Leptogenesis with Dirac Neutrinos and Dark Matter

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This talk is based on my works: Pei-Hong Gu, 1209.4579; 1410.5753; 1410.5759.

中国科学技术大学 交叉学科理论研究中心　2014.10.23

Outline

Introduction

‡ No primordial antimatter significantly exists in the present universe.

‡ All particles should come in particle-antiparticle pairs.

‡ An initial matter-antimatter asymmetry cannot survive after inflation.

We need a dynamical baryogenesis mechanism!

The matter-antimatter asymmetry is as same as a baryon asymmetry, which has been precisely measured by the cosmological observations (e.g. Planck collaboration, 1303.5076.),

$$\Omega_B h^2 = \frac{\rho_B}{\rho_c} h^2 = 0.02205 \pm 0.00028\,, \quad \rho_B = n_B m_p\,,$$

$$\eta_B = \frac{n_B}{n_\gamma} \simeq 2.68 \times 10^{-8} (\Omega_B h^2) = (5.91 \pm 0.08) \times 10^{-10}\,.$$

If CPT (C – charge conjugation, P – parity, T – time reversal.) is invariant, any successful baryogenesis mechanisms should satisfy the Sakharov conditions (Sakharov 67'):

‡ baryon number nonconservation,

‡ C and CP violation,

‡ departure from equilibrium.

$$\left.\begin{array}{l} B \xrightarrow{C} -B \ \text{ for } \ q_{L(R)} \xrightarrow{C} q^c_{L(R)} \\ B \xrightarrow{CP} -B \ \text{ for } \ q_L \xrightarrow{CP} q^c_R \end{array}\right\} \Longrightarrow n_B \equiv n_b - n_{\bar{b}} = \frac{1}{3}(n_{q_L} - n_{\bar{q}_L} + n_{q_R} - n_{\bar{q}_R}) \xrightarrow{C,CP} 0 \,.$$

$$\langle B \rangle = \mathrm{Tr}(e^{-\frac{H}{T}} B) = \mathrm{Tr}[e^{-\frac{H}{T}} (CPT)^{-1} B (CPT)] = \mathrm{Tr}[e^{-\frac{H}{T}} (-B)] = -\langle B \rangle \Rightarrow \langle B \rangle = 0 \,.$$

Both of the baryon (B) and lepton (L) numbers are violated by quantum effects in the standard model ('t Hooft, 76'.). The transition of the baryon and lepton numbers from one vacuum to the next vacuum is

$$\partial_\mu J^\mu_B = \partial_\mu J^\mu_L = N_f \frac{g_2^2}{32\pi^2} \epsilon_{\mu\nu\rho\sigma} \mathrm{Tr}\,(W^{\mu\nu} W^{\rho\sigma}) \Rightarrow \Delta B = \Delta L = N_f\,, \quad \Delta(B-L) = 0\,.$$

At zero temperature, the baryon and lepton number violating processes via a tunneling between the topologically distinct vacua are highly suppressed. However, such sphaleron processes could become efficient during the temperatures near and above the electroweak phase transition (Kuzmin, Rubakov, Shaposhnikov, 85'.),

$$100\,\mathrm{GeV} < T < 10^{12}\,\mathrm{GeV}\,.$$

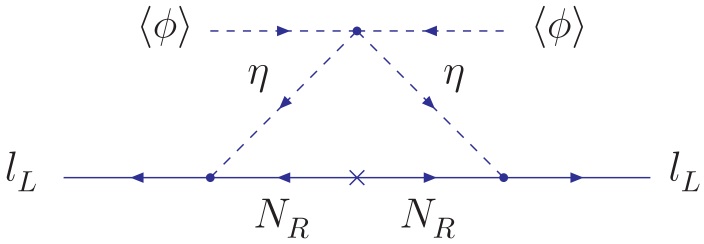
In the standard model, the sphaleron processes, the CKM phase and the electroweak phase transition can fulfill all of the three Sakharov conditions to realize an electroweak baryogenesis scenario (e.g. Morrissey, Ramsey-Musolf, 12'.).

Unfortunately, the baryon asymmetry induced by the electroweak baryogenesis in the standard model is too small to explain the observed value.

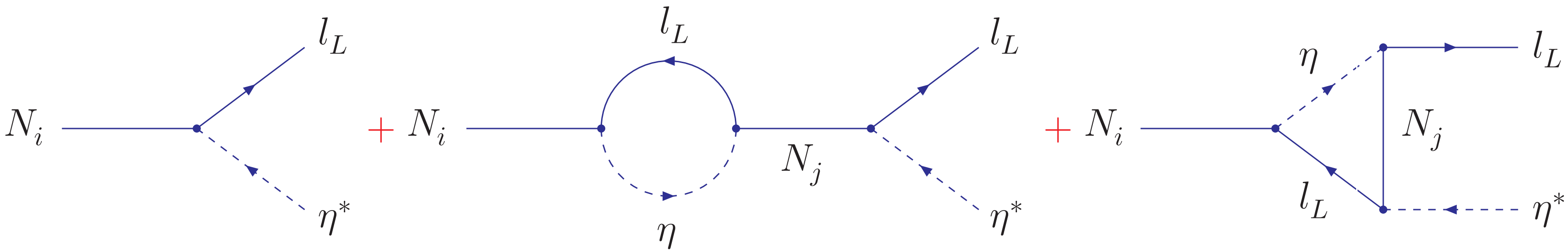
‡ The electroweak phase transition should be strongly first-order to avoid the washout of the induced baryon asymmetry. This requires the Higgs boson lighter than about $m_H < 40$ GeV, which is much lower than the experimental value $m_H = 125$ GeV.

‡ Even if the electroweak phase transition is strongly first-order, the induced baryon asymmetry can only arrive at the order of $\eta_B = \mathcal{O}(10^{-20})$.

We need a baryogenesis beyond the standard model!



$$N_i = N_{Ri} + N_{Ri}^c$$



The leptogenesis (Fukugita, Yanagida, 86'.) mechanism within various seesaw models (Minkowski, 77'; Yanagida, 79'; Gell-Mann, Ramond, Slansky, 79'; Glashow, 80'; Mohapatra, Senjanović, 80'; Magg, Wetterich, 80'; Schechter, Valle, 80'; Cheng, Li, 80'; Lazarides, Shafi, Wetterich, 81'; Mohapatra, Senjanović, 81'; Foot, Lew, He, 89'.) can simultaneously explain the observed baryon asymmetry and the small neutrino masses. The leptogenesis can even include dark matter (Ma, 06'.).

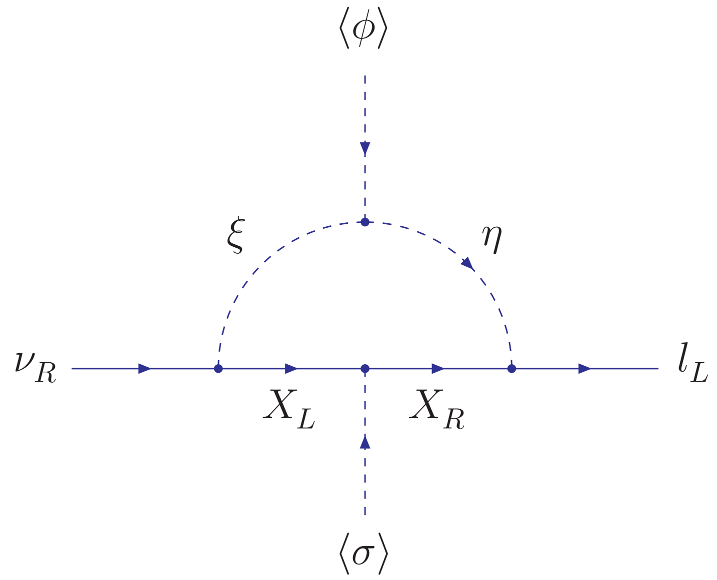
In the popular leptogenesis-seesaw context, the lepton number is explicitly violated and the neutrinos are Majorana particles.

‡ The Majorana nature of neutrinos is just a theoretical assumption and has not been experimentally confirmed so far. Even if the running and coming experiments observed a neutrinoless double beta decay, they could not distinguish the Majorana neutrinos from the pseudo-Dirac neutrinos.

‡ The neutrinos can be Dirac particles just like the charged fermions.

It is worth exploring the picture of Dirac neutrinos (Roncadelli, Wyler, 83'; Roy, Shanker, 84'; PHG, He, 06'.)!

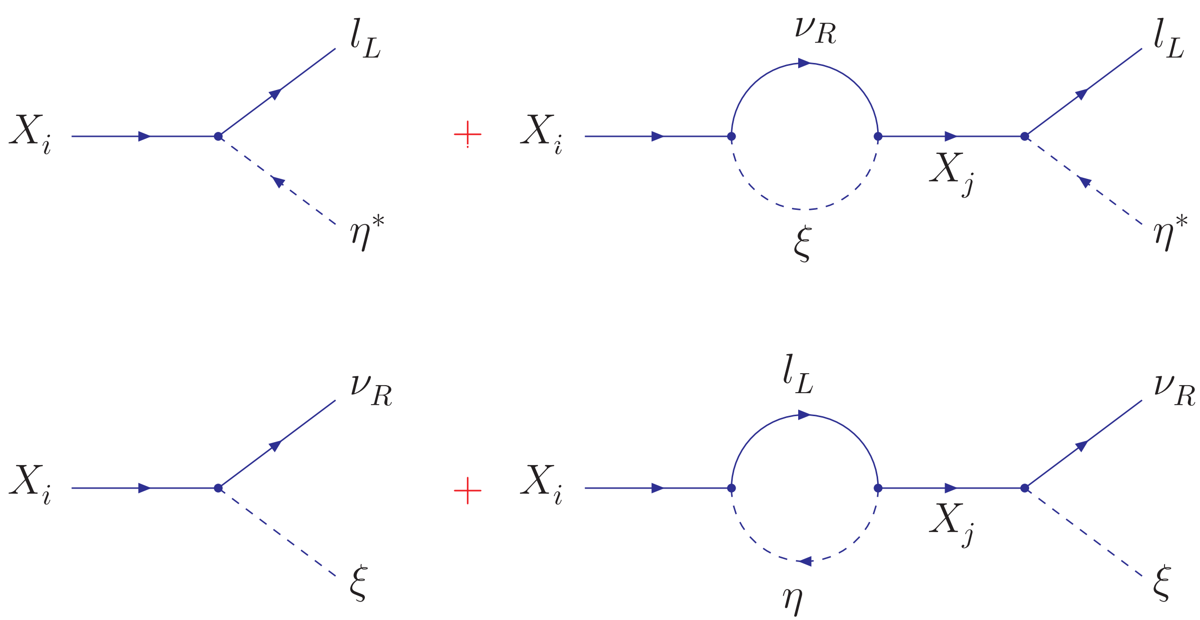
A lepton-number-conserving leptogenesis can be realized in the presence of Dirac neutrinos (Dick, Lindner, Ratz, Wright, 99'; Murayama, Pierce, 02'; Thomas, Toharia, 06'; Abel, Page, 06'; PHG, He, 06'; PHG, He, Sarkar, 07'.).



$$X_i = X_{Li} + X_{Ri}$$

The dark matter field can play an essential role in the leptogenesis with Dirac neutrinos

(PHG, Sarkar, 07'.).



