Dirac Neutrinos, Dark Energy and Baryon Asymmetry

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We explore a new origin of neutrino dark energy and baryon asymmetry in the universe. The neutrinos acquire small masses through the Dirac seesaw mechanism. The pseudo-Nambu-Goldstone boson associated with neutrino mass-generation provides a candidate for dark energy. The puzzle of cosmological baryon asymmetry is resolved via neutrinogenesis.

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I. INTRODUCTION

Strong evidence from cosmological observations [1] indicates that our universe is expanding with an accelerated rate at the present. This acceleration can be attributed to the dark energy. The dark energy may be a dynamical scalar field, such as the quintessence [2] with an extremely flat potential. The quintessence can be realized by a pseudo-Nambu-Goldstone boson (pNGB) arising from spontaneous breaking of certain global symmetry near the Planck scale [3].

On the other hand, various neutrino oscillation experiments [4] have confirmed that the neutrinos have tiny but nonzero masses, of the order $10^{-2}\,$eV. The smallness of neutrino masses can be naturally explained by the seesaw mechanism [5]. In the original seesaw scenario, the neutrinos are of Majorana nature which, however, has not been experimentally verified so far. In fact, the ultralight Dirac neutrinos were discussed many years ago [6, 7]. Recently some interesting models were proposed [8, 9], in which the neutrinos can naturally acquire small Dirac masses, meanwhile, the observed baryon asymmetry in the universe can be produced by a new type of leptogenesis [10], called neutrinogenesis [11].

It is striking that the scale of dark energy ($\sim(3\times 10^{-3}\,\mathrm{eV})^4$) is far lower than all the known scales in particle physics except that of the neutrino masses. The intriguing coincidence between the neutrinos mass scale and the dark energy scale inspires us to consider them in a unified scenario, as in the neutrino dark energy model [12, 13]. Recently a number of works studied the possible connection between the pNGB dark energy and the Majorana neutrinos [14].

In this paper, we propose a novel model to unify the mass-generation of Dirac neutrinos and the origin of dark energy. In particular, a pNGB associated with the neutrino mass-generation provides the candidate for dark energy while the neutrino masses depending on the dark energy field are generated through the Dirac seesaw [9]. Furthermore, our model also resolves the puzzle of cosmological baryon asymmetry via the neutrinogenesis [11].

II. THE MODEL

We extend the standard model (SM) gauge symmetry $SU(2)_L\otimes U(1)_Y$ with an approximate global symmetry $U(1)^3 \equiv U(1)_1\otimes U(1)_2\otimes U(1)_3$ as well as a discrete symmetry Z_2. The quantum number assignment is shown in Table I, where $i,j=1,2,3$ denote the family indices, x_i is the $U(1)_i$ charge, ψ_{Li} is the left-handed lepton doublet, ν_{Ri} is the right-handed neutrino, H and η_{ij} are the Higgs doublets, $\xi_{ij}\equiv\xi^*_{ji}\,(i\neq j)$ is the Higgs singlet, χ is a real scalar. Since all η_{ii}'s carry zero $U(1)_i$ charge, we only need to introduce one such doublet-field by defining $\eta_{ii}\equiv\eta_0$. As for the other SM fields, which carry even parity under the Z_2, they are all singlets under the $U(1)_i$ except that the right-handed charged leptons ℓ_{Ri} have the same $U(1)_i$ charge as ψ_{Li}. Thus H plays a role of the SM Higgs.

The phase transformations of the three Higgs singlets, $\xi_{ij}\equiv\xi^*_{ji}$ $(i\neq j)$, are supposed to be independent and hence will result in a global $U(1)^3$ symmetry. Subsequently, the transformations of the Higgs doublets η_{ij} $(i\neq j)$ under this $U(1)^3$ are determined by requiring the invariance of the following scalar interactions,

$$\xi_{ij}\chi\eta^\dagger_{ij}H+\mathrm{h.c.}\,. \tag{1}$$

However, the six Higgs doublets η_{ij} $(i\neq j)$ only have two independent phase transformations to keep the following Yukawa interactions invariant,

$$\overline{\psi_{Li}}\eta_{ij}\nu_{Rj}+\mathrm{h.c.}\,, \tag{2}$$

which explicitly break the $U(1)_1\otimes U(1)_2\otimes U(1)_3$ down to a $U(1)^{'}_1\otimes U(1)^{'}_2$. So, in the presence of Eqs. (1) and (2), we will have two massless Nambu-Goldstone bosons (NGBs) and one pNGB after this global symmetry is spontaneously broken by the vacuum expectation values (*vevs*) of three Higgs singlets ξ_{ij}.

