



PUBLISHED FOR SISSA BY SPRINGER

RECEIVED: May 3, 2014
ACCEPTED: June 14, 2014
PUBLISHED: July 9, 2014

Neutrino universe

Tetsutaro Higaki,^a Ryuichiro Kitano^{a,b} and Ryosuke Sato^a

^a*Institute of Particle and Nuclear Studies,
High Energy Accelerator Research Organization (KEK),
Tsukuba 305-0801, Japan*

^b*The Graduate University for Advanced Studies (Sokendai),
Tsukuba 305-0801, Japan*

E-mail: thigaki@post.kek.jp, ryuichiro.kitano@kek.jp,
rsato@post.kek.jp

ABSTRACT: The Standard Model of particle physics fails to explain the important pieces in the standard cosmology, such as inflation, baryogenesis, and dark matter of the Universe. We consider the possibility that the sector to generate small neutrino masses is responsible for all of them; the inflation is driven by the Higgs field to break $B - L$ gauge symmetry which provides the Majorana masses to the right-handed neutrinos, and the reheating process by the decay of the $B - L$ Higgs boson supplies the second lightest right-handed neutrinos whose CP violating decays produce $B - L$ asymmetry, *à la*, leptogenesis. The lightest right-handed neutrinos are also produced by the reheating process, and remain today as the dark matter of the Universe. In the minimal model of the inflaton potential, one can set the parameter of the potential by the data from CMB observations including the BICEP2 and the Planck experiments. In such a scenario, the mass of the dark matter particle is predicted to be of the order of PeV. We find that the decay of the PeV right-handed neutrinos can explain the high-energy neutrino flux observed at the IceCube experiments if the lifetime is of the order of 10^{28} s.

KEYWORDS: Cosmology of Theories beyond the SM, Neutrino Physics

ARXIV EPRINT: [1405.0013](https://arxiv.org/abs/1405.0013)

Contents

1	Introduction	1
2	Model	2
3	Inflation with the $B - L$ Higgs field	6
4	Reheating by the inflaton decay	7
4.1	Leptogenesis	8
4.2	Dark matter abundance	9
5	PeV neutrinos as a signal of decaying N_1	9
5.1	The branching fractions of N_1	10
5.2	Neutrino flux from decay of N_1	11
6	Summary	14

1 Introduction

Various cosmological observations are telling us that the Standard Model of particle physics needs some extension. The observation of the cosmic microwave background (CMB) and its anisotropy strongly supports the inflationary cosmology [1–3], which requires a process to generate the Standard Model particles after the inflation era. The particle-antiparticle asymmetry should also be generated after or during the reheating process. Also, the dark matter of the Universe must also be produced in the course of the cosmological history. The Standard Model should be extended to accommodate the inflation, baryogenesis and dark matter of the Universe.

One of the clues towards the mysteries of the Universe may be the finite neutrino masses, which are another evidence to go beyond the Standard Model. Once three kinds of right-handed neutrinos are introduced in the same way as other fermions, the global $U(1)_{B-L}$ symmetry becomes non-anomalous, and thus can be promoted to a Higgsed gauge symmetry. It seems that all the ingredients to accommodate the realistic cosmology are present in this $U(1)_{B-L}$ extended Standard Model.

The inflation can be driven by the Higgs field to break $U(1)_{B-L}$ gauge symmetry [4, 5] by assuming an appropriate form of the potential based on the idea of the chaotic inflation [6]. After the inflation, the $B - L$ Higgs field oscillates about the minimum of the potential where $U(1)_{B-L}$ is broken. The spontaneous breaking of the $B - L$ symmetry can give Majorana masses to the right-handed neutrinos through the Yukawa coupling, explaining the smallness of the neutrino masses by the seesaw mechanism [7–12]. The very

same coupling allows the decay of the inflaton oscillation into the right-handed neutrinos to reheat the Universe. The subsequent decay of the right-handed neutrinos can provide the baryon asymmetry of the Universe by the leptogenesis mechanism [13]. The lightest right-handed neutrino should also be produced by the inflaton decay. If it is long-lived, this non-thermal component is a good candidate of the dark matter of the Universe.

There have been other minimalistic approaches to the connection between particle physics and cosmology. An realistic model with the minimal particle content has been constructed in ref. [14], where the inflaton and the dark matter particle are both introduced as new scalar fields. The possibility of the inflaton as the Higgs-like field, thus playing important roles both in particle physics and cosmology, has been considered in refs. [15–29]. The dark matter of the Universe as the right-handed neutrino has also been considered in refs. [30–39] where the mass range of $\mathcal{O}(\text{keV})$ are assumed.

In this paper, we consider the $U(1)_{B-L}$ extended Standard Model which covers the shortages in the Standard Model including the small neutrino masses as well as cosmological observations. We find that this minimalistic scenario is consistent with various observations such as tensor-to-scalar ratio, spectral index of the CMB fluctuations, the neutrino masses, baryon asymmetry of the Universe, and the energy density of the dark matter. We find, in the case where the reheating process is dominated by the decay of inflaton into the second lightest right-handed neutrinos, the mass of the dark matter particle is predicted to be of the order of PeV.

Since there is no reason to assume that the dark matter particle, the lightest right-handed neutrino, to be absolutely stable, we expect the decay of the dark matter to happen occasionally somewhere in the Universe. Through the dimension-four Yukawa interactions, the main decay mode would be into a lepton and a W boson, or a neutrino and a Z/h boson. We demonstrate that the PeV neutrino events found at the IceCube experiment [40, 41] can be explained by the decaying right-handed neutrinos if the lifetime is of the order of 10^{28} s.¹

In the following sections, based on the above scenario with $U(1)_{B-L}$ extended Standard Model, we discuss the neutrino flavor structure, an inflation model with the $B-L$ Higgs, the non-thermal leptogenesis, the dark matter abundance produced by the decay of the inflaton, and the signals of decaying right-handed neutrinos at the IceCube experiment.

2 Model

We extend the gauge group of the Standard Model into,

$$SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)_{B-L},$$

and introduce the right-handed neutrinos, N_i ($i = 1, 2, 3$), and the $U(1)_{B-L}$ Higgs field ϕ_{B-L} which is neutral under the Standard Model gauge group and has charge -2 under $U(1)_{B-L}$. The $U(1)_{B-L}$ symmetry is gauged, and thus the spontaneous breaking of $U(1)_{B-L}$ would not leave the massless Nambu-Goldstone boson. The following interaction

¹See, e.g., refs. [42, 43] for studies on PeV decaying dark matter.

terms are added to the Standard Model:

$$\mathcal{L}_{\text{int}} = y_{\nu}^{ij} \bar{N}_i P_L (\ell_j \cdot \tilde{H}) + \text{h.c.} + \frac{\lambda_i}{2} \phi_{B-L} \bar{N}_i P_L N_i^c + \text{h.c.}, \quad (2.1)$$

where ℓ_i and N_i are four-component Weyl fermions, *i.e.*, $P_L \ell_i = \ell_i$, $P_R N_i = N_i$. The coupling constant λ_i can be taken to be real and positive without loss of generality, and the components of y_{ν}^{ij} , in general, are complex valued. The potential terms for ϕ_{B-L} field can be written as,

$$V(\phi) = \frac{\kappa}{4} (|\phi_{B-L}|^2 - v_{B-L}^2)^2 = \frac{\kappa v_{B-L}^4}{4} \left(\frac{|\phi_{B-L}|^2}{v_{B-L}^2} - 1 \right)^2. \quad (2.2)$$

There can also be an interaction term such as,

$$\mathcal{L}_{\phi H} = \kappa' |\phi_{B-L}|^2 |H|^2. \quad (2.3)$$

For $v_{B-L} \gtrsim 5M_{\text{Pl}}$ which we assume later, the coupling constant κ' is extremely small if we demand this term would not contribute significantly to the Higgs potential.

