cos 2\theta_{12} c_{13}^2 - s_{13}^2} \tag{40}$$

and

$$m_{\rm NS}^{\rm low} \simeq m_{ee}^{\rm min}, \tag{41}$$

which show that $m_{\rm NS}^{\rm up}$ strongly depends on the θ_{12} angle while $m_{\rm NS}^{\rm low}$ is only affected by the value of $m_{ee}^{\rm min}$. Equivalent expressions can be found for the INMS considering that

$$m_{ee}^{\rm max} \geq m_{ee}^{\rm IS} \geq \left| c_{13}^2 (1-2s_{12}^2)\sqrt{m_3^2+|\Delta m_{32}^2|} - m_3 s_{13}^2 \right|, \tag{42}$$

$$m_{ee}^{\rm min} \leq m_{ee}^{\rm IS} \leq \sqrt{m_3^2+|\Delta m_{32}^2|}, \tag{43}$$

for $m_{ee}^{\rm max} \gg \sqrt{\Delta m_{21}^2} s_{12}^2$. These inequalities imply

$$m_{\rm IS}^{\rm up} \simeq \frac{m_{ee}^{\rm max} s_{13}^2 + \cos 2\theta_{12} c_{13}^2 \sqrt{(m_{ee}^{\rm max})^2 - |\Delta m_{32}^2|(\cos^2 2\theta_{12} c_{13}^4 - s_{13}^4)}}{\cos^2 2\theta_{12} c_{13}^4 - s_{13}^4}, \tag{44}$$

$$m_{\rm IS}^{\rm low} \simeq \sqrt{(m_{ee}^{\rm min})^2 - |\Delta m_{32}^2|}, \tag{45}$$

which reduce to the expressions given in Eq. (40) in the limit where the terms proportional to $|\Delta m_{32}^2|$ can be neglected. It can easily be seen from the above expression that for $m_{ee}^{\rm min} \lesssim \sqrt{|\Delta m_{32}^2|_{\rm max}} \simeq 0.06$ eV there is no lower bound on m_3 in the INMS case.

Let us now discuss the implications of the Heidelberg-Moscow experimental result given in Eq. (12) on the values of neutrino masses. In this case $m_{ee} \lesssim (0.35-1.24)$ eV and therefore only an upper bound on the neutrino masses can be set. In order to compute it we can use the first expression in Eq. (40) for both NNMS and INMS since for this range of $m_{ee}^{\rm max}$ the first inequality in Eq. (39) is verified. Therefore, we will denote both $m_{\rm NS}^{\rm up}$ and $m_{\rm IS}^{\rm up}$ just by $m^{\rm up}$. In Fig. 2(a) we show the allowed region for $m^{\rm up}$ as a function of s_{12}^2, obtained for $s_{13}^2=0.05$. Taking into account the Troitsk and Mainz bound on $m^{\rm up}$ and the 99% C.L. allowed range for s_{12}^2 one can conclude that $m_{ee}^{\rm max} \lesssim 1$ eV. This can be seen putting $s_{13}^2=0.05$ and $m_{\rm NS}^{\rm up}=2.2$ eV in Eq. (40). Using the best-fit value $s_{12}^2=0.3$ and the bound (9) one gets $m_{ee}^{\rm max} \lesssim 0.9$ eV.

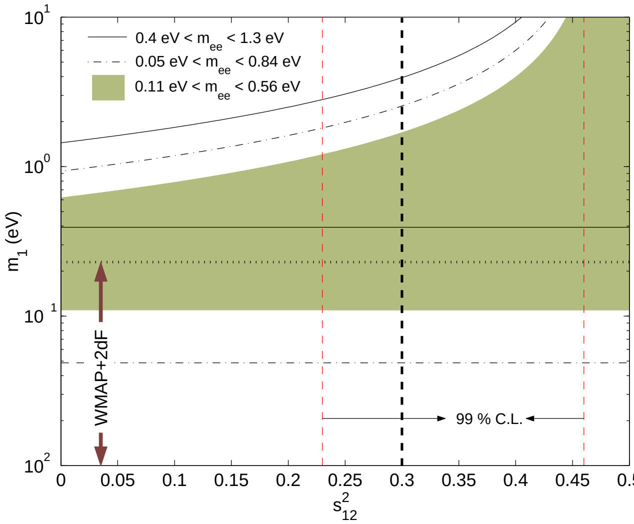

FIG. 3. Allowed ranges for the lightest neutrino mass m_1 in the NS case, when the latest results of the Heidelberg-Moscow experiment are considered. The shaded region corresponds to the interval of m_{ee} given in Eq. (14) and the dash-dotted and solid lines to the ranges (15) and (16), respectively.

