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Large invisible decay of a Higgs boson to neutrinos

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We show that the standard model (SM)-like Higgs boson may decay into neutrinos with a sizable decay branching ratio in one well-known two Higgs doublet model, so-called neutrinophilic Higgs model. This could happen if the mass of the lighter extra neutral Higgs boson is smaller than one half of the SM-like Higgs boson mass. The definite prediction of this scenario is that the rate of the SM-like Higgs boson decay into diphoton normalized by the SM value is about 0.9. In the case that a neutrino is Majorana particle, a displaced vertex of right-handed neutrino decay would be additionally observed. This example indicates that a large invisible Higgs boson decay could be irrelevant to dark matter.

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I. INTRODUCTION

The newly discovered particle at the Large Hadron Collider (LHC) is now identified as a Higgs boson [1,2]. Its measured properties, such as spin, parity, and couplings, are consistent with the Higgs boson h in the standard model (SM) of particle physics [3–6] within uncertainties, which are not very small yet. Possible deviations from the SM prediction on the Higgs boson also have been examined.

One of those is an invisible Higgs boson decay. Actually, an invisible Higgs boson decay occurs even in the SM through an off-shell Z boson Z pair into four neutrinos ν, as $h \to Z^* Z^* \to 2\nu 2\bar{\nu}$. Its branching ratio in the SM is of the order of 10^{-3}. If once it is found with a larger branching ratio than that due to SM processes, this must be a sign of a beyond the SM (BSM). Such BSM models include, for instance, a light neutralino in supersymmetric models [7], a Majoron [8], graviscalars [9], fourth generation neutrino [10], and Higgs portal dark matter [11]. Searches of invisible decays of the Higgs boson h have been carried out and to date only an upper bound on the branching ratio of the invisible decay has been obtained [12,13].

In this paper, we show that the Higgs boson h would decay into four neutrinos through an extra Higgs boson, which can be seen as the invisible decay, in a class of the two Higgs doublet model (THDM). The remarkable feature in this scenario is that the invisible final states are a SM particle, neutrinos, compared with other BSM models mentioned above where final invisible states are new hypothetical particles, such as a supersymmetric particle or dark matter. We consider the so-called neutrinophilic THDM [14–16], where one Higgs doublet provides the mass of the SM fermions, while the other generates neutrino Dirac masses with its small vacuum expectation value (VEV). Phenomenology of the charged Higgs boson was studied in Refs. [17,18]. In this paper, we will study a possible phenomenology of neutral Higgs bosons in those models. Because of these Yukawa couplings, both extra CP-even and extra CP-odd neutral Higgs bosons, H and A, respectively, couple mostly with neutrinos. Thus, through interactions between the SM-like Higgs boson h and the extra Higgs boson H (A), as the Z boson makes, new decay processes

$$h \to H^{(*)} H^{(*)} \quad \text{or} \quad A^{(*)} A^{(*)} \to 2\nu 2\bar{\nu} \tag{1}$$

arise.[1] If the intermediate H or A is off shell, the resultant contribution is comparable to the SM contribution by the Z boson and is not so large. However, if either H or A is on shell, the resultant invisible decay width is large.

II. NEUTRINOPHILIC TWO HIGGS DOUBLET MODEL

The Higgs sector is of the so-called neutrinophilic THDM, where one Higgs doublet Φ_1 with its VEV v_1 generates the mass of the SM fermions, while the other Φ_2 generates neutrino Dirac masses through its VEV $v_2 \ll v_1$. Such a Yukawa coupling is realized by introducing the softly broken Z_2-parity charge assigned as in Table I. The Yukawa interaction is given by

$$\begin{aligned}\mathcal{L}_Y = & -y_{\ell_\alpha} \overline{L_\alpha} \Phi_1 \ell_{R_\alpha} - y_{u_\alpha} \overline{Q_\alpha} \tilde{\Phi}_1 u_{R_\alpha} - y_{d_\alpha} \overline{Q_\alpha} \Phi_1 d_{R_\alpha} \\ & - y_{\alpha i} \overline{L_\alpha} \tilde{\Phi}_2 \nu_{R_i} + \text{H.c.},\end{aligned} \tag{2}$$

where $\tilde{\Phi} = i\sigma_2 \Phi^*$, Q (L) is the left-handed $SU(2)$ doublet quark (lepton), and u_R, d_R, e_R, and ν_R are the right-handed (RH) $SU(2)$ singlet fermions, respectively. α denotes flavor where we neglect mixing in quarks and i represents the generation index of RH neutrinos. If we admit lepton