The spontaneous breaking of $U(1)_{B-L}$ by $\langle \phi_{B-L} \rangle = v_{B-L}$ generates masses of N_i :

$$M_i = \lambda_i v_{B-L}. \quad (2.4)$$

The neutrino masses are, in turn, generated by the seesaw mechanism:

$$m_{\nu}^{ij} = y_{\nu}^{ki} M_k^{-1} y_{\nu}^{kj} \langle H \rangle^2. \quad (2.5)$$

We assume that the lightest right-handed neutrino, N_1 , to be long-lived, and it serves as the dark matter of the Universe. That means,

$$|y_{\nu}^{1i}| \ll 1. \quad (2.6)$$

As we will see in section 5, in the scenario where the PeV neutrino events at the IceCube experiment to be explained by the decay of N_1 , the lifetime of N_1 has to be around 10^{28} s. This lifetime corresponds to $y_{\nu}^{1i} \sim 10^{-29}$. In fact, this model has various unexplained small numbers such as the Higgs mass parameter, the θ parameter in QCD, the cosmological constant, κ' , κ as well as y_{ν}^{1i} . Although we do not look for particular reasons for such small numbers here, a very small y_{ν}^{1i} is somewhat special since it can be protected by a Z_2 symmetry, $N_1 \leftrightarrow -N_1$. If such a symmetry is only violated by some non-perturbative effects of gauge or gravity interactions at high scales, the size may be understood as a natural value.² In such a scenario, it is likely that the non-perturbative effects respect the flavor symmetry, and thus the effective operator to break the Z_2 symmetry, for example, takes the form of

$$\mathcal{L}_{NP} = \frac{1}{\Lambda^{14}} (\ell_1 \cdot \ell_2) (\ell_2 \cdot \ell_3) (\ell_3 \cdot \ell_1) e_1^c e_2^c e_3^c N_1^c N_2^c N_3^c + \text{h.c.} \quad (2.7)$$

²The quantum theory of gravity may give natural ground for such considerations [44–47].

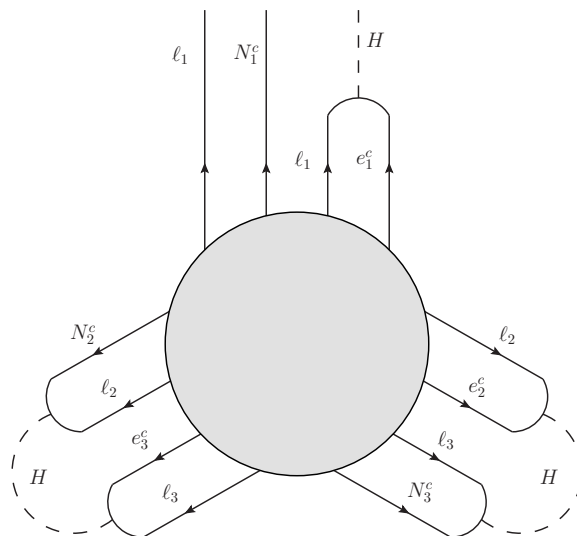


Figure 1. One of the diagrams generating y_ν^{1k} with the interaction in eq. (2.7).

Here, Λ is expected to be the scale which characterizes the non-perturbative effects such as, $\mu e^{-8\pi^2/g^2(\mu)}$, in the case of a gauge theory. This is analogous to the interaction considered in QCD [48, 49]. Together with the Yukawa interactions of the charged lepton sector y_e^{ij} in the Standard Model and $y_\nu^{\alpha i}$ ($\alpha = 2, 3$) in eq. (2.1), y_ν^{1i} is generated as in the diagram in figure 1:

$$y_\nu^{1k} \propto (\det y_e) \epsilon^{ijk} y_\nu^{2i} y_\nu^{3j}. \quad (2.8)$$

One can also consider interactions such as $\mathcal{L}_{NP} = (q_1 \cdot \ell_2)(q_2 \cdot \ell_3)(q_3 \cdot \ell_1) d_1^c d_2^c d_3^c N_1^c N_2^c N_3^c / \Lambda^{14}$. From this operator, we obtain $y_\nu^{1k} \propto (\det y_d) \epsilon^{ijk} y_\nu^{2i} y_\nu^{3j}$. In any case, the flavor symmetry implies an interesting proportionality:

$$y_\nu^{1k} \propto \epsilon^{ijk} y_\nu^{2i} y_\nu^{3j}. \quad (2.9)$$

We will see in section 5 that if this type of contribution is dominated, the branching ratio of the N_1 decay is directly related to the neutrino mixing parameters. By introducing a small parameter c , eq. (2.9) is explicitly written as

$$y_\nu^{1e} = c(y_\nu^{2\mu} y_\nu^{3\tau} - y_\nu^{3\mu} y_\nu^{2\tau}), \quad y_\nu^{1\mu} = c(y_\nu^{2\tau} y_\nu^{3e} - y_\nu^{3\tau} y_\nu^{2e}), \quad y_\nu^{1\tau} = c(y_\nu^{2e} y_\nu^{3\mu} - y_\nu^{3e} y_\nu^{2\mu}). \quad (2.10)$$

Because of tiny $y_\nu^{1\ell}$'s, N_1 provides very little contribution to the neutrino masses. In this case, the neutrino sector is essentially that of the model with only two right-handed neutrinos [50, 51]. Here, we define the following Yukawa matrix \tilde{y} and mass matrix \tilde{M} :

$$\tilde{y} = \begin{pmatrix} y_\nu^{2e} & y_\nu^{2\mu} & y_\nu^{2\tau} \\ y_\nu^{3e} & y_\nu^{3\mu} & y_\nu^{3\tau} \end{pmatrix}, \quad \tilde{M} = \begin{pmatrix} M_2 & 0 \\ 0 & M_3 \end{pmatrix}. \quad (2.11)$$

