Direct search for right-handed neutrinos and neutrinoless double beta decay

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We consider an extension of the Standard Model by two right-handed neutrinos, especially with masses lighter than the charged \(K\) meson. This simple model can realize the seesaw mechanism for neutrino masses and also baryogenesis by flavor oscillations of right-handed neutrinos. We summarize the constraints on right-handed neutrinos from direct searches as well as big bang nucleosynthesis. It is then found that the possible range for the quasi-degenerate mass of right-handed neutrinos is \(M_N \geq 163\) MeV for the normal hierarchy of neutrino masses, while \(M_N = 188–269\) MeV and \(M_N \geq 285\) MeV for the inverted hierarchy case. Furthermore, we find in the latter case that the possible value of the Majorana phase is restricted for \(M_N = 188–350\) MeV, which leads to the fact that the rate of neutrinoless double beta decay is also limited.

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1. Introduction

Various oscillation experiments have revealed non-zero masses of neutrinos. Observation shows that there exist two mass scales of neutrinos, the differences of mass-squared \(\Delta m^2_{\text{atm}} \simeq 2.43 \times 10^{-3}\) eV\(^2\) and \(\Delta m^2_{\text{sol}} \simeq 7.54 \times 10^{-5}\) eV\(^2\) (see, e.g., Ref. [1]), related to the so-called atmospheric and solar neutrinos, respectively. In the (canonical) Standard Model neutrinos are exactly massless, and new physics beyond the Standard Model is indicated. The crucial questions are then (i) what is the origin of neutrino masses? and (ii) how do we verify it experimentally?

One of the simplest and most attractive ways to generate neutrino masses is to introduce right-handed neutrinos \(\nu_R\) into the Standard Model. In this case neutrinos can obtain Dirac masses as quarks and charged leptons. Furthermore, since these neutral fermions are singlets under the gauge group of the Standard Model, the Majorana masses of right-handed neutrinos are also allowed. Notice that we should require at least two right-handed neutrinos in order to explain \(\Delta m^2_{\text{atm}}\) and \(\Delta m^2_{\text{sol}}\). When the Majorana masses are much heavier than the Dirac ones, the smallness of the neutrino masses can be explained by the seesaw mechanism [2–7]. The mass eigenstates are then separated into three lighter and two heavier ones. The former ones are responsible for oscillation phenomena while the latter ones have not been confirmed by experiments.

In this framework all the neutrinos are Majorana particles, and the (total) lepton number is violated, which may be tested by the neutrinoless double beta (0\(\nu\)2\(\beta\)) decay. In particular, when only two right-handed neutrinos are present, the lightest (active) neutrino becomes massless and the predicted range
for the rate of $0\nu2\beta$ decay is very limited. Interestingly, future $0\nu2\beta$ experiments may probe such a range in the inverted hierarchy of neutrino masses. (See, e.g., Ref. [8] for a review.)

On the other hand, right-handed neutrinos can also play important roles in generating the baryon asymmetry of the universe. The concrete scenario for this baryogenesis strongly depends on the mass scales of right-handed neutrinos. In the canonical leptogenesis scenario [9] the lightest right-handed neutrino should be heavier than $\mathcal{O}(10^9)$ GeV (see, e.g., Ref. [10], and Refs. [11,12] for reviews) when they have hierarchical masses. If $\nu_R$ are produced non-thermally at the reheating of the inflation, the required mass can be as small as $\mathcal{O}(10^6)$ GeV [13–16]. Moreover, the resonant leptogenesis [17] by quasi-degenerate right-handed neutrinos is possible even for smaller $\nu_R$ masses. Interestingly, the successful scenario of baryogenesis can be realized even if the masses are smaller than the electroweak scale [18,19]. In these extensions of the Standard Model, therefore, a detailed examination of right-handed neutrinos is crucial in order to elucidate the mechanism to generate neutrino masses as well as the cosmic baryon asymmetry. For this purpose, the scenario with lighter $\nu_R$ is more promising.

