An SO(10) × SO(10)′ model for common origin of neutrino masses, ordinary and dark matter-antimatter asymmetries

Pei-Hong Gu

Department of Physics and Astronomy, Shanghai Jiao Tong University, 800 Dongchuan Road, Shanghai 200240, China

E-mail: peihong.gu@sjtu.edu.cn

Received October 27, 2014
Revised November 24, 2014
Accepted November 27, 2014
Published December 22, 2014

Abstract. We propose an SO(10) × SO(10)′ model to simultaneously realize a seesaw for Dirac neutrino masses and a leptogenesis for ordinary and dark matter-antimatter asymmetries. A (16 × 16′)H scalar crossing the SO(10) and SO(10)′ sectors plays an essential role in this seesaw-leptogenesis scenario. As a result of lepton number conservation, the lightest dark nucleon as the dark matter particle should have a determined mass around 15 GeV to explain the comparable fractions of ordinary and dark matter in the present universe. The (16 × 16′)H scalar also mediates a U(1)em × U(1)′em kinetic mixing after the ordinary and dark left-right symmetry breaking so that we can expect a dark nucleon scattering in direct detection experiments and/or a dark nucleon decay in indirect detection experiments. Furthermore, we can impose a softly broken mirror symmetry to simplify the parameter choice.

Keywords: dark matter theory, baryon asymmetry, neutrino theory

ArXiv ePrint: 1410.5759
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1 Introduction

In the most popular grand unified theories, we can naturally obtain the extremely light Majorana neutrinos through the famous seesaw mechanism [1–11]. The lepton-number-violating interactions for the Majorana neutrinos can also accommodate a leptogenesis [12–28] mechanism to explain the cosmic matter-antimatter asymmetry. However, the Majorana nature of the neutrinos is just a theoretical assumption and has not been confirmed experimentally. Meanwhile, all of the other observed fermions are the Dirac particles rather than the Majorana particles. Therefore, it is worth exploring the possibility of the Dirac neutrinos [29–46] in the grand unification framework.

On the other hand, the dark and ordinary matter contribute comparable energy densities in the present universe [47]. This coincidence can be understood in a nature way if the dark matter relic density is a dark matter-antimatter asymmetry [48–93] and has a common origin with the ordinary matter-antimatter asymmetry. The mirror world based on the gauge groups $[\text{SU}(3)_c \times \text{SU}(2)_L \times \text{U}(1)_Y] \times [\text{SU}(3)'_c \times \text{SU}(2)'_L \times \text{U}(1)'_Y]$ is a very attractive asymmetric dark matter scenario [33, 94–143]. The mirror models can contain a tiny $\text{U}(1)_Y \times \text{U}(1)'_Y$ kinetic mixing input by hand to open a window for dark matter direct detections.

In this paper we shall propose an $\text{SO}(10) \times \text{SO}(10)'$ model with a $(16 \times 10')^H$ scalar to simultaneously realize a seesaw for Dirac neutrino masses and a leptogenesis for ordinary and dark matter-antimatter asymmetries. After the ordinary and dark left-right symmetry breaking, the $(16 \times 10'^I)^H$ scalar can acquire an induced vacuum expectation value. The ordinary right-handed neutrinos and the dark left-handed neutrinos then can form three heavy Dirac fermions to highly suppress the masses between the ordinary left-handed neutrinos and the dark right-handed neutrinos. Meanwhile, these heavy Dirac fermions can decay to generate a lepton asymmetry in the ordinary leptons and an opposite lepton asymmetry in the dark leptons. The $\text{SU}(2)_L$ and $\text{SU}(2)'_R$ sphaleron processes respectively can transfer such lepton asymmetries to an ordinary baryon asymmetry and a dark baryon asymmetry. With calculable lepton-to-baryon conversations in the ordinary and dark sectors, the lightest dark nucleon as the dark matter particle should have a predictive mass about 15 GeV to explain the ordinary and dark matter in the present universe as the ordinary proton has a known
mass about 1 GeV. Benefited from the $U(1)_{\text{em}} \times U(1)'_{\text{em}}$ kinetic mixing mediated by the $(16 \times \overline{16}')_H$ scalar, the dark proton as the dark matter particle can scatter off the ordinary nucleons at a testable level while the dark proton/neutron as the dark matter particle can decay to produce the ordinary fermion pairs. Furthermore, a softly broken mirror symmetry can be imposed to simplify the parameter choice.