A $G_{\rm SM} \times G_{\rm SM'}$ Scenario

It is intriguing that the baryonic and dark matter contribute comparable energy densities to the present universe (e.g. Planck collaboration, 1303.5076.),

$$\Omega_{DM}h^2 : \Omega_B h^2 = (0.1199 \pm 0.0027) : (0.02205 \pm 0.00028) \simeq 5\,.$$

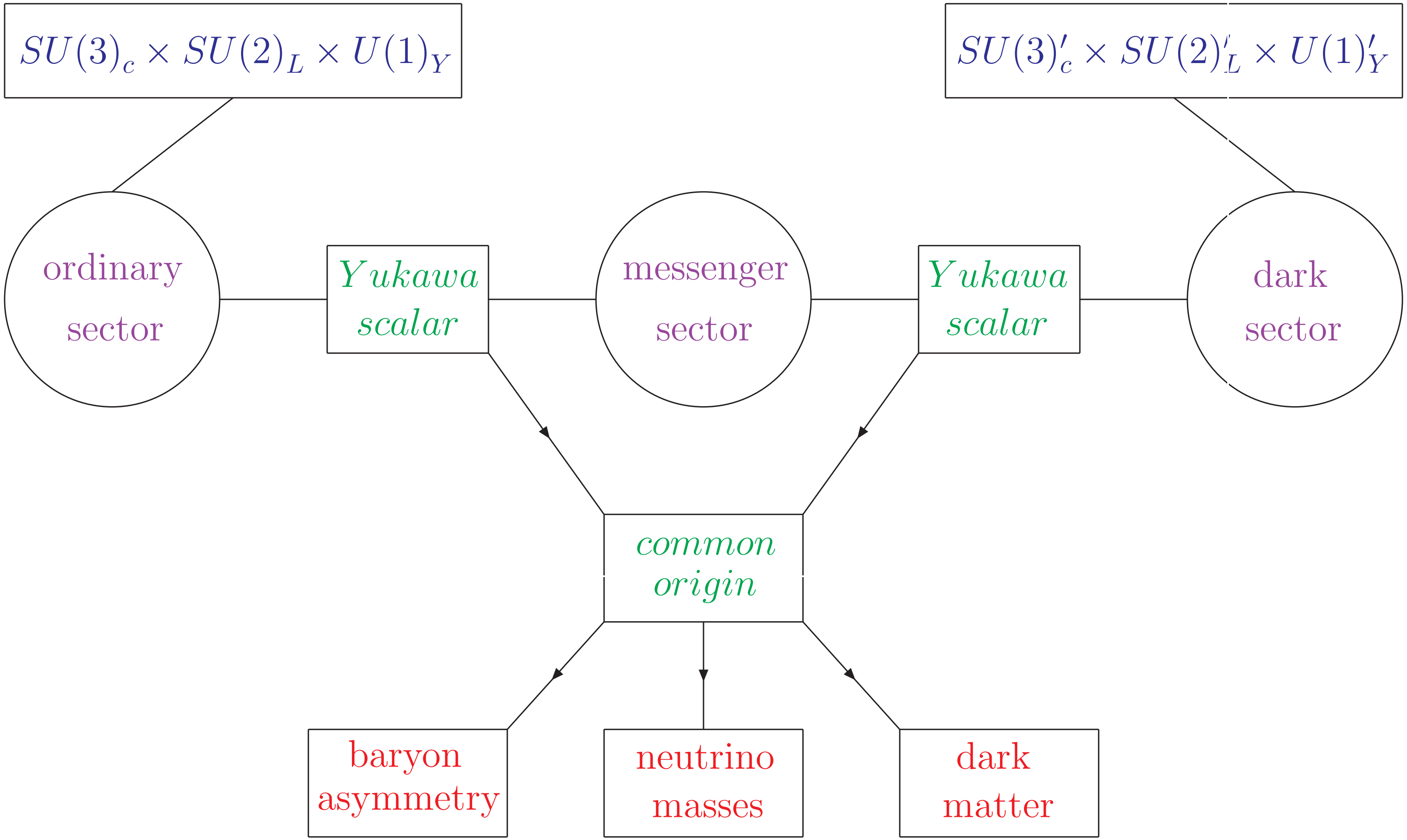
This coincidence can be elegantly explained in the asymmetric dark matter scenario.

‡ The dark matter relic density is an asymmetry between dark matter and antimatter.

‡ The dark matter asymmetry is produced together with the baryon asymmetry.

The Dirac neutrino masses, the baryon asymmetry and the dark matter relic density can have a common origin (PHG, 1209.4579.).

The model



$SU(3)_c \times SU(2)_L \times U(1)_Y$
$SU(3)'_c \times SU(2)'_L \times U(1)'_Y$
ordinary sector
Yukawa scalar
messenger sector
Yukawa scalar
dark sector
common origin
baryon asymmetry
neutrino masses
dark matter

Ordinary and dark fields

$$q_L(\mathbf{3},\mathbf{2},+\tfrac{1}{6}) = \begin{bmatrix} u_L \\ d_L \end{bmatrix} \quad \Longleftrightarrow \quad q'_L(\mathbf{3},\mathbf{2},-\tfrac{1}{6}) = \begin{bmatrix} d'_L \\ u'_L \end{bmatrix}\,,$$

$$d_R(\mathbf{3},\mathbf{1},-\tfrac{1}{3}) \quad \Longleftrightarrow \quad d'_R(\mathbf{3},\mathbf{1},+\tfrac{1}{3})\,,$$

$$u_R(\mathbf{3},\mathbf{1},+\tfrac{2}{3}) \quad \Longleftrightarrow \quad u'_R(\mathbf{3},\mathbf{1},-\tfrac{2}{3})\,;$$

$$l_L(\mathbf{1},\mathbf{2},-\tfrac{1}{2}) = \begin{bmatrix} \nu_L \\ e_L \end{bmatrix} \quad \Longleftrightarrow \quad l'_L(\mathbf{3},\mathbf{2},+\tfrac{1}{2}) = \begin{bmatrix} e'_L \\ \nu'_L \end{bmatrix}\,,$$

$$e_R(\mathbf{1},\mathbf{1},-1) \quad \Longleftrightarrow \quad e'_R(\mathbf{1},\mathbf{1},+1)\,;$$

$$\phi(\mathbf{1},\mathbf{2},-\tfrac{1}{2}) = \begin{bmatrix} \phi^0 \\ \phi^-_0 \end{bmatrix} \quad \Longleftrightarrow \quad \phi'(\mathbf{3},\mathbf{2},+\tfrac{1}{2}) = \begin{bmatrix} \phi'^+ \\ \phi'^0 \end{bmatrix}\,,$$

$$\delta(\mathbf{1},\mathbf{1},-2) \quad \Longleftrightarrow \quad \delta'(\mathbf{1},\mathbf{1},+2)\,;$$

$$\cdots \quad \Longleftrightarrow \quad \cdots\,.$$

Ordinary and dark symmetry breaking

$$\langle \phi' \rangle \equiv \begin{bmatrix} 0 \\ \langle \phi'^0 \rangle \end{bmatrix} \neq \langle \phi \rangle \equiv \begin{bmatrix} \langle \phi^0 \rangle \\ 0 \end{bmatrix} \simeq 174\,\text{GeV}\,,$$

$$\langle \delta' \rangle \neq \langle \delta \rangle \equiv 0\,,$$

$$SU(3)_c \times SU(2)_L \times U(1)_Y \xrightarrow{\langle \phi \rangle} SU(3)_c \times U(1)_{em}\,,$$

$$SU(3)'_c \times SU(2)'_L \times U(1)'_Y \xrightarrow{\langle \phi \rangle} SU(3)'_c \times U(1)'_{em} \xrightarrow{\langle \delta' \rangle} SU(3)'_c\,.$$

Note the dark photon will become massive while the ordinary photon will keep massless.

Ordinary and dark fermion masses

$$\mathcal{L} \supset -(y_d)_{ij}\bar{q}_{Li}\tilde{\phi}d_{Rj} - (y_u)_{ij}\bar{q}_{Li}\phi u_{Rj} - (y_e)_{ij}\bar{l}_{Li}\tilde{\phi}e_{Rj} - (y_\delta)_{ij}\delta\bar{e}_{Ri}e^c_{Rj} + \text{H.c.}\,.$$

$$\mathcal{L} \supset -(y_{d'})_{ij}\bar{q}'_{Li}\tilde{\phi}'d'_{Rj} - (y_{u'})_{ij}\bar{q}'_{Li}\phi' u'_{Rj} - (y_{e'})_{ij}\bar{l}'_{Li}\tilde{\phi}'e'_{Rj} - (y_{\delta'})_{ij}\delta'\bar{e}'_{Ri}e'^c_{Rj} + \text{H.c.}\,.$$

$$m_d = y_d\langle\phi\rangle \quad \Longleftrightarrow \quad m_{d'} = y_{d'}\langle\phi'\rangle\,,$$

$$m_u = y_u\langle\phi\rangle \quad \Longleftrightarrow \quad m_{u'} = y_{u'}\langle\phi'\rangle\,,$$

$$m_e = y_e\langle\phi\rangle \quad \Longleftrightarrow \quad m_{e'} = y_{e'}\langle\phi'\rangle\,.$$

Note the dark charged leptons should be the quasi-Dirac particles since they have a small Majorana mass term,

$$\delta m_{e'} = y_{\delta'}\langle\delta'\rangle \ll m_{e'}\,.$$

Messenger fields

Guage-singlet fermions: $\begin{cases} N_R(\mathbf{1},\mathbf{1},0)(\mathbf{1},\mathbf{1},0) & \text{with a lepton number} \quad L=+1\,, \\ N'_R(\mathbf{1},\mathbf{1},0)(\mathbf{1},\mathbf{1},0) & \text{with a lepton number} \quad L=-1\,. \end{cases}$

$$\begin{aligned} \mathcal{L} &\supset -(y_N)_{\alpha i}\bar{l}_{L\alpha}\phi N_{Ri} - (y_{N'})_{\alpha i}\bar{l}'_{L\alpha}\phi' N'_{Ri} - (M_N)_{ij}\bar{N}'^{c}_{Ri}N_{Rj} + \text{H.c.} \\ &= -(y_N)_{\alpha i}\bar{l}_{L\alpha}\phi N_i - (y_{N'})_{\alpha i}\bar{l}'_{L\alpha}\phi' N_i + \text{H.c.} - M_{N_i}\bar{N}_i N_i \quad \text{with} \quad N_i = N_{Ri} + N'^{c}_{Ri}\,. \end{aligned}$$

$[SU(2)_L \times SU(2)'_L]$-bidoublet scalars: $\Sigma_a(\mathbf{1},\mathbf{2},-\frac{1}{2})(\mathbf{1},\mathbf{2},+\frac{1}{2}) = \begin{bmatrix} \bar{\sigma}^+ & \sigma^0 \\ \bar{\sigma}^0 & \sigma^- \end{bmatrix}_a$ with a zero lepton number .

$$\mathcal{L} \supset -f_a\bar{l}_L\Sigma_a l'^{c}_L - \rho_a\phi^\dagger\Sigma_a\phi'^* + \text{H.c.} - M^2_{\Sigma_a}\text{Tr}(\Sigma^\dagger_a\Sigma_a)\,.$$

Neutrino mass generation

If the messenger sector only contains the gauge-singlet fermions, the mass matrix involving the ordinary and dark neutrinos should be

$$\mathcal{L} \supset -\left[\begin{array}{cc} \bar{\nu}_L & \bar{N}'^c_R \end{array}\right] \left[\begin{array}{cc} 0 & y_N \langle \phi \rangle \\ y^T_{N'} \langle \phi' \rangle & M_N \end{array}\right] \left[\begin{array}{c} \nu'^c_L \\ N_R \end{array}\right] + \text{H.c.}\,.$$

We can block diagonalize the above mass matrix to be

$$\mathcal{L} \supset -\left[\begin{array}{cc} \bar{\nu}_L & \bar{N}'^c_R \end{array}\right] \left[\begin{array}{cc} m_\nu & 0 \\ 0 & M_N \end{array}\right] \left[\begin{array}{c} \nu'^c_L \\ N_R \end{array}\right] + \text{H.c.} \text{ with } m_\nu = -y_N \frac{\langle \phi \rangle \langle \phi' \rangle}{M_N} y^T_{N'} \equiv m^{\text{I}}_\nu\,,$$

as long as the seesaw condition is satisfied, i.e.

$$y_N \langle \phi \rangle\,,\; y_{N'} \langle \phi' \rangle \ll M_N\,.$$

The ordinary and dark neutrinos thus will form three light Dirac neutrinos. This mechanism for generating the Dirac neutrino masses may be named as the type-I Dirac seesaw.

If the messenger sector only contains the $[SU(2)_L \times SU(2)'_L]$-bidoublet scalars Σ_a, we can also obtain the light Dirac neutrinos in a natural way. Specifically, the bidoublet scalars can pick up their seesaw-suppressed vacuum expectation values:

$$\langle \Sigma_a \rangle = \begin{bmatrix} 0 & \langle \sigma_a^0 \rangle \\ 0 & 0 \end{bmatrix} \simeq -\frac{\rho_a \langle \phi \rangle \langle \phi' \rangle}{M_{\Sigma_a}^2} \ll \langle \phi \rangle \, , \langle \phi' \rangle \, ,$$

and then the ordinary and dark neutrinos can acquire their Dirac masses,

$$\mathcal{L} \supset -m_\nu \bar{\nu}_L \nu_L'^c + \text{H.c.} \;\; \text{with} \;\; m_\nu = \sum_a f_a \langle \Sigma_a \rangle \equiv \sum_a m_\nu^{\text{II}a} \equiv m_\nu^{\text{II}} \, .$$

This mechanism for generating the Dirac neutrino masses may be named as the type-II Dirac seesaw.