[1] This possibility was briefly mentioned in Ref. [17] for a very heavy SM-like Higgs boson in a different type of neutrinophilic Higgs model.

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TABLE I. The assignment of Z_2 parity and lepton number.

Fields	Z_2 parity	Lepton number
First Higgs doublet, Φ_1	+	0
Second Higgs doublet, Φ_2	−	0
Lepton doublet, L	+	1
Right-handed neutrino, ν_R	−	1
Right-handed charged lepton, ℓ_R	+	1
Others	+	0

number violation in theory, the lepton number violating Majorana mass term

$$\mathcal{L}_M = -\frac{1}{2}\overline{\nu^c_{R_i}} M_i \nu_{R_i} \tag{3}$$

also can be introduced [14]. The scalar potential is given by

$$\begin{aligned} V &= \mu_1^2|\Phi_1|^2 + \mu_2^2|\Phi_2|^2 - (\mu_{12}^2\Phi_1^\dagger\Phi_2 + \mathrm{H.c.}) \\ &\quad + \lambda_1|\Phi_1|^4 + \lambda_2|\Phi_2|^4 + \lambda_3|\Phi_1|^2|\Phi_2|^2 + \lambda_4|\Phi_1^\dagger\Phi_2|^2 \\ &\quad + \left\{\frac{\lambda_5}{2}(\Phi_1^\dagger\Phi_2)^2 + \mathrm{H.c.}\right\}, \end{aligned} \tag{4}$$

where μ_{12} is the soft breaking parameter of the Z_2 parity, as introduced above. Conditions that the potential (4) is bounded from below and a stable vacuum are given by [19]

$$\begin{aligned} &\lambda_1 > 0, \qquad \lambda_2 > 0, \\ &2\sqrt{\lambda_1\lambda_2} + \lambda_3 + \min[0, \lambda_4 - |\lambda_5|] > 0. \end{aligned} \tag{5}$$

Components in two Higgs doublets, each with a VEV, are parametrized as

$$\Phi_1 = \begin{pmatrix} \phi_1^+ \\ \frac{v_1 + h_1 + i a_1}{\sqrt{2}} \end{pmatrix}, \qquad \Phi_2 = \begin{pmatrix} \phi_2^+ \\ \frac{v_2 + h_2 + i a_2}{\sqrt{2}} \end{pmatrix}. \tag{6}$$

Following the concept of neutrinophilic Higgs model, we take $v_1 \simeq v \simeq 246$ GeV and $v_2 \ll v_1$. The smallness of v_2 is due to the small μ_{12}^2 [14]. We define $\tan\beta = v_1/v_2$ as usual; this corresponds to $\tan\beta \gg 1$. The states h_1 and h_2 are diagonalized to the mass eigenstates (h and H) as

$$\begin{pmatrix} h_1 \\ h_2 \end{pmatrix} = \begin{pmatrix} \cos\alpha & -\sin\alpha \\ \sin\alpha & \cos\alpha \end{pmatrix} \begin{pmatrix} h \\ H \end{pmatrix}. \tag{7}$$