Considering $m_{ee}\lesssim 0.35$ eV one can see that the value of $m^{\rm up}$ extracted from $(\beta\beta)_{0\nu}$ decay is smaller than the Troitsk and Mainz bound for $s_{12}^2\lesssim 0.4$. Alternatively, with $s_{12}^2=0.3$ and $s_{13}^2=0.01$ we get $m^{\rm up}\simeq 0.9$ eV for $m_{ee}^{\rm max}=0.35$ eV. From Fig. 2(a) we can also see that the bounds on $m^{\rm up}$ are always more conservative than the WMAP bound.

The sensitivity to m_{ee} is expected to be improved by several future experiments like the NEMO3 experiment [38], which has started to take data last year, and the CUORICINO project [39] to be operative this year. Other experimental setups are being planned or under construction with the goal of reaching sensitivities of $m_{ee}\simeq 0.01$ eV [40]. Taking this into consideration, it is interesting to analyze how these upcoming experiments can improve the present bounds on $m^{\rm up}$. It can be shown that in order for this to happen the following condition has to be verified:

$$s_{12}^2\lesssim \frac{1}{2}\left(1-\frac{m_{ee}^{\rm max}+\sqrt{|\Delta m_{32}^2|+(m^{\rm up})^2}s_{13}^2}{m^{\rm up}(1-s_{13}^2)}\right)$$
$$\simeq \frac{1}{2}\left(1-\frac{m_{ee}^{\rm max}/m^{\rm up}+s_{13}^2}{(1-s_{13}^2)}\right)\equiv (s_{12}^2)_{\rm max}, \qquad (46)$$

where the last expression corresponds to the limit $m_{\rm up} \gg \sqrt{|\Delta m_{32}^2|}$. In Fig 2(b) we show the dependence of $(s_{12}^2)_{\rm max}$ on the values of $m_{ee}^{\rm max}$ for $m^{\rm up}=2.2$ eV (Troitsk+Mainz) and $m^{\rm up}=0.23$ eV (WMAP+2dF). The vertical dotted lines indicate the sensitivities of some future $(\beta\beta)_{0\nu}$-decay experiments.

Turning now our attention to the evidence for $(\beta\beta)_{0\nu}$ decay reported by the Heidelberg-Moscow Collaboration and taking into account the m_{ee} ranges given in Eqs. (14)–(16), we can use Eqs. (35) and (36) to compute the values of $m_{\rm NS}^{\rm up}$ and $m_{\rm NS}^{\rm low}$ which are shown in Fig. 3. The horizontal lines correspond to the values of $m_{\rm NS}^{\rm low}$ for $m_{ee}^{\rm min}=0.05,0.11,0.4$ eV. We can see that the lower bounds on the lightest neutrino mass in each case obey the relation $m_{\rm NS}^{\rm low}\simeq m_{ee}^{\rm min}$, as already shown in Eq. (41). Consequently, the lower bound of m_{ee} given in Eq. (16) is in conflict with the WMAP result. Considering the results (14) and (15), together with the cosmological bound on neutrino masses, one has for the NNMS

$$0.05\ \ {\rm eV}\lesssim m_1\lesssim 0.23\ \ {\rm eV}. \qquad (47)$$

Regarding the inverted neutrino mass spectrum, we note that the above results obtained for $m^{\rm up}$ remain valid, as already discussed before. However, the situation changes for $m_{\rm IS}^{\rm low}$ in the sense that now the lower bound on the lightest neutrino mass may not exist. Considering $m_{ee}^{\rm min}=0.1,0.4$ eV we obtain $m_{\rm IS}^{\rm low}\simeq 0.1,0.4$ eV since for these values of $m_{ee}^{\rm min}$ we can neglect $|\Delta m_{32}^2|$ in Eq. (45). On the other hand, from Eq. (45) and the SK+K2K allowed ranges for Δm_{32}^2 given in Eq. (6) one has $m_{\rm IS}^{\rm low}\lesssim 1.1\times 10^{-3}$ eV for $m_{ee}^{\rm min}=0.05$.