Neutrino masses are given by,

$$m_\nu \equiv \text{diag}(m_1, m_2, m_3) = (U_{\text{PMNS}}^T \tilde{y}^T \tilde{M}^{-1} \tilde{y} U_{\text{MNS}}) \langle H \rangle^2, \quad (2.12)$$

where U_{MNS} is the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix [52, 53]:

$$U_{\text{PMNS}} = \begin{pmatrix} U_{e1} & U_{e2} & U_{e3} \\ U_{\mu1} & U_{\mu2} & U_{\mu3} \\ U_{\tau1} & U_{\tau2} & U_{\tau3} \end{pmatrix} = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix} \times \text{diag}(1, e^{i\alpha/2}, 1). \quad (2.13)$$

Eq. (2.12) tells us that the lightest neutrino is massless (up to $O((y_\nu^{1i})^2)$ contributions) because the rank of \tilde{y} and \tilde{M} is two. There is only one Majorana phase in eq. (2.13) in this effectively two-generation model. We can parametrize \tilde{y} which satisfies eq. (2.12) by using a 3×2 complex matrix R [54, 55]:

$$\tilde{y} = \frac{1}{\langle H \rangle} \tilde{M}^{1/2} R m_\nu^{1/2} U_{\text{PMNS}}^\dagger, \quad (2.14)$$

where R can be expressed in terms of a complex parameter z ,

$$R = \begin{pmatrix} 0 & \cos z & -\sin z \\ 0 & \sin z & \cos z \end{pmatrix}, \quad (2.15)$$

for normal hierarchy, and,

$$R = \begin{pmatrix} \cos z & -\sin z & 0 \\ \sin z & \cos z & 0 \end{pmatrix}, \quad (2.16)$$

for inverted hierarchy.

By using the above parametrization and eqs. (2.10), (2.14), we can determine the structure of the Yukawa coupling y_ν . For normal hierarchy, we obtain,

$$y_\nu^{1\ell} = c \frac{\sqrt{M_2 M_3 m_2 m_3}}{\langle H \rangle^2} \det U_{\text{PMNS}}^* \times U_{\ell 1}, \quad (2.17)$$

$$y_\nu^{2\ell} = \frac{\sqrt{M_2}}{\langle H \rangle} (\sqrt{m_2} U_{\ell 2}^* \cos z - \sqrt{m_3} U_{\ell 3}^* \sin z), \quad (2.18)$$

$$y_\nu^{3\ell} = \frac{\sqrt{M_3}}{\langle H \rangle} (\sqrt{m_2} U_{\ell 2}^* \sin z + \sqrt{m_3} U_{\ell 3}^* \cos z). \quad (2.19)$$

For inverted hierarchy,

$$y_\nu^{1\ell} = c \frac{\sqrt{M_2 M_3 m_1 m_3}}{\langle H \rangle^2} \det U_{\text{PMNS}}^* \times U_{\ell 3}, \quad (2.20)$$

$$y_\nu^{2\ell} = \frac{\sqrt{M_2}}{\langle H \rangle} (\sqrt{m_1} U_{\ell 1}^* \cos z - \sqrt{m_2} U_{\ell 2}^* \sin z), \quad (2.21)$$

$$y_\nu^{3\ell} = \frac{\sqrt{M_3}}{\langle H \rangle} (\sqrt{m_1} U_{\ell 1}^* \sin z + \sqrt{m_2} U_{\ell 2}^* \cos z). \quad (2.22)$$

Here, we used the unitarity of U_{PMNS} for the calculation of $y_\nu^{1\ell}$. These structures are important for the discussion of the flavor of the decay products of N_1 . We will discuss their effects on the energy spectrum of the neutrino flux from the decay of N_1 in section 5.

3 Inflation with the $B - L$ Higgs field

In this section, we consider an inflation model with the $B - L$ Higgs field. The potential for ϕ_{B-L} in eq. (2.2) can drive inflation of the Universe. By defining $\phi = \sqrt{2}|\phi_{B-L}|$, the potential is recast in the form of,

$$V(\phi) = \Lambda^4 \left(\frac{\phi^2}{\mu^2} - 1 \right)^2, \quad (3.1)$$

where $\mu^2 = 2v_{B-L}^2$ and $\Lambda^4 = \kappa v_{B-L}^4/4$, and we define $\mu > 0$. The phase direction can be gauged away.

The inflaton field can slow roll when $\mu \gg M_{\text{Pl}}$, either from the $|\phi| > \mu$ or $|\phi| < \mu$ region towards the minimum at $\phi = \mu$. In both cases, the slow-roll parameters at $\phi = \phi_0$ are given by [56],

$$\epsilon = \frac{M_{\text{Pl}}^2}{2} \left(\frac{V'}{V} \right)^2 = \frac{M_{\text{Pl}}^2}{2\mu^2} \left(\frac{\frac{4\phi_0}{\mu}}{\frac{\phi_0^2}{\mu^2} - 1} \right)^2, \quad \eta = M_{\text{Pl}}^2 \frac{V''}{V} = \frac{4M_{\text{Pl}}^2}{\mu^2} \frac{\frac{3\phi_0^2}{\mu^2} - 1}{\left(\frac{\phi_0^2}{\mu^2} - 1 \right)^2}. \quad (3.2)$$

The field value ϕ_0 at the pivot scale $k_0 = 0.002 \text{ Mpc}^{-1}$ is expressed in terms of the number of e -folds N :

$$N \simeq \frac{1}{M_{\text{Pl}}^2} \int_{\phi_{\text{end}}}^{\phi_0} \frac{V}{V'} d\phi = \frac{\mu}{M_{\text{Pl}}^2} \int_{\phi_{\text{end}}}^{\phi_0} \frac{\frac{\phi^2}{\mu^2} - 1}{\frac{4\phi}{\mu}} d\phi, \quad (3.3)$$

where the field value at the end of the inflation, ϕ_{end} , is obtained from,

$$1 \simeq M_{\text{Pl}}^2 \left(\frac{V'}{V} \right)^2 \Big|_{\phi=\phi_{\text{end}}} = \frac{M_{\text{Pl}}^2}{\mu^2} \left(\frac{\frac{4\phi_{\text{end}}}{\mu}}{\frac{\phi_{\text{end}}^2}{\mu^2} - 1} \right)^2. \quad (3.4)$$

The tensor-to-scalar ratio, r , and the spectral index, n_s , is expressed in terms of the slow-roll parameters as,

$$r = 16\epsilon, \quad n_s = 1 - 6\epsilon + 2\eta. \quad (3.5)$$

The Planck normalization sets the overall scale [57],

$$\left(\frac{V}{\epsilon} \right)^{1/4} \Big|_{\phi_0} = 6.4 \times 10^{16} \text{ GeV}, \quad (3.6)$$

and the observed spectral index is given by,

$$n_s = 0.9603 \pm 0.0073. \quad (3.7)$$

The results from the BICEP2 experiment prefer,

$$r = 0.20_{-0.05}^{+0.07}, \quad V^{1/4} = 2.0 \times 10^{16} \text{ GeV} \cdot \left(\frac{r}{0.16} \right)^{1/4}, \quad (3.8)$$