In this paper, we consider the Standard Model with two right-handed neutrinos, which is probably the minimal extension to explain the neutrino oscillation results and the baryon asymmetry. In particular, we assume that the masses of these neutral leptons are smaller than the $K^\pm$ meson. Such light $\nu_R$ are good targets for search experiments using $K^\pm$ and $\pi^\pm$ decays (see Refs. [20–24]). It should be noted that, even when the $\nu_R$ are lighter than $m_{K^\pm}$, $\Delta m^2_{\text{atm}}$ and $\Delta m^2_{\text{sol}}$ from the oscillation experiments can be explained via the seesaw mechanism by requiring that the Yukawa coupling constants of neutrinos are sufficiently small, and also that sufficient baryon asymmetry can be generated via the mechanism [18,19] by requiring that the $\nu_R$ are sufficiently degenerate.

Under the above situation, we shall summarize the constraints on such light $\nu_R$ from direct search experiments as well as cosmological observations, and then identify the allowed region in the parameter space of the model. In particular, possible values of $\nu_R$ masses are presented, as pointed out in Refs. [21,25].

In addition, we shall discuss the implications of the allowed parameter space. In particular, we find that the Majorana phase (which is the one of the CP-violating parameters in the lepton sector) is restricted from the search and cosmological constraints when the inverted hierarchy of neutrino masses is considered. It will then be discussed that this will have an important impact on future $0\nu2\beta$ experiments.

### 2. Extension by two right-handed neutrinos

Let us start with the framework of the present analysis. We consider the Standard Model extended by two right-handed neutrinos $\nu_R$\(^1\) together with

$$
\mathcal{L} = i \bar{\nu}_R \gamma^\mu \partial_\mu \nu_R - \left( F \bar{L} \Phi \nu_R + \frac{M_M}{2} \bar{\nu}_R \nu_R + \text{h.c.} \right),
$$

where $\Phi$ and $L_\alpha = (e_L, \nu_L)^T$ are Higgs and lepton doublets, respectively. Here and hereafter, the indices of flavor are implicit unless otherwise mentioned. The $3 \times 2$ Yukawa matrix of neutrinos is denoted by $F$ and the $2 \times 2$ matrix of Majorana masses by $M_M$. Notice that we choose the basis in which the mass matrix of charged leptons and $M_M$ are diagonal.

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\(^1\)By adding the keV right-handed neutrino, it can play the role of cosmic dark matter [26]. The results in the present analysis can be applied to the neutrino Minimal Standard Model ($\nu$MSM) [19,26] if the dark matter physics requires no limitation on the parameters of the two right-handed neutrinos considered.
We assume that the Dirac masses of neutrinos \( M_D = F\langle \Phi \rangle \) are much smaller than the Majorana ones \( M_M \) for the seesaw mechanism. In this case, the mass matrix of light neutrinos participating in the observed flavor oscillation is given by \( M_\nu = -M_D M_M^{-1} M_D^T \). By using this relation, we can parametrize, without loss of generality, the Yukawa matrix \( F \) as follows [27, 28]:

\[
F = \frac{i}{\langle \Phi \rangle} U (M_\nu^{\text{diag}})^{1/2} \Omega (M_M)^{1/2},
\]

where \( M_\nu^{\text{diag}} = \text{diag}(m_1, m_2, m_3) \) with masses of light neutrinos \( m_i \), and \( M_\nu^{\text{diag}} = U^\dagger M_\nu U^* \). This parametrization is the same as that in Ref. [29]. The mixing matrix of light neutrinos, called the Pontecorvo–Maki–Nakagawa–Sakata (PMNS) matrix [30, 31], is written as

\[
U = \begin{pmatrix}
c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta}
\end{pmatrix}
\begin{pmatrix}
-c_{23}s_{12} - s_{23}s_{13}e^{i\delta} & c_{23}c_{12} - s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\
-s_{23}s_{12} - c_{23}s_{13}e^{i\delta} & -c_{23}c_{12} - c_{23}s_{13}e^{i\delta} & c_{23}c_{13}
\end{pmatrix} \times \text{diag}(1, \ e^{i\eta}, 1),
\]

with \( s_{ij} = \sin \theta_{ij} \) and \( c_{ij} = \cos \theta_{ij} \).

It is then found that the flavor neutrino \( \nu_L \) is written in terms of mass eigenstates \( \nu \) and \( N \) as

\[
\nu_L = U \nu + \Theta N^C,
\]

where each heavy neutrino \( N \) almost coincides with \( \nu_R \), \( N \sim \nu_R \). The mixing elements of \( N \) are given by \( \Theta = M_D / M_M \).

