



# On neutrinoless double beta decay in the $\nu$ MSM



Takehiko Asaka<sup>a</sup>, Shintaro Eijima<sup>b</sup>, Hiroyuki Ishida<sup>c,\*</sup>

<sup>a</sup> Department of Physics, Niigata University, Niigata 950-2181, Japan

<sup>b</sup> Institute of Physics, École Polytechnique Fédérale de Lausanne, CH-1015 Lausanne, Switzerland

<sup>c</sup> Graduate School of Science and Engineering, Shimane University, Matsue 690-8504, Japan

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## ABSTRACT

We consider the neutrinoless double beta ( $0\nu\beta\beta$ ) decay in the  $\nu$ MSM, in which three right-handed neutrinos with masses below the electroweak scale are additionally introduced to the Standard Model. In this model there appear three heavy neutral leptons  $N_1$ ,  $N_2$ , and  $N_3$  corresponding to right-handed neutrinos. It has been known that the lightest one  $N_1$  with keV mass, which is a candidate for dark matter, gives a negligible contribution to the  $0\nu\beta\beta$  decay. By contrast, the heavier ones  $N_2$  and  $N_3$ , which are responsible to the seesaw mechanism of neutrino masses and baryogenesis, give the destructive contribution (compared with one from active neutrinos). This is because their mass degeneracy at high precision has been assumed, which is expected by analytical studies of baryogenesis. In this analysis, we find that the effective mass of the  $0\nu\beta\beta$  decay becomes larger than one from active neutrinos due to the  $N_2$  and  $N_3$  constructive contribution when the mass difference becomes larger and the mass ordering of active neutrinos is inverted. Such a possibility will be explored by the current and near future experiments of the  $0\nu\beta\beta$  decay.

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## 1. Introduction

The fate of the lepton number is an interesting question of particle physics. It is an accidental global symmetry in the Standard Model (SM) and then physics beyond the SM can violate the symmetry. One of the most interesting possibilities is the seesaw mechanism [1] for generating masses of active neutrinos. In this case right-handed neutrinos are introduced with Majorana masses which break the lepton number. Thus, the non-zero masses of active neutrinos observed in oscillation experiments may lead to the violation if neutrinos are Majorana particles.

The lepton number violation induces striking processes which are absent in the SM. One famous example is the neutrinoless double beta ( $0\nu\beta\beta$ ) decay:  $(A, Z) \rightarrow (A, Z + 2) + 2e^-$ , which changes the lepton number by two units [2]. The decay occurs if the neutrinos are massive Majorana particles and its rate is proportional to the squared of the effective neutrino mass  $m_{\text{eff}}$ . As for the  $0\nu\beta\beta$  decay of  $^{136}\text{Xe}$ , the lower bound on the half-life is  $T_{1/2} > 1.1 \times 10^{26}$  yr at 90% C.L. [3] which leads to the upper bound on the effective mass  $|m_{\text{eff}}| < m_{\text{eff}}^{\text{UB}} = (61\text{--}161)$  meV, where the uncertainties in Ref. [4] are taken into account. The bound on the

$0\nu\beta\beta$  decay of  $^{76}\text{Ge}$  is  $T_{1/2} > 3.0 \times 10^{25}$  yr at 90% C.L. [5]. In this case the bound becomes  $|m_{\text{eff}}| < m_{\text{eff}}^{\text{UB}} = (213\text{--}308)$  meV.

The contribution to  $m_{\text{eff}}$  from active neutrinos  $\nu_i$  ( $i = 1, 2, 3$ ),  $m_{\text{eff}}^\nu$ , is given by

$$m_{\text{eff}}^\nu = \sum_{i=1,2,3} U_{ei}^2 m_i, \quad (1)$$

where  $m_i$  are masses of active neutrinos and their ordering is  $m_3 > m_2 > m_1$  in the normal hierarchy (NH) case and  $m_2 > m_1 > m_3$  in the inverted hierarchy (IH) case. The mixing matrix of active neutrinos [6] is denoted by  $U_{\alpha i}$  ( $\alpha = e, \mu, \tau$ ), which is represented by