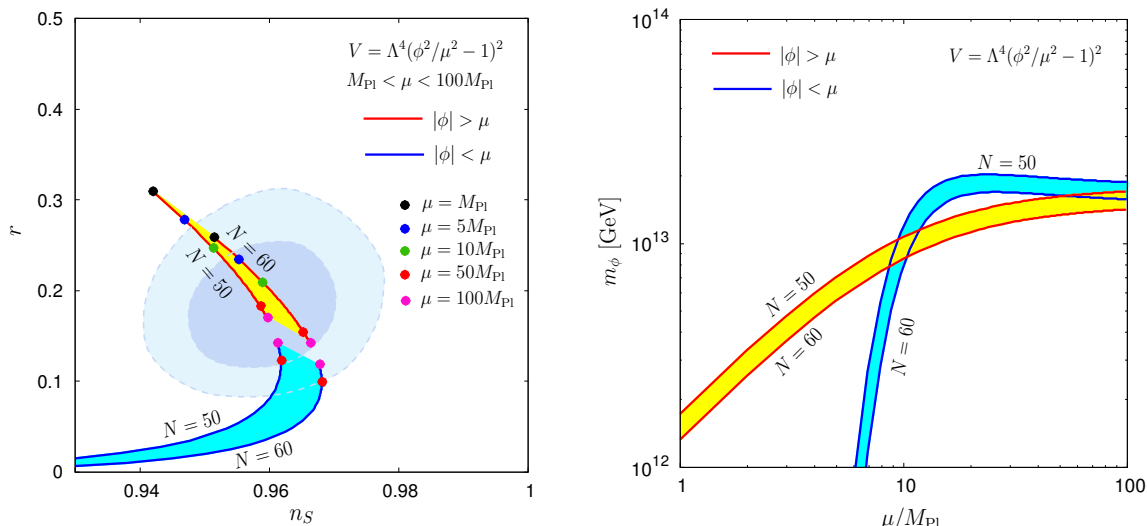


Figure 2. Predictions of the inflation model in the $r-n_s$ plane (left) and $m_\phi-\mu/M_{\text{Pl}}$ plane (right). The region favored by CMB observations (Planck+WP+highL+BICEP2) [58] are also shown in the left figure. The dark blue region corresponds to the region consistent with the BICEP2 at 1σ level whereas the light blue region does to that at 2σ level. The left figure is consistent with the result obtained in ref. [5].

when one combines the data from the Planck experiment.³ Here, the preferred range of r will be modified to $r = 0.16^{+0.06}_{-0.05}$ after subtracting the best available estimate for foreground dust [58].

The predictions for r and n_s is shown in figure 2 with varying μ . The region favored by the CMB observations are also shown. We see that for $N = 60$, $|\phi| > \mu$ and $\mu \gtrsim 5M_{\text{Pl}}$ is favored. The inflaton mass, $m_\phi = 2\sqrt{2}\Lambda^2/\mu$, as a function of μ/M_{Pl} is also shown in figure 2. For $|\phi| > \mu$ and $\mu \gtrsim 5M_{\text{Pl}}$, we find,

$$m_\phi \sim 10^{13} \text{ GeV}. \tag{3.9}$$

This value corresponds to a very small value of κ such as $\kappa \sim 10^{-12}$ for $\mu \sim 5M_{\text{Pl}}$. In the following discussion, we fix the inflaton mass at this value, and will see that the correct amount of the baryon asymmetry and the dark matter can be obtained after the decay of the inflaton fields.

4 Reheating by the inflaton decay

After the inflation, the decay of ϕ can produce the Standard Model particles. The dominant decay mode can either be into two right-handed neutrinos via the interaction term in eq. (2.1) or two Higgs fields (including the Goldstone modes) via the term in eq. (2.3).

³ We will not consider the tension between the data from the Planck satellite ($r < 0.11$) [57] and that from the BICEP2 experiment ($r \sim 0.2$). The tension can be relaxed if one considers a running spectral index, an extra relativistic component, non-zero neutrino mass [58, 59], an anti-correlation between tensor and scalar modes [60] or between tensor and isocurvature modes [61, 62]. See also [63–69] for other solutions.

In the case where the $\phi \rightarrow N_i N_i$ mode is dominated and for $\lambda_1 \ll \lambda_2$ and $M_3 > m_\phi$ which are justified later, the decay width is given by,

$$\Gamma_\phi = \frac{1}{2} \frac{m_\phi}{16\pi} \frac{M_2^2}{v_{B-L}^2} \left(1 - \frac{4M_2^2}{m_\phi^2}\right)^{3/2}. \quad (4.1)$$

By equating Γ_ϕ with the Hubble parameter $H(T_R)$ at the reheating temperature, T_R , we obtain,

$$T_R \simeq 2 \times 10^7 \text{ GeV} \left(\frac{M_2}{10^{12} \text{ GeV}}\right) \left(\frac{m_\phi}{10^{13} \text{ GeV}}\right)^{1/2} \left(\frac{v_{B-L}}{5M_{\text{Pl}}}\right)^{-1} \left(1 - \frac{4M_2^2}{m_\phi^2}\right)^{3/4}. \quad (4.2)$$

Here we used $T_R = (90/\pi^2 g_*(T_R))^{1/4} \sqrt{\Gamma_\phi M_{\text{Pl}}}$ and $g_*(T_R) = 106.75$, where $g_*(T_R)$ is the relativistic degrees of freedom in plasma at the temperature T_R . If the Higgs mode $\phi \rightarrow hh, WW, ZZ$ is the dominant decay channel through eq. (2.3), the reheating temperature can be arbitrarily higher than the above estimate. If T_R is higher than m_ϕ , the perturbative analysis of the reheating process becomes unreliable [70, 71]. Therefore, we restrict ourselves to the region of $T_R < m_\phi \sim 10^{13}$ GeV.

4.1 Leptogenesis

For the case where $\phi \rightarrow N_2 N_2$ is the dominant decay mode, the decay of N_2 can generate $B - L$ asymmetry by leptogenesis. The baryon-to-entropy ratio obtained from the non-thermal leptogenesis is [72, 73],

$$\frac{n_B}{s} = -\frac{28}{79} \cdot \frac{3}{2} \cdot \epsilon \cdot \frac{T_R}{m_\phi}, \quad (4.3)$$

where $(3/2)T_R/m_\phi$ is the number density $n_{N_2} \simeq n_\phi/2$ divided by the entropy density produced by the decay of ϕ . The ϵ factor is the magnitude of the CP violation [74]:

$$\epsilon \simeq -\frac{3}{16\pi} \frac{\text{Im}(y_\nu y_\nu^\dagger)_{23}^2}{(y_\nu y_\nu^\dagger)_{22}} \frac{M_2}{M_3}, \quad (4.4)$$

for $M_2 \ll M_3$. It is bounded by [51, 75],

$$|\epsilon| \lesssim \begin{cases} \frac{3}{16\pi} \frac{M_2}{\langle H \rangle^2} (m_3 - m_2) \sim 8 \times 10^{-5} \left(\frac{M_2}{10^{12} \text{ GeV}}\right), & \text{(Normal)} \\ \frac{3}{16\pi} \frac{M_2}{\langle H \rangle^2} (m_2 - m_1) \sim 2 \times 10^{-6} \left(\frac{M_2}{10^{12} \text{ GeV}}\right). & \text{(Inverted)} \end{cases} \quad (4.5)$$