In the case under consideration, there is one Majorana phase \( \eta \) in addition to the Dirac phase \( \delta \), and the lightest neutrino becomes exactly massless. Thus, the possible patterns of neutrino masses are (i) the normal hierarchy (NH) with \( m_3 > m_2 > m_1 = 0 \), and (ii) the inverted hierarchy (IH) with \( m_2 > m_1 > m_3 = 0 \). We express the Majorana mass matrix as \( M_M = \text{diag}(M_N - \Delta M / 2, M_N + \Delta M / 2) \) and the \( 3 \times 2 \) matrix \( \Omega \) as

\[
\Omega = \begin{pmatrix}
0 & 0 \\
\cos \omega & -\sin \omega \\
\xi \sin \omega & \xi \cos \omega
\end{pmatrix},
\]

in the NH case, and

\[
\Omega = \begin{pmatrix}
\cos \omega & -\sin \omega \\
\xi \sin \omega & \xi \cos \omega \\
0 & 0
\end{pmatrix},
\]

in the IH case. Here \( \omega \) is an arbitrary complex number and \( \xi = \pm 1 \) is the sign parameter. Here we apply the convention \( \text{Im} \omega \geq 0, \xi = \pm 1 \) and \( \eta = 0 - \pi \).

The observational data for the mixing angles are \( s_{12}^2 = 0.307^{+0.052}_{-0.048}, s_{13}^2 = 0.386^{+0.251}_{-0.055} (0.392^{+0.271}_{-0.057}), \) and \( s_{13}^2 = 0.0241^{+0.0072}_{-0.0072} (0.0244^{+0.0071}_{-0.0073}) \), respectively, and the masses are \( \delta m^2 = m_2^2 - m_1^2 = (7.54^{+0.64}_{-0.55}) \times 10^{-5} \text{ eV}^2 \) and \( \Delta m^2 = |m_3^2 - (m_1^2 + m_2^2)|/2 = (2.43^{+0.19}_{-0.24}) \times 10^{-3} \text{ eV}^2 \) for the NH (IH) case (at the 3\sigma level) [1]. Hereafter, we shall adopt the central values of these observables. It should be stressed that these observational data from the oscillation experiments can be reproduced, being independent of the parameters of right-handed neutrinos, i.e., \( M_N, \Delta M, \omega, \) and \( \xi \). In practice, we shall consider the case when \( M_N < m_{K^\pm} \) and \( \Delta M \ll M_N \) (see the discussions below).

It is interesting to note that the two right-handed neutrinos with \( M_N < m_{K^\pm} \) introduced above can be responsible for the baryon asymmetry of the universe by invoking the mechanism proposed...
in Ref. [19]. By originating the CP violations in the flavor oscillation as well as the production/destruction processes of right-handed neutrinos, asymmetries of left-handed leptons are generated, which are partially converted into baryon asymmetry due to the rapid sphaleron transition [32]. See the analysis in Refs. [19,25,33–37].

We perform a numerical study of the generation of baryon asymmetry\(^2\) and identify the parameter region in which the observed value (the baryon to entropy ratio) \(n_B/s = 8.8 \times 10^{-11}\) [38] can be explained. Here, we use the central values for the neutrino oscillation parameters and vary over all possible values for the unknown parameters.

The region of successful baryogenesis in the \(M_N-\Delta M\) plane is shown in Fig. 1. It is then found that sufficient baryon asymmetry can be generated if the masses of right-handed neutrinos are

\[
M_N \geq \begin{cases} 
2.1 \text{ MeV} & \text{for the NH case} \\
0.7 \text{ MeV} & \text{for the IH case}
\end{cases}
\]

(7)

It is also found that the mass difference of the two right-handed neutrinos should be \(\Delta M \ll M_N\). In these regions, we can explain the neutrino masses in the oscillation experiments and the baryon asymmetry at the same time by using the two right-handed neutrinos only. Such a mass bound has already been obtained, as shown in Fig. 1 of Ref. [25]. The region that we have obtained is wider than theirs; in particular, the lower bound on \(M_N\) in the NH case is smaller by a factor of five. This is mainly because different values for \(\theta_{ij}\) and \(\Delta m_{ij}\) are chosen.