Because of $\tan\beta \gg 1$, ϕ_1^+ and a_1 are mostly eaten by the W and Z bosons, while we can identify the physical states as $H^+ \simeq \phi_2^+$, $A \simeq a_2$, $H \simeq h_2$, and $h \simeq h_1$. Then, the mixing angle is found to be

$$\sin\alpha \simeq \frac{v_2}{v_1}. \tag{8}$$

Automatically almost, the so-called SM limit $\sin(\beta - \alpha) = 1$ is realized. Because of $\tan\beta \gg 1$, $\alpha \simeq 0$ is realized. The Higgs boson h with the mass of 125 GeV is also composed as $h \simeq h_1$. From Eq. (2), the Yukawa interactions of extra neutral Higgs bosons are written as

$$\begin{aligned} \mathcal{L}_Y \supset &-\sum_{f=u_i,d_i,\ell_i} \frac{m_f}{v}\frac{\sin\alpha}{\sin\beta}\bar{f}Hf - i\frac{m_{u_i}}{v}(-\cot\beta)\bar{u}_i A\gamma_5 u_i \\ &- i\sum_{f=d_i,\ell_i}\frac{m_f}{v}\cot\beta\bar{f}A\gamma_5 f - \frac{y_{\alpha i}}{\sqrt{2}}\cos\alpha\overline{\nu_\alpha}HP_R\nu_i \\ &+ i\frac{y_{\alpha i}}{\sqrt{2}}\sin\beta\overline{\nu_\alpha}AP_R\nu_i + \mathrm{H.c.} \end{aligned} \tag{9}$$

We find that H or A decays into mostly neutrinos for

$$y_{\alpha i} \gg \frac{\sqrt{2}m_f}{v\tan\beta}. \tag{10}$$

Masses of extra Higgs bosons are given by

$$m_H^2 = \mu_2^2 + \frac{\lambda_3 + \lambda_4 + \lambda_5}{2}v^2, \tag{11}$$

$$m_A^2 = \mu_2^2 + \frac{\lambda_3 + \lambda_4 - \lambda_5}{2}v^2, \tag{12}$$

$$m_{H^\pm}^2 = \mu_2^2 + \frac{\lambda_3}{2}v^2. \tag{13}$$

To be consistent with the electroweak precision test, one neutral Higgs boson mass should be close to the charged Higgs boson mass as

$$m_{H^+} \simeq m_A \quad \text{or} \quad m_{H^+} \simeq m_H. \tag{14}$$

Interactions of extra Higgs bosons with h is

$$\mathcal{L} \supset -v\left(\frac{1}{2}(\lambda_3 + \lambda_4 - \lambda_5)A^2 + \frac{1}{2}(\lambda_3 + \lambda_4 + \lambda_5)H^2 + \lambda_3|H^+|^2\right)h. \tag{15}$$

III. EXOTIC SM-LIKE HIGGS BOSON DECAY

Now we consider a case where either H or A is light enough to be produced on shell by the h decay. There are two mass spectra of Higgs bosons that are as consistent with the electroweak precision test:

$$m_H < m_h/2 \ll m_{H^+} \simeq m_A \tag{16}$$

and

$$m_A < m_h/2 \ll m_{H^+} \simeq m_H. \tag{17}$$

From the mass formulas (11), (12), and (13), we find that mass spectra (16) and (17) can be realized for

$$0 > \lambda_4 \simeq \lambda_5 \quad \text{and} \quad 0 > \lambda_4 \simeq -\lambda_5, \tag{18}$$

respectively.

With couplings (9) and (15), if h decays into H or A, which decays into neutrinos, then this fraction is measured as its invisible decay. The decay width of $h \to HH$ is given by

$$\Gamma(h \to HH) = \lambda_H^2 \frac{v^2}{32\pi m_h} \left(1 - \frac{4m_H^2}{m_h^2}\right)^{1/2}, \tag{19}$$

with $\lambda_H = \lambda_3 + \lambda_4 + \lambda_5$. For the case of $h \to AA$, we obtain the same result just by replacing λ_H and m_H with $\lambda_A = \lambda_3 + \lambda_4 - \lambda_5$ and m_A, respectively. The LHC constraints on exotic decay modes, $h \to HH$ or $h \to AA$, indicate that $\lambda_H(\lambda_A) \lesssim \mathcal{O}(10^{-2})$ is allowed. We define