Here, we take $\Delta m_\odot^2 = (0.0086 \text{ eV})^2$ and $\Delta m_A^2 = (0.048 \text{ eV})^2$ [76]. Therefore,

$$\frac{n_B}{s} \Big|_{\text{max}} \simeq \left(\frac{M_2}{10^{12} \text{ GeV}}\right)^2 \left(\frac{m_\phi}{10^{13} \text{ GeV}}\right)^{-1/2} \left(\frac{v_{B-L}}{5M_{\text{Pl}}}\right)^{-1} \times \begin{cases} 1 \times 10^{-10} \text{ (Normal)} \\ 2 \times 10^{-12} \text{ (Inverted)} \end{cases}. \quad (4.6)$$

For normal hierarchy, compared with the observed baryon-to-entropy ratio, $n_B/s \simeq 10^{-10}$ [77], we need $M_2 \gtrsim 10^{12}$ GeV. On the other hand, for inverted hierarchy, we need

$M_2 \gtrsim 10^{13}$ GeV which is on the border of the constraint: $m_\phi > 2M_2$. In any case, these result justify $M_3 > m_\phi$ which we assumed before.

If the Higgs mode is important, the branching ratio into M_2 is suppressed, and thus non-thermal leptogenesis becomes difficult. With fixed m_ϕ from the CMB observations, there is no freedom to make M_2 larger since the decay into N_2 becomes kinematically forbidden. Instead, if the reheating temperature is high enough, it is possible to produce N_2 thermally. The thermal leptogenesis is possible for 10^9 GeV $\lesssim M_2 \lesssim T_R$ [75, 78].

4.2 Dark matter abundance

The inflaton also decays into two N_1 's. The assumption that N_1 is long-lived makes it possible to identify this component to be the dark matter of the Universe.

The partial decay width is given by,

$$\Gamma(\phi \rightarrow N_1 N_1) = \frac{1}{2} \frac{m_\phi}{16\pi} \frac{M_1^2}{v_{B-L}^2} \left(1 - \frac{4M_1^2}{m_\phi^2}\right)^{3/2}. \quad (4.7)$$

By using the relation $H(T_R) \sim \Gamma_\phi \sim T_R^2/M_{Pl}$, and $n_{N_1}/s \simeq (3/2)(T_R/m_\phi)\text{Br}(\phi \rightarrow N_1 N_1)$, we find,

$$\Omega_{N_1}^{\text{NT}} \simeq 0.2 \left(\frac{M_1}{4 \text{ PeV}}\right)^3 \left(\frac{T_R}{3 \times 10^7 \text{ GeV}}\right)^{-1} \left(\frac{v_{B-L}}{5M_{Pl}}\right)^{-2}. \quad (4.8)$$

Here, we used $\Omega_{N_1}^{\text{NT}} = (M_1 n_{N_1}/s)/(\rho_c/s)_0$, where $(\rho_c/s)_0 \simeq 1.8 \times 10^{-9}$ GeV is the critical density divided by the entropy density today. The contribution from the thermal production from the scattering processes by the $U(1)_{B-L}$ gauge interaction is much smaller such as [38, 39],

$$\Omega_{N_1}^{\text{TH}} \sim 10^{-23} \left(\frac{M_1}{4 \text{ PeV}}\right) \left(\frac{T_R}{5 \times 10^7 \text{ GeV}}\right)^3 \left(\frac{v_{B-L}}{5M_{Pl}}\right)^{-4}. \quad (4.9)$$

This is estimated with the interaction between N_1 and the Standard Model fermions in plasma through the s-channel exchange of the $U(1)_{B-L}$ gauge boson. We summarize the allowed regions in figure 3. We see that one obtains the correct amount of the baryon asymmetry and the dark matter abundance at $M_1 \sim 1$ PeV and $M_2 \sim 10^{12}$ GeV within the region consistent with the BICEP2 results at the 1σ level. The PeV dark matter opens up an interesting possibility that the high energy neutrinos observed at the IceCube experiment [40, 41] are explained by the decay of N_1 , which will be studied in the next section. For a heavier N_1 , we see a region where the thermal leptogenesis works. There, a high enough reheating temperature is realized by the inflaton decay into Higgs fields through the coupling in eq. (2.3).

5 PeV neutrinos as a signal of decaying N_1

In this section, we discuss observational signatures of the right-handed neutrino dark matter. As discussed so far, the inflation, the baryon asymmetry and the correct amount of

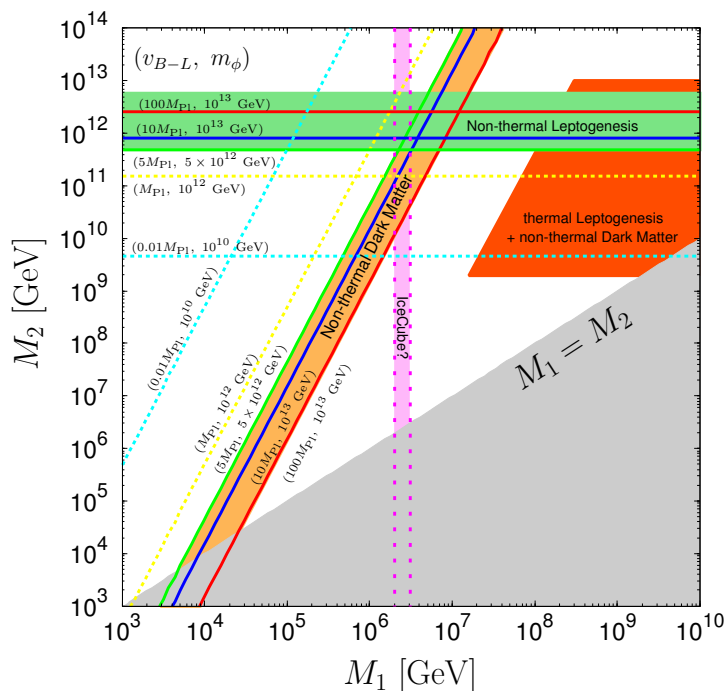


Figure 3. Consistent regions with neutrino masses and cosmological observations. The two shaded regions (green and light orange) are consistent with the BICEP2 at 1σ level respectively, and imply that the non-thermal leptogenesis works (green) and the dark matter is explained via the inflaton decay (light orange). Here, we assume normal hierarchy. We also show the mass range of N_1 favored by the IceCube experiment (pink shaded region). In the dark orange region where thermal leptogenesis is viable, the reheating temperature is treated as a free parameter satisfying $M_2 \leq T_R \leq m_\phi = 10^{13}$ GeV with $5M_{\text{Pl}} \leq v_{B-L}$. A high reheating temperature is realized by the decay into hh , WW and ZZ via the coupling in eq. (2.3).

dark matter can be explained for $M_1 = \mathcal{O}(1)$ PeV. Since there is no reason to expect that N_1 is absolutely stable, we have a chance to see high energy cosmic rays produced from the decay of N_1 . It is interesting that the PeV is indeed the energy region where an excess of high energy neutrinos events are observed at the IceCube experiment. In this section, we discuss the possibility that neutrino excess which is observed at IceCube experiment [40, 41] is explained by the decay products of N_1 .