3. Constraints from direct search and cosmology

Heavy neutrinos \((N \simeq \nu_R)\) have weak interaction due to the mixing induced by the seesaw mechanism, whose strength is suppressed by the mixing elements \(\Theta\) compared with ordinary neutrinos. Thereby, they can be produced by meson decays (e.g. \(\pi \to Ne, K \to N\mu\)), and can decay to

\(^2\)To estimate BAU, we solve numerically the kinetic equations (5.2) and (5.3) in Ref. [37] by neglecting the momentum dependence in the density matrix of the \(\nu_R\), for simplicity. The details are given in Ref. [37].
Fig. 2. The allowed region of $M_N$ and $\tau_N$ in the NH (left panel) and IH (right panel) cases, respectively. The (red) solid line shows the lower bound on $\tau_N$ derived from the search experiments, while the (blue) dashed line shows the upper bound on $\tau_N$ by BBN.

Standard Model particles (e.g. $N \rightarrow \nu\nu\bar{\nu}$, $N \rightarrow \nu e^+ e^-$, $N \rightarrow e\pi$). Using such processes, various experiments have been conducted to search for heavy neutrinos directly. Since such neutrinos have not been discovered, upper bounds have been imposed on the mixing elements. In the mass range under discussion, the beam dump experiment, the PS191 experiment [39,40], has placed the strongest bounds on $\Theta_{1\alpha}$. In this experiment heavy neutrinos are produced by $\pi$ and/or $K$ decays and charged particles from $N$ decays in the far detector are searched for as signal events. These processes are induced as the $\Theta^4$ effect, therefore the experiment sets the upper bounds of the mixing elements in the form $|\Theta_{\alpha I}|^2 (a|\Theta_{e I}|^2 + b|\Theta_{\mu I}|^2 + c|\Theta_{\tau I}|^2)$, where $\alpha$ and $I$ are the flavor indices of left-handed neutrinos and heavy neutrinos, respectively. $a$, $b$, and $c$ are coefficients depending on the mass and decay channel. In Refs. [39,40], such bounds were derived by assuming that the heavy neutrino is a Dirac particle and that it has only charged current interaction. As pointed out in Ref. [20] (see also Ref. [41]), however, the bounds in the scenario under consideration must be evaluated with two Majorana (heavy) neutrinos including the neutral current interaction.

On the other hand, heavy neutrinos are also restricted from cosmological observation. To maintain the success of Big Bang nucleosynthesis (BBN), the lifetime of $N$, $\tau_N$, is required to be shorter than 0.1 sec for $M_N > m_\pi$, and $\tau_N/\sec < t_1 (M_N/\text{MeV})^\beta + t_2$ with $t_1 = 128.7$, $t_2 = 0.04179$, and $\beta = -1.828$ for $M_N \leq m_\pi$ [42]. Note that the lifetime bound for $M_N \leq m_\pi$ was discussed recently in Ref. [43]. Here we apply the bounds in Ref. [42] to make a conservative analysis.

From these constraints we evaluate the possible region of heavy neutrinos, as performed in previous work [21,33,41]. The direct search experiments set upper bounds on $|\Theta|$, which result in the lower bound on $\tau_N$. On the other hand, BBN sets the upper bound on $\tau_N$. Thus, we can identify the possible region of $M_N$ and $\tau_N$. Figure 2 shows the results of this work. It is found that the following $N$ masses are allowed:

$$M_N \geq 163 \text{ MeV} \quad \text{for the NH case},$$

$$M_N = 188-269 \text{ MeV and } M_N \geq 285 \text{ MeV} \quad \text{for the IH case}. \quad (8)$$

It should be noted that a specific element of the mixing matrix $\Theta$ can be very suppressed compared with other elements by choosing the parameters carefully [29,41]. The possible component of this suppression is $|\Theta_{e I}|^2$ in the NH case, and $|\Theta_{\mu I}|^2$ or $|\Theta_{\tau I}|^2$ in the IH case. For the current data from
The allowed region of $M_N$ and $\eta$ in the IH case. The shaded region is excluded by the bounds from the search experiments and BBN. The left side of the (red) solid line is excluded by using the search mode $K \rightarrow eN \rightarrow e(\pi\pi)$, while the region inside the (blue) dashed line is excluded by $K \rightarrow \mu N \rightarrow \mu(\pi\pi)$. To the left of the (green) dotted line only the sign parameter $\xi = +1$ is allowed, while both possibilities $\xi = \pm 1$ are allowed on the opposite side.