$$U = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -c_{23}s_{12} - s_{23}c_{12}s_{13}e^{i\delta} & c_{23}c_{12} - s_{23}s_{12}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{23}s_{12} - c_{23}c_{12}s_{13}e^{i\delta} & -s_{23}c_{12} - c_{23}s_{12}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix} \times \text{diag}(1, e^{i\eta}, e^{i\eta'}), \quad (2)$$

where  $s_{ij} = \sin\theta_{ij}$  and  $c_{ij} = \cos\theta_{ij}$ .  $\delta$  is the Dirac phase and  $\eta$  and  $\eta'$  are the Majorana phases. By using the central values of mixing angles  $\theta_{ij}$  in Ref. [7] and applying the cosmological bound  $\sum m_i < 0.23$  eV [8], the effective mass due to active neutrinos is the range  $|m_{\text{eff}}^\nu| = (1.49\text{--}72.0)$  meV for the NH case or  $(18.5\text{--}82.4)$  meV for the IH case, respectively. Thus, the predicted range, especially in

\* Corresponding author.

E-mail address: [hiroyuki403@gmail.com](mailto:hiroyuki403@gmail.com) (H. Ishida).

the IH case, will begin to be tested by the near future experiments of the  $0\nu\beta\beta$  decay.

We revisit the  $0\nu\beta\beta$  decay in the neutrino Minimal Standard Model ( $\nu$ MSSM) [9,10], which is the extension of the SM by three right-handed neutrinos  $\nu_{RI}$  with masses below  $\mathcal{O}(10^2)$  GeV. The model realizes the seesaw mechanism of neutrino masses and the baryogenesis via oscillation of right-handed neutrinos [11,10], and offers a candidate of dark matter (known as sterile neutrino dark matter [12,13]). The  $0\nu\beta\beta$  decay of the model has been investigated in Refs. [14–17]. In this model  $m_{\text{eff}}$  is induced not only by active neutrinos but also by heavy neutral leptons (HNLs) associated with right-handed neutrinos. It has been shown that the contributions from HNLs can be comparable to the one from active neutrinos and, importantly, it is a destructive contribution due to the strong mass degeneracy of HNLs [15].

In this analysis we will show that HNLs can give a constructive contribution to  $m_{\text{eff}}$  in a certain parameter region, which should be contrast to previous results. This enhances the rate of the  $0\nu\beta\beta$  decay and hence it is a good target for the current and near future experiments.

## 2. $0\nu\beta\beta$ decay in the $\nu$ MSSM

Let us consider the  $\nu$ MSSM, in which three right-handed neutrinos  $\nu_{RI}$  ( $I = 1, 2, 3$ ) are introduced to the SM. Its Lagrangian is

$$\mathcal{L}_{\nu\text{MSSM}} = \mathcal{L}_{\text{SM}} + i\overline{\nu_{RI}} \gamma^\mu \partial_\mu \nu_{RI} - \left( F_{\alpha I} \bar{L}_\alpha \Phi \nu_{RI} + \frac{M_I}{2} \overline{\nu_{RI}^c} \nu_{RI} + h.c. \right), \quad (3)$$

where  $\mathcal{L}_{\text{SM}}$  is the SM Lagrangian.  $L_\alpha = (\nu_{L\alpha}, e_{L\alpha})^T$  and  $\Phi = (\phi^0, \phi^-)^T$  are lepton and Higgs weak-doublets.  $M_I$  are Majorana masses of  $\nu_{RI}$  (which are taken to be real and positive) and  $F_{\alpha I}$  are Yukawa coupling constants. The two types of neutrino masses,  $[M_D]_{\alpha I} = F_{\alpha I} \langle \Phi \rangle$  and  $M_I$ , are assumed to be  $|[M_D]_{\alpha I}| \ll M_I \lesssim \mathcal{O}(10^2)$  GeV. The model realizes the seesaw mechanism and the mass eigenstates in the neutrino sector are active neutrinos  $\nu_i$  and heavy neutral leptons (HNLs)  $N_I$ . The left-handed neutrinos are then written as

$$\nu_{L\alpha} = U_{\alpha i} \nu_i + \Theta_{\alpha I} N_I^c, \quad (4)$$

where the mixing matrix of HNLs is given by  $\Theta_{\alpha I} = [M_D]_{\alpha I}/M_I$ .