2 Fields and symmetry breaking

In the ordinary SO(10) sector, we have the fermions and scalars including

$$q_L^c(3, 2, 1, -1/3) \oplus q_R^c(3, 1, 2, +1/3) \oplus l_L^c(1, 2, 1, +1) \oplus l_R^c(1, 1, 2, -1) = 16_F,$$

$$\chi_L^{\ast}(1, 2, 1, +1) \oplus \chi_R(1, 1, 2, -1) \in 16_H,$$

$$\Delta_L^\ast(1, 3, 1, -2) \oplus \Delta_R(1, 1, 3, +2) \oplus \Omega_L^0(3, 3, 1, -2/3) \oplus \Omega_R^0(3, 1, 3, +2/3) \in 126_H,$$

$$\Phi(1, 2, 0) \in 10_H \text{ and/or others}, \tag{2.1}$$

where the brackets following the fields describe the transformations under the SU(3)$_c$ × SU(2)$_L$ × SU(2)$_R$ × U(1)$_{B-L}$ gauge groups. Accordingly, the fermions and scalars in the dark SO(10)' sector contain

$$q_R^c(3, 2, 1, -1/3) \oplus q_L^c(3, 1, 2, +1/3) \oplus l_R^c(1, 2, 1, +1) \oplus l_L^c(1, 1, 2, -1) = 16_F,$$

$$\chi_R^{\ast}(1, 2, 1, +1) \oplus \chi_L(1, 1, 2, -1) \in 16_H',$$

$$\Delta_R^\ast(1, 3, 1, -2) \oplus \Delta_L^\ast(1, 1, 3, +2) \oplus \Omega_L^0(3, 3, 1, -2/3) \oplus \Omega_R^0(3, 1, 3, +2/3) \in 126_H',$$

$$\Phi^\prime(1, 2, 0) \in 10_H' \text{ and/or others}, \tag{2.2}$$

where the brackets give the SU(3)$_c$' × SU(2)$_R'$ × SU(2)$_L'$ × U(1)$_{B-L}'$ quantum numbers. There is also a $(16 \times \overline{16}')_H$ scalar crossing the SO(10) and SO(10)' sectors,

$$(16 \times \overline{16}')_H = \Sigma_{L^{L\prime}_R}(1, 2, 1, -1)(1, 1, 2, +1)\Sigma_{L^{R\prime}_{L\prime}}(1, 1, 2, +1)(1, 1, 2, -1) + \ldots. \tag{2.3}$$

For simplicity, we shall not consider the details of the SO(10) and SO(10)' symmetry breaking. Instead, we shall demonstrate at the left-right level. The ordinary and dark left-right symmetries are expected to have the breaking patterns as below,

$$\begin{align*}
\text{SU}(3)_c \times \text{SU}(2)_L \times \text{SU}(2)_R \times \text{U}(1)_{B-L} \\
^{(x_R)} = \frac{1}{\sqrt{2}}(v_R, 0)^T \\
\text{SU}(3)_c \times \text{SU}(2)_L \times \text{U}(1)_{\text{em}}, \tag{2.4a}
\end{align*}$$

$$\begin{align*}
\text{SU}(3)'_c \times \text{SU}(2)'_R \times \text{SU}(2)'_L \times \text{U}(1)'_{B-L} \\
^{(x_L')} = \frac{1}{\sqrt{2}}(v_L', 0)^T \\
\text{SU}(3)'_c \times \text{SU}(2)'_R \times \text{U}(1)'_{\text{em}},
\end{align*}$$

$$(2.4b)$$
We further impose a $U(1)_L$ global symmetry under which $(\chi_L^*, \chi_R)$ and $(\chi_R^*, \chi_L)$ carry a same charge. This means the following cubic terms

\[ V \supset \rho_\Phi \chi_L^\dagger \Phi \chi_R + \tilde{\rho}_\Phi \chi_L^\dagger \tilde{\Phi} \chi_R + \rho_\Phi \chi_R^\dagger \Phi \chi_L + \tilde{\rho}_\Phi \chi_R^\dagger \tilde{\Phi} \chi_L \]