If the messenger sector contains not only the gauge-singlet fermions but also the bidoublet scalars, the masses involving the ordinary and dark neutrinos should be

$$\mathcal{L} \supset -\left[\begin{array}{cc} \bar{\nu}_L & \bar{N}'^c_R \end{array}\right] \left[\begin{array}{cc} \sum_a f_a \langle \Sigma_a \rangle & y_N \langle \phi \rangle \\ y^T_{N'} \langle \phi' \rangle & M_N \end{array}\right] \left[\begin{array}{c} \nu'^c_L \\ N_R \end{array}\right] + \text{H.c.}\,.$$

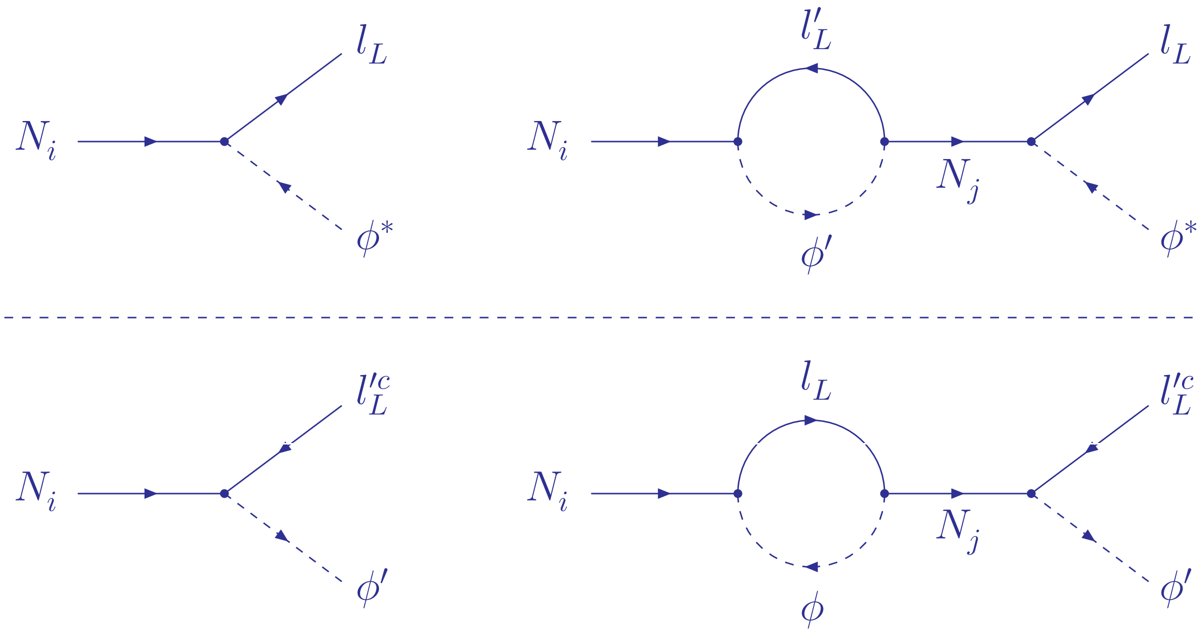
Under the seesaw condition,

$$\sum_a f_a \langle \Sigma_a \rangle \ll y_N \langle \phi \rangle\,,\; y_{N'} \langle \phi' \rangle \ll M_N\,,$$

the masses between the ordinary and dark neutrinos should be a sum of the type-I Dirac seesaw and the type-II Dirac seesaw, i.e.

$$\mathcal{L} \supset -m_\nu \bar{\nu}_L \nu'^c_L + \text{H.c.} \quad \text{with} \quad m_\nu = m^{\mathrm{I}}_\nu + m^{\mathrm{II}}_\nu = -y_N \frac{\langle \phi \rangle \langle \phi' \rangle}{M_N} y^T_{N'} - \sum_a f_a \frac{\rho_a \langle \phi \rangle \langle \phi' \rangle}{M^2_{\Sigma_a}}\,.$$

Ordinary and dark lepton asymmetries



$$\eta_{l_L} = -\eta_{l'_L} \propto \varepsilon_{N_i}\,.$$

The CP asymmetry ε_{N_i} can be calculated at one-loop order,

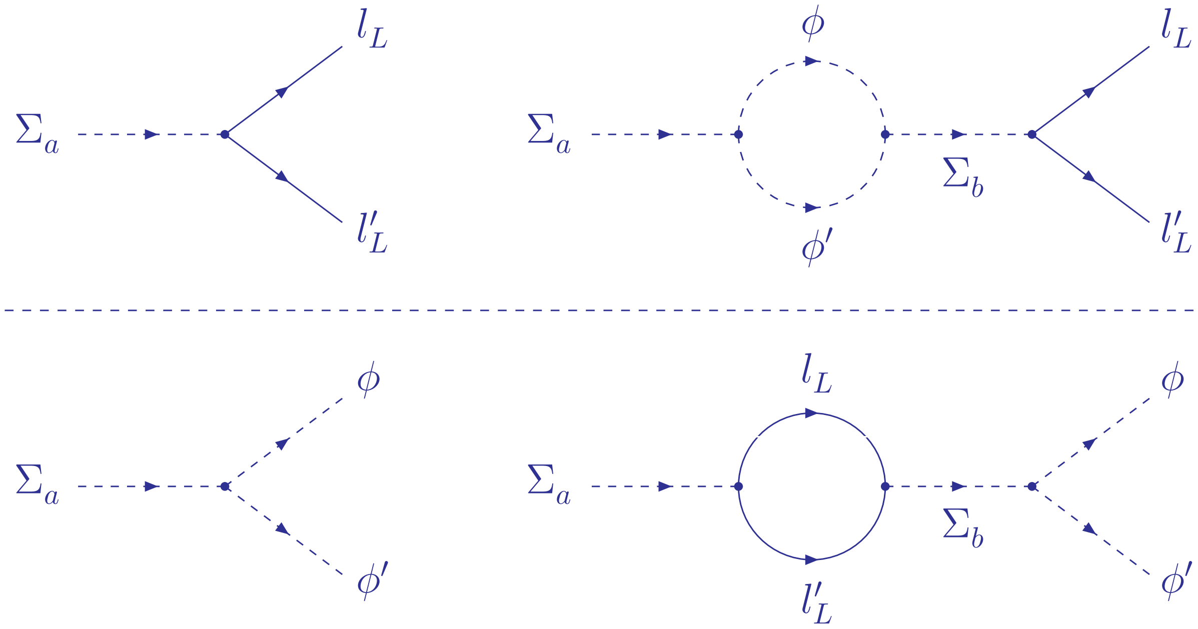
$$\varepsilon_{N_i} = \frac{\Gamma_{N_i \to l_L \phi^*} - \Gamma_{N_i^c \to l_L^c \phi}}{\Gamma_{N_i}} = \frac{\Gamma_{N_i^c \to l'_L \phi'^*} - \Gamma_{N_i \to l'^c_L \phi'}}{\Gamma_{N_i}}$$

$$= -\frac{1}{4\pi} \sum_{j \neq i} \frac{\mathrm{Im}\left[(y_N^\dagger y_N)_{ij} (y_{N'}^\dagger y_{N'})_{ij}\right]}{(y_N^\dagger y_N)_{ii} + (y_{N'}^\dagger y_{N'})_{ii}} \frac{M_{N_i} M_{N_j}}{M_{N_j}^2 - M_{N_i}^2},$$

with Γ_{N_i} being the decay width,

$$\Gamma_{N_i} = \Gamma_{N_i \to l_L \phi^*} + \Gamma_{N_i \to l'^c_L \phi'} = \Gamma_{N_i^c \to l_L^c \phi} + \Gamma_{N_i^c \to l'_L \phi'^*} = \frac{1}{16\pi} [(y_N^\dagger y_N)_{ii} + (y_{N'}^\dagger y_{N'})_{ii}] M_{N_i}.$$

Note at least two gauge-singlet Dirac fermions should be introduced to give a nonzero CP asymmetry ε_{N_i}.



$$\eta_{l_L} = -\eta_{l'_L} \propto \varepsilon_{\Sigma_a} .$$

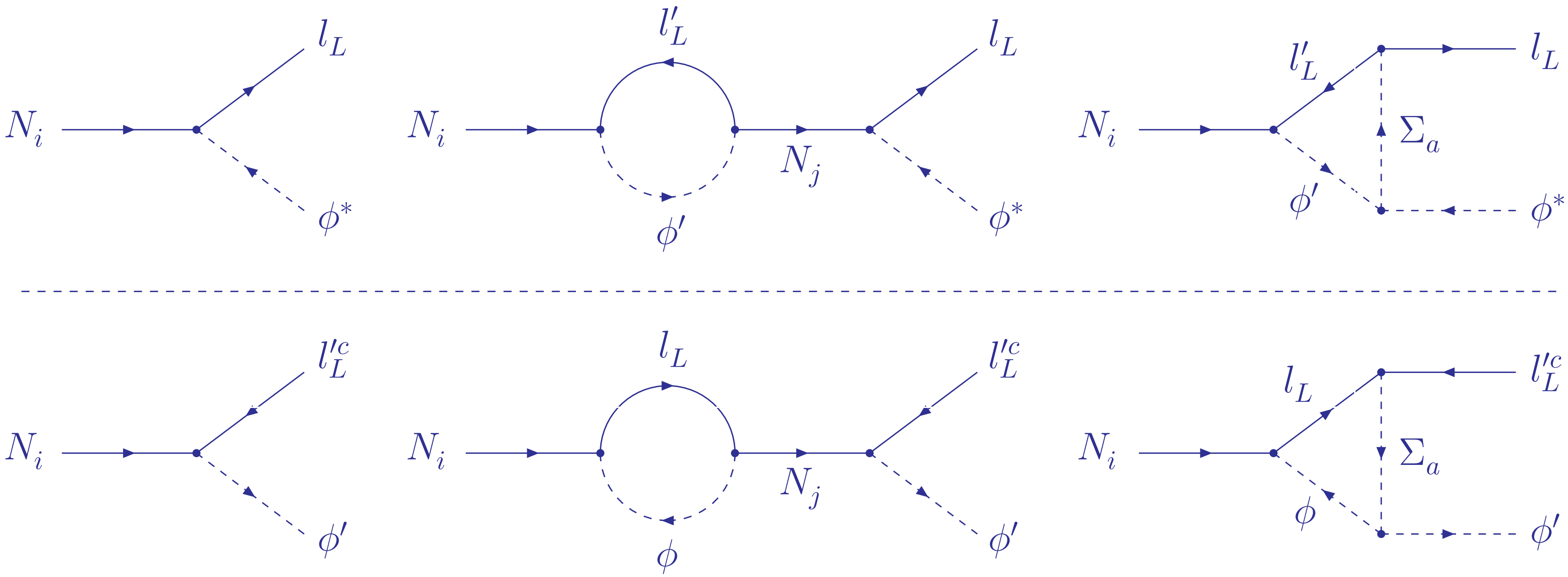
The CP asymmetry ε_{Σ_a} can be calculated at one-loop order,

$$\begin{aligned}\varepsilon_{\Sigma_a} &= \frac{\Gamma_{\Sigma_a \to l_L l'_L} - \Gamma_{\Sigma_a^* \to l_L^c l'^c_L}}{\Gamma_{\Sigma_a}} = \frac{\Gamma_{\Sigma_a^* \to \phi^* \phi'^*} - \Gamma_{\Sigma_a \to \phi \phi'}}{\Gamma_{\Sigma_a}} \\ &= -\frac{1}{4\pi} \sum_{b \neq a} \frac{\mathrm{Im}\left[\mathrm{Tr}\left(f_b^\dagger f_a\right)\right]}{\mathrm{Tr}\left(f_a^\dagger f_a\right) + \frac{\rho_a^2}{M_{\Sigma_a}^2}} \frac{\rho_b \rho_a}{M_{\Sigma_b}^2 - M_{\Sigma_a}^2}\,,\end{aligned}$$