The lightest HNL  $N_1$  with keV scale mass is a candidate of dark matter. In order to avoid the constraints for dark matter its mixing  $\Theta_{\alpha 1}$  must be sufficiently small [18]. Consequently,  $N_1$  decouples from the seesaw mechanism and from the baryogenesis. The heavier ones  $N_2$  and  $N_3$  are then responsible to these two mechanism. It has been shown that they should be quasi-degenerate  $\Delta M = (M_3 - M_2)/2 \ll M_N = (M_3 + M_2)/2$  and  $M_N = \mathcal{O}(0.1) - \mathcal{O}(10^2)$  GeV for the successful scenario. (See the analyses in Refs. [19,20,16,21].)

The effective mass in the  $\nu$ MSSM is given by

$$m_{\text{eff}}^{\nu\text{MSSM}} = m_{\text{eff}}^\nu + m_{\text{eff}}^N, \quad (5)$$

where the first term is the contribution from active neutrinos in Eq. (1). Since  $N_1$  gives a suppressed contribution to the masses of active neutrinos, the lightest one is smaller than  $\mathcal{O}(10^{-4})$  eV. In this case, the range of  $m_{\text{eff}}^\nu$  is limited as  $|m_{\text{eff}}^\nu| = (1.49 - 3.66)$  meV for the NH case or (18.5–47.9) meV for the IH case, respectively.

On the other hand, the second term in Eq. (5) denotes the contributions from HNLs given by

$$m_{\text{eff}}^N = \sum_{I=1,2,3} \Theta_{\alpha I}^2 f_\beta(M_I) M_I, \quad (6)$$

where the  $f_\beta(M_I)$  represents the suppression of the nuclear matrix element of the  $0\nu\beta\beta$  decay for  $M_I \gg \mathcal{O}(0.1)$  GeV. In this analysis,

we follow the results in Ref. [4] (see also Ref. [22]) and take the expression

$$f_\beta(M_I) = \frac{\langle p^2 \rangle}{\langle p^2 \rangle + M_I^2}, \quad (7)$$

with  $\langle p^2 \rangle \simeq (200 \text{ MeV})^2$ . As shown in Ref. [14] the contribution from  $N_1$  is much smaller than  $m_{\text{eff}}^\nu$  and those from  $N_2$  and  $N_3$  (see also Refs. [15,23]). Thus, we shall neglect its contribution and set its Yukawa couplings  $F_{\alpha 1} = 0$  from now on. In this case, the Yukawa couplings of  $N_2$  and  $N_3$  are parameterized as presented in Eq. (2.8) of Ref. [15].

The contributions from  $N_2$  and  $N_3$  is decomposed into two parts [15]

$$m_{\text{eff}}^N = \sum_{I=2,3} \Theta_{\alpha I}^2 f_\beta(M_I) M_I = \bar{m}_{\text{eff}}^{N_{2,3}} + \delta m_{\text{eff}}^{N_{2,3}}, \quad (8)$$

where

$$\bar{m}_{\text{eff}}^{N_{2,3}} = f_\beta(M_N) \sum_{I=2,3} \Theta_{\alpha I}^2 M_I = -f_\beta(M_N) m_{\text{eff}}^\nu, \quad (9)$$

$$\delta m_{\text{eff}}^{N_{2,3}} = \sum_{I=2,3} [f_\beta(M_I) - f_\beta(M_N)] \Theta_{\alpha I}^2 M_I, \quad (10)$$

where we have used in Eq. (9) the seesaw relation of masses and mixings of active neutrinos and HNLs

$$0 = \sum_{i=1,2,3} U_{\alpha i}^2 m_i + \sum_{I=1,2,3} \Theta_{\alpha I}^2 M_I. \quad (11)$$

In the previous analysis we have neglected  $\delta m_{\text{eff}}^{N_{2,3}}$  which vanishes for  $\Delta M = 0$ . This is because the baryogenesis via neutrino oscillations requires  $\Delta M/M_N \ll 1$  in order to generate the sufficient amount of the baryon asymmetry of the universe (BAU). (See, however, the discussion below.) In this case the effective mass in the  $\nu$ MSSM is given by a simple expression [15]