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Fields	$SU(2)_L$	$U(1)_Y$	$U(1)_1\otimes U(1)_2\otimes U(1)_3$	Z_2
ψ_{Li}	**2**	$-1/2$	$x_i\times(\delta_{i1},\delta_{i2},\delta_{i3})$	$+$
ν_{Ri}	**1**	0	$x_i\times(\delta_{i1},\delta_{i2},\delta_{i3})$	$-$
H	**2**	$-1/2$	0	$+$
η_{ij}	**2**	$-1/2$	$x_i\times(\delta_{i1},\delta_{i2},\delta_{i3})-x_j\times(\delta_{j1},\delta_{j2},\delta_{j3})$	$-$
$\xi_{ij}\,(i\neq j)$	**1**	0	$x_i\times(\delta_{i1},\delta_{i2},\delta_{i3})-x_j\times(\delta_{j1},\delta_{j2},\delta_{j3})$	$+$
χ	**1**	0	0	$-$

TABLE I: The field content and quantum number assignment.

We then write down the relevant Lagrangian,

$$\begin{aligned}-\mathcal{L}\supset&\sum_{ij}\big(\rho_{ij}^2+\sum_{k\neq\ell}\lambda_{ij,k\ell}\xi_{k\ell}^\dagger\,\xi_{k\ell}\big)\,\eta_{ij}^\dagger\eta_{ij}\\&+\sum_{i\neq j,k\neq\ell,ij\neq k\ell}\lambda_{ij,k\ell}^{'}\xi_{ij}^\dagger\xi_{k\ell}\eta_{ij}^\dagger\eta_{k\ell}+\Big(-\mu_0\chi\eta_0^\dagger H\\&+\sum_{i\neq j}h_{ij}\xi_{ij}\chi\eta_{ij}^\dagger H+\sum_{ij}y_{ij}\overline{\psi_{Li}}\eta_{ij}\nu_{Rj}+\text{h.c.}\Big),\end{aligned}\quad(3)$$

where ρ_{ij} and μ_0 have the mass-dimension one while $\lambda^{(')}_{ij,kl}$, h_{ij} and y_{ij} are dimensionless. For convenience, we will denote $\rho_{ii}\equiv\rho_0$ and $\lambda_{ii,k\ell}\equiv\lambda_{0,k\ell}$ corresponding to $\eta_{ii}\equiv\eta_0$.

After the three Higgs singlets ξ_{ij} acquire their *vevs*, $\langle\xi_{ij}\rangle\equiv\frac{1}{\sqrt{2}}f_{ij}$, we can write

$$\xi_{ij}=\frac{1}{\sqrt{2}}\left(\sigma_{ij}+f_{ij}\right)\exp\left(i\varphi_{ij}/f_{ij}\right),\quad(i\neq j),\quad(4)$$

with σ_{ij} , φ_{ij} $(i,j=1,2,3)$ being the three neutral Higgs and the three NGBs, respectively. Here $f_{ij}\equiv f_{ji}$, $\sigma_{ij}\equiv\sigma_{ji}$ and $\varphi_{ij}\equiv-\varphi_{ji}$ since $\xi_{ij}\equiv\xi_{ji}^*$. In this approach, due to the explicit breaking of $U(1)_1\otimes U(1)_2\otimes U(1)_3\to U(1)_1^{'}\otimes U(1)_2^{'}$, one of these three NGBs will acquire a finite mass via the Coleman-Weinberg potential and thus become a pNGB, while the other two remain massless, as a result of spontaneous breaking of the subgroup $U(1)_1^{'}\otimes U(1)_2^{'}$.