5.1 The branching fractions of N_1

The partial decay widths of N_1 at tree level are,

$$\Gamma(N_1 \rightarrow \ell^- W^+) = \Gamma(N_1 \rightarrow \ell^+ W^-) = \frac{|y_\nu^{1\ell}|^2 M_1}{16\pi} \left(1 - \frac{m_W^2}{M_1^2}\right)^2 \left(1 + \frac{2m_W^2}{M_1^2}\right), \quad (5.1)$$

$$\Gamma(N_1 \rightarrow \nu_\ell Z) = \Gamma(N_1 \rightarrow \bar{\nu}_\ell Z) = \frac{|y_\nu^{1\ell}|^2 M_1}{32\pi} \left(1 - \frac{m_Z^2}{M_1^2}\right)^2 \left(1 + \frac{2m_Z^2}{M_1^2}\right), \quad (5.2)$$

$$\Gamma(N_1 \rightarrow \nu_\ell h) = \Gamma(N_1 \rightarrow \bar{\nu}_\ell h) = \frac{|y_\nu^{1\ell}|^2 M_1}{32\pi} \left(1 - \frac{m_h^2}{M_1^2}\right)^2. \quad (5.3)$$

For $M_1 \gg m_W, m_Z, m_h$, we can see that $\Gamma(N_1 \rightarrow \ell^\mp W^\pm) : \Gamma(N_1 \rightarrow \nu Z, \bar{\nu} Z) : \Gamma(N_1 \rightarrow \nu h, \bar{\nu} h) \simeq 2 : 1 : 1$ due to the equivalence theorem [79–82]. The lifetime of N_1 for $M_1 \gg m_W, m_Z, m_h$ is calculated as,

$$\tau_{N_1} = \left(\frac{M_1}{4\pi} \sum_{\ell} |y_{1\ell}|^2 \right)^{-1} \sim 8 \times 10^{28} \text{ s} \left(\frac{M_1}{1 \text{ PeV}} \right)^{-1} \left(\sum_{\ell} \left| \frac{y_{1\ell}}{10^{-29}} \right|^2 \right)^{-1}. \quad (5.4)$$

The branching fractions for each lepton family $\text{Br}(\ell) \equiv \text{Br}(N_1 \rightarrow \ell^\mp W^\pm, \nu_\ell Z, \bar{\nu}_\ell Z, \nu_\ell h, \bar{\nu}_\ell h)$ are determined by $y_\nu^{1\ell}$'s. For each neutrino mass hierarchy, by the assumption of eq. (2.10), $\text{Br}(\ell)$'s are completely determined by the PMNS matrix,

$$(\text{Br}(e), \text{Br}(\mu), \text{Br}(\tau)) = (|U_{e1}|^2, |U_{\mu 1}|^2, |U_{\tau 1}|^2), \quad (\text{Normal}) \quad (5.5)$$

$$(\text{Br}(e), \text{Br}(\mu), \text{Br}(\tau)) = (|U_{e3}|^2, |U_{\mu 3}|^2, |U_{\tau 3}|^2). \quad (\text{Inverted}) \quad (5.6)$$

We take $\sin^2 \theta_{12} = 0.31$, $\sin^2 \theta_{23} = 0.39$ and $\sin^2 \theta_{13} = 0.02$ [76], then, the numerical values of the branching fraction are given by,

$$(\text{Br}(e), \text{Br}(\mu), \text{Br}(\tau)) = (0.68, 0.24 + 0.02 \cos \delta, 0.08 - 0.02 \cos \delta), \quad (\text{Normal}) \quad (5.7)$$

$$(\text{Br}(e), \text{Br}(\mu), \text{Br}(\tau)) = (0.02, 0.38, 0.60). \quad (\text{Inverted}) \quad (5.8)$$

The branching fractions for normal hierarchy has small dependence on CP-violating phase δ . On the other hand, the branching fractions for inverted hierarchy is completely determined independent of δ .

5.2 Neutrino flux from decay of N_1

We have calculated the energy spectrum of neutrinos dN_ν/dE_ν from decay of N_1 by using PYTHIA 8.1 [83]. The neutrino spectrum for $M_1 = 2.3 \text{ PeV}$ is shown in figure 4. We have a sharp peak in the neutrino energy spectrum at $E_\nu = M_1/2$. In the case of inverted hierarchy, since the fractions of muon and tau are large compared to the normal hierarchy, the number of neutrinos is slightly larger around $E_\nu \sim 10^{5-6} \text{ GeV}$ due to the decay products of the muons and taus.

As the neutrino travels towards the Earth, the neutrinos change their flavors by the neutrino oscillation according to the following probabilities:

$$P(\nu_\ell \rightarrow \nu_{\ell'}) = P(\bar{\nu}_\ell \rightarrow \bar{\nu}_{\ell'}) = \sum_{i=1}^3 |U_{\ell i} U_{\ell' i}|^2 \simeq \begin{pmatrix} 0.55 & 0.27 + 0.02 \cos \delta & 0.18 - 0.02 \cos \delta \\ 0.27 + 0.02 \cos \delta & 0.36 - 0.02 \cos \delta & 0.37 + 0.00 \cos \delta \\ 0.18 - 0.02 \cos \delta & 0.37 + 0.00 \cos \delta & 0.45 + 0.02 \cos \delta \end{pmatrix}. \quad (5.9)$$

In figure 5, we show the energy spectrum of the neutrinos after the oscillation.