with Γ_{Σ_a} being the decay width,

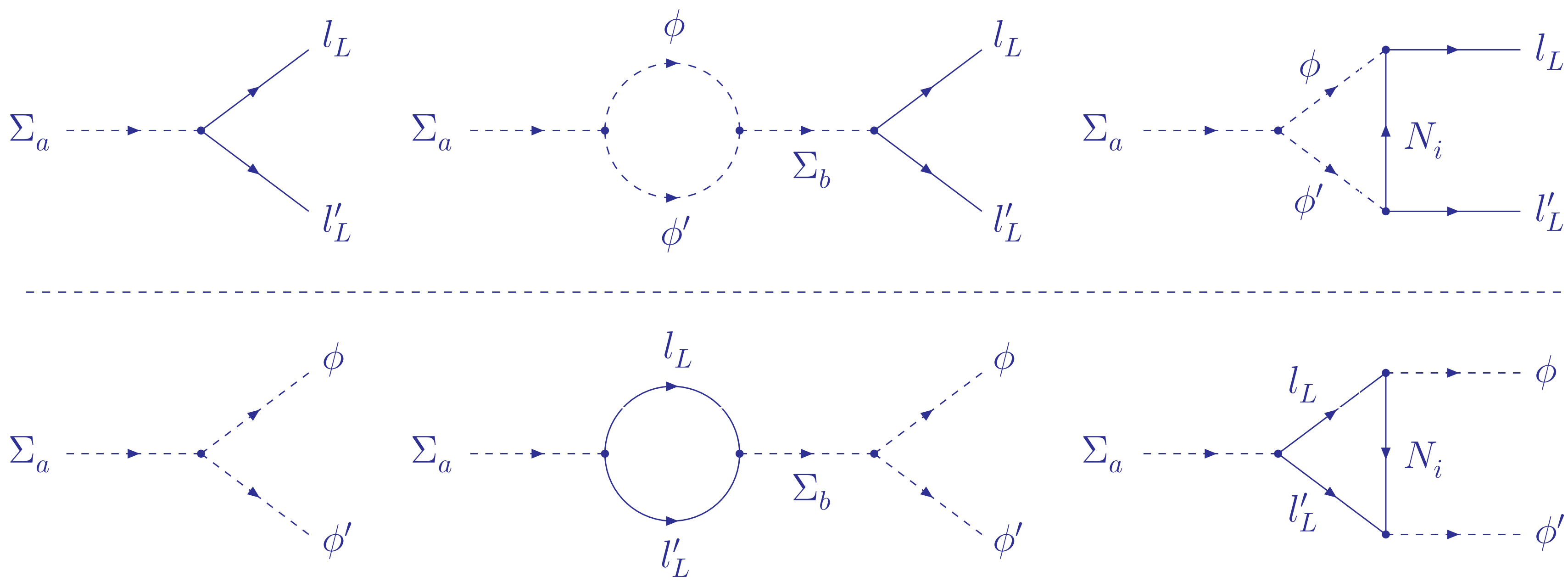
$$\Gamma_{\Sigma_a} = \Gamma_{\Sigma_a \to l_L l'_L} + \Gamma_{\Sigma_a \to \phi\phi'} = \Gamma_{\Sigma_a^* \to l_L^c l'^c_L} + \Gamma_{\Sigma_a^* \to \phi^* \phi'^*} = \frac{1}{16\pi}\left[\mathrm{Tr}\left(f_a^\dagger f_a\right) + \frac{\rho_a^2}{M_{\Sigma_a}^2}\right] M_{\Sigma_a}\,.$$

Note at least two $[SU(2)_L \times SU(2)'_L]$-bidoublet scalars should be introduced to give a nonzero CP asymmetry ε_{Σ_a}.



$$\varepsilon_{N_i} = -\frac{1}{4\pi}\frac{1}{(y_N^\dagger y_N)_{ii}+(y_{N'}^\dagger y_{N'})_{ii}}\left\{\sum_{j\neq i}\mathrm{Im}\left[(y_N^\dagger y_N)_{ij}(y_{N'}^\dagger y_{N'})_{ij}\right]\frac{M_{N_i}M_{N_j}}{M_{N_j}^2-M_{N_i}^2}\right.$$

$$\left.+2\,\mathrm{Im}\left[(y_N^\dagger f_a y_{N'}^*)_{ii}\right]\frac{\rho_a}{M_{N_i}}\left[1-\frac{M_{\Sigma_a}^2}{M_{N_i}^2}\ln\left(1+\frac{M_{N_i}^2}{M_{\Sigma_a}^2}\right)\right]\right\}.$$



$$\varepsilon_{\Sigma_a} = -\frac{1}{4\pi}\frac{1}{\mathrm{Tr}(f_a^\dagger f_a)+\frac{\rho_a^2}{M_{\Sigma_a}^2}}\left\{\sum_{b\neq a}\frac{\mathrm{Im}\left[\mathrm{Tr}\left(f_b^\dagger f_a\right)\right]\rho_b\rho_a}{M_{\Sigma_b}^2-M_{\Sigma_a}^2}\right.$$

$$\left.+\sum_i \mathrm{Im}\left[(y_N^\dagger f_a y_{N'}^*)_{ii}\right]\frac{\rho_a M_{N_i}}{M_{\Sigma_a}^2}\ln\left(1+\frac{M_{\Sigma_a}^2}{M_{N_i}^2}\right)\right\}.$$

Ordinary and dark baryon asymmetries

The ordinary lepton asymmetry η_{l_L} and the dark lepton asymmetry $\eta_{l'_L}$ will decouple each other after they are produced by the decays of the messenger fields.

The dark $SU(2)'_L$ sphaleron processes will partially transfer the dark lepton asymmetry $\eta_{l'_L}$ to a dark baryon asymmetry η'_B,

$$\eta'_B = -C'\eta_{l'_L} \,.$$

while the ordinary $SU(2)_L$ sphaleron processes will partially transfer the ordinary lepton asymmetry η_{l_L} to an ordinary baryon asymmetry η_B,

$$\eta_B = -C\eta_{l_L} \,.$$

Note there will be a lepton-number-violating Majorana mass term of the dark charged leptons after the $U(1)'_{em}$ symmetry is broken. The $U(1)'_{em}$ symmetry breaking and then the lepton number violation should appear at a lower temperature, where the $SU(2)'_L$ sphaleron processes have become very weak, to avoid the washout of the dark baryon asymmetry.

The exact relation between the baryon and lepton numbers can be calculated by means of the analysis of the chemical potentials. The excess of a particle i over its antiparticle can be described by the net number density (e.g. Kolb, Turner, *The Early Universe.*),

$$n_i = n_i^+ - n_i^- = \begin{cases} \frac{1}{6} g_i T^3 \left(\frac{\mu_i}{T}\right) & \text{for fermion}\,, \\ \frac{1}{3} g_i T^3 \left(\frac{\mu_i}{T}\right) & \text{for boson}\,, \end{cases}$$

$$n_B = \frac{1}{3}\sum_{i=1}^{3}(n_{q_i} + n_{u_i} + n_{d_i})\,,$$

$$n_L = \sum_{i=1}^{3}(n_{l_i} + n_{e_i})\,.$$

A $B-L$ asymmetry can be partially converted into a baryon asymmetry (Kuzmin, Rubakov, Shaposhnikov, 85'.).

$$\left.\begin{array}{ll} \text{Yukawa}: & -\mu_q+\mu_u+\mu_\phi=0\,,\ -\mu_q+\mu_d-\mu_\phi=0\,,\ -\mu_l+\mu_e-\mu_\phi=0\,, \\ \text{Hypercharge}: & 3(\mu_q+2\mu_u-\mu_d-\mu_l-\mu_e)-2\mu_\phi=0\,, \\ \text{Sphalerons}: & 3(3\mu_q+\mu_l)=0\,, \end{array}\right\}$$

$$\Rightarrow \left\{\begin{array}{l} n_B=\frac{1}{6}T^2\sum_{i=1}^{3}(2\mu_q+\mu_u+\mu_d)=\frac{2}{3}\mu_l T^2 \\ n_L=\frac{1}{6}T^2\sum_{i=1}^{3}(2\mu_l+\mu_e)=-\frac{17}{14}\mu_l T^2 \end{array}\right\} \Rightarrow n_B=C(n_B-n_L)\ \text{ with }\ C=\frac{28}{79}\,.$$

$$\left.\begin{array}{ll} \text{Yukawa}: & -\mu_{q'}+\mu_{u'}+\mu_{\phi'}=0\,,\ -\mu_{q'}+\mu_{d'}-\mu_{\phi'}=0\,,\ -\mu_{l'}+\mu_{e'}-\mu_{\phi'}=0\,,\ \mu_{\delta'}-2\mu_{e'}=0\,; \\ \text{Hypercharge}: & 3(-\mu_{q'}-2\mu_{u'}+\mu_{d'}+\mu_{l'}+\mu_{e'})+2\mu_{\phi'}+4\mu_{\delta'}=0\,; \\ \text{Sphalerons}: & 3(3\mu_{q'}+\mu_{l'})=0\,; \end{array}\right\}$$

$$\Rightarrow \left\{\begin{array}{l} n'_B=-\frac{1}{6}T^2\sum_{i=1}^{3}(2\mu_{q'}+\mu_{u'}+\mu_{d'})=\frac{2}{3}\mu_l T^2\,, \\ n'_L=-\frac{1}{6}T^2\left[\sum_{i=1}^{3}(2\mu_{l'}+\mu_{e'})+4\mu_{\delta'}\right]=-\frac{3}{2}\mu_l T^2 \end{array}\right\} \Rightarrow n'_B=C'(n'_B-n'_L)\ \text{ with }\ C'=\frac{4}{13}\,.$$

Predictive dark matter mass

The $U(1)_Y$ and $U(1)'_Y$ gauge fields can have a kinetic mixing (Foot, He, 91'.),

$$\mathcal{L} \supset -\frac{1}{4}B_{\mu\nu}B^{\mu\nu} - \frac{1}{4}B'_{\mu\nu}B'^{\mu\nu} - \frac{\epsilon}{2}B_{\mu\nu}B'^{\mu\nu} = -\frac{1}{4}\tilde{B}_{\mu\nu}\tilde{B}^{\mu\nu} - \frac{1}{4}\tilde{B}'_{\mu\nu}\tilde{B}'^{\mu\nu}$$

$$\text{with} \quad \tilde{B}_\mu = B_\mu + \epsilon B'_\mu\,, \quad \tilde{B}'_\mu = \sqrt{1-\epsilon^2}B'_\mu\,.$$

We then can define the orthogonal fields,

$$A_\mu = W^3_\mu s_W + \tilde{B}_\mu c_W\,, \quad Z_\mu = W^3_\mu c_W - \tilde{B}_\mu s_W\,, \quad A'_\mu = W'^3_\mu s_W + \tilde{B}'_\mu c_W\,, \quad Z'_\mu = W'^3_\mu c_W - \tilde{B}'_\mu s_W\,,$$

among which the field A is exactly massless and is the ordinary photon, while the dark photon A' is massive and mixes with the Z and Z' bosons.