V. CONCLUSIONS

In this paper we have focused on the implications of the available data from $(\beta\beta)_{0\nu}$-decay experiments in the light of the latest neutrino oscillation and WMAP results. We have briefly commented on the possible occurrence of cancellations in the effective Majorana neutrino mass parameter taking into account the allowed ranges for the neutrino oscillation parameters at 99% C.L. given in Refs. [23,24]. We conclude that cancellations are only possible for the HI neutrino mass spectrum. However, this is no longer true if one relies on the best-fit values given in Eqs. (7), (9) and Table I since in this case the condition $m_{ee}\simeq 0$ cannot be fulfilled for any of the neutrino mass schemes considered here. As for the extraction of neutrino mass bounds from the presently available $(\beta\beta)_{0\nu}$-decay data we have seen that, while the establishment of an upper bound on the mass of the lightest neutrino strongly depends on the value of s_{12}^2, the lower bound is only sensitive to $m_{ee}^{\rm min}$, for the NNMS. In particular, the range 0.05 eV$\lesssim m_1\lesssim 0.23$ eV is obtained if one considers the intervals for m_{ee} given in Eqs. (14)–(16) together with the WMAP bound. In the INMS case, the knowledge of $|\Delta m_{32}^2|$ may be relevant for the lower bound on the lightest neutrino mass since in this case the lower bound on m_3 may not exist if a $\pm 50\%$ uncertainty in the nuclear matrix elements is taken into account. We have also discussed how the upcoming $(\beta\beta)_{0\nu}$-decay experiments may improve the WMAP result and concluded that, again, the key point relies on the measurement of s_{12}^2 which is expected to be improved by the future KamLAND and BOREXINO [41] data. Nevertheless, a sensitivity of $m_{ee}\simeq 0.09$ eV is required if the best-fit value $s_{12}^2=0.3$ is considered. Finally, we would like to remark that the cosmological bound on the absolute neutrino mass scale has important consequences for the future prospects in the study of $(\beta\beta)_{0\nu}$-decay searches. On the other hand, the bound $m_i<0.23$ eV is not encouraging for the KATRIN experiment which will be sensitive to m_i

$\gtrsim$0.35 eV. In any case, one should remember that the kind of analysis presented here is based on the assumption that $(\beta\beta)_{0\nu}$ decays are mediated by the exchange of massive Majorana neutrinos. However, other mechanisms can give rise to these processes [42]. This possibility opens a new challenging question, namely, how one can identify the physics behind $(\beta\beta)_{0\nu}$ decays.

ACKNOWLEDGMENTS

The author thanks G. C. Branco and R. González Felipe for the careful reading of the manuscript and valuable suggestions. This work was supported by *Fundação para a Ciência e a Tecnologia* (FCT, Portugal) under the Grant No. PRAXISXXI/BD/18219/98.

[1] KamLAND Collaboration, K. Eguchi *et al.*, Phys. Rev. Lett. **90**, 021802 (2003).

[2] Super-Kamiokande Collaboration, Y. Fukuda *et al.*, Phys. Rev. Lett. **81**, 1562 (1998).

[3] K2K Collaboration, M.H. Ahn *et al.*, Phys. Rev. Lett. **90**, 041801 (2003).

[4] The possibilities of distinguishing Dirac from Majorana neutrinos have been discussed in M. Zralek, Acta Phys. Pol. B **28**, 2225 (1997).

[5] M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity*, edited by P. van Nieuwenhuizen and D. Freedman (North-Holland, Amsterdam, 1979), p. 315; S.L. Glashow, in *Quarks and Leptons*, Cargèse, 1979, edited by M. Levy *et al.* (Plenum, New York, 1980), p. 707; T. Yanagida, in *Proceedings of the Workshop on the Unified Theory and Baryon Number in the Universe*, edited by O. Sawada and A. Sugamoto (KEK, Tsukuba, 1979), p. 95, KEK Report No. 79-18; R.N. Mohapatra and G. Senjanović, Phys. Rev. Lett. **44**, 912 (1980).

[6] For a complete review on the status and prospects for measuring neutrino masses see, e.g., S.M. Bilenky, C. Giunti, J.A. Grifols, and E. Masso, Phys. Rep. **379**, 69 (2003).

[7] J. Bonn *et al.*, Prog. Part. Nucl. Phys. **48**, 133 (2002).