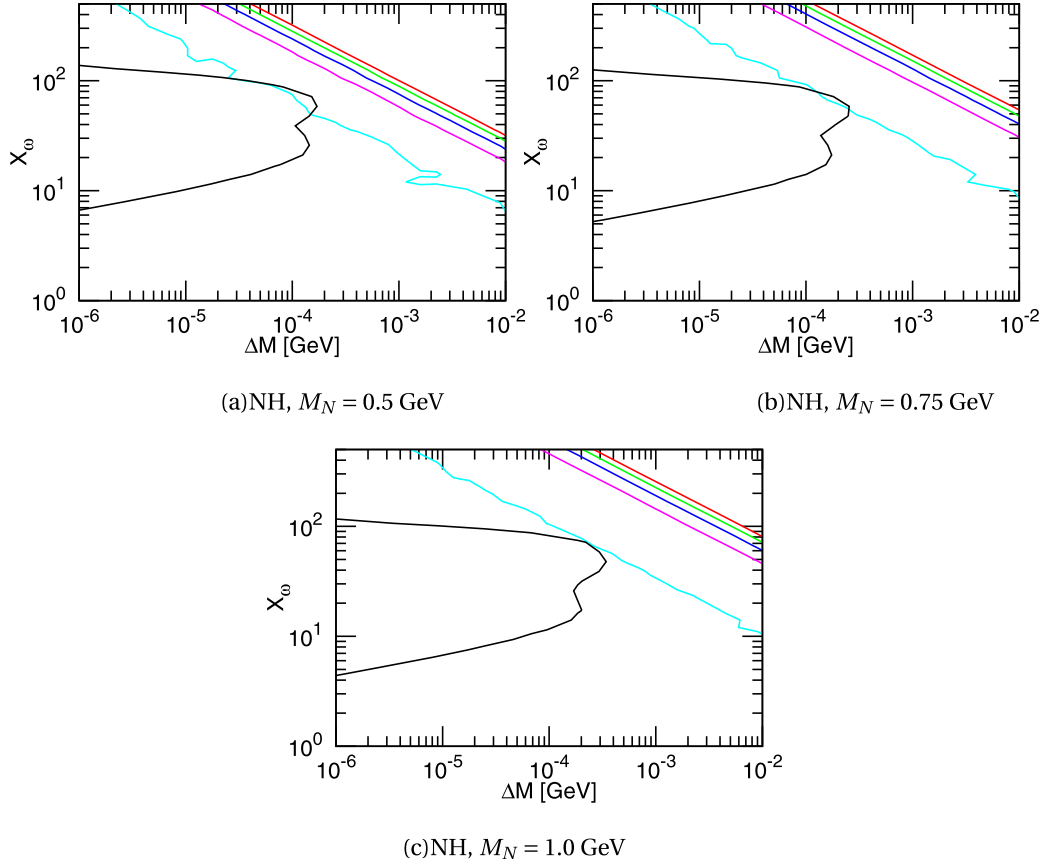
$$m_{\text{eff}}^{\nu\text{MSSM}} = [1 - f_\beta(M_N)] m_{\text{eff}}^\nu = \frac{M_N^2}{\langle p^2 \rangle + M_N^2} m_{\text{eff}}^\nu. \quad (12)$$

It is seen that the contributions from HNLs can be neglected when  $M_N^2 \gg \langle p^2 \rangle$  [14]. On the other hand, when  $M_N^2 \lesssim \langle p^2 \rangle$ ,  $N_2$  and  $N_3$  give sizable, destructive contributions to  $m_{\text{eff}}^\nu$ . Especially, the effective mass in the  $\nu$ MSSM becomes zero as  $M_N \rightarrow 0$ . This is a general consequence of the seesaw relation (11) when HNLs are much lighter than  $\langle p^2 \rangle^{1/2}$ .