For $\epsilon \ll 1$ and $\langle\Sigma\rangle$, $\langle\delta'\rangle \ll \langle\phi\rangle \ll \langle\phi'\rangle$, the mass eigenstates of the A', Z and Z' bosons can approximate to

$$\hat{A}' \simeq A' - \epsilon s_W c_W Z\,,\ \ \hat{Z} \simeq Z + \epsilon s_W c_W A\,,\ \ \hat{Z}' \simeq Z'\,.$$

The dark photon can couple to the ordinary fermions besides the dark fermions,

$$\mathcal{L} \supset g_2 s_W \hat{A}'_\mu \left[\left(\frac{1}{3}\bar{d}'\gamma^\mu d' - \frac{2}{3}\bar{u}'\gamma^\mu u' + \bar{e}'\gamma^\mu e'\right) + \epsilon c_W^2 \left(-\frac{1}{3}\bar{d}\gamma^\mu d + \frac{2}{3}\bar{u}\gamma^\mu u - \bar{e}\gamma^\mu e\right)\right].$$

As long as the dark photon is heavy enough,

$$m_{\hat{A}'}^2 \simeq m_{A'}^2 = 8 g_2^2 s_W^2 \langle\delta'\rangle^2 \simeq (400\,\text{MeV})^2 \left(\frac{\langle\delta'\rangle}{470\,\text{MeV}}\right)^2,$$

it can efficiently decay into the ordinary fermions,

$$\begin{aligned}\Gamma_{A'} &= \Gamma_{A'\to e^- e^+} + \Gamma_{A'\to u\bar{u}} + \Gamma_{A'\to d\bar{d}} + \Gamma_{A'\to s\bar{s}} \simeq \frac{15}{16}\epsilon^2 m_{A'} \alpha c_W^4 \\ &\simeq \frac{1}{5\times 10^{-10}\,\text{sec}} \left(\frac{\epsilon}{10^{-6}}\right)^2 \left(\frac{m_{A'}}{400\,\text{MeV}}\right).\end{aligned}$$

Since the dark charged leptons are the quasi-Dirac fermions, their lepton asymmetries cannot survive (Buckley, Profumo, 11'.). The lightest dark charged lepton (the dark electron) thus should have a thermally produced relic density, which depends on its pair annihilation into the dark photon,

$$\langle \sigma_{e'^+ e'^- \to A' A'} v_{\rm rel} \rangle \simeq \frac{\pi \alpha'^2}{m_{e'}^2} \simeq 6.3 \times 10^5 \,\mathrm{pb} \left(\frac{\alpha'}{\alpha} \right)^2 \left(\frac{322 \,\mathrm{MeV}}{m_{e'}} \right)^2 .$$

If the dark electron is at the GeV scale, it will have a frozen temperature far below its mass and give a negligible contribution to the dark matter relic density.

We now have known

$$\eta'_B = -\frac{C'}{C}\eta_B \text{ for } \eta_{l_L} = -\eta_{l'_L}\,.$$

If the dark matter relic density is dominated by the lightest dark nucleon N', the cosmological observations will imply

$$\Omega_B h^2 : \Omega_{DM} h^2 = m_p \eta_B : m_{N'}(-\eta'_B) \Rightarrow m_{N'} = \frac{C}{C'}\frac{\Omega_{DM} h^2}{\Omega_B h^2} m_p\,,$$

from which the dark matter mass can be determined by

$$m_{N'} \simeq \begin{cases} 5.67\,\mathrm{GeV} & \text{for} \quad \frac{\Omega_{DM}h^2}{\Omega_B h^2} = \frac{0.1172}{0.02233}\,, \\ 5.88\,\mathrm{GeV} & \text{for} \quad \frac{\Omega_{DM}h^2}{\Omega_B h^2} = \frac{0.1199}{0.02205}\,, \\ 6.09\,\mathrm{GeV} & \text{for} \quad \frac{\Omega_{DM}h^2}{\Omega_B h^2} = \frac{0.1226}{0.02177}\,. \end{cases}$$

Effective neutrino number

The BBN stringently restricts the existence of the new relativistic degrees of freedom. The constraint on the new degrees of freedom is conventionally quoted as $\Delta N_\nu = N_\nu - 3.046$, the effective number of additional light neutrinos (e.g. Planck collaboration, 1304.5076.),

$$N_\nu = 3.30 \pm 0.27 \Rightarrow \Delta N_\nu = 0.25 \pm 0.27\,.$$

So, the gauge interactions of the dark neutrinos should decouple before the BBN. We can estimate

$$\Gamma \sim \left(\frac{g_2'/\cos\theta_W'}{g_2/\cos\theta_W}\right)^4 \left(\frac{\langle\phi\rangle}{\langle\phi'\rangle}\right)^4 G_F^2 T^5 = H(T) = \left[\frac{8\pi^3 g_*(T)}{90}\right]^{\frac{1}{2}} \frac{T^2}{M_{\mathrm{Pl}}}\Big|_{T=T_D}$$

$$\Rightarrow T_D \simeq 12\,\mathrm{GeV}\left(\frac{g_*}{100}\right)^{\frac{1}{6}} \left(\frac{\langle\phi'\rangle}{630\,\langle\phi\rangle}\right)^{\frac{4}{3}} \left(\frac{g_2/\cos\theta_W}{g_2'/\cos\theta_W'}\right)^{\frac{4}{3}}\,.$$

The temperature of the dark neutrinos at the BBN epoch $T_{\mathrm{BBN}} \sim 1\,\mathrm{MeV}$ then should be about

$$\left(\frac{T_{\nu'}}{T_{\mathrm{BBN}}}\right)^4 = \left[\frac{g_*(T_{\mathrm{BBN}})}{g_*(T_D)}\right]^{\frac{4}{3}} \sim \left(\frac{10.75}{80+44.5}\right)^{\frac{4}{3}} \simeq 0.038 \Rightarrow \Delta N_\nu = \frac{4}{7}\left(\frac{T_{\nu'}}{T_{\mathrm{BBN}}}\right)^4 < 0.02\,.$$

Discrete mirror symmetry

We can impose a discrete mirror symmetry (Lee, Yang, 56'; Kobzarev, Okun, Pomeranchuk, 66'; Pavsic, 74'; Blinnikov, Khlopov, 82; Glashow, 86'; Foot, Lew, Volkas, 91'.) under which the fields transform as

$$G'^a_\mu \leftrightarrow G^a_\mu\,,\ W'^a_\mu \leftrightarrow W^a_\mu\,,\ B'_\mu \leftrightarrow B_\mu\,,\ q'_L \leftrightarrow q_L\,,\ d'_R \leftrightarrow d_R\,,\ u'_R \leftrightarrow u_R\,,\ l'_L \leftrightarrow l_L\,,\ e'_R \leftrightarrow e_R\,,\ \phi' \leftrightarrow \phi\,,$$
$$\delta' \leftrightarrow \delta\,,\ N'_R \leftrightarrow N_R\,,\ \Sigma_a \leftrightarrow \Sigma^T_a\,,$$

to simplify the dimensionless parameters,

$$g'_3 = g_3\,,\ g'_2 = g_2\,,\ g'_1 = g_1\,,\ y_{d'} = y_d\,,\ y_{u'} = y_u\,,\ y_{e'} = y_e\,,\ y_{\delta'} = y_\delta\,,\ y_{N'} = y_N\,,\ f_a = f^T_a\,.$$

Benefited from the discrete mirror symmetry, the neutrino mass matrix induced by the type-I, type-II or type-I+II Dirac seesaw will have a symmetric structure:

Type-I Dirac seesaw: $m_\nu = -y_N \frac{\langle\phi\rangle\langle\phi'\rangle}{M_N} y^T_N$ with $M_N = \mathrm{diag}\{M_{N_1}, M_{N_2}, ...\}\,,$

Type-II Dirac seesaw: $m_\nu = -\sum_a f_a \frac{\rho_a\langle\phi\rangle\langle\phi'\rangle}{M^2_{\Sigma_a}}$ with $f_a = f^T_a\,,$

Type-I+II Dirac seesaw: $m_\nu = -y_N \frac{\langle\phi\rangle\langle\phi'\rangle}{M_N} y^T_N - \sum_a f_a \frac{\rho_a\langle\phi\rangle\langle\phi'\rangle}{M^2_{\Sigma_a}}$ with $M_N = \mathrm{diag}\{M_{N_1}, ...\}\,,\ f_a = f^T_a\,.$

As the dark vacuum expectation value $\langle\phi'\rangle$ is determined by the dark matter mass, our Dirac seesaw will not contain new parameters compared with the traditional type-I, type-II or type-I+II Majorana seesaw.

Another implication of the mirror symmetry is that the beta functions of the QCD in the ordinary and dark sectors, i.e.

$$\alpha_s^{-1}(\mu) = \frac{11 - \frac{2}{3}n_f}{2\pi} \ln\left(\frac{\mu}{\Lambda}\right), \quad \alpha_{s'}^{-1}(\mu) = \frac{11 - \frac{2}{3}n_{f'}}{2\pi} \ln\left(\frac{\mu}{\Lambda'}\right),$$

can yield a relation between the electroweak and hadronic scales in the dark sector. Here $n_f(n_{f'})$ counts the number of the ordinary(dark) quarks involved at a given scale μ. By matching $\alpha_s(\mu)$ at the scale $\mu = m_t$ with $n_f = 6$ and $n_f = 5$, at the scale $\mu = m_b$ with $n_f = 5$ and $n_f = 4$, at the scale $\mu = m_c$ with $n_f = 4$ and $n_f = 3$, respectively, we can deduce

$$\Lambda_{(5)} = m_t^{\frac{2}{23}} \Lambda_{(6)}^{\frac{21}{23}}, \quad \Lambda_{(4)} = m_b^{\frac{2}{25}} \Lambda_{(5)}^{\frac{23}{25}}, \quad \Lambda_{(3)} = m_c^{\frac{2}{27}} \Lambda_{(4)}^{\frac{25}{27}} \Rightarrow \Lambda_{\textrm{QCD}} = \Lambda_{(3)} = (m_c m_b m_t)^{\frac{2}{27}} \Lambda_{(6)}^{\frac{21}{27}}.$$

Similarly, we can have

$$\Lambda'_{(5)} = m_{t'}^{\frac{2}{23}} \Lambda'^{\frac{21}{23}}_{(6)}, \quad \Lambda'_{(4)} = m_{b'}^{\frac{2}{25}} \Lambda'^{\frac{23}{25}}_{(5)}, \quad \Lambda'_{(3)} = m_{c'}^{\frac{2}{27}} \Lambda'^{\frac{25}{27}}_{(4)}, \quad \Lambda'_{(2)} = m_{s'}^{\frac{2}{29}} \Lambda'^{\frac{27}{29}}_{(4)}, \quad \Lambda'_{(1)} = m_{d'}^{\frac{2}{31}} \Lambda'^{\frac{29}{31}}_{(2)}, \quad \Lambda'_{(0)} = m_{u'}^{\frac{2}{33}} \Lambda'^{\frac{31}{33}}_{(1)},$$

$$\Rightarrow \Lambda_{\textrm{QCD}'} = \Lambda'_{(0)} = (m_{u'} m_{d'} m_{s'} m_{c'} m_{b'} m_{t'})^{\frac{2}{33}} \Lambda'^{\frac{21}{33}}_{(6)} \ \textrm{for} \ \Lambda_{\textrm{QCD}'} < m_{u'}.$$

Due to

$$\alpha_s(\mu \gg m_t) = \alpha'_s(\mu \gg m_{t'}) \Rightarrow \Lambda_{(6)} = \Lambda'_{(6)}\,, \quad \frac{\langle\phi'\rangle}{\langle\phi\rangle} = \frac{m_{u'}}{m_u} = \frac{m_{d'}}{m_d} = \frac{m_{s'}}{m_s} = \frac{m_{c'}}{m_c} = \frac{m_{b'}}{m_b} = \frac{m_{t'}}{m_t}\,,$$

the dark hadronic scale $\Lambda_{\mathrm{QCD}'}$ should arrive at

$$\Lambda_{\mathrm{QCD}'} = \left(\frac{\langle\phi'\rangle}{\langle\phi\rangle}\right)^{\frac{4}{11}} (m_u m_d m_s)^{\frac{2}{33}} \Lambda_{\mathrm{QCD}}^{\frac{9}{11}} \quad \text{for} \quad \Lambda_{\mathrm{QCD}'} < m_{u'}\,.$$

So, by inputting $\langle\phi'\rangle = 630\,\langle\phi\rangle$, we can determine

$$\Lambda_{\mathrm{QCD}'} = 1.2\,\mathrm{GeV} \;\; \text{for} \;\; \Lambda_{\mathrm{QCD}} = 200\,\mathrm{MeV}\,, \;\; m_u = 2.3\,\mathrm{MeV}\,, \;\; m_d = 4.8\,\mathrm{MeV}\,, \;\; m_s = 95\,\mathrm{MeV}\,.$$

Consequently, the dark proton and neutron masses can be given by

$$m_{p'} \simeq 2m_{u'} + m_{d'} \simeq 6.0\,\mathrm{GeV}\,, \; m_{n'} \simeq m_{u'} + 2m_{d'} \simeq 7.5\,\mathrm{GeV} \;\; \text{with} \;\; m_{u'} = 1.5\,\mathrm{GeV}\,, \;\; m_{d'} = 3.0\,\mathrm{GeV}\,.$$