\[ + \rho_\Delta (\chi_L^\dagger i\tau_2 \Delta_L \chi_L + \chi_R^\dagger i\tau_2 \Delta_R \chi_R^*) + \rho_\Delta (\chi_R^\dagger i\tau_2 \Delta_R \chi_R + \chi_L^\dagger i\tau_2 \Delta_L \chi_L^*) + \text{H.c.}, \]

should be absent from the scalar potential. Therefore the neutral components of the scalars $\chi_L, \Delta_L, \chi_R, \Delta_R$ will not acquire any induced vacuum expectation values. Accordingly, we can give a nonzero $\langle \Sigma_{LR} \rangle$ and a zero $\langle \Sigma_{LR'} \rangle$ from the scalar interactions as follows,

\[ V \supset \rho_\Sigma (\chi_L^\dagger \Sigma_{LR} \chi_R^* + \chi_R^\dagger \Sigma_{RL} \chi_L^*) + \text{H.c.}. \]

### 3 Dirac neutrinos and lepton asymmetries

We write down the Yukawa couplings relevant for the fermion mass generation,

\[
\mathcal{L} \supset -y_{qL} q_L \Phi q_R - \tilde{y}_{qL} \tilde{q}_L \tilde{\Phi} q_R - y_{tL} t_L \Phi t_R - \tilde{y}_{tL} \tilde{t}_L \tilde{\Phi} t_R \\
- y_{qR} q_R^\dagger \Phi^\dagger q_L - \tilde{y}_{qR} \tilde{q}_R^\dagger \tilde{\Phi}^\dagger q_L - y_{tR} t_R^\dagger \Phi^\dagger t_L - \tilde{y}_{tR} \tilde{t}_R^\dagger \tilde{\Phi}^\dagger t_L \\
- \frac{1}{2} f_\Delta (\tilde{t}_L^\dagger i\tau_2 \Delta_L \tilde{t}_L^* + \tilde{t}_R^\dagger i\tau_2 \Delta_R \tilde{t}_R^* - f_\Delta (\tilde{t}_R^\dagger i\tau_2 \Delta_R \tilde{t}_R^* + \tilde{t}_L^\dagger i\tau_2 \Delta_L \tilde{t}_L^*) \\
- f_\Sigma (\tilde{t}_L^\dagger \Sigma_{LR} \tilde{t}_L^* + \tilde{t}_R^\dagger \Sigma_{RL} \tilde{t}_R^* + \text{H.c.}).
\]

When the left-right symmetries are broken down to the electroweak symmetries, we can derive

\[
\mathcal{L} \supset -y_{uL} u_L \phi u_R - y_{dL} d_L \phi d_R - y_{eL} e_L \phi e_R \\
- y_{uR} u_R^\dagger \phi^\dagger u_L - y_{dR} d_R^\dagger \phi^\dagger d_L - y_{eR} e_R^\dagger \phi^\dagger e_L - y_{eL} e_L^\dagger \phi^\dagger e_R \\
- \frac{1}{2} f_\Delta (\tilde{t}_L^\dagger i\tau_2 \Delta_L \tilde{t}_L^* - f_\Delta (\tilde{t}_R^\dagger i\tau_2 \Delta_R \tilde{t}_R^* - M_N \tilde{u}_R \tilde{v}_L^* + \text{H.c.})
\]

With

\[
\begin{align*}
y_u &= \frac{v_1 y_u + v_2 \tilde{y}_u}{\sqrt{v_1^2 + v_2^2}}, & y_{u'} &= \frac{v_1' y_u + v_2' \tilde{y}_u}{\sqrt{v_1'^2 + v_2'^2}}, \\
y_d &= \frac{v_2 y_d + v_1 \tilde{y}_d}{\sqrt{v_1^2 + v_2^2}}, & y_{d'} &= \frac{v_2 y_d + v_1' \tilde{y}_d}{\sqrt{v_1'^2 + v_2'^2}}, \\
y_e &= \frac{v_1 y_e + v_2 \tilde{y}_e}{\sqrt{v_1^2 + v_2^2}}, & y_{e'} &= \frac{v_1' y_e + v_2' \tilde{y}_e}{\sqrt{v_1'^2 + v_2'^2}}, \\
M_N &= \frac{1}{\sqrt{2}} f_\Sigma v_L.
\end{align*}
\]
Figure 1. The heavy masses between the ordinary right-handed neutrinos $\nu_R$ and the dark left-handed neutrinos $\nu'_L$ are responsible for suppressing the masses between the ordinary left-handed neutrinos $\nu_L$ and the dark right-handed neutrinos $\nu'_R$.