Let us summarize the results of the  $0\nu\beta\beta$  in the  $\nu$ MSSM so far: (1) HNLs  $N_2$  and  $N_3$  gives a destructive contribution and then  $m_{\text{eff}}^{\nu\text{MSSM}}$  is given by Eq. (12), and hence  $|m_{\text{eff}}^{\nu\text{MSSM}}| < |m_{\text{eff}}^\nu|$ . (2) The upper bound on  $|m_{\text{eff}}^{\nu\text{MSSM}}|$  is the same as the upper bound on  $|m_{\text{eff}}^\nu|$  for the lightest active neutrino with  $\mathcal{O}(10^{-4})$ . It is achieved for  $M_N \gg \langle p^2 \rangle^{1/2}$ . (3) There is the lower bound on  $|m_{\text{eff}}^{\nu\text{MSSM}}|$  which is obtained by the smallest value of  $|m_{\text{eff}}^\nu|$  (for the lightest active neutrino with  $\mathcal{O}(10^{-4})$ ) multiplied by the factor  $[1 - f_\beta(M_N)]$  with the smallest value of  $M_N$ . Notice that the lower bound on  $M_N \simeq 150$  MeV to avoid the cosmological constraint as well as the bounds from the direct search experiments (see the recent analysis [16]).

## 3. Enhancement of $0\nu\beta\beta$ decay in the $\nu$ MSSM

We are now at the position to discuss the main point of this analysis. It will be shown that there is a certain parameter region in which the effective mass  $|m_{\text{eff}}^{\nu\text{MSSM}}|$  can be larger than  $|m_{\text{eff}}^\nu|$ . This is because of  $\delta m_{\text{eff}}^{N_{2,3}}$  in Eq. (10) which was neglected in the previous analysis. First, we note that



**Fig. 1.** Contours of the maximal value of  $|m_{\text{eff}}^{\nu\text{MSM}}|_{\text{max}}$  in  $\Delta M$ - $X_\omega$  plane for the NH case. Here contour lines correspond to  $|m_{\text{eff}}^{\nu\text{MSM}}|_{\text{max}}/|m_{\text{eff}}^\nu|_{\text{max}} = 1$  (cyan line), 2 (magenta line), 3 (blue line), 4 (green line) and 5 (red line). Here we take  $|m_{\text{eff}}^{\text{SM}}|_{\text{max}} = 3.66$  meV. In the region enclosed by the black line the BAU can be explained. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

$$\begin{aligned} \delta m_{\text{eff}}^{N_{2,3}} &\simeq f'_\beta(M_N)\Delta M \left( \Theta_{e3}^2 M_3 - \Theta_{e2}^2 M_2 \right) \\ &\simeq -\frac{2\langle p^2 \rangle \langle \Phi^2 \rangle \Delta M}{(\langle p^2 \rangle + M_N^2)^2} (F_{e3}^2 - F_{e2}^2), \end{aligned} \quad (13)$$

where we have neglected the terms suppressed by  $\mathcal{O}(\Delta M^2/M_N^2)$ . Now we use the parameterization of the Yukawa coupling constants in Ref. [15], and then obtain the expression

$$\delta m_{\text{eff}}^{N_{2,3}} = -\frac{\langle p^2 \rangle M_N^2}{(\langle p^2 \rangle + M_N^2)^2} \frac{\Delta M}{M_N} X_\omega^2 m_*, \quad (14)$$

where we have assumed  $X_\omega \gg 1$ . Here the mass parameter  $m_*$  is given by the mixing angles and phases as

$$m_* = e^{-2i(\text{Re}\omega + \delta)} \left( e^{i(\delta + \eta)} \sqrt{m_2} \sin \theta_{12} \cos \theta_{13} + i\xi \sqrt{m_3} \sin \theta_{13} \right)^2, \quad (15)$$

for the NH case and

$$m_* = e^{-2i\text{Re}\omega} \cos^2 \theta_{13} \left( \sqrt{m_1} \cos \theta_{12} + i\xi e^{i\eta} \sqrt{m_2} \sin \theta_{12} \right)^2, \quad (16)$$