Before closing this section, it should be mentioned that we have numerically confirmed that sufficient baryon asymmetry can be generated in all the allowed regions shown in Fig. 2. As a result, we have found the parameter space of heavy neutrinos, which can explain the neutrino masses as well as the observed values of BAU without conflicting with the experimental and cosmological constraints.

4. Implications for neutrinoless double beta decay

In the allowed parameter space, we find a distinctive feature in the IH case; namely, the possible value of the Majorana phase $\eta$ is restricted by the present search experiments and BBN, as shown in Fig. 3. In the figure, the shaded region is excluded, and the (red) solid and (blue) dashed lines are derived by the experimental bound from the search mode $K \rightarrow eN \rightarrow e(\pi\pi)$ and that from the mode $K \rightarrow \mu N \rightarrow \mu(\pi\pi)$ in PS191, respectively. The former essentially sets the upper bound on $|\Theta_e|^2$, while the latter sets the bound on $|\Theta_\mu|^2$. To the left of the (green) dotted line the sign parameter $\xi$ is allowed to be only $+1$, while both signs $\xi = \pm 1$ are allowed to the right. It can be seen that, when the mass is $M_N = 188–269$ MeV, the Majorana phase close to $\pi/2$ is possible and $\eta \approx 0$ and $\pi$ are forbidden. When the mass becomes heavy, i.e. $M_N = 285–350$ MeV, the possible range of $\eta$ is changed, and $\eta \approx 0$ and $\pi$ as well as $\eta \approx \pi/2$ are disfavored. If the $N$ are heavy enough, i.e. $M_N > 350$ MeV, the full range of $\eta$ is consistent with the constraints.

oscillation experiments [1], the suppression condition in the NH case cannot be satisfied exactly. However, $|\Theta_e|^2$ can be smaller by a few orders of magnitude compared with $|\Theta_\mu|^2$ and $|\Theta_\tau|^2$ near the suppression point. In the IH case the suppression condition can be satisfied even for the current data. The suppression is, however, forbidden to escape the constraints from search experiments and BBN for $M_N \leq 350$ MeV. On the other hand, the experimental bounds on $|\Theta_e|^2$ are stronger than those on $|\Theta_\mu|^2$ in almost all mass ranges of interest. Due to the feature of mixing elements $|\Theta|^2$ and the present experimental bounds, the allowed range of $M_N$ in the NH case is wider than that in the IH case, as shown in Fig. 2.
To understand these outcomes, we present the formulæ of the mixing elements for $X_\omega \equiv e^{i\omega_\nu} \gg 1$ [29]:

$$|\Theta_e|^2 \simeq 1.2 \times 10^{-8} \left( \frac{\text{MeV}}{M_N} \right) [1 - 0.93 \xi \sin \eta]X_{\omega}^2,$$

$$|\Theta_\mu|^2 \simeq 7.6 \times 10^{-9} \left( \frac{\text{MeV}}{M_N} \right) [1 + 0.90 \xi \sin \eta - 0.25 \xi \cos \eta \sin \delta + 0.09 \xi \sin \eta \cos \delta]X_{\omega}^2,$$

$$|\Theta_\tau|^2 \simeq 5.0 \times 10^{-9} \left( \frac{\text{MeV}}{M_N} \right) [1 + 0.86 \xi \sin \eta + 0.38 \xi \cos \eta \sin \delta - 0.14 \xi \sin \eta \cos \delta]X_{\omega}^2.$$

Notice that the typical value of $X_\omega$ is $\mathcal{O}(10)$ in the allowed region of Fig. 2. We can see from Eq. (9) that the mixing elements depend crucially on the combination $\xi \sin \eta$. For $M_N \lesssim 332$ MeV, the constraint on $|\Theta_e|^2$ favors $\xi \sin \eta \sim 1$; hence, $\xi = -1$ and $\eta \sim 0, \pi$ are excluded. In addition, when $M_N = 269 - 350$ MeV, the constraint on $|\Theta_\mu|^2$ favors $\xi \sin \eta \sim -1$, which excludes the possibility $\eta \sim \pi/2$.

We have found that the possible range of $\eta$ is restricted in the IH case when $M_N \lesssim 350$ MeV. We should comment that a similar analysis can be done for the NH case. However, the present observational data of neutrino oscillation and search experiments cannot constrain the value of the Majorana phase. If future experiments improve the data, we may also obtain a limitation on $\eta$ even in the NH case.