For convenience we redefine the Higgs doublets η_{ij} $(i\neq j)$ as

$$\exp\left(i\varphi_{ij}/f_{ij}\right)\eta_{ij}\;\to\;\eta_{ij}\,,\quad(5)$$

and then express the Lagrangian (3) in a new form,

$$\begin{aligned}-\mathcal{L}\;\supset\;&M_0^2\eta_0^\dagger\eta_0+\sum_{i\neq j,k\neq\ell}\left(M^2\right)_{ij,kl}\eta_{ij}^\dagger\eta_{k\ell}\\&+\Big\{-\mu_0\chi\eta_0^\dagger H+y_{ii}\overline{\psi_{Li}}\eta_0\nu_{Ri}+\sum_{i\neq j}\Big[-\mu_{ij}\chi\eta_{ij}^\dagger H\\&+\;y_{ij}\exp\left(-i\varphi_{ij}/f_{ij}\right)\overline{\psi_{Li}}\eta_{ij}\nu_{Rj}\Big]+\text{h.c.}\Big\}\end{aligned}\quad(6)$$

with the definitions,

$$M_0^2\;\equiv\;\rho_0^2+\sum_{k\neq\ell}\lambda_{0,k\ell}f_{k\ell}^2\,,\quad(7)$$

$$\begin{aligned}(M^2)_{ij,k\ell}\;\equiv\;&\big(\rho_{ij}^2+\frac{1}{2}\sum_{m\neq n}\lambda_{ij,mn}\,f_{mn}^2\big)\,\delta_{ij,k\ell}\\&+\frac{1}{2}\lambda_{ij,k\ell}^{'}f_{ij}f_{k\ell}(1-\delta_{ij,k\ell})\,,\end{aligned}\quad(8)$$

$$\mu_{ij}\;\equiv\;-\frac{1}{\sqrt{2}}h_{ij}f_{ij}\,,\quad(9)$$

At this stage, the last Yukawa term in (6) depends on all three fields φ_{ij}. However, by making the further phase rotations on the left-handed lepton doublets and the right-handed neutrinos, we can find that except one combination of φ_{ij} still remains in the Yukawa interaction, the other two disappear from (6), so they only have derivative interactions and stay as the massless NGBs. For instance, we can make the following rotations,

$$\exp(-i\varphi_{12}/f_{12})\psi_{L2}\;\to\;\psi_{L2}\,,\quad(10)$$
$$\exp(-i\varphi_{12}/f_{12})\nu_{R2}\;\to\;\nu_{R2}\,,\quad(11)$$
$$\exp(+i\varphi_{31}/f_{31})\psi_{L3}\;\to\;\psi_{L3}\,,\quad(12)$$
$$\exp(+i\varphi_{31}/f_{31})\nu_{R3}\;\to\;\nu_{R3}\,,\quad(13)$$

and then obtain

$$-\mathcal{L}_Y\;=\;\sum_{ij}Y_{ij}\overline{\psi_{Li}}\eta_{ij}\nu_{Ri}+\text{h.c.}\,,\quad(14)$$

where

$$Y\;\equiv\;\left(\begin{array}{ccc}y_{11}&y_{12}&y_{13}\\y_{21}&y_{22}&y_{23}e^{-i\phi/f}\\y_{31}&y_{32}e^{+i\phi/f}&y_{33}\end{array}\right)\quad(15)$$

with the definition

$$\phi/f\;\equiv\;\varphi_{12}/f_{12}+\varphi_{23}/f_{23}+\varphi_{31}/f_{31}\,.\quad(16)$$

Here f should be of the order of the $U(1)_1^{'}\otimes U(1)_2^{'}$ breaking scales, i.e., $f\sim f_{ij}$. It is impossible to remove ϕ from the Yukawa interactions by further transformation. We will show later that this ϕ is a pNGB with a tiny mass and can naturally serve as the candidate of dark energy.

III. NEUTRINO MASSES

We consider the case that after the $U(1)_1^{'} \otimes U(1)_2^{'}$ breaking, the mass-square (7) of η_0 and the eigenvalues of the mass-square matrix (8) for η_{ij}'s $(i \neq j)$ are all positive. So, these Higgs doublets can develop nonzero *vevs* only after the SM Higgs-doublet H and the real scalar χ both acquire their *vevs* [9],

$$\langle \eta_{ij} \rangle \simeq \begin{cases} \langle H \rangle \langle \chi \rangle \sum\limits_{k \neq \ell} \left(M^{-2}\right)_{ij,k\ell} \mu_{k\ell}\,, & \text{for } i \neq j\,, \\ \langle H \rangle \langle \chi \rangle M_0^{-2} \mu_0\,, & \text{for } i = j\,. \end{cases} \qquad (17)$$