Here the Higgs scalars $\phi$ and $\phi'$ with the vacuum expectation values,

$$
\langle \phi \rangle = \begin{bmatrix} \frac{1}{\sqrt{2}} v \\ 0 \end{bmatrix} \quad \left( v = \sqrt{v_1^2 + v_2^2} \approx 246 \text{ GeV} \right),
$$

$$
\langle \phi' \rangle = \begin{bmatrix} \frac{1}{\sqrt{2}} v' \\ 0 \end{bmatrix} \quad \left( v' = \sqrt{v_1'^2 + v_2'^2} \right),
$$

are responsible for spontaneously breaking the ordinary and dark electroweak symmetries.

According to the symmetry breaking pattern (2.4), the fermion masses thus should be

$$
\mathcal{L} \supset -m_u \bar{u}_L u_R - m_d \bar{d}_L d_R - m_e \bar{e}_L e_R - m'_u \bar{u}'_R u'_L - m'_d \bar{d}'_R d'_L - m'_e \bar{e}'_R e'_L - \frac{1}{2} \bar{\nu} e'_R e'_L + \text{H.c. (3.4)}
$$

with

$$
m_f = \frac{1}{\sqrt{2}} y_f v, \quad m'_f = \frac{1}{\sqrt{2}} y'_f v', \quad \bar{m}_{e'} = \frac{1}{\sqrt{2}} \Delta v'_{em},
$$

$$
m_{LR} = \frac{1}{\sqrt{2}} y_{\nu} v, \quad m_{R'L'} = \frac{1}{\sqrt{2}} y_{\nu'} v'.
$$

Note the dark charged leptons should be the so-called pseudo-Dirac particles for $v'_{em} \ll v'$.

As for the ordinary and dark neutrinos, their mass matrix can be block-diagonalized if the off-diagonal blocks are much lighter than the diagonal block,

$$
\mathcal{L} \supset -m_\nu \bar{\nu}_L \nu'_R - M_N \bar{\nu}_R \nu'_L + \text{H.c. (3.5)}
$$

with

$$
m_\nu = -m_{LR} \frac{1}{M_N} m_{R'L'}.
$$

Clearly, the ordinary left-handed neutrinos and the dark right-handed neutrinos can form the extremely light Dirac neutrinos as their masses are highly suppressed by the masses between the ordinary right-handed neutrinos and the dark left-handed neutrinos. This Dirac seesaw is definitely a variation of the canonical Majorana seesaw, see figure 1. For the following discussions we can conveniently define the mass eigenstates by a proper phase rotation,

$$
N_i = \nu_{Ri} + \nu'_{Li} \quad \text{with} \quad M_N = \text{diag}\{M_{N_1}, M_{N_2}, M_{N_3}\}. \quad (3.7)
$$
As long as the CP is not conserved, the heavy Dirac fermions composed of the ordinary right-handed neutrinos and the dark left-handed neutrinos can have the lepton-number-conserving decays to generate a lepton asymmetry $\bar{\eta}_L$ stored in the ordinary leptons and an opposite lepton asymmetry $\bar{\eta}'_L$ stored in the dark leptons,

$$\bar{\eta}_L = -\bar{\eta}'_L \propto \varepsilon_{N_i}. \quad (3.8)$$

Here $\varepsilon_{N_i}$ is the CP asymmetry defined as below,

$$\varepsilon_{N_i} = \frac{\Gamma(N_i \to l_L \phi^*) - \Gamma(N_i \to l'_L \phi)}{\Gamma_{N_i}} = \frac{\Gamma(N_i \to l'_R \phi') - \Gamma(N_i \to l'_R \phi'^*)}{\Gamma_{N_i}} \quad (3.9)$$

with $\Gamma_{N_i} = \Gamma(N_i \to l_L \phi^*) + \Gamma(N_i \to l'_L \phi') = \Gamma(N_i \to l'_R \phi^*) + \Gamma(N_i \to l'_R \phi')$.

