Dirac neutrinos, dark energy and baryon asymmetry

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Abstract. 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 dark energy candidate. The puzzle of cosmological baryon asymmetry is resolved via neutrino-generation.

Keywords: dark energy theory, neutrino properties, baryon asymmetry

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1. 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 quintessence [2] with an extremely flat potential. 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 non-zero masses, of the order of $10^{-2} \text{ eV}$. The smallness of neutrino masses can be naturally explained using 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 ($\sim (3 \times 10^{-3} \text{ eV})^4$) 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 for unifying the mass generation of Dirac neutrinos and the origin of dark energy. In particular, a pNGB associated with the neutrino mass generation provides the dark energy candidate 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].
denote the family indices, one pNGB after this global symmetry is spontaneously broken by the vacuum expectation value of the Higgs singlet, \( \eta \). The quantum number assignment is shown in Table 1, where \( i, j = 1, 2, 3 \) denote the family indices, \( x \) is the \( U(1)_i \) charge, \( \psi_{Li} \) is the left-handed lepton doublet, \( \nu_{Ri} \) is the right-handed neutrino, \( H \) and \( \eta_{ij} \) are the Higgs doublets, \( \xi_{ij} \equiv \xi_{ji}^\dagger (i \neq j) \) is the Higgs singlet, \( \chi \) is a real scalar. Since all of the \( \eta_{ii} \) carry zero \( U(1)_i \) charge, we only need to introduce one such doublet field by defining \( \eta_0 \equiv \eta_0 \). As for the other SM fields, which carry even parity under \( Z_2 \), they are all singlets under \( U(1)_i \). So, in the presence of equations (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 \( \xi_{ij} \).

We then write down the relevant Lagrangian:

\[
-L = \sum_{ij} \left[ \rho_{ij}^2 + \sum_{k \neq l} \lambda_{ij,kl}^\dagger \xi_{kl} \xi_{kl}^\dagger \eta_{ij} \eta_{ij} + \sum_{i \neq j, k \neq l, i \neq j \neq k} \lambda_{ij,kl}^\dagger \xi_{kl} \xi_{kl}^\dagger \eta_{ij} \eta_{kl} + \left( -\mu_0 \chi_{ij}^\dagger H + \sum_{i \neq j} h_{ij} \xi_{ij} \chi_{ij}^\dagger H + \sum_{ij} y_{ij} \psi_{Li} \eta_{ij} \nu_{Rj} + \text{h.c.} \right) \right],
\]

where \( \rho_{ij} \) and \( \mu_0 \) have the mass dimension 1 while \( \lambda_{ij,kl} \), \( h_{ij} \) and \( y_{ij} \) are dimensionless. For convenience, we will define \( \rho_{ii} \equiv 0 \) and \( \lambda_{ii,kl} \equiv \lambda_{0,kl} \) corresponding to \( \eta_{ii} \equiv \eta_0 \).
After the three Higgs singlets $\xi_{ij}$ acquire their vevs, $\langle \xi_{ij} \rangle \equiv (1/\sqrt{2})f_{ij}$, we can write

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

with $\sigma_{ij}, \varphi_{ij} (i, j = 1, 2, 3)$ 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 $\xi_{ij} \equiv \xi_{ji}^\dagger$. 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 $\eta_{ij} (i \neq j)$ as

$$\exp(i\varphi_{ij}/f_{ij})\eta_{ij} \to \eta_{ij},$$

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

$$-\mathcal{L} \supset M_0^2\eta_0^\dagger\eta_0 + \sum_{i \neq j, k \neq \ell} (M^2)_{ij,kl}\eta_{ij}^\dagger\eta_{kl} + \left\{-\mu_0\chi_0^\dagger H + y_{ii}\overline{\psi}_{Li}\eta_{0i}\nu_{Ri} + \sum_{i \neq j} \left[-\mu_{ij}\chi_{ij}^\dagger H + y_{ij}\overline{\psi}_{Li}\eta_{ij}\nu_{Rj} \right] + \text{h.c.} \right\},$$

with the definitions

$$M_0^2 \equiv \rho_0^2 + \sum_{k \neq \ell} \lambda_{0,kl} f_{k\ell}^2,$$

$$(M^2)_{ij,kl} \equiv \left(\rho_{ij}^2 + \frac{1}{2} \sum_{m \neq n} \lambda_{ij,mn} f_{mn}^2 \right) \delta_{ij,kl} + \frac{1}{2} \lambda'_{ij,kl} f_{ij} f_{kl} (1 - \delta_{ij,kl}),$$

$$\mu_{ij} \equiv -\frac{1}{\sqrt{2}}h_{ij} f_{ij}.$$

At this stage, the last Yukawa term in (6) depends on all three fields $\varphi_{ij}$. However, by making the further phase rotations on the left-handed lepton doublets and the right-handed neutrinos, we can find that just one combination of $\varphi_{ij}$ 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:

$$\exp(-i\varphi_{12}/f_{12})\psi_{L2} \to \psi_{L2},$$

$$\exp(-i\varphi_{12}/f_{12})\nu_{R2} \to \nu_{R2},$$

$$\exp(+i\varphi_{31}/f_{31})\psi_{L3} \to \psi_{L3},$$

$$\exp(+i\varphi_{31}/f_{31})\nu_{R3} \to \nu_{R3},$$

and then obtain

$$-\mathcal{L}_Y = \sum_{ij} Y_{ij} \overline{\psi}_{Li}\eta_{ij}\nu_{Ri} + \text{h.c.},$$
where

\[
Y \equiv \begin{pmatrix}
y_{11} & y_{12} & y_{13} e^{-i\phi/f} \\
y_{21} & y_{22} & y_{23} e^{-i\phi/f} \\
y_{31} & y_{32} e^{+i\phi/f} & y_{33}
\end{pmatrix}
\]  

(15)

with the definition

\[
\phi/f \equiv \frac{\varphi_{12}}{f_{12}} + \frac{\varphi_{23}}{f_{23}} + \frac{\varphi_{31}}{f_{31}}.
\]  

(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 \(\phi\) from the Yukawa interactions by further transformation. We will show later that this \(\phi\) is a pNGB with a tiny mass and can naturally serve as the dark energy candidate.

3. Neutrino masses

We consider the case where after the \(U(1)'_1 \otimes U(1)'_2\) breaking, the mass square (7) of \(\eta_0\) and the eigenvalues of the mass square matrix (8) for the \(\eta_{ij}(i \neq j)\) are all positive. So, these Higgs doublets can develop non-zero vevs only after the SM Higgs doublet \(H\) and the real scalar \(\chi\) both acquire their vevs [9],

\[
\langle \eta_{ij} \rangle \simeq \begin{cases} 
\langle H \rangle \langle \chi \rangle \sum_{k \neq \ell} (M^{-2})_{ij,k\ell} \mu_{k\ell}, & \text{for } i \neq j, \\
\langle H \rangle \langle \chi \rangle M_0^{-2} \mu_0, & \text{for } i = j.
\end{cases}
\]  

(17)

In consequence, the neutrinos obtain small Dirac masses,

\[
(m_{\nu})_{ij} \simeq Y_{ij} \langle \eta_{ij} \rangle.
\]  

(18)

The discrete \(Z_2\) symmetry is expected to be broken at the TeV scale by the vev of the real scalar \(\chi\), so we will set \(\langle \chi \rangle\) around \(\mathcal{O}(\text{TeV})\). Furthermore, it is reasonable to take \(\mu\) less than \(M\) in (17). Under this set-up, 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 \(\langle H \rangle \simeq 174\text{ GeV}, M \sim 10^{14}\text{ GeV}, \mu \sim 10^{13}\text{ GeV}\) and \(Y \sim \mathcal{O}(1)\), the neutrino masses can be naturally around \(\mathcal{O}(0.1\text{ eV})\). We see that this mechanism of 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 a Dirac seesaw [9].

4 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 \(\nu_R, \eta_{ij}\) and \(\chi\) carry the \(U(1)\) charges \(-\frac{1}{2}, +\frac{1}{2}\) and \(+\frac{1}{2}\), respectively. This \(U(1)_X\) symmetry is spontaneously broken at TeV scale by the vev of \(\chi\). So \(\langle \chi \rangle\) is fixed by the \(U(1)_X\) symmetry breaking scale around \(\mathcal{O}(\text{TeV})\). In the case of a local \(U(1)_X\) symmetry, we need add three massless left-handed singlet fermions, \(s_{ij}(i = 1, 2, 3)\) 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 \(s_L, \nu_R, \eta_{ij}\) and \(\chi\) rather than the SM particles, so it is expected to escape the detection at the LHC and ILC.

5 Here we are not concerned with the naturalness issue of scalar masses as in any non-supersymmetric model.
4. 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 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 get their vevs, one NGB $\phi$ (as shown in (15) and (16)) will remain in the neutrino Yukawa interactions (which explicitly breaks 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 $\phi$ at one-loop order,

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

(19)

where $m_k$ (as a function of $\phi$) is the $k$th eigenvalue of the neutrino mass matrix $m_\nu$ and $\Lambda$ is the ultraviolet cut-off. Note that there is an irrelevant quadratic term in $V$, $(\Lambda^2/16\pi^2) \sum_k m_k^2$, which has no dependence on $\phi$ and is thus omitted in (19). A typical term in $V$ contributing to the potential of a pNGB field $Q$ has the form

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

(20)

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

5. 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 non-zero 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 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

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

(21)

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

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

$$\eta_{ij} \rightarrow \begin{cases} \psi_{Li}^c \nu_{Rj}, \\ \chi H. \end{cases}$$

(22)
Figure 1. The Higgs doublets decay into the leptons at one-loop order. Here $i \neq j$, $k \neq \ell$ and $ij \neq k\ell$.