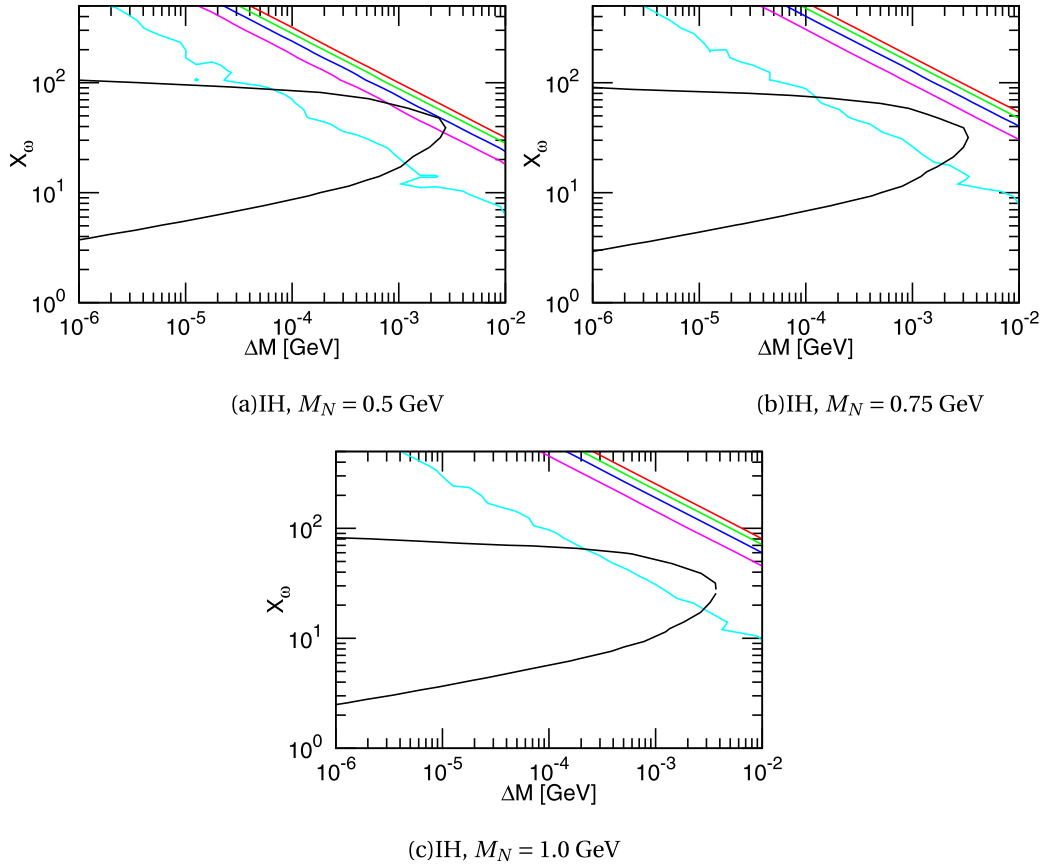
for the IH case. The mass parameter takes its maximal value  $|m_*| = 7.0$  meV if  $\delta + \eta = \pi/2$  and  $\xi = +1$  (or  $\pi/3$  and  $\xi = -1$ ) for the NH case, whereas  $|m_*| = 92$  meV if  $\eta = 3\pi/2$  and  $\xi = +1$  (or  $\eta = \pi/2$  and  $\xi = -1$ ) for the IH case. Therefore, we find that  $|\delta m_{\text{eff}}^{N_{2,3}}|$  can be significantly large, when  $M_N \simeq \langle p^2 \rangle^{1/2}$  and  $X_\omega \Delta M / M_N \gtrsim 1$ , as

$$\left| \delta m_{\text{eff}}^{N_{2,3}} \right| \simeq X_\omega^2 \left( \frac{\Delta M}{M_N} \right) |m_*|. \quad (17)$$

In order to discuss an impact of  $\delta m_{\text{eff}}^{N_{2,3}}$ , hereafter we restrict CP phases and  $\xi$  as derived above. It should be noted that  $\delta m_{\text{eff}}^{N_{2,3}}$  can always be a constructive contribution by choosing  $\text{Re}\omega$ . As we will demonstrate below, it can overcome  $|m_{\text{eff}}^\nu|$ . This point was missed in the previous analysis.