In consequence, the neutrinos obtain small Dirac masses,

$$(m_\nu)_{ij} \;\simeq\; Y_{ij} \langle \eta_{ij} \rangle\,. \qquad (18)$$

The discrete Z_2 symmetry is expected to break at the TeV scale by the *vev* of the real scalar χ [1], so we will set $\langle \chi \rangle$ around $\mathcal{O}(\mathrm{TeV})$ [2]. Furthermore, it is reasonable to take μ less than M in (17). Under this setup, it is straightforward to see that the Dirac neutrino masses will be efficiently suppressed by the ratio of the electroweak scale over the heavy masses. For instance, we find that, for $\langle H \rangle \simeq 174\,\mathrm{GeV}$, $M \sim 10^{14}\,\mathrm{GeV}$, $\mu \sim 10^{13}\,\mathrm{GeV}$ and $Y \sim \mathcal{O}(1)$, the neutrino masses can be naturally around $\mathcal{O}(0.1\,\mathrm{eV})$. We see that this mechanism of the neutrino mass generation has two essential features: (i) it generates Dirac masses for neutrinos, and (ii) it retains the essence of the conventional seesaw [5] by making the neutrino masses tiny via the small ratio of the electroweak scale over the heavy mass scale. This is a realization of Dirac Seesaw [9].

IV. DARK ENERGY

So far cosmological observations [1] strongly support the existence of dark energy which accelerates the expansion of our universe. One plausible explanation for the dark energy has its origin in a dynamical scalar field, such as the quintessence [2] with an extremely flat potential. It was shown [3] that the pNGB provides an attractive realization of the quintessence field.

We have pointed out that after the Higgs singlets getting their *vevs*, one NGB ϕ [as shown in (15)-(16)] will remain in the neutrino Yukawa interactions (which explicitly breaks the global $U(1)^3$). Therefore, this NGB will develop a finite mass from the Coleman-Weinberg effective potential via these neutrino Yukawa interactions, and thus become a pNGB. We can explicitly compute the Coleman-Weinberg potential for ϕ at one-loop order,

$$V(\phi) \;=\; -\frac{1}{16\pi^2} \sum_{k=1}^{3} m_k^4 \ln \frac{m_k^2}{\Lambda^2}\,, \qquad (19)$$

where m_k (as a function of ϕ) is the kth eigenvalue of the neutrino mass matrix m_ν and Λ is the ultraviolet cutoff. Note that there is an irrelevant quadratical term in V, $\frac{\Lambda^2}{16\pi^2}\sum_k m_k^2$, which has no ϕ-dependence and is thus omitted in (19). A typical term in V contributing to the potential of a pNGB field Q has the form,

$$V(Q) \;\simeq\; V_0 \cos(Q/f)\,, \qquad (20)$$

with $V_0 = \mathcal{O}(m_\nu^4)$. It is well-known that with f of the order of Planck scale M_{Pl}, Q obtains a mass of $\mathcal{O}(m_\nu^2/M_{\mathrm{Pl}})$ and is a consistent candidate for the quintessence dark energy.

V. BARYON ASYMMETRY

We now demonstrate how to generate the observed baryon asymmetry in our model. We make use of the neutrinogenesis mechanism [11]. Since the sphalerons [16] have no direct effect on the right-handed fields, a nonzero lepton asymmetry stored in the right-handed fields could survive above the electroweak phase transition and then produce the baryon asymmetry in the universe, although the lepton asymmetry stored in the left-handed fields had been destroyed by the sphalerons. For all the SM species, the Yukawa couplings are sufficiently strong to rapidly cancel the stored left- and right-handed lepton asymmetry. But the effective Yukawa interactions of the Dirac neutrinos are exceedingly weak, and the equilibrium between the left-handed lepton doublets and the right-handed neutrinos will not be realized until temperatures fall well below the electroweak scale. At that time the lepton asymmetry stored in the left-handed lepton doublets has already been converted to the baryon asymmetry by the sphalerons. In particular, the final baryon asymmetry should be

$$B \;=\; \frac{28}{79}\,(B - L_{SM}) \;=\; \frac{28}{79} L_{\nu_R}\,, \qquad (21)$$

for the SM with three generation fermions and one Higgs doublet.