Dark matter detection

The dark proton as the dark matter particle can scatter off the ordinary nucleons,

$$\sigma_{p'N\to p'N}(Z,A) \simeq \epsilon^2\pi\alpha^2 c_W^4 \frac{\mu_r^2}{m_{A'}^4}\left(\frac{Z}{A}\right)^2$$

$$\simeq 4\times 10^{-45}\,\mathrm{cm}^2\left(\frac{\epsilon}{10^{-7}}\right)^2\left(\frac{\mu_r}{0.81\,\mathrm{GeV}}\right)^2\left(\frac{400\,\mathrm{MeV}}{m_{A'}}\right)^4\left(\frac{Z}{A}\right)^2,$$

where μ_r is the reduced mass,

$$\mu_r = \frac{m_{p'}m_p}{m_{p'}+m_p} \simeq 0.81\,\mathrm{GeV} \quad \text{for} \quad m_{p'} \simeq 6\,\mathrm{GeV}\,.$$

An $SO(10) \times SO(10)'$ Scenario

(PHG, 1410.5759.)

$$q^c_L(\mathbf{3},\mathbf{2},\mathbf{1},-\tfrac{1}{3}) \oplus q_R(\mathbf{3},\mathbf{1},\mathbf{2},+\tfrac{1}{3}) \oplus l^c_L(\mathbf{1},\mathbf{2},\mathbf{1},+1) \oplus l_R(\mathbf{1},\mathbf{1},\mathbf{2},-1) = \mathbf{16}_F\,,$$

$$\chi^*_L(\mathbf{1},\mathbf{2},\mathbf{1},+1) \oplus \chi_R(\mathbf{1},\mathbf{1},\mathbf{2},-1) \in \mathbf{16}_H\,,$$

$$\Delta^*_L(\mathbf{1},\mathbf{3},\mathbf{1},-2) \oplus \Delta_R(\mathbf{1},\mathbf{1},\mathbf{3},+2) \oplus \Omega^*_L(\mathbf{3},\mathbf{3},\mathbf{1},-\tfrac{2}{3}) \oplus \Omega_R(\mathbf{3},\mathbf{1},\mathbf{3},+\tfrac{2}{3}) \in \mathbf{126}_H\,,$$

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

$$q'^c_R(\mathbf{3},\mathbf{2},\mathbf{1},-\tfrac{1}{3}) \oplus q'_L(\mathbf{3},\mathbf{1},\mathbf{2},+\tfrac{1}{3}) \oplus l'^c_R(\mathbf{1},\mathbf{2},\mathbf{1},+1) \oplus l'_L(\mathbf{1},\mathbf{1},\mathbf{2},-1) = \mathbf{16}'_F\,,$$

$$\chi'^*_R(\mathbf{1},\mathbf{2},\mathbf{1},+1) \oplus \chi'_L(\mathbf{1},\mathbf{1},\mathbf{2},-1) \in \mathbf{16}'_H\,,$$

$$\Delta'^*_R(\mathbf{1},\mathbf{3},\mathbf{1},-2) \oplus \Delta'_L(\mathbf{1},\mathbf{1},\mathbf{3},+2) \oplus \Omega'^*_R(\mathbf{3},\mathbf{3},\mathbf{1},-\tfrac{2}{3}) \oplus \Omega'_L(\mathbf{3},\mathbf{1},\mathbf{3},+\tfrac{2}{3}) \in \mathbf{126}'_H\,,$$

$$\Phi'(\mathbf{1},\mathbf{2},\mathbf{2},0) \in \mathbf{10}'_H \text{ and/or others}\,.$$

$$(\mathbf{16} \times \overline{\mathbf{16}'})_H = \Sigma_{l_L l'_R}(\mathbf{1},\mathbf{2},\mathbf{1},-1)(\mathbf{1},\mathbf{2},\mathbf{1},+1) \oplus \Sigma_{l_R l'_L}(\mathbf{1},\mathbf{1},\mathbf{2},+1)(\mathbf{1},\mathbf{1},\mathbf{2},-1) \oplus \ldots\,.$$

Symmetry breaking

$$
\begin{aligned}
& SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \\
& \xrightarrow{\langle \chi_R \rangle = \frac{1}{\sqrt{2}}(v_R,\, 0)^T} SU(3)_c \times SU(2)_L \times U(1)_Y \\
& \xrightarrow{\langle \Phi \rangle = \frac{1}{\sqrt{2}}\mathrm{diag}\{v_1,\, v_2\}} SU(3)_c \times U(1)_{em}\,.
\end{aligned}
$$

$$
\begin{aligned}
& SU(3)'_c \times SU(2)'_R \times SU(2)'_L \times U(1)'_{B-L} \\
& \xrightarrow{\langle \chi'_L \rangle = \frac{1}{\sqrt{2}}(v'_L,\, 0)^T} SU(3)'_c \times SU(2)'_R \times U(1)'_Y \\
& \xrightarrow{\langle \Phi' \rangle = \frac{1}{\sqrt{2}}\mathrm{diag}\{v'_1,\, v'_2\}} SU(3)'_c \times U(1)'_{em} \\
& \xrightarrow{\langle \Delta'_R \rangle = \frac{1}{\sqrt{2}}\left[\begin{array}{cc} 0 & 0 \\ v'_{em} & 0 \end{array}\right]} SU(3)'_c\,.
\end{aligned}
$$

There is a $U(1)_G$ global symmetry under which the scalars (χ_L^*, χ_R) and $(\chi_R'^*, \chi_L')$ carry a same charge.

$$\text{Forbid} \quad V \supset \rho_\Phi \chi_L^\dagger \Phi \chi_R + \tilde{\rho}_\Phi \chi_L^\dagger \widetilde{\Phi} \chi_R + \rho_{\Phi'} \chi_R'^\dagger \Phi' \chi_L' + \tilde{\rho}_{\Phi'} \chi_R'^\dagger \widetilde{\Phi}' \chi_L'$$
$$+\rho_\Delta (\chi_L^T i\tau_2 \Delta_L \chi_L + \chi_R^\dagger i\tau_2 \Delta_R^* \chi_R^*)$$
$$+\rho_{\Delta'} (\chi_R'^T i\tau_2 \Delta_R' \chi_R' + \chi_L'^\dagger i\tau_2 \Delta_L'^* \chi_L'^*) + \text{H.c.}\,,$$

$$\text{Allow} \quad V \supset \rho_\Sigma (\chi_L^T \Sigma_{LR'}^* \chi_R'^* + \chi_R^\dagger \Sigma_{RL'} \chi_L') + \text{H.c.}\,.$$

The neutral components of the scalars χ_L, $\Delta_{L,R}$, χ_R', $\Delta_{R,L}'$ will not acquire any induced vacuum expectation values.

$$\langle \Sigma_{RL'} \rangle \leq \langle \chi_{R,L'} \rangle \neq 0\,, \quad \langle \Sigma_{LR'} \rangle = 0\,.$$

Dirac neutrinos, lepton asymmetry and dark matter mass

At the left-right level,

$$
\begin{aligned}
\mathcal{L} \supset\; & -y_q\bar{q}_L\Phi q_R - \tilde{y}_q\bar{q}_L\widetilde{\Phi}q_R - y_l\bar{l}_L\Phi l_R - \tilde{y}_l\bar{l}_L\widetilde{\Phi}l_R \\
& -y_{q'}\bar{q}'_R\Phi' q'_L - \tilde{y}_{q'}\bar{q}'_R\widetilde{\Phi}' q'_L - y_{l'}\bar{l}'_R\Phi' l'_L - \tilde{y}_{l'}\bar{l}'_R\widetilde{\Phi}' l'_L \\
& -\frac{1}{2}f_\Delta(\bar{l}_L i\tau_2\Delta_L^* l_L^c + \bar{l}_R^c i\tau_2\Delta_R l_R) - \frac{1}{2}f_{\Delta'}(\bar{l}'_R i\tau_2\Delta_R'^* l_R'^c + \bar{l}_L'^c i\tau_2\Delta'_L l'_L) \\
& -f_\Sigma(\bar{l}_L^c\Sigma_{l_L l'_R} l_R'^c + \bar{l}_R\Sigma_{l_R l'_L} l'_L) + \mathrm{H.c.}\,.
\end{aligned}
$$

After the left-right symmetry breaking,

$$\begin{aligned}\mathcal{L} \supset & -y_u\bar{q}_L\phi u_R - y_d\bar{q}_L\tilde{\phi}d_R - y_\nu\bar{l}_L\phi\nu_R - y_e\bar{l}_L\tilde{\phi}e_R \\ & -y_{u'}\bar{q}'_R\phi' u'_L - y_{d'}\bar{q}'_R\widetilde{\phi'}d'_L - y_{\nu'}\bar{l}'_R\phi'\nu'_L - y_{e'}\bar{l}'_R\widetilde{\phi'}e'_L \\ & -\frac{1}{2}f_\Delta\bar{l}_L i\tau_2\Delta_L^* l_L^c - \frac{1}{2}f_{\Delta'}\bar{l}'_R i\tau_2\Delta'^*_R l'^c_R - M_N\bar{\nu}_R\nu'_L + \text{H.c.}\end{aligned}$$

$$\text{with}\left\{\begin{array}{l} y_u = \frac{v_1 y_q + v_2\tilde{y}_q}{\sqrt{v_1^2+v_2^2}}\,,\; y_d = \frac{v_2 y_q + v_1\tilde{y}_q}{\sqrt{v_1^2+v_2^2}}\,, \\ y_{u'} = \frac{v'_1 y_{q'} + v'_2\tilde{y}_{q'}}{\sqrt{v'^2_1+v'^2_2}}\,,\; y_{d'} = \frac{v'_2 y_{q'} + v'_1\tilde{y}_{q'}}{\sqrt{v'^2_1+v'^2_2}}\,, \\ M_N = \frac{1}{\sqrt{2}}f_\Sigma v'_L\,. \end{array}\right.$$

The Higgs scalars ϕ and ϕ' are responsible for spontaneously breaking the ordinary and dark electroweak symmetries.

After the electromagnetic symmetry breaking,

$$\mathcal{L} \supset -m_u\bar{u}_L u_R - m_d\bar{d}_L d_R - m_e\bar{e}_L e_R - m_{u'}\bar{u}'_R u'_L$$
$$-m_{d'}\bar{d}'_R d'_L - m_{e'}\bar{e}'_R e'_L - \frac{1}{2}\bar{m}_{e'}\bar{e}'_R e'^c_R$$
$$-\begin{bmatrix} \bar{\nu}_L & \bar{\nu}'_L \end{bmatrix}\begin{bmatrix} 0 & m_{LR} \\ m^\dagger_{R'L'} & M^\dagger_N \end{bmatrix}\begin{bmatrix} \nu'_R \\ \nu_R \end{bmatrix} + \text{H.c. with}$$
$$m_f = \frac{1}{\sqrt{2}}y_f v\,,\; m_{f'} = \frac{1}{\sqrt{2}}y_{f'}v'\,,\; \bar{m}_{e'} = \frac{1}{\sqrt{2}}f_{\Delta'}v'_{em}\,,$$
$$m_{LR} = \frac{1}{\sqrt{2}}y_\nu v\,,\; m_{R'L'} = \frac{1}{\sqrt{2}}y_{\nu'}v'\,.$$

$$\mathcal{L} \supset -m_\nu\bar{\nu}_L\nu'_R - M_N\bar{\nu}_R\nu'_L + \text{H.c. with } m_\nu = -m_{LR}\frac{1}{M^\dagger_N}m^\dagger_{R'L'}\,.$$

Mass eigenstates:

$$N_i = \nu_{Ri} + \nu'_{Li} \ \text{ with } \ M_N = \text{diag}\{M_{N_1}, M_{N_2}, M_{N_3}\} \, .$$

Decay width:

$$\begin{aligned}\Gamma_{N_i} &= \Gamma(N_i \to l_L \phi^*) + \Gamma(N_i \to l'_R \phi'^*) = \Gamma(N_i \to l^c_L \phi) + \Gamma(N_i \to l'^c_R \phi') \\ &= \frac{1}{16\pi}[(y^\dagger_\nu y_\nu)_{ii} + (y^\dagger_{\nu'} y_{\nu'})_{ii}] M_{N_i} \, .\end{aligned}$$