We can calculate the decay width at tree level,

$$\Gamma_{N_i} = \frac{1}{16\pi} [(y^L_i y^\nu_{ij})_{ii} + (y^L_{i'} y^\nu_{ij})_{ij}] M_{N_i}, \quad (3.10)$$

and the CP asymmetry at one-loop level,

$$\varepsilon_{N_i} = \frac{1}{4\pi} \sum_{j \neq i} \frac{\text{Im}[(y^L_{i} y^\nu_{ij})(y^L_{i'} y^\nu_{ij})_{ij}]}{(y^L_{i} y^\nu_{ij})_{ii} + (y^L_{i'} y^\nu_{ij})_{ij}} \frac{M_{N_i} M_{N_j}}{M_{N_i}^2 - M_{N_j}^2}. \quad (3.11)$$

The relevant diagrams are shown in figure 2.

### 4 Dark matter mass

In the absence of other baryon asymmetries, the produced ordinary lepton asymmetry $\bar{\eta}_L$ is equivalent to an ordinary $B - L$ asymmetry $\eta_{B-L} = -\bar{\eta}_L$ while the dark lepton asymmetry $\bar{\eta}'_L$ is equivalent to a dark $B - L$ asymmetry $\eta'_{B-L} = -\bar{\eta}'_L$. The ordinary SU(2)$_L$ sphaleron processes and the dark SU(2)$_R$ sphaleron processes then will partially transfer the ordinary and dark $B - L$ asymmetries to an ordinary baryon asymmetry $\eta_B$ and a dark baryon asymmetry $\eta'_B$, respectively,

$$\eta_B = C \eta_{B-L} = -C \bar{\eta}_L \quad \text{with} \quad C = \frac{28}{79} \quad (4.1a)$$
Note when computing the dark lepton-to-baryon conversation factor $C'$ we should take the $[\text{SU}(2)']^c_R$-triplet scalar $\Delta'_R$ into account since this scalar drives the dark electromagnetic symmetry breaking much below the dark electroweak scale.

After the dark electromagnetic symmetry breaking, the dark charged leptons acquire a lepton-number-violating Majorana mass term so that the final dark charged lepton asymmetry cannot survive at all [144]. The lightest dark charged lepton denoted as the dark electron will only give a negligible relic density if its mass is at the GeV scale. Furthermore, we will show later the dark electron will only leave a thermally produced relic density.

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Here $\alpha'$ is the dark fine-structure constant. It is easy to check the dark electron will only give a negligible relic density if its mass is at the GeV scale. Furthermore, we will show later the dark electron will only leave a thermally produced relic density. Therefore, if the lightest dark nucleon $N'$ is expected to serve as the dark matter particle, its mass should be determined by

$$m_{N'} = C' \frac{\Omega_{\text{DM}} h^2}{\langle \sigma_{\nu e}^- \rangle v_{\text{vel}}} \approx 14.79 \text{GeV} \left( \frac{\Omega_{\text{DM}} h^2/0.1199}{\Omega_B h^2/0.2205} \right) .$$

5 Dark matter detection

We can calculate the $U(1)_{\text{em}} \times U(1)_{\text{em}}'$ kinetic mixing at one-loop level,

$$\mathcal{L} \supset -\frac{\epsilon}{2} A_{\mu} A^{\mu} \quad \text{with} \quad \epsilon = \frac{\sqrt{\alpha \alpha'}}{12 \pi} \sum_{QQ'} \sum_{QQ'} \sigma_{QQ'} C_{QQ'} C_{QQ'} \ln \left[ \frac{M_{QQ'}^2}{\mu^2} \right] .$$

Here $Q, Q' = \pm 1, \pm \frac{1}{2}, \pm \frac{3}{2}$ are the ordinary and dark electric charges of the scalars $\sigma(Q, Q') \in (16 \times \overline{16})_R$, $M_{QQ'}$ denotes the $\sigma(Q, Q')$'s mass, $\mu$ is a renormalizable scale, while $C_{QQ'} = 1$ for $Q, Q' = \pm 1$ and $C_{QQ'} = 3$ for $Q, Q' = \pm \frac{1}{2}, \pm \frac{3}{2}$ are the color factors. Clearly, we have $\epsilon = 0$ at the GUT scale. However, such kinetic mixing can appear after the left-right symmetry breaking,