$$Br(h \to HH) = \frac{\Gamma(h \to HH)}{\Gamma(h \to all)}, \tag{20}$$

which is shown in Fig. 1. By combining the constraint on $\lambda_H(\lambda_A)$ and Eq. (18), we find that, for the light H (A),

$$\lambda_3 \simeq -\lambda_4 - (+)\lambda_5 \tag{21}$$

should be positive and of $\mathcal{O}(1)$, which leads to the deviation in the diphoton decay rate of h from the SM value [20], as shown in Fig. 2. Here, the brown and magenta shaded regions correspond to the region where the extra light Higgs boson is tachyonic and its on shell production is kinematically forbidden, respectively. One can find that this scenario predicts $R_{\gamma\gamma} \simeq 0.9$. The signal strength of $h \to \gamma\gamma$ has been reported as 1.17 ± 0.27 by ATLAS [21] and $1.14^{+0.26}_{-0.23}$ by CMS [22].

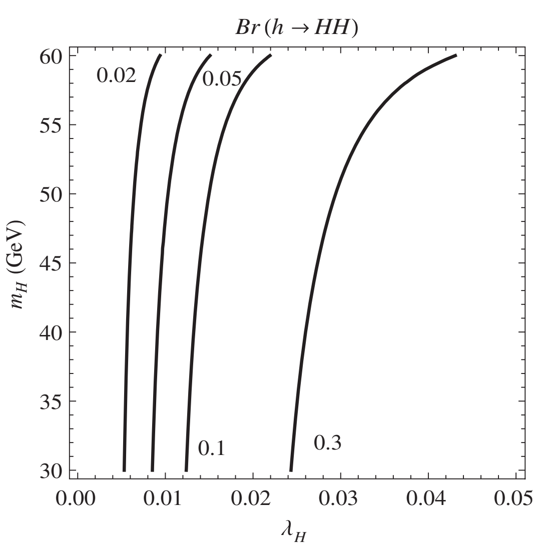

FIG. 1. Contours of the decay branching ratio of h into HH in the $\lambda_H - m_H$ plane.

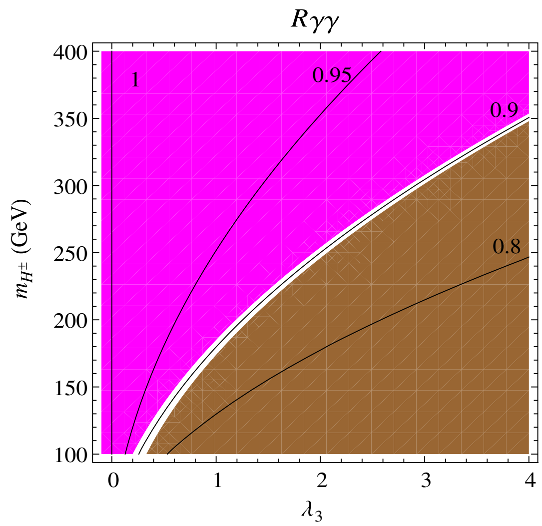

FIG. 2 (color online). Contours of $h \to \gamma\gamma$ rate normalized by the SM value on the $\lambda_3 - m_{H^\pm}$ plane. In the magenta shaded region, the lighter extra Higgs boson is too heavy to be produced by the h decay. In the brown shaded region, the extra light Higgs boson is tachyonic. Here, the extra Higgs boson mass is evaluated with $\lambda_4 = -\lambda_5$.

In the following sections, we discuss more detailed phenomenology which depends on neutrino mass nature.

IV. DIRAC NEUTRINO CASE

The condition (10) is satisfied by a large $\tan\beta$ for the Dirac neutrino case. Thus, the light H is, in practice, invisible and we have $Br(h \to inv) = Br(h \to HH)$ (shown in Fig. 1). The same is true for a light A as well.