[8] V.M. Lobashev *et al.*, Nucl. Phys. (Proc. Suppl.) **91**, 280 (2001).

[9] KATRIN Collaboration, A. Osipowicz *et al.*, hep-ex/0109033.

[10] For a review on several aspects related with $(\beta\beta)_{0\nu}$ decay see, e.g., S.M. Bilenky, S. Pascoli, and S.T. Petcov, Phys. Rev. D **64**, 053010 (2001), and references therein.

[11] J. Schechter and J.W. Valle, Phys. Rev. D **25**, 2951 (1982); S.T. Petcov and A.Y. Smirnov, Phys. Lett. B **322**, 109 (1994); M. Hirsch, H.V. Klapdor-Kleingrothaus, and S.G. Kovalenko, Phys. Rev. D **53**, 1329 (1996); R.N. Mohapatra, hep-ph/9507234; S.M. Bilenky, A. Bottino, C. Giunti, and C.W. Kim, Phys. Rev. D **54**, 1881 (1996); F. Vissani, J. High Energy Phys. **06**, 022 (1999); H.V. Klapdor-Kleingrothaus, H. Pas, and A.Y. Smirnov, Phys. Rev. D **63**, 073005 (2001); K. Matsuda, T. Kikuchi, T. Fukuyama, and H. Nishiura, Mod. Phys. Lett. A **17**, 2597 (2002); H. Minakata and H. Sugiyama, hep-ph/0212240; G.C. Branco, R. Gonzalez Felipe, F.R. Joaquim, and T. Yanagida, Phys. Lett. B **562**, 265 (2003); G. Bhattacharyya, H. Pas, L.G. Song, and T.J. Weiler, *ibid.* **564**, 175 (2003); K. Matsuda, T. Fukuyama, and H. Nishiura, hep-ph/0302254; K. Cheung, Phys. Lett. B **562**, 97 (2003); H.V. Klapdor-Kleingrothaus and U. Sarkar, hep-ph/0304032; A. Abada and G. Bhattacharyya, hep-ph/0304159.

[12] O. Elgaroy *et al.*, Phys. Rev. Lett. **89**, 061301 (2002).

[13] C.L. Bennett *et al.*, astro-ph/0302207; D.N. Spergel *et al.*, astro-ph/0302209.

[14] E. W. Kolb and M. S. Turner, *The Early Universe* (Addison-Wesley, Reading, 1990).

[15] A. Pierce and H. Murayama, hep-ph/0302131.

[16] LSND Collaboration, A. Aguilar *et al.*, Phys. Rev. D **64**, 112007 (2001).

[17] M. Maltoni, T. Schwetz, M.A. Tortola, and J.W.F. Valle, Nucl. Phys. B (Proc. Suppl.) **114**, 203 (2003).

[18] KARMEN Collaboration, W. Kretschmer, Acta Phys. Pol. B **33**, 1775 (2002).

[19] MiniBooNE Collaboration, I. Stancu *et al.*, FERMILAB Report No. TM-2207.

[20] S.M. Bilenky, J. Hosek, and S.T. Petcov, Phys. Lett. **94B**, 495 (1980); I.Y. Kobzarev, B.V. Martemyanov, L.B. Okun, and M.G. Shchepkin, Yad. Fiz. **32**, 1590 (1980) [Sov. J. Nucl. Phys. **32**, 823 (1980)]; M. Doi, T. Kotani, H. Nishiura, K. Okuda, and E. Takasugi, Phys. Lett. **102B**, 323 (1981); L. Wolfenstein, *ibid.* **107B**, 77 (1981); B. Kayser and A.S. Goldhaber, Phys. Rev. D **28**, 2341 (1983).

[21] See, e.g., M. Lindner, hep-ph/0209083, and references therein.

[22] W. Rodejohann, hep-ph/0203214; V. Barger, S.L. Glashow, P. Langacker, and D. Marfatia, Phys. Lett. B **540**, 247 (2002); S. Pascoli, S.T. Petcov, and W. Rodejohann, *ibid.* **549**, 177 (2002).

[23] G.L. Fogli, E. Lisi, A. Marrone, and D. Montanino, Phys. Rev. D **67**, 093006 (2003).

[24] G.L. Fogli, E. Lisi, A. Marrone, D. Montanino, A. Palazzo, and A.M. Rotunno, Phys. Rev. D **67**, 073002 (2003).