Next, we discuss the impact on the $0\nu2\beta$ decay from the result in the IH case. The decay rate of $0\nu2\beta$ decay is characterized by the effective neutrino mass $m_{\text{eff}}$ (see, e.g., Refs. [8,44,45]). In the model under consideration, it is given by [29]

$$m_{\text{eff}} = m_{\text{eff}}^\nu + \sum_I M_I \Theta_{ei}^2 f_\beta(M_I),$$

(10)

where $m_{\text{eff}}^\nu = \sum_i m_i U_{ei}^2$ and $M_I$ denotes the mass eigenvalue of the $I$th heavy neutrino. $f_\beta(M_I)$ is a function that represents the suppression of the nuclear matrix element of the contribution of heavy neutrinos. The function is unity for $M_I \ll 1$ GeV, and decreases as $1/M_I^2$ for $M_I \gg 1$ GeV. The details of $f_\beta$ are described in Ref. [29] (see also Refs. [46,47]).

To evaluate $m_{\text{eff}}$, we can neglect the mass difference of heavy neutrinos because the observed value of baryon asymmetry requires a very small mass difference. In this case, we can approximate $m_{\text{eff}}$ in Eq. (10) as [29]

$$|m_{\text{eff}}| = \left[ 1 - f_\beta(M_N) \right] |m_{\text{eff}}^\nu|,$$

(11)

where

$$|m_{\text{eff}}^\nu| = \cos^2 \theta_{13}(m_1^2 \cos^4 \theta_{12} + m_2^2 \sin^4 \theta_{12} + 2m_1m_2 \cos^2 \theta_{12} \sin^2 \theta_{12} \cos 2\eta)^{1/2}.$$  

(12)

As shown in this equation, $|m_{\text{eff}}^\nu|$ has a significant dependence on $\eta$. By varying $\eta$ from 0 to $\pi$ we find $|m_{\text{eff}}^\nu| = (1.82 - 4.79) \times 10^{-2}$ eV. This range is indeed for the conventional case when only light neutrinos contribute to the $0\nu2\beta$ decay.

There are two important impacts on $|m_{\text{eff}}|$ from heavy neutrinos when the masses are smaller than about 500 MeV. The first is the destructive contribution from these particles [29], which can be easily seen from Eq. (11). The second is the impact from the Majorana phase restricted by experimental and cosmological constraints. Figure 4 shows the predicted range of $|m_{\text{eff}}|$ which is consistent with the constraints on heavy neutrinos.
Fig. 4. The allowed region of \(M_N\) and \(|m_\text{eff}|\) in the IH case. The region between the (green) dashed lines is allowed when we take the full range of the Majorana phase \(\eta\). The region between the (red) solid lines is allowed by the constraints on \(\eta\) from the search experiment and BBN.

We can see the correlation between the allowed region of \(\eta\) in Fig. 3 and the predicted range of \(|m_\text{eff}|\) in Fig. 4. For \(M_N = 188–269\) MeV, \(\eta \sim \pi/2\) is favored and then the lower value of \(|m_\text{eff}|\) is predicted, while \(\eta \sim \pi/2\) is disfavored for \(M_N = 285–350\) MeV and then the larger value of \(|m_\text{eff}|\) is predicted. When \(M_N > 350\) MeV, there is no limitation of \(\eta\) but the presence of heavy neutrinos induces a smaller value of \(|m_\text{eff}|\) compared with \(|m_\nu^{\text{eff}}|\).

5. Conclusions

We have considered right-handed neutrinos that are responsible for neutrino masses and BAU. In particular, we have discussed the case when they are quasi-degenerate and lighter than charged kaons. It has been found that the constraints from BBN and direct search experiments can be avoided when the degenerate mass is \(M_N \geq 163\) MeV for the NH case, and \(M_N = 188–269\) MeV and \(M_N \geq 285\) MeV for the IH case.

Interestingly, we have found that the possible value of the Majorana phase is limited by the present constraints from cosmology and search experiments if the neutrino masses obey the IH and \(M_N = 188–350\) MeV. In such a case, the rate of 0\(\nu\)2\(\beta\) decay is also limited, which may be different from the prediction of the conventional scenario with three active neutrinos. Thus, future experiments of 0\(\nu\)2\(\beta\) decay may give us an important hint for the right-handed neutrinos considered in this analysis.

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