CP asymmetry:

$$\begin{aligned}\varepsilon_{N_i} &= \frac{\Gamma(N_i \to l_L \phi^*) - \Gamma(N_i \to l^c_L \phi)}{\Gamma_{N_i}} = \frac{\Gamma(N_i \to l'^c_R \phi') - \Gamma(N_i \to l'_R \phi'^*)}{\Gamma_{N_i}} \\ &= \frac{1}{4\pi} \sum_{j \neq i} \frac{\text{Im}[(y^\dagger_\nu y_\nu)_{ij} (y^\dagger_{\nu'} y_{\nu'})_{ji}]}{(y^\dagger_\nu y_\nu)_{ii} + (y^\dagger_{\nu'} y_{\nu'})_{ii}} \frac{M_{N_i} M_{N_j}}{M^2_{N_i} - M^2_{N_j}} \, .\end{aligned}$$

Lepton asymmetry:

$$\bar{\eta}_L = -\bar{\eta}'_L \propto \varepsilon_{N_i} \,.$$

Baryon asymmetry:

$$\eta_B = C\eta_{B-L} = -C\bar{\eta}_L \ \text{with} \ C = \frac{28}{79} \,,$$

$$\eta'_B = C\eta'_{B-L} = -C'\bar{\eta}'_L \ \text{with} \ C' = \frac{28}{229} \,.$$

Dark matter mass:

$$\Omega_{\mathrm{B}} h^2 : \Omega_{\mathrm{DM}} h^2 = \eta_B m_p : (-\eta'_B) m_{N'} \Rightarrow$$

$$m_{N'} = \frac{C'}{C} \frac{\Omega_{\mathrm{DM}} h^2}{\Omega_{\mathrm{B}} h^2} m_p = 14.79 \, \mathrm{GeV} \left(\frac{\Omega_{\mathrm{DM}} h^2 / 0.1199}{\Omega_{\mathrm{B}} h^2 / 0.02205} \right) \,.$$

$U(1)_{em} \times U(1)'_{em}$ kinetic mixing and dark matter detection

The $(\mathbf{16} \times \overline{\mathbf{16}'})_H$ scalar can contribute to the $U(1)_{em} \times U(1)'_{em}$ kinetic mixing at one-loop level,

$$\mathcal{L} \supset -\frac{\epsilon}{2} A'_{\mu\nu} A^{\mu\nu} \quad \text{with} \quad \epsilon = \frac{\sqrt{\alpha\alpha'}}{12\pi} \sum_{Q,Q'} QQ' C_Q C_{Q'} \ln\left[\frac{M^2_{(Q,Q')}}{\mu^2}\right] .$$

$$Q, Q' = \pm 1, \pm\frac{1}{3}, \pm\frac{2}{3} .$$

$$C_{Q,Q'} = \begin{cases} 1 & \text{for} \quad Q, Q' = \pm 1 \,, \\ 3 & \text{for} \quad Q, Q' = \pm\frac{1}{3}, \pm\frac{2}{3} . \end{cases}$$

Before the $SO(10) \times SO(10)'$ symmetry breaking, the contributions from the scalars $\sigma(Q, Q') \in (\mathbf{16} \times \overline{\mathbf{16}'})_H$ will be canceled exactly.

At the left-right level,

$$V \supset \lambda \left(\chi_R^\dagger \widetilde{\Sigma}_{f_R f'_L} \widetilde{\Sigma}^\dagger_{f_R f'_L} \chi_R + \chi_L'^\dagger \widetilde{\Sigma}^\dagger_{f_R f'_L} \widetilde{\Sigma}_{f_R f'_L} \chi'_L + \chi_R^\dagger \Sigma_{f_R f'_L} \Sigma^\dagger_{f_R f'_L} \chi_R + \chi_L'^\dagger \Sigma^\dagger_{f_R f'_L} \Sigma_{f_R f'_L} \chi'_L \right)$$

$$+ M_1^2 \mathrm{Tr} \left(\Sigma^\dagger_{l_{R,L} l'_{L,R}} \Sigma_{l_{R,L} l'_{L,R}} \right) + M_3^2 \sum_{f_{R,L} \neq l_{R,L}}^{f'_{L,R} \neq l'_{L,R}} \mathrm{Tr} \left(\Sigma^\dagger_{f_{R,L} f'_{L,R}} \Sigma_{f_{R,L} f'_{L,R}} \right) ,$$

the $U(1)_{em} \times U(1)'_{em}$ kinetic mixing can arrive at a nonzero value,

$$\epsilon = \frac{\sqrt{\alpha \alpha'}}{12\pi} \left[\ln \frac{\left(1 + \frac{1}{2}\lambda v_L'^2 / M_1^2\right)\left(1 + \frac{1}{2}\lambda v_R^2 / M_1^2\right)}{1 + \frac{1}{2}\lambda (v_L'^2 + v_R^2)/M_1^2} + \ln \frac{\left(1 + \frac{1}{2}\lambda v_L'^2 / M_3^2\right)\left(1 + \frac{1}{2}\lambda v_R^2 / M_3^2\right)}{1 + \frac{1}{2}\lambda (v_L'^2 + v_R^2)/M_3^2} \right]$$

$$\simeq \frac{\sqrt{\alpha \alpha'}}{48\pi} \frac{\lambda^2 v_L'^2 v_R^2}{M_1^4} \quad \text{for} \quad M_3^2 \gg M_1^2 \gg \lambda v_L'^2 , \lambda v_R^2$$

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

The $U(1)_{em} \times U(1)'_{em}$ kinetic mixing can mediate a testable dark matter scattering if the dark proton is the dark matter particle.

Furthermore, when the dark leptoquark scalars $\Omega'_{R,L}$ is allowed much lighter than the ordinary ones $\Omega_{L,R}$, 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, in the dark matter decay chains,

$$p' \to \pi'^0 e'^+ \ (\text{or } n' \to \pi'^0 \bar{\nu}'_R)\,, \ \ \pi'^0 \to \gamma'\gamma'\,, \ \ \gamma' \to e^+ e^-\,,$$

the induced positrons/electrons can have a distinct energy,

$$\begin{aligned} E_{e^\pm} &\simeq \frac{m_{N'}^2 + m_{\pi'^0}^2}{8 m_{N'}} \ (\text{for } \ m_{e'} \ll m_{\pi'^0} < m_{N'}) \\ &\simeq 1.9\,\text{GeV} \left(\frac{m_{N'}}{15\,\text{GeV}} \right) \left[1 + \left(\frac{m_{\pi'^0}}{m_{N'}} \right)^2 \right] \in (1.9\text{GeV},\ 3.8\text{GeV})\,. \end{aligned}$$

A Left-Right Symmetric Scenario

In QCD, the axial baryon number current has an anomaly (Callan, Dashen, Gross, 76'; Jackiw, Rebbi, 76'.),

$$J_{5\mu} = \sum_{a=1}^{N_q} \bar{q}_a \gamma_\mu \gamma_5 q_a \,, \quad \partial^\mu J_{5\mu} = N_q \frac{g_3^2}{32\pi^2} \epsilon_{\mu\nu\rho\sigma} \mathrm{Tr}(G^{\mu\nu} G^{\rho\sigma}) \,.$$

This is equivalent to add a term to the QCD Lagrangian,

$$\mathcal{L}_{\rm QCD} \supset -\bar{q}_L m_q q_R - \bar{q}_R m_q^\dagger q_L + \theta_{\rm QCD} \frac{g_3^2}{32\pi^2} \epsilon_{\mu\nu\rho\sigma} \mathrm{Tr}(G^{\mu\nu} G^{\rho\sigma}) \quad \text{with} \quad (m_q)_{ab} = m_a \delta_{ab} e^{i\theta_{\rm QFD}/N_q} \,.$$

It should be noted that only the sum of $\theta_{\rm QCD}$ and $\theta_{\rm QFD} = \mathrm{ArgDet}(m_q)$, i.e.

$$\bar{\theta} = \theta_{\rm QCD} + \theta_{\rm QFD} \,,$$

is physical. Actually, we can remove either $\theta_{\rm QCD}$ or $\theta_{\rm QFD}$ by doing a chiral transformation,

$$\theta_{\rm QCD} \to \theta_{\rm QCD} - \theta_{\rm QCD} = 0 \,, \ \ \theta_{\rm QFD} \to \theta_{\rm QCD} + \theta_{\rm QFD} = \bar{\theta} \quad \text{for} \quad q_a \to e^{i\theta_{\rm QCD}/(2N_q)\gamma_5} q_a \,,$$

$$\theta_{\rm QCD} \to \theta_{\rm QCD} + \theta_{\rm QFD} = \bar{\theta} \,, \ \ \theta_{\rm QFD} \to \theta_{\rm QFD} - \theta_{\rm QFD} = 0 \quad \text{for} \quad q_a \to e^{-i\theta_{\rm QFD}/(2N_q)\gamma_5} q_a \,.$$

The physical $\bar{\theta}$ violates P and CP. Due to this CP violation, the pion-nucleon interactions now can be generalized to

$$\mathcal{L}_{\pi NN} = \bar{N}\tau \cdot \pi(i\gamma_5 g_{\pi NN} + \overline{g_{\pi NN}})N \quad \text{with} \quad \overline{g_{\pi NN}} = -\bar{\theta}\frac{2(m_\Xi - m_\Sigma)m_u m_d}{f_\pi(m_u + m_d)(2m_s - m_u - m_d)} \simeq -0.023\bar{\theta}\,.$$

The neutron electric dipole moment then can be estimated as

$$\frac{d_n}{\mathrm{e}} = \frac{g_{\pi NN}\overline{g_{\pi NN}}}{4\pi^2 m_N} \ln\left(\frac{m_N}{m_\pi}\right)\,.$$

The recent upper bound on the neutron electric dipole moment is

$$|d_n| < 2.9 \times 10^{-26}\,\text{e-cm} \Rightarrow |\bar{\theta}| < 0.7 \times 10^{-11}\,.$$

The extremely small upper bound on the strong CP phase $\bar{\theta}$ has led to the so-called strong CP problem (e.g. Kim, Carosi, 08'.).

Currently, the most popular solution to the strong CP problem is to introduce a continuous Peccei-Quinn (Peccei, Quinn, 77'.) symmetry which predicts a pseudo Goldstone boson, the axion (Peccei, Quinn, 77'; Weinberg, 78; Wilczek, 78'; Kim, 79'; Shifman, Vainshtein, Zakharov, 80'; Zhitnitsky, 80'; Dine, Fischler, Srednicki, 81'.), and hence a dynamical strong CP phase. However, the axion has not been experimentally seen so far.

Alternatively, we can consider certain discrete symmetries such as P or CP (Mohapatra, Senjanović, 78'; Beg, Tsao, 78'; Georgi, 78'.) to suppress or remove the strong CP phase. For example (Babu, Mohapatra, 89'.), we can realize the parity symmetry for solving the strong CP problem in the $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ left-right symmetric theories (Pati, Salam, 74'; Mohapatra, Pati, 75'; Mohapatra, Senjanović, 75'.).

In the left-right symmetric context with a parity symmetry for solving the strong CP problem, we can realize a leptogenesis with predictions on Dirac neutrinos' CP violation and mass scale (PHG, 1410.5753.). In this new leptogenesis scenario, some dark matter scalars play an essential role.

The model

$$SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_L \times U(1)_R \times U(1)_N \times Z_2 \times Z_2'$$

$SU(2)_L \times U(1)_L$: $q_L(2, +\frac{1}{6})$, $d_R(1, -\frac{1}{3})$, $u_R(1, +\frac{2}{3})$, $l_L(2, -\frac{1}{2})$, $e_R(1, -1)$; $\phi(2, -\frac{1}{2})$, $\eta(2, -\frac{1}{2})$.

$SU(2)_R \times U(1)_R$: $q_R'(2, +\frac{1}{6})$, $d_L'(1, -\frac{1}{3})$, $u_L'(1, +\frac{2}{3})$, $l_R'(2, -\frac{1}{2})$, $e_L'(1, -1)$; $\phi'(2, -\frac{1}{2})$, $\eta'(2, -\frac{1}{2})$.

$U(1)_L \times U(1)_R$: $\omega(-3, +3)$, $\chi_d(-\frac{1}{3}, +\frac{1}{3})$, $\chi_u(+\frac{2}{3}, -\frac{2}{3})$, $\chi_e(-1, +1)$.