In the  $\nu\text{MSM}$ , the production mechanism of the lepton asymmetry is completely different from the ordinary leptogenesis. For the whole system the total lepton asymmetry is conserved at the early universe due to the smallness of the Majorana masses, and then the asymmetries are stored separately in left- and right-handed lepton sectors. Since the decay processes of right-handed neutrinos are irrelevant at the temperature above sphaleron freeze-out, the lepton asymmetry is produced through CP violating right-handed neutrino oscillation [11,10] which starts at the typical temperature  $T_{\text{osc}} = (M_0 \Delta M M_N / 6)^{1/3}$  ( $M_0 = 7.12 \times 10^{17}$  GeV).

In the baryogenesis via neutrino oscillations the upper bounds on  $X_\omega$  and  $\Delta M$  can be obtained in order to generate the observed baryon asymmetry [20]. When  $X_\omega \gg 1$ , the Yukawa coupling constants become larger which leads to the washout of the produced asymmetry. On the other hand, when  $\Delta M$  becomes larger, the oscillation of right-handed neutrinos begins at earlier epoch, which again leads to the suppression of the yield of baryon asymmetry. In addition, when  $M_N$  becomes smaller than  $K$ -meson mass,  $N_2$  and  $N_3$  can be produced in  $K$ -meson decay and then receive the stringent constraints from direct search experiments. In this case, the mixing angles  $|\Theta_{\alpha 2}|$  and  $|\Theta_{\alpha 3}|$  cannot be large and an extremely large  $X_\omega$  is disfavored. See, for example, the recent analysis in Ref. [24].



**Fig. 2.** Contours of the maximal value of  $|m_{\text{eff}}^{\text{vMSM}}|_{\text{max}}$  in  $\Delta M$ - $X_\omega$  plane for the IH case. Here contour lines correspond to  $|m_{\text{eff}}^{\text{vMSM}}|_{\text{max}}/|m_{\text{eff}}^{\text{v}}|_{\text{max}} = 1$  (cyan line), 2 (magenta line), 3 (blue line), 4 (green line) and 5 (red line). Here we take  $|m_{\text{eff}}^{\text{vMSM}}|_{\text{max}} = 47.9$  meV. In the region enclosed by the black line the BAU can be explained. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

Based on these arguments, we shall illustrate our idea in Figs. 1 and 2. Here we take  $M_N = 0.5, 0.75,$  and  $1$  GeV for both NH and IH cases. First, the successful baryogenesis is possible in the region enclosed by the black line for each figure.

In this analysis, we use the kinetic equations of the case (ii) given in Ref. [25] for the estimation of the BAU.

$$\begin{aligned} \frac{dR_N}{dt} = & -i\langle [H_N], R_N \rangle - \frac{3\langle \gamma_N^d \rangle}{2} \{F^\dagger F, R_N - \mathbf{1}\} \\ & + 2\langle \gamma_N^d \rangle F^\dagger (A - \mathbf{1}) F - \frac{\langle \gamma_N^d \rangle}{2} \{F^\dagger (A^{-1} - \mathbf{1}) F, R_N\}, \end{aligned} \quad (18)$$

$$\begin{aligned} \frac{d\mu_{\nu_\alpha}}{dt} = & -\frac{3\gamma_\nu^d(T)}{2} [FF^\dagger]_{\alpha\alpha} \tanh \mu_{\nu_\alpha} \\ & + \frac{\gamma_\nu^d(T)}{2} [FR_N F^\dagger - F^* R_N^* F^T]_{\alpha\alpha} \frac{1}{\cosh \mu_{\nu_\alpha}} \\ & + \frac{\gamma_\nu^d(T)}{4} \left\{ [F(R_N - \mathbf{1})F^\dagger]_{\alpha\alpha} (1 - \tanh \mu_{\nu_\alpha}) \right. \\ & \left. - [F^*(R_N^* - \mathbf{1})F^T]_{\alpha\alpha} (1 + \tanh \mu_{\nu_\alpha}) \right\}, \end{aligned} \quad (19)$$

where  $R_N$  ( $R_{\bar{N}}$ ) and  $\mu_{\nu_\alpha}$  are the density matrix of right-handed (anti-)neutrino and the chemical potential of left-handed lepton,  $A = \text{diag}(e^{\mu_{\nu_e}}, e^{\mu_{\nu_\mu}}, e^{\mu_{\nu_\tau}})$ ,  $\gamma_\nu^d$  and  $\gamma_N^d$  are destruction rates for left- and right-handed neutrinos, respectively. The equations for  $\bar{N}$  can be obtained by the CP conjugation of Eq. (18).  $T$  is the temperature of the universe. (See the details in Ref. [25].) In Figs. 1 and 2, the yield of the baryon asymmetry from these kinetic equa-