The constraint on the charged Higgs boson, which decays into a lepton and a neutrino, is, in fact, stringent. This decay mode is similar to that of a slepton in supersymmetric models [17] and masses of the first and the second generation slepton is constrained as $m_{\tilde{l}} \gtrsim 300$ GeV by ATLAS [23] or $m_{\tilde{l}} \gtrsim 260$ GeV by CMS [24]. Although some differences due to its decay branching ratio exist [17], roughly speaking, there is a similar bound on H^+. Referring to Fig. 2, we find that a rather large coupling $\lambda_3 \gtrsim 3(2)$ for the ATLAS (CMS) bound is required.

V. MAJORANA NEUTRINO CASE

With the presence of the Majorana mass term (3), the neutrino mass matrix is given as

$$\mathcal{M} = \begin{pmatrix} m_\nu^{(\text{loop})} & \frac{y}{\sqrt{2}} v_2 \\ \frac{y}{\sqrt{2}} v_2 & M_k \end{pmatrix}, \tag{22}$$

with the radiative generated mass $m_\nu^{(\text{loop})}$ [25],

$$m_\nu^{(\mathrm{loop})} = \sum_k \frac{y_{\alpha k} y_{k\beta}^T M_k}{16\pi^2}\left(\frac{m_H^2}{m_H^2 - M_k^2}\ln\frac{m_H^2}{M_k^2} - \frac{m_A^2}{m_A^2 - M_k^2}\ln\frac{m_A^2}{M_k^2}\right). \tag{23}$$

We obtain a light neutrino mass

$$(m_\nu)_{\alpha\beta} = m_\nu^{(\mathrm{loop})} - \sum_k \frac{y_{\alpha k} y_{k\beta}^T v_2^2/2}{M_k}, \tag{24}$$

the mass of a heavier RH-like neutrino $m_{N_R} \simeq M_k$ and the left-right mixing angle θ [26],

$$\sin\theta \simeq \frac{y v_2}{\sqrt{2} M_k}. \tag{25}$$

$m_\nu^{(\mathrm{loop})}$ could be indeed dominant—or at least comparable with tree level seesaw mass—because of $\lambda_5 = \pm\mathcal{O}(1)$ from Eqs. (18) and (21).

The charged Higgs boson decays into cb or tb, depending on the mass in the neutrinophilic Higgs model [18]. The results for the $H^\pm \to tb$ mode with a normalizing production cross section of 1 pb can be found in Ref. [27]. However, this constraint is not so stringent because the actual production cross section is not so large. The LHC data constrains the mass of $H^\pm$ decaying into bc between 90 and 150 GeV [28] and $H^\pm$ decaying into $\tau\nu$ [29].

We note here one cosmological argument on the Majorana neutrino case. The lepton number violation by the Majorana nature of neutrino plays an important role in cosmology. Several cosmological discussions on neutrinophilic Higgs model were held in Refs. [30–32]. One of them is an enhancement $\Delta L = 2$ washout process by large Yukawa couplings and relatively light RH neutrinos in a neutrinophilic Higgs model [30]. Although a discussion of baryogenesis is beyond the scope and purpose of this paper, as a necessary condition to have nonvanishing baryon asymmetry in our Universe, we roughly evaluate the condition of no strong washout of lepton asymmetry,[2] provided a nonvanishing lepton asymmetry has been generated by any means of a higher energy physics process. If this condition were violated, it would be difficult to explain nonvanishing baryon asymmetry in our Universe, because any generated lepton asymmetry is washed out. The washout rate is given by $\Gamma_{\Delta L=2} \simeq y^4 T$. The condition $\Gamma_{\Delta L=2} < H(T)$ at $T \simeq 100$ GeV, with H being the cosmic expansion rate, is rewritten as

$$y \lesssim 10^{-4}, \tag{26}$$

which would be regarded as a cosmologically favored region.[3]