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[1] It is also possible to replace the Z_2 symmetry by a global or local $U(1)_X$ symmetry [9, 15], under which all SM particles transform as singlets, while ν_{Ri}, η_{ij} and χ carry the $U(1)$ charge $-\frac{1}{2}$, $+\frac{1}{2}$ and $+\frac{1}{2}$, respectively. This $U(1)_X$ symmetry is spontaneously broken at TeV by the *vev* of χ. So $\langle \chi \rangle$ is fixed by the $U(1)_X$ symmetry breaking scale around $O(\mathrm{TeV})$. In the case of a local $U(1)_X$ symmetry, we need add three massless left-handed singlet fermions, s_{Li} $(i = 1, 2, 3)$ with the $U(1)$ charge $+\frac{1}{2}$, which decouple from everything and make the theory anomaly free. In this case, the new gauge boson couples to s_{Li}, ν_{Ri}, η_{ij} and χ rather than the SM particles, so it is expected to escape the detection at the LHC and ILC.

[2] Here we are not concerned with the naturalness issue of scalar masses as in any non-supersymmetric model.

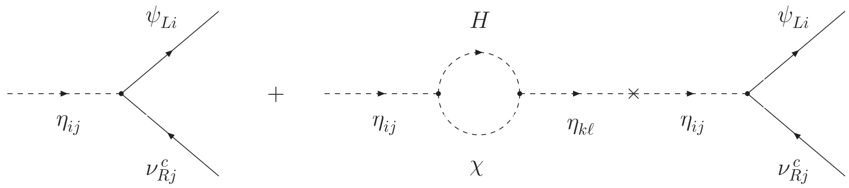

FIG. 1: The Higgs doublets decay into the leptons at one-loop order. Here $i \neq j$, $k \neq \ell$ and $ij \neq k\ell$.

There are two types of final states coexisting in the decays of every heavy Higgs doublet,

$$\eta_{ij} \rightarrow \begin{cases} \psi_{Li}\, \nu^c_{Rj}\,, \\ \chi\, H\,. \end{cases} \tag{22}$$

The channels of $\eta \rightarrow \psi_L\, \nu^c_R$ and $\eta^* \rightarrow \psi^c_L\, \nu_R$ can provide the expected asymmetry between the right-handed neutrinos and anti-neutrinos if the CP is not conserved and the decays are out of thermal equilibrium. As shown in Fig. 1, the mixing (8) among $\eta_{ij}\,(i \neq j)$ help to generate the interference between the tree-level decay and the irreducible loop-correction. For convenience, we adopt the following definitions,

$$\widehat{\eta}_a \;\equiv\; \sum_{i\neq j} U_{a,ij}\eta_{ij}\,, \tag{23}$$

$$\widehat{M}^2_a \;\equiv\; \sum_{i\neq j, k\neq l} U_{a,ij}\left(M^2\right)_{ij,kl} U_{a,kl}\,, \tag{24}$$

$$\widehat{\mu}_a \;\equiv\; \sum_{i\neq j} U_{a,ij}\mu_{ij}\,, \tag{25}$$

$$\widehat{Y}_{a,ij} \;\equiv\; U_{a,ij}Y_{ij}\,, \qquad (\text{for} \quad i \neq j)\,, \tag{26}$$

where U is the orthogonal rotation matrix to diagonalize $\eta_{ij}\,(i \neq j)$ in their mass-eigenbasis $\widehat{\eta}_a$. We then derive the relevant CP-asymmetry,

$$\begin{aligned} \varepsilon_a \;&\equiv\; \frac{\sum_{ij}\left[\Gamma\left(\widehat{\eta}^*_a \rightarrow \psi^c_{Li}\nu_{Rj}\right) - \Gamma\left(\widehat{\eta}_a \rightarrow \psi_{Li}\nu^c_{Rj}\right)\right]}{\Gamma_a} \\ &=\; \frac{1}{4\pi}\sum_{b\neq a}\frac{\mathrm{Im}\left[(\widehat{Y}^\dagger\widehat{Y})_{ba}\widehat{\mu}^*_b\widehat{\mu}_a\right]}{(\widehat{Y}^\dagger\widehat{Y})_{aa}\widehat{M}^2_a + |\widehat{\mu}_a|^2}\frac{\widehat{M}^2_a}{\widehat{M}^2_a - \widehat{M}^2_b} \end{aligned} \tag{27}$$