We estimate the observed flux of neutrinos on the Earth in the following way. We have two classes of contribution from the decaying dark matter; one is from halo of our galaxy,

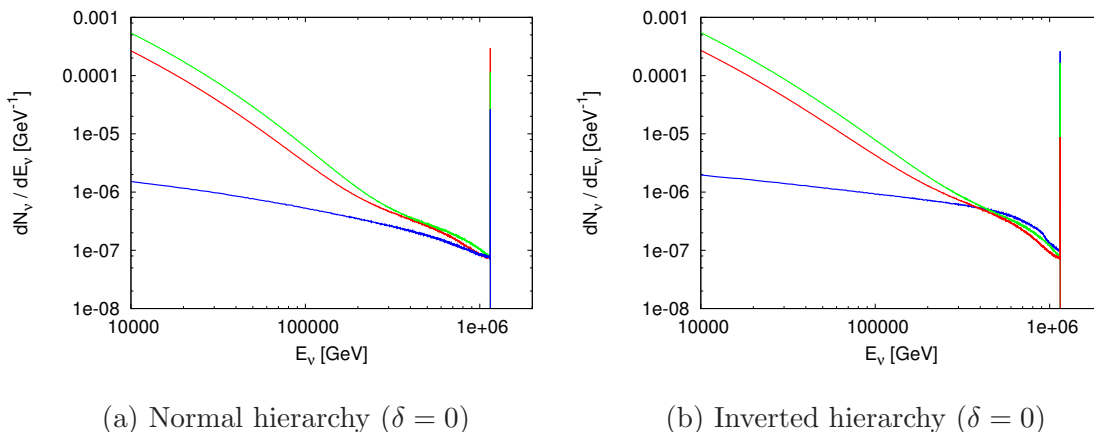


Figure 4. dN_ν/dE_ν for $M_1 = 2.3$ PeV when produced by the decay of N_1 . We take normal hierarchy with $\delta = 0$ in left figure and inverted hierarchy in right figure. Red, green and blue lines show the spectrum of $\nu_e + \bar{\nu}_e$, $\nu_\mu + \bar{\nu}_\mu$ and $\nu_\tau + \bar{\nu}_\tau$, respectively.

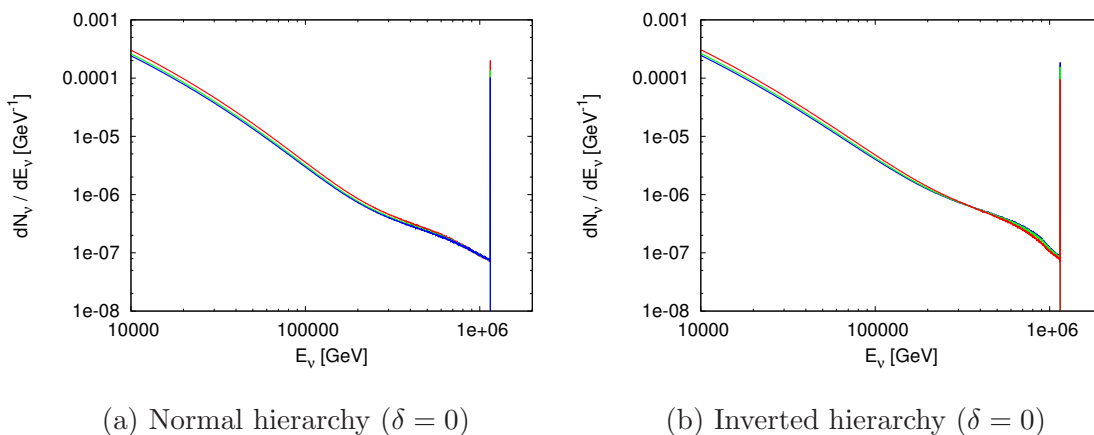


Figure 5. dN_ν/dE_ν for $M_1 = 2.3$ PeV which takes into account the effect of the neutrino oscillation (See eq. (5.9)). We take normal hierarchy in left figure and inverted hierarchy in right figure. In both figure, we take $\delta = 0$. Red, green and blue lines show the spectrum of $\nu_e + \bar{\nu}_e$, $\nu_\mu + \bar{\nu}_\mu$ and $\nu_\tau + \bar{\nu}_\tau$, respectively.

and another is from extra-galactic. For a review of the calculation of the neutrino flux, e.g., see ref. [84]. The halo contribution which is averaged over the full sky is proportional to dN_ν/dE_ν :

$$\frac{d\Phi_{\text{halo}}}{dE_\nu} = D_{\text{halo}} \frac{dN_\nu}{dE_\nu}, \tag{5.10}$$

where D_{halo} is determined by the halo density profile $\rho_{\text{halo}}(r)$,

$$D_{\text{halo}} = \frac{1}{4\pi} \int_{-1}^1 d\cos\theta \int_0^{2\pi} d\phi \left(\frac{1}{4\pi M_1 \tau_{N_1}} \int_0^\infty ds \rho_{\text{halo}}(r(s, \theta, \phi)) \right). \tag{5.11}$$

The parameter s in the integral of eq. (5.11) is the distance from the Earth, and it is related to the distance r from the galactic center as, $r(s, \theta, \phi) = \sqrt{s^2 + R_\odot^2 - 2sR_\odot \cos \theta \cos \phi}$. Here, R_\odot is the distance of the Sun to the galactic center, and we take its value as 8.0 kpc [85]. For the calculation of D_{halo} , we adopt the Navarro-Frenk-White (NFW) density profile [86],

$$\rho_{\text{halo}}(r) = \rho_\odot \frac{(R_\odot/r_c)(1 + R_\odot/r_c)^2}{(r/r_c)(1 + r/r_c)^2}, \quad (5.12)$$

and take $r_c = 20$ kpc and $\rho_\odot = 0.4$ GeV cm⁻³ [87]. Then, D_{halo} is calculated as,

$$D_{\text{halo}} = 1.7 \times 10^{-13} \left(\frac{1 \text{ PeV}}{M_1} \right) \left(\frac{10^{28} \text{ s}}{\tau_{N_1}} \right) \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}. \quad (5.13)$$

Extra galactic contribution is redshifted because of the expansion of the Universe. Their contribution is written by,

$$\frac{d\Phi_{\text{eg}}}{dE_\nu} = \frac{\Omega_{\text{DM}} \rho_c c}{4\pi M_1 \tau_{N_1}} \int_0^\infty \frac{dz}{H(z)} e^{-s(E_\nu, z)} \frac{dN_\nu}{dE} \Big|_{E=(1+z)E_\nu}, \quad (5.14)$$

where we estimate the integrand just from $z = 0$ to z_{eq} for simplicity and hence also neglect the contribution from the dark matter which had decayed at the radiation dominated era, because we assume the dark matter mass is around PeV and the energy of neutrino from early universe is too low to explain the IceCube excess. In eq. (5.14), $H(z) = H_0 \sqrt{\Omega_\Lambda + \Omega_m(1+z)^3}$ is the Hubble expansion rate at the redshift z . $c = 3.0 \times 10^{10}$ cm s⁻¹ is the speed of light. $s(E_\nu, z)$ is neutrino opacity, which is estimated as $s(E_\nu, z) \sim 10^{-17} (1+z)^{7/2} (E_\nu/1\text{TeV})$ for $z < z_{\text{eq}}$ [88]. However, in the present situation, this effect is negligibly small. Then, we take $s(E_\nu, z)$ to be zero for an approximation. For the cosmological parameters, we take $\Omega_\Lambda = 0.68$, $\Omega_m = 0.32$, $\Omega_{\text{DM}} = 0.27$, $H_0 = 67$ km s⁻¹ Mpc⁻¹, $\rho_c = 3H_0^2 M_{\text{Pl}}^2 \simeq 4.7 \times 10^{-6}$ GeV cm⁻³ and $z_{\text{eq}} = 3.4 \times 10^3$. These values are derived from the Planck data [77].