tions can be consistent with the observed value for the regions inside the black colored line. The behavior of this line is interpreted as follows: The value of  $X_\omega$  is bounded from both below and above as explained above. In addition, in the interesting parameter space where  $\Delta M$  is relatively large, the generated asymmetry at the sphaleron freeze-out temperature decreases as  $\Delta M$  increases due to the suppression of the oscillation effect, and then  $\Delta M$  is bounded from above.

In these figures we also plot contours of the ratio  $|m_{\text{eff}}^{\text{vMSM}}|_{\text{max}}/|m_{\text{eff}}^{\text{v}}|_{\text{max}}$ , where the maximal value of the contribution from active neutrinos is  $|m_{\text{eff}}^{\text{v}}|_{\text{max}} = 3.66$  and  $47.9$  meV for the NH and IH cases, respectively. It is then found that the ratio can be large as unity at most in the NH case within our setup. This means that the effective neutrino mass in the  $\nu$ MSM cannot exceed the active neutrino one when the mass ordering is NH. On the other hand, it should be noted that the ratio can be large as three for the IH case, which means that  $|m_{\text{eff}}^{\text{vMSM}}|$  can be large as  $140$  meV. This should be contrast the results in the previous works. Such a large value is realized when  $M_N \simeq 500$  MeV,  $\Delta M \simeq 10^{-3}$  GeV, and  $X_\omega \simeq 50$ . Here the choice of these parameters gives the mixing angles of heavy neutral leptons as [24]

$$|\Theta|^2 = \sum_{\alpha=e,\mu,\tau} \sum_{l=2,3} |\Theta_{\alpha l}|^2 = \frac{\sum_{i=1,2,3} m_i}{2M_N} X_\omega^2 \simeq 1.3 \times 10^{-7}. \quad (20)$$

Note that this parameter set gives rise to the large effective mass and the sufficient amount of the BAU at the same time. The whole of such possibilities can be derived by the numerical scan of the full parameter space of the model taking into account the exper-

imental and cosmological constraints. This issue is, however, beyond the scope of the present analysis.

#### 4. Summary and outlook

We have investigated the  $0\nu\beta\beta$  decay in the  $\nu$ MSM. Especially, we have estimated the contribution  $\delta m_{\text{eff}}^{N_{2,3}}$  proportional to the mass difference  $\Delta M$  between HNLs  $N_2$  and  $N_3$ , which was not fully taken into account in the previous analysis.

It has been shown that the effective neutrino mass  $|m_{\text{eff}}^{\nu\text{MSM}}|$  can be large owing to this effect if HNLs have a large mass difference and strong mixing and the common mass is close to the typical momentum of this process ( $(p^2)^{1/2} \simeq 200$  MeV). Actually, we have shown that the maximal value of  $|m_{\text{eff}}^{\nu\text{MSM}}|$  is 140 meV when  $M_N \simeq 500$  MeV,  $\Delta M \simeq 10^{-3}$  GeV, and  $X_\omega \simeq 50$ , if the mass ordering of active neutrinos is the inverted one. Such a large value is comparable to the current upper bound by using  $0\nu\beta\beta$  decay of  $^{136}\text{Xe}$  from KamLAND-Zen experiment. It should be stressed that large  $|m_{\text{eff}}^{\nu\text{MSM}}|$  is possible only if the CP phases including Majorana phase and the mixing angle of HNLs are aligned appropriately. Our analysis, therefore, indicates that these unexplored parameters of the  $\nu$ MSM start to be revealed by the  $0\nu\beta\beta$  decay experiments.<sup>1</sup>

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<sup>1</sup> While writing up this manuscript, one of the authors (S.E.) and M. Drewes presented numerical results of this issue in Ref. [26] where the parameter scan including cosmological and experimental constraints is performed. In addition, after we submitted this work to arXiv, we noticed that a paper (arXiv:1606.06719), which discusses the similar issue, appeared on the same day. These results are consistent with ours.