$$\epsilon = \frac{\sqrt{\alpha \alpha'}}{12 \pi} \left[ \ln \left( \frac{1 + \frac{1}{2} \lambda v^2_L / M_1^2}{1 + \frac{1}{2} \lambda (v^2_L + v^2_R) / M_1^2} \right) \ln \left( \frac{1 + \frac{1}{2} \lambda v^2_L / M_1^2}{1 + \frac{1}{2} \lambda (v^2_L + v^2_R) / M_1^2} \right) \right]$$

$$\approx \frac{\sqrt{\alpha \alpha'}}{48 \pi} \frac{\lambda^2 v^2_L v^2_R}{M_1^4} \quad \text{for} \quad M_3^2 \gg M_1^2 \gg \lambda v^2_L , \lambda v^2_R$$

$$\approx \frac{\sqrt{\alpha \alpha'}}{48 \pi} \lambda^2 = 10^{-9} \left( \frac{\lambda}{0.0046} \right)^2 \sqrt{\frac{\alpha'}{\alpha}} \quad \text{for} \quad M_1^2 \sim v^2_L \sim v^2_R .$$

(5.2)
In the above calculation we have simplified the left-right level interactions as
\[
V \supset \lambda \left( \chi_R \tilde{\Sigma}_{f_R f_L} \Sigma_{f_R f_L} \chi_R + \chi_L \tilde{\Sigma}_{f_R f_L} \Sigma_{f_R f_L} \chi_L + \chi_R \Sigma_{f_R f_L} \Sigma_{f_R f_L} \chi_R + \chi_L \Sigma_{f_R f_L} \Sigma_{f_R f_L} \chi_L \right) + M^2 f_{f_R \neq f_L} \left( \Sigma_{f_{R,L} f_{R,L}} \right) + M^2 \sum_{f_{R,L} \neq f_{R,L}} \text{Tr} \left( \Sigma_{f_{R,L} f_{R,L}} \right) .
\]  

(5.3)

Due to the U(1)$_{em}$ × U(1)′$_{em}$ kinetic mixing, the physically dark photon will couple to not only the dark charged fermions but also the ordinary charged fermions although the physically ordinary photon doesn’t couple to the dark charged fermions,
\[
\mathcal{L} \supset e \left( \hat{A}_\mu - \frac{e}{\sqrt{1-\epsilon^2}} \hat{A}'_\mu \right) \left( - \bar{e} \gamma^\mu e - \frac{1}{3} \bar{d} \gamma^\mu d + \frac{2}{3} \bar{u} \gamma^\mu u \right)
\]
\[+ e \hat{A}'_\mu \left( - \bar{e}' \gamma^\mu e' - \frac{1}{3} \bar{d}' \gamma^\mu d' + \frac{2}{3} \bar{u}' \gamma^\mu u' \right) ,
\]

(5.4)

where the physical photons have been defined by [145]
\[
\hat{A}_\mu = A_\mu + \epsilon A'_\mu , \quad \hat{A}'_\mu = \sqrt{1-\epsilon^2} A'_\mu .
\]

(5.5)

Once the kinematics is allowed, the dark photon can efficiently decay into the ordinary charged fermion pairs,
\[
\Gamma_{A' \rightarrow ff} = \frac{\epsilon^2 e^2 Q_f^2 C_f m_{A'}}{12\pi} \left( 1 - \frac{m_f^2}{m_{A'}^2} \right) \sqrt{1 - 4 \frac{m_f^2}{m_{A'}^2}} ,
\]

(5.6)

with the dark photon mass $m_{A'}^2 = 16\pi \alpha' \nu_{em}^2$ and the ordinary electric charges $Q_{e,\mu,\tau} = -1$, $Q_{d,s,b} = -\frac{1}{3}$ and $Q_{u,c,t} = +\frac{2}{3}$.

The dark photon can mediate an elastic scattering of the dark nucleons off the ordinary nucleons. If the dark proton is the dark matter particle, its scattering will have a spin-independent cross section,
\[
\sigma_{p'N \rightarrow p'N}(Z, A) \simeq \frac{\epsilon^2 \alpha_\alpha}{2} \frac{[m_{p'} m_p/(m_{p'} + m_p)]^2}{m_{A'}^2} \left( \frac{Z}{A} \right)^2 \simeq 5.1 \times 10^{-46} \text{cm}^2 \left( \frac{Z}{A} \right)^2 \left( \frac{\epsilon}{10^{-9}} \right)^2 \left( \frac{100 \text{ MeV}}{m_{A'}} \right)^4 .
\]