Now we discuss the decay of H or A. For $m_{N_R} < m_{H/A}$, an extra neutral Higgs boson H (A) decay produces one light left-handed-like neutrino ν and the other heavy RH-like neutrino N_R. The amplitude is calculated as

$$\begin{aligned}\overline{|\mathcal{M}(H/A\bar{\nu}N_R)|^2} &= 2|y|^2(p_1 \cdot p_2)\\ &= |y|^2((p_1+p_2)^2 - m_{N_R}^2),\end{aligned} \tag{27}$$

where p_1 and p_2 are outgoing momentum of ν and N_R, respectively. In addition, m_ν is neglected and indexes of y are omitted. The decay width is given by

$$\begin{aligned}&\Gamma(H/A \to \nu N_R)\\ &\quad= \frac{1}{16\pi {m_{H/A}}^3}\sum |y|^2({m_{H/A}}^2 - m_{N_R}^2)^2.\end{aligned} \tag{28}$$

Here, the summation $\sum$ is taken for all kinematically possible modes. An extra Higgs boson decays into SM fermions, mostly the bottom quark, through a tiny mixing α. Thus, its decay width is strongly suppressed by a large $\tan\beta$ as

$$\Gamma(H/A \to b\bar{b}) \simeq \frac{3}{8\pi}\left(\frac{m_b}{v\tan\beta}\right)^2 m_{H/A}. \tag{29}$$

Here, we define

$$Br(H/A \to inv) = \frac{\Gamma(H/A \to \nu N_R)}{\Gamma(H/A \to b\bar{b}) + \Gamma(H/A \to \nu N_R)}, \tag{30}$$

and

$$\begin{aligned}&Br(h \to 2H/2A \to 2\nu 2N_R)\\ &\quad= Br(h \to HH/AA)Br(H/A \to inv).\end{aligned} \tag{31}$$

Figure 3 shows the contour plot of the invisible decay branching ratio with thick black lines as well as the contour of Eq. (25) with thin blue thin lines of H and A, respectively. In both cases, the invisible decay branching ratio is large for $v_2 < 0.1$ GeV. The dashed green (thick) lines are contours of the typical size of Yukawa coupling $y = 10^{-5}(10^{-4})$ estimated from Eq. (24) with the atmospheric neutrino mass. As discussed above, $y \simeq 10^{-4}$ would be critical when we consider nonvanishing baryon asymmetry in our Universe. In both panels, neutrino masses dominantly come from the tree level seesaw at the upper left region, and do from m_ν^{loop} at the lower right region.

[2]This is because lepton asymmetry is a potential source of the baryon asymmetry in our Universe in the large class of baryogenesis scenario [33].

[3]One known mechanism of baryogenesis which works without any lepton asymmetry is "baryogenesis via neutrino oscillation" [34,35].

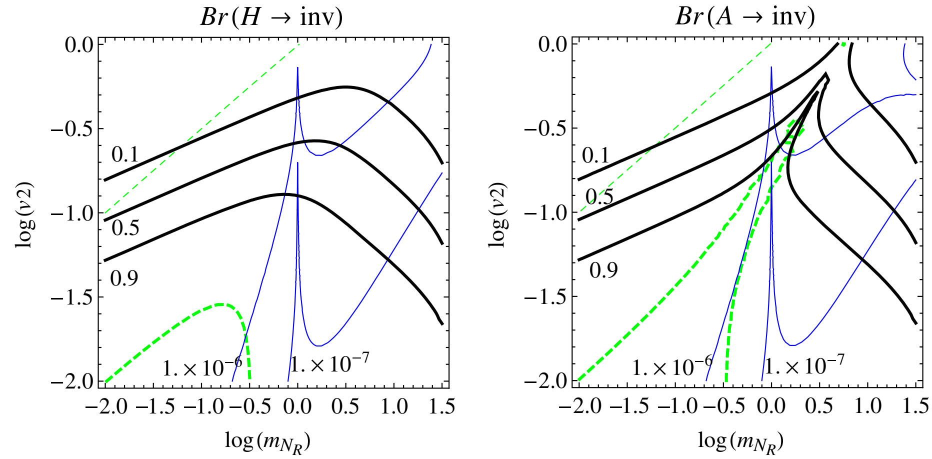

FIG. 3 (color online). Contours of the invisible decay branching ratio of the lighter extra neutral Higgs boson, H or A, for $m_\nu = 0.05$ eV on the $\log_{10}(m_{N_R}/\mathrm{GeV}) - \log_{10}(v_2/\mathrm{GeV})$ plane. The thin blue lines are contours of $\sin\theta$. For the dashed lines, see the text. Here we take $m_H, m_A = 60$ and 200 GeV (left panel) and 200 and 60 GeV (right panel).