The channels of $\eta \rightarrow \psi_L \nu_R^c$ and $\eta^* \rightarrow \psi_R^c \nu_L$ can provide the expected asymmetry between the right-handed neutrinos and antineutrinos if $CP$ is not conserved and the decays are out of thermal equilibrium. As shown in figure 1, the mixing (8) among $\eta_{ij}$ ($i \neq j$) helps to generate the interference between the tree-level decay and the irreducible loop correction.

For convenience, we adopt the following definitions:

$$\hat{\eta}_a \equiv \sum_{i \neq j} U_{a,ij} \eta_{ij},$$

$$\hat{M}_a^2 \equiv \sum_{i \neq j,k \neq l} U_{a,ij} (M^2)_{ij,kl} U_{a,kl},$$

$$\hat{\mu}_a \equiv \sum_{i \neq j} U_{a,ij} \mu_{ij},$$

$$\hat{Y}_{a,ij} \equiv U_{a,ij} Y_{ij},$$

(for $i \neq j$),

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

$$\varepsilon_a \equiv \frac{\sum_{ij} \left[ \Gamma (\hat{\eta}_a^* \rightarrow \psi_L^c \nu_R) - \Gamma (\hat{\eta}_a \rightarrow \psi_L \nu_R^c) \right]}{\Gamma_x}$$

$$= \frac{1}{4\pi} \sum_{b \neq a} \frac{\text{Im} \left[ (\hat{Y}^\dagger \hat{Y})_{ba} \hat{\mu}_a \hat{\mu}_b \right]}{(\hat{Y}^\dagger \hat{Y})_{aa} \hat{M}_a^2 + |\hat{\mu}_a|^2 \hat{M}_a^2 - \hat{M}_b^2}$$

(27)

with

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

(28)

being the total decay width of $\hat{\xi}_a$ or $\hat{\xi}_a^*$. For illustration, we will use $\hat{\xi}_a$ to denote the lightest one among all heavy Higgs doublets (including $\eta_0$), and hence the contribution of $\hat{\xi}_a$ is expected to dominate the final baryon asymmetry, which is given by the approximate relation [17]

$$Y_B \equiv \frac{n_B}{s} \simeq \frac{28}{19} \times \left\{ \begin{array}{ll}
\frac{\varepsilon_a}{g_s}, & \text{for } K \ll 1, \\
0.3 \varepsilon_a g_s K (\ln K)^{0.6}, & \text{for } K \gg 1.
\end{array} \right. $$

(29)
Here the parameter $K$ is defined as

$$K \equiv \frac{\Gamma_a}{2H(T)} \bigg|_{T=\tilde{M}_a} = \left( \frac{45}{16\pi^3g_*} \right)^{1/2} \frac{M_{Pl}\Gamma_a}{\hat{M}_a^2}$$

(30)

which characterizes the deviation from equilibrium. For instance, inputting $\hat{M}_a = 0.1\hat{M}_b = 10^{14} \text{ GeV}, |\hat{\mu}_a| = |\hat{\mu}_b| = 10^{13} \text{ GeV}, |\sum_{b\neq a}(\hat{Y}^\dagger \hat{Y})_{ba}| = 1.5, (\hat{Y}^\dagger \hat{Y})_{aa} = 1$, 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 $M_{Pl} \sim 10^{19} \text{ GeV}$. In consequence, we deduce $n_B/s \simeq 10^{-10}$, as desired.

6. Summary and discussion

In this paper, we have proposed a new model for realizing Dirac neutrinos, dark energy and baryon asymmetry. In particular, the heavy Higgs doublets develop small vevs 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 dark energy candidate. 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 evolve 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 experiments, such as observations on the cosmic microwave background and large scale structures [18], measurement of extremely high energy cosmic neutrinos [19] and analysis of the neutrino oscillation data [20].

Finally, we note that the real scalar $\chi$ has a vev around the electroweak scale; it can mix with and couple to the SM Higgs boson via the quartic interaction

$$\kappa_{\text{eff}}\chi^2 H^\dagger H \equiv \left[ \kappa - \left( \sum_a |\hat{\mu}_a|^2 \frac{M_a^2}{M^2_0} + |\hat{\mu}_0|^2 \frac{M^2_0}{M^2_0} \right) \right] \chi^2 H^\dagger H,$$

(31)

with $\kappa$ 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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