with

$$\Gamma_a = \frac{1}{16\pi}\left[(\widehat{Y}^\dagger\widehat{Y})_{aa} + \frac{|\widehat{\mu}_a|^2}{\widehat{M}^2_a}\right]\widehat{M}_a \tag{28}$$

being the total decay width of $\widehat{\xi}_a$ or $\widehat{\xi}^*_a$.

For illustration, we will use $\widehat{\xi}_a$ to denote the lightest one among all heavy Higgs doublets (including η_0), and hence the contribution of $\widehat{\xi}_a$ is expected to dominate the final baryon asymmetry, which is given by the approximate relation [17],

$$Y_B \;\equiv\; \frac{n_B}{s} \;\simeq\; \frac{28}{79}\times\begin{cases} \dfrac{\varepsilon_a}{g_*}\,, & \text{for} \;\; K \ll 1, \\ \dfrac{0.3\,\varepsilon_a}{g_* K\,(\ln K)^{0.6}}\,, & \text{for} \;\; K \gg 1. \end{cases} \tag{29}$$

Here the parameter K is defined as

$$K \equiv \left.\frac{\Gamma_a}{2H(T)}\right|_{T=\widehat{M}_a} = \left(\frac{45}{16\pi^3 g_*}\right)^{\frac{1}{2}}\frac{M_{\mathrm{Pl}}\Gamma_a}{\widehat{M}^2_a} \tag{30}$$

which characterizes the deviation from equilibrium. For instance, inputting $\widehat{M}_a = 0.1\widehat{M}_b = 10^{14}\,\mathrm{GeV}$, $|\widehat{\mu}_a| = |\widehat{\mu}_b| = 10^{13}\,\mathrm{GeV}$, $\left|\sum_{b\neq a}(\widehat{Y}^\dagger\widehat{Y})_{ba}\right| = 1.5$, $(\widehat{Y}^\dagger\widehat{Y})_{aa} = 1$, and the maximum CP-phase, we derive the sample predictions: $K \simeq 60$ and $\varepsilon_a \simeq 8.0\times 10^{-6}$, where we have used $g_* \sim 100$ and $M_{\mathrm{Pl}} \sim 10^{19}\,\mathrm{GeV}$. In consequence, we deduce, $n_B/s \simeq 10^{-10}$, as desired.

VI. SUMMARY AND DISCUSSIONS

In this paper, we have proposed a new model to realize Dirac neutrinos, dark energy and baryon asymmetry. In particular, the heavy Higgs doublets develop small *vev*s which make the neutrinos acquire small masses through the Dirac seesaw. Furthermore, the pNGB associated with the Dirac neutrino mass-generation can be the quintessence field and thus provide an attractive candidate for dark energy. Finally, our model generates the matter-antimatter asymmetry in the universe via the out-of-equilibrium decays of the heavy Higgs doublets with CP-violation.

In our model, the Dirac neutrino masses are functions of the dark energy field. The dark energy is a dynamical component and will evolute with time and/or in space. In consequence, the Dirac neutrino masses are variable, rather than constant. The prediction of the neutrino-mass variation could be verified in the experiments, such as the observation on the cosmic microwave background and the large scale structures [18], the measurement of the extremely high-energy cosmic neutrinos [19] and the analysis of the neutrino oscillation data [20].

Finally, we note that the real scalar χ has a *vev* around the electroweak scale, it can mix with and couple to the

SM Higgs boson via the quartic interaction,

$$\kappa_{\rm eff}\chi^2 H^\dagger H \equiv \left[\kappa - \left(\sum_a \frac{|\widehat{\mu}_a|^2}{\widehat{M}_a^2} + \frac{|\mu_0|^2}{M_0^2}\right)\right]\chi^2 H^\dagger H, \quad (31)$$

with κ being a dimensionless parameter. Hence, the SM Higgs boson is no longer a mass-eigenstate, and its collider signatures will be modified [21]. Further phenomenological analyses for such non-standard Higgs boson will be given elsewhere.

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