In the light A case shown in the right panel of Fig. 3, the destructive cancellation between m_ν^{loop} and the seesaw term takes place. This cancellation makes cuspy curves of the invisible branching ratio and Yukawa couplings in contours. For the H decay, $m_H = 60$ GeV and $m_A = 200$ GeV are taken. For the A decay, $m_H = 200$ GeV and $m_A = 60$ GeV are taken.

Figure 4 is the contour plot of the invisible decay branching ratio of h. Here, $m_H = 60$ GeV and $m_{N_R} = 10$ GeV are taken.

The produced RH neutrino N_R decays as $N_R \to Z^*\nu$, $h^*\nu$ or $W^*\ell$ through a tiny left-right mixing of $\sin\theta \lesssim \mathcal{O}(10^{-6})$. Here, a sign of inequality becomes more appropriate as $m_\nu^{(\mathrm{loop})}$ becomes sizable. For such a left-right mixing of the order of 10^{-6} or less, the displaced vertex of N_R decay could be generated, and the decay length of the RH neutrino becomes $c\tau_{N_R} \gtrsim 1$ cm for $m_{N_R} = \mathcal{O}(10)$ GeV [36]. For a further lighter m_{N_R} or a much smaller $\sin\theta$, N_R would not decay inside the detector. One can see that such a small mixing is realized in Fig. 3.

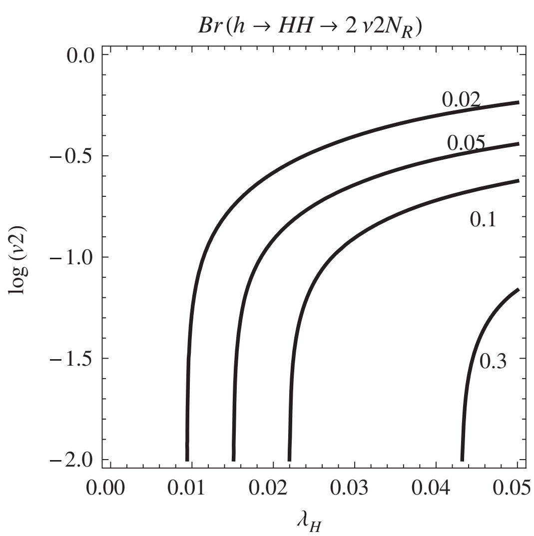

FIG. 4. Contours of the decay branching ratio of the SM-like Higgs boson, $Br(h \to 2H \to 2\nu 2N_R)$, in the $\lambda_H - v_2$ plane. Here, we take $m_{N_R} = 10$ GeV.

On the other hand, for $m_{N_R} > m_{H/A}$, H or A decays into SM fermions through a tiny mixing of a Higgs boson (8), as in the usual type-I THDM.

VI. SUMMARY

We have shown that the SM-like Higgs boson could have a sizable invisible decay branching ratio such as $\mathcal{O}(10)\%$ with four neutrinos final states, $h \to 2\nu 2\bar{\nu}$ for a Dirac neutrino and $h \to 2\nu 2N_R$ for a Majorana neutrino, in neutrinophilic Higgs doublet models, if one of the extra Higgs bosons is light enough to be produced by the SM-like Higgs boson decay. For the Majorana neutrino, this becomes a case in the parameter region $v_2 \lesssim 0.1$ GeV. Because of this mass spectrum of Higgs bosons, the SM normalized decay rate of $h \to \gamma\gamma$ is predicted to be 0.9. In the Majorana neutrino case, the displaced vertex of a N_R decay also would be observed. Although the invisible decay of the Higgs boson was recently widely discussed in [11] or applied to dark matter physics [12,13], we emphasize that such a size of invisible decay can be realized just within simple THDM without a dark matter candidate.

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