Finally, the expected number of events at the IceCube detector per 662 days with given energy is calculated as,

$$N(E_0 \leq E \leq E_1) = 4\pi \times 662 \text{ days} \times \sum_{\ell=e, \mu, \tau} \int_{E_0}^{E_1} dE_\nu \left(\frac{d\Phi_{\text{halo}}^{(\nu_\ell + \bar{\nu}_\ell)}}{dE_\nu} + \frac{d\Phi_{\text{eg}}^{(\nu_\ell + \bar{\nu}_\ell)}}{dE_\nu} \right) \sigma_{\text{eff}}^{(\nu_\ell)}(E_\nu), \quad (5.15)$$

where $\sigma_{\text{eff}}^{(\nu_\ell)}$ is the neutrino effective area for each flavor which is given in refs. [41, 89]. The IceCube experiment observed 28 events with deposited energies between 30 and 1200 TeV, and the expected number of events from atmospheric muons and neutrinos is $10.6_{-3.6}^{+5.0}$ [41]. For 2.3 PeV dark matter, the total expected number of events for each pattern of the neutrino mass hierarchy is,

$$N(30 \text{ TeV} \leq E_\nu) = 10.8 \times (\tau_{N_1}/10^{28} \text{ s})^{-1}, \quad (\text{Normal}) \quad (5.16)$$

$$N(30 \text{ TeV} \leq E_\nu) = 13.7 \times (\tau_{N_1}/10^{28} \text{ s})^{-1}. \quad (\text{Inverted}) \quad (5.17)$$

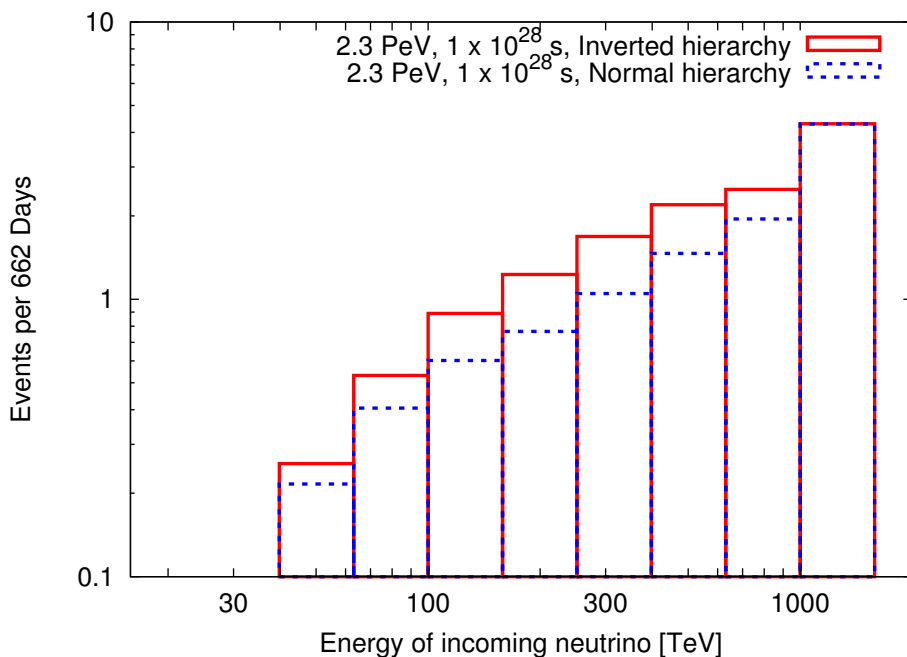


Figure 6. Number of events per 662 days at the IceCube experiment from neutrinos with given energy. For the parameter of N_1 , we take $M_1 = 2.3$ PeV, $\tau_{N_1} = 10^{28}$ s and CP-violating phase $\delta = 0$. We assume normal hierarchy for dotted blue boxes and inverted hierarchy for solid red boxes.

From this estimate, we see that the total excess can be explained for $\tau_{N_1} \simeq 1 \times 10^{28}$ s for both normal and inverted hierarchy. We also show the energy distribution of the neutrinos in figure 6. In this figure, we take $M_1 = 2.3$ PeV and $\tau_{N_1} = 10^{28}$ s.

The IceCube experiment provides the data of the event rate per the deposited energies in the detector in figure 4 in ref. [41]. Note that our results in figure 6 are, in contrast, those for incoming neutrino energies, and thus the deposited ones should be smaller due to escaping neutrinos and muons. One needs to take into account the correction when the shape of the distribution is compared. For $M_{N_1} = 2.3$ PeV, the expected number of neutrinos with the energy higher than 1 PeV is,

$$N(1000 \text{ TeV} \leq E_\nu) = 4.3 \times (\tau_{N_1}/10^{28} \text{ s})^{-1}, \quad (5.18)$$

for both normal and inverted hierarchy. Thus, by assuming that the deposited energy is equal to that of incoming neutrinos, the two observed neutrino events around PeV energies can be explained for $\tau_{N_1} \simeq 2 \times 10^{28}$ s. We can expect more sub-PeV events for the inverted hierarchy than the normal hierarchy. Implications from the IceCube experiment will be important to distinguish neutrino models.

6 Summary

In this paper, we considered a minimalistic cosmological scenario based on the $U(1)_{B-L}$ extended Standard Model. The model consistently explains the neutrino masses, the in-

flation, the baryon asymmetry of the Universe and dark matter abundance, which are left unexplained in the Standard Model. If both the baryon asymmetry and the dark matter abundance are explained directly by the inflaton decay, we obtain the mass of N_1 and the second lightest right-handed neutrino N_2 to be 1 PeV and 10^{12} GeV, respectively.

Interestingly, the mass of N_1 , PeV, turns out to be the energy scale of the excess of the neutrino events at the IceCube experiment. We see that the PeV neutrino events can be explained by the decaying N_1 with its lifetime being $\mathcal{O}(10^{28})$ s. Predicted number of neutrinos with sub-PeV energies depends on the neutrino mass hierarchy and the CP violating phase. Further observations of high energy neutrino events may, in principle, provide information on the flavor structures in the neutrino sector.

If the coupling between the Standard Model Higgs and the $B - L$ Higgs field is significant, the reheating temperature can be higher than the second lightest right-handed neutrino N_2 , depending on the coupling. In such a case, thermal leptogenesis is possible whereas the dark matter should be heavier than $\mathcal{O}(10)$ PeV to explain the abundance by the inflaton decay.

Acknowledgments

We would like to thank Aya Ishihara for the explanation of the IceCube experiments. This work is supported by JSPS Grant-in-Aid for Young Scientists (B) (No. 23740165 [RK], No. 25800169 [TH]), MEXT Grant-in-Aid for Scientific Research on Innovative Areas (No. 25105011 [RK]) and JSPS Research Fellowships for Young Scientists [RS].

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