[25] M. Maltoni, T. Schwetz, and J.W. Valle, Phys. Rev. D **67**, 093003 (2003); J.N. Bahcall, M.C. Gonzalez-Garcia, and C. Peña-Garay, J. High Energy Phys. **02**, 009 (2003); P.C. de Holanda and A.Y. Smirnov, J. Cosmol. Astropart. Phys. **02**, 001 (2003); A.B. Balantekin and H. Yuksel, J. Phys. G **29**, 665 (2003).

[26] CHOOZ Collaboration, M. Apollonio *et al.*, Phys. Lett. B **466**, 415 (1999).

[27] F. Boehm *et al.*, Phys. Rev. Lett. **84**, 3764 (2000).

[28] S. Pascoli, S.T. Petcov, and L. Wolfenstein, Phys. Lett. B **524**, 319 (2002).

[29] P. Lipari, Phys. Rev. D **61**, 113004 (2000).

[30] H.V. Klapdor-Kleingrothaus, Nucl. Phys. B (Proc. Suppl.) **100**, 309 (2001).

[31] C.E. Aalseth *et al.*, Yad. Fiz. **63**, 1341 (2000) [Phys. At. Nucl. **63**, 1268 (2000)].

[32] H.V. Klapdor-Kleingrothaus, A. Dietz, H.L. Harney, and I.V. Krivosheina, Mod. Phys. Lett. A **16**, 2409 (2001).

[33] Particle Data Group, K. Hagiwara *et al.*, Phys. Rev. D **66**, 010001 (2002), pp. 389–392.

[34] H.V. Klapdor-Kleingrothaus, hep-ph/0302237; Heidelberg-

Moscow Collaboration, H.V. Klapdor-Kleingrothaus, hep-ph/0303217.

[35] F. Feruglio, A. Strumia, and F. Vissani, Nucl. Phys. **B637**, 345 (2002); C.E. Aalseth *et al.*, Mod. Phys. Lett. A **17**, 1475 (2002); Y.G. Zdesenko, F.A. Danevich, and V.I. Tretyak, Phys. Lett. B **546**, 206 (2002).

[36] H.V. Klapdor-Kleingrothaus and U. Sarkar, Mod. Phys. Lett. A **16**, 2469 (2001); E. Ma, *ibid.* **17**, 289 (2002); V. Barger, S.L. Glashow, D. Marfatia, and K. Whisnant, Phys. Lett. B **532**, 15 (2002); H. Minakata and H. Sugiyama, *ibid.* **532**, 275 (2002); G.C. Branco, R. González Felipe, F.R. Joaquim, and M.N. Rebelo, Nucl. Phys. **B640**, 202 (2002); Z.Z. Xing, Phys. Rev. D **65**, 077302 (2002); S. Pakvasa and P. Roy, Phys. Lett. B **535**, 181 (2002); G. Barenboim, J.F. Beacom, L. Borissov, and B. Kayser, *ibid.* **537**, 227 (2002); M. Frigerio and A.Y. Smirnov, Nucl. Phys. **B640**, 233 (2002).

[37] S. Pascoli, S.T. Petcov, and W. Rodejohann, Phys. Lett. B **558**, 141 (2003).

[38] NEMO Collaboration, C. Marquet, Nucl. Phys. B (Proc. Suppl.) **87**, 298 (2000).

[39] E. Fiorini, Phys. Rep. **307**, 309 (1998).

[40] H.V. Klapdor-Kleingrothaus, J. Hellmig, and M. Hirsch, J. Phys. G **24**, 483 (1998); H. Ejiri, J. Engel, R. Hazama, P. Krastev, N. Kudomi, and R.G. Robertson, Phys. Rev. Lett. **85**, 2917 (2000); G. Bellini *et al.*, Eur. Phys. J. C **19**, 43 (2001); K. Zuber, Phys. Lett. B **519**, 1 (2001); Majorana Collaboration, C.E. Aalseth *et al.*, hep-ex/0201021; Y.G. Zdesenko, O.A. Ponkratenko, and V.I. Tretyak, J. Phys. G **27**, 2129 (2001); EXO Collaboration, K. Wamba, hep-ph/0210186.

[41] Borexino Collaboration, T. Shutt, Nucl. Phys. B (Proc. Suppl.) **110**, 323 (2002).

[42] For a review see, e.g., J.D. Vergados, Phys. Rep. **361**, 1 (2002), and references therein.