(5.7)

Such dark matter scattering can be verified in the direct detection experiments [146]. If the dark neutron is the dark matter particle, its scattering off the ordinary nucleons will be further suppressed by its dark magnetic moment [136]. In the present SO(10) × SO(10)' framework, we can expect a dark nucleon decay according to the ordinary proton decay. It should be noted the dark leptoquark scalars $\Omega_{R,L}$ can be allowed much lighter than the ordinary ones $\Omega_{L,R}$. This means the dark nucleon decay can be fast enough to open a window for the indirect detection experiments although the ordinary proton decay is extremely slow. For example, we can have the dark matter decay chains $p' \rightarrow \pi^0 e^+$ (or $n' \rightarrow \pi^0 \nu'_p$), $\pi^0 \rightarrow \gamma \gamma'$, $\gamma' \rightarrow e^+ e^-$, $u\bar{u}, d\bar{d}, \mu^+ \mu^-$, . . . . Clearly, if the dark photon mass is about 1–2 MeV, the dark matter should mostly decay into the positron/electron pairs.
The dark electromagnetic interactions will lead to a dark matter self-interaction. For example, if the dark proton is the dark matter particle, we can have the self-interacting cross section as below,

\[
\sigma_{p'p'\rightarrow p'p'} = \frac{\pi\alpha'^2 m_{p'}^2}{2m_{A'}^4}
\]

\[
= \left(\frac{\alpha'}{\alpha}\right)^2 \left(\frac{m_{p'}}{15 \text{ GeV}}\right)^2 \left(\frac{100 \text{ MeV}}{m_{A'}}\right)^4 \times 7.4 \times 10^{-26} \text{ cm}^2. \tag{5.8}
\]

In the case the dark neutron is the dark matter particle, its self-interaction should be determined by a dark magnetic moment and hence should be further suppressed. The dark strong interactions will also result in the dark matter self-interaction. We have known the scattering of the ordinary neutrons off the ordinary protons should have a cross section \(\sigma_{np} \sim 10^{-24} \text{ cm}^2\). The isospin symmetry then can give \(\sigma_{pp} \simeq \sigma_{nn} \simeq \sigma_{np}\). We hence can estimate the cross sections of the dark nucleons’ self-interactions to be

\[
\sigma_{N'N'} \sim \left(\frac{\Lambda_{\text{QCD}}}{\Lambda_{\text{QCD}'}}\right)^2 \sigma_{np} \sim 10^{-26} \text{ cm}^2 \text{ for } \Lambda_{\text{QCD}'} \sim 10 \Lambda_{\text{QCD}}, \tag{5.9}
\]

with \(\Lambda_{\text{QCD}}\) and \(\Lambda_{\text{QCD}'}\) being the ordinary and dark hadronic scales. It is easy to see that the self-interactions (5.8) and (5.9) can be consistent with the constraints from simulations and observations [147–150],

\[
\sigma_{\text{self}} \lesssim 1 \text{ cm}^2 \text{ gram}^{-1} m_{N'}
\]

\[= 2.6 \times 10^{-23} \text{ cm}^2 \text{ for } m_{N'} \simeq 15 \text{ GeV}. \tag{5.10}
\]

### 6 Discrete mirror symmetry

We can impose a softly or spontaneously broken mirror symmetry under which the ordinary and dark fields transform as

\[
16_F \leftrightarrow 16_F', \quad 16_H \leftrightarrow 16_H', \quad \ldots \tag{6.1}
\]

to simplify the parameter choice,

\[
y_f = y_{f'}', \quad \tilde{y}_f = \tilde{y}_{f'}', \quad f_\Sigma = f_\Sigma', \quad \ldots \tag{6.2}
\]

By further assuming

\[
\frac{v_1'}{v_2'} = \frac{v_1}{v_2}, \tag{6.3}
\]

we can read

\[
\frac{\langle v' \rangle}{\langle v \rangle} = \frac{m_{u'}}{m_u} = \frac{m_{d'}}{m_d} = \frac{m_{s'}}{m_s} = \frac{m_{c'}}{m_c} = \frac{m_{b'}}{m_b} = \frac{m_t'}{m_t}
\]

\[= \frac{m_{e'}}{m_e} = \frac{m_{\mu'}}{m_{\mu}} = \frac{m_{\tau'}}{m_{\tau}}. \tag{6.4}
\]
We then can make use of the beta functions of the ordinary and dark QCDs to determine
\[
\Lambda_{\text{QCD}'} = \left( \frac{v'}{v} \right)^{-\frac{11}{3}} (m_u m_d m_s \Lambda_{\text{QCD}})^{\frac{2}{3}} \Lambda_{\text{QCD}}^{\frac{2}{3}} \quad \text{for} \quad \Lambda_{\text{QCD}'} < m_{u'}.
\] (6.5)

Since the dark hadronic scale is lighter than the dark quark masses, we can simply ignore the dark QCD contributions to the masses of the dark baryons and mesons such as
\[
m_{p'} \simeq 2m_{u'} + m_{d'}, \quad m_{n'} \simeq m_{u'} + 2m_{d'}, \quad m_{e'} \simeq m_{e^0} \simeq m_{e^\pm} \simeq m_{u'} + m_{d'}.
\] (6.6)

From eqs. (6.4)–(6.6), we can obtain
\[
m_{e'} = 1.5 \text{ GeV}, \quad m_{u'} = 3.75 \text{ GeV}, \quad m_{d'} = 7.5 \text{ GeV}, \\
\Lambda_{\text{QCD}'} = 2 \text{ GeV}, \quad m_{p'} = 15 \text{ GeV}, \quad m_{n'} = 18.75 \text{ GeV}, \\
m_{e'} = 11.25 \text{ GeV},
\] (6.7)
by inputting
\[
v' = 3000 v, \quad m_e = 0.511 \text{ MeV}, \quad m_{u} = 1.25 \text{ MeV}, \\
m_{d} = 2.5 \text{ MeV}, \quad m_s = 100 \text{ MeV}, \quad \Lambda_{\text{QCD}} = 200 \text{ MeV}.
\] (6.8)

In this case, the dark proton is the lightest dark nucleon and hence is the dark matter particle.

Another interesting consequence of the above parameter choice is that the Dirac seesaw (3.5) now can be given by
\[
m_{\nu} = -\frac{v'}{v} m_{LR}^T T_{LR} = -3000 m_{LR}^T M_{N}^T T_{LR},
\] (6.9)
which doesn’t contain unknown parameters compared with the canonical Majorana seesaw.

7 Summary

In this paper we have proposed an SO(10) × SO(10)′ model to simultaneously explain the smallness of the Dirac neutrino masses and the coincidence between the ordinary and dark matter. Specifically we introduced a (16 × 16′) \( \tilde{H} \) scalar crossing the ordinary SO(10) sector and the dark SO(10)′ sector. This (16 × 16′) \( \tilde{H} \) scalar can acquire an induced vacuum expectation value after the 16 \( \tilde{H} \) and 16′ \( \tilde{H} \) scalars drive the spontaneous breaking of the ordinary and dark left-right symmetries. Consequently the ordinary right-handed neutrinos and the dark left-handed neutrinos can form the heavy Dirac fermions to highly suppress the masses between the ordinary left-handed neutrinos and the dark right-handed neutrinos. The decays of such heavy Dirac fermions can generate an ordinary lepton asymmetry and an opposite dark lepton asymmetry. We hence can obtain an ordinary baryon asymmetry and a dark baryon asymmetry due to the SU(2)\(_L\) and SU(2)\(_R\) sphaleron processes. By taking into account the difference between the ordinary and dark lepton-to-baryon conversations, we can expect the lightest dark nucleon as the dark matter particle to have a determined mass around 15 GeV. Furthermore, the (16 × 16′) \( \tilde{H} \) scalar can mediate a small U(1)\(_{\text{em}}\) × U(1)′\(_{\text{em}}\) kinetic mixing after the ordinary and dark left-right symmetry breaking. Therefore, the dark proton as the dark matter particle can be verified by the direct and indirect detection experiments. Alternatively, if the dark neutron is the dark matter particle, it can be only found by the indirect detection experiments. Our model can accommodate a softly broken mirror symmetry to simplify the parameters.
Acknowledgments

This work was supported by the Shanghai Jiao Tong University under Grant No. WF220407201 and the Shanghai Laboratory for Particle Physics and Cosmology under Grant No. 11DZ2260700.

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