$U(1)_N$: $\nu_R(1)$, $\nu_L'(1)$; $\eta(-1)$, $\eta'(-1)$, $\sigma(-1)$.

$Z_2 \times Z_2'$: fermions$(-,-)$; $\chi_d(-,-)$, $\chi_u(-,-)$, $\chi_e(-,-)$.

$$U(1)_L \times U(1)_R \xrightarrow{\langle\omega\rangle} U(1)_{B-L}\,.$$

$$
\begin{aligned}
\mathcal{L}_K &\supset (D_\mu\phi)^\dagger D^\mu\phi + (D_\mu\phi')^\dagger D^\mu\phi' + (D_\mu\eta)^\dagger D^\mu\eta + (D_\mu\eta')^\dagger D^\mu\eta' \\
&+ i\bar{q}_L \not{D} q_L + i\bar{q}'_R \not{D} q'_R + i\bar{d}_R \not{D} d_R + i\bar{d}'_L \not{D} d'_L + i\bar{u}_R \not{D} u_R + i\bar{u}'_L \not{D} u'_L \\
&+ i\bar{l}_L \not{D} l_L + i\bar{l}'_R \not{D} l'_R + i\bar{e}_R \not{D} e_R + i\bar{e}'_L \not{D} e'_L \quad \text{with} \quad g_L = g_R = g_2\,. \\
V &\supset \mu_\phi^2 \phi^\dagger\phi + \mu_{\phi'}^2 \phi'^\dagger\phi + \mu_\eta^2 \xi^\dagger\xi + \mu_{\xi'}^2 \xi'^\dagger\xi') + \rho(\sigma\eta^\dagger\phi + \text{H.c.}) + \rho'(\sigma\eta'^\dagger\phi' + \text{H.c.}) \\
&+ \lambda_3(\chi_e^*\chi_d^3 + \text{H.c.}) + \lambda_2(\chi_e^*\chi_d^*\chi_u^2 + \text{H.c.}) \quad \text{with} \quad \mu_\phi^2 \neq \mu_{\phi'}^2\,, \;\; \mu_\eta^2 \neq \mu_{\eta'}^2\,, \;\; \rho \neq \rho'\,. \\
\mathcal{L}_Y &= -y_d(\bar{q}_L\tilde{\phi} d_R + \bar{q}'_R\tilde{\phi}' d'_L) - y_u(\bar{q}_L\phi u_R + \bar{q}'_R\phi' u'_L) - y_e(\bar{l}_L\tilde{\phi} e_R + \bar{l}'_R\tilde{\phi}' e'_L) - y_\nu(\bar{l}_L\eta\nu_R + \bar{l}'_R\eta'\nu'_L) \\
&- f_d\chi_d\bar{d}_R d'_L + f_u\chi_u\bar{u}_R u'_L - f_e\chi_e\bar{e}_R e'_L + \text{H.c.} \quad \text{with} \quad f_d = f_d^\dagger\,, \;\; f_u = f_u^\dagger\,, \;\; f_e = f_e^\dagger\,.
\end{aligned}
$$

Symmetry breaking

$$
\begin{array}{c}
SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times U(1)_N \\
\Big\Downarrow \langle\sigma\rangle \\
SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \\
\Big\Downarrow \langle\eta'\rangle \\
SU(3)_c \times SU(2)_L \times U(1)_Y \\
\Big\Downarrow \langle\phi\rangle \\
SU(3)_c \times U(1)_{em}\,.
\end{array}
$$

$$\langle\eta'\rangle \equiv \begin{bmatrix} \langle\eta'^0\rangle \\ 0 \end{bmatrix} \geq \langle\phi'\rangle \equiv \begin{bmatrix} \langle\phi'^0\rangle \\ 0 \end{bmatrix} \simeq -\frac{\rho'\langle\sigma\rangle\langle\eta'\rangle}{M_{\eta'}^2}\,.$$

$$\langle\phi\rangle \equiv \begin{bmatrix} \langle\phi^0\rangle \\ 0 \end{bmatrix} \simeq 174\,\mathrm{GeV} \gg \langle\eta\rangle \equiv \begin{bmatrix} \langle\eta^0\rangle \\ 0 \end{bmatrix} \simeq -\frac{\rho\langle\sigma\rangle\langle\phi\rangle}{M_\eta^2} = \mathcal{O}(0.01\text{-}0.1\,\mathrm{eV})\,.$$

Masses of fermions and charged gauge bosons

$$\left.\begin{aligned} m_f &= y_f\langle\phi\rangle \\ M_{f'} &= y_f\langle\phi'\rangle \end{aligned}\right\} \Rightarrow \frac{M_{d'}}{m_d} = \frac{M_{s'}}{m_s} = \frac{M_{b'}}{m_b} = \frac{M_{u'}}{m_u} = \frac{M_{c'}}{m_c} = \frac{M_{t'}}{m_t} = \frac{M_{e'}}{m_e} = \frac{M_{\mu'}}{m_\mu} = \frac{M_{\tau'}}{m_\tau} \equiv \frac{\langle\phi'\rangle}{\langle\phi\rangle}\,,$$

$$\left.\begin{aligned} m_\nu &= y_\nu\langle\eta\rangle \\ M_{\nu'} &= y_\nu\langle\eta'\rangle \end{aligned}\right\} \Rightarrow \frac{M_{\nu'_1}}{m_{\nu_1}} = \frac{M_{\nu'_2}}{m_{\nu_2}} = \frac{M_{\nu'_3}}{m_{\nu_3}} \equiv \frac{\langle\eta'\rangle}{\langle\eta\rangle}\,, \quad \Longrightarrow$$

$$\begin{cases} V'_{\text{CKM}} = V_{\text{CKM}} \equiv V\,, \\ U'_{\text{PMNS}} = U_{\text{PMNS}} \equiv U = \begin{bmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{bmatrix}. \end{cases}$$

$$m_{W_L} = \frac{1}{\sqrt{2}}g\sqrt{\langle\phi\rangle^2 + \langle\eta\rangle^2} \simeq \frac{1}{\sqrt{2}}g\langle\phi\rangle \quad \text{for} \quad \langle\eta\rangle \ll \langle\phi\rangle\,,$$

$$M_{W_R} = \frac{1}{\sqrt{2}}g\sqrt{\langle\phi'\rangle^2 + \langle\eta'\rangle^2} = \frac{1}{\sqrt{2}}g\langle\eta'\rangle/\sin\beta' \quad \text{with} \quad \sin\beta' = \frac{\langle\phi'\rangle}{\sqrt{\langle\phi'\rangle^2+\langle\eta'\rangle^2}} \leq \frac{1}{\sqrt{2}} \quad \text{for} \quad \langle\phi'\rangle \leq \langle\eta'\rangle\,.$$

Solution to the strong CP problem

The strong CP phase contains two parts,

$$\bar{\theta} = \theta_{\text{QCD}} + \theta_{\text{QFD}}\,.$$

First of all, the original QCD phase θ_{QCD} can be removed as a result of the parity invariance.

Furthermore, the contribution from the ordinary and mirror quark masses will arrive at zero since the ordinary and mirror quarks have the mass matrices as below,

$$\mathcal{L} \supset -[\bar{d}_L \;\; \bar{d}'_L] M_d \begin{bmatrix} d_R \\ d'_R \end{bmatrix} - [\bar{u}_L \;\; \bar{u}'_L] M_u \begin{bmatrix} u_R \\ u'_R \end{bmatrix} + \text{H.c. with } M_d = \begin{bmatrix} y_{d(u)}\langle\phi\rangle & 0 \\ 0 & y^{\dagger}_{d(u)}\langle\phi'\rangle \end{bmatrix},$$

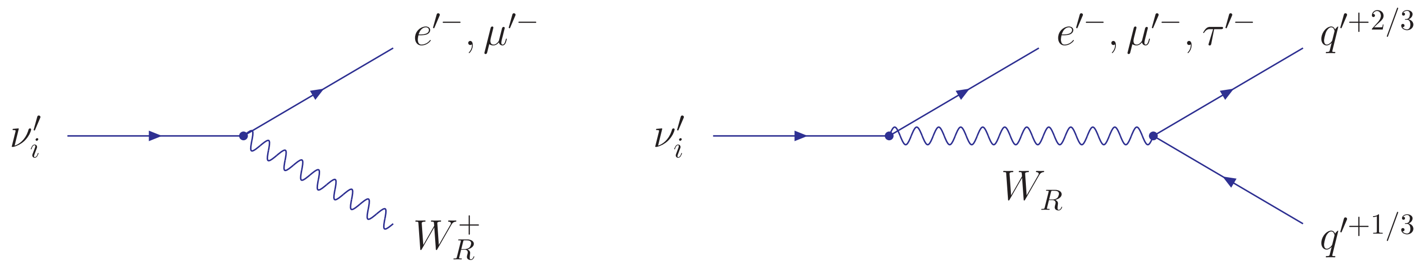
which means

$$\theta_{\text{QFD}} = \text{Arg}[\text{Det}(M_u M_d)] = \text{Arg}[\text{Det}(M_u)\text{Det}(M_d)] = 0\,,$$

for $\text{Det}(M_d)$ and $\text{Det}(M_u)$ being real.

Therefore, we can obtain a vanishing strong CP phase $\bar{\theta} = 0$.

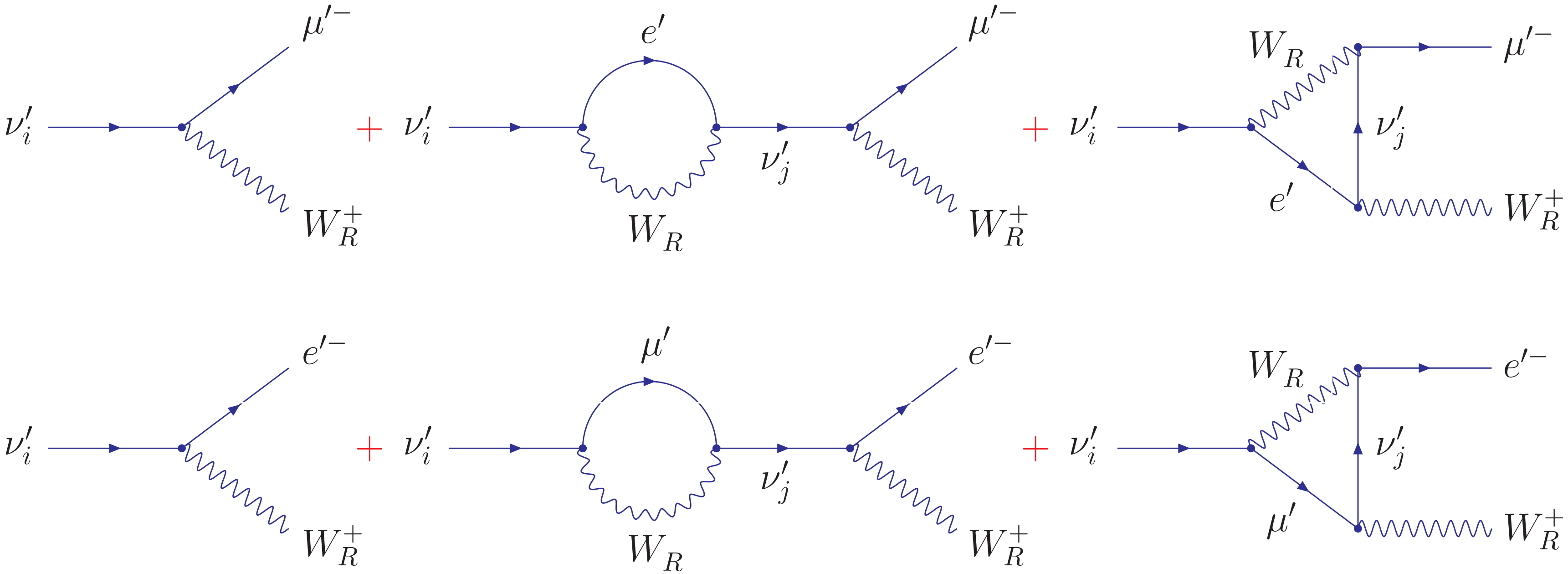
Mirror lepton asymmetry



$$\Gamma_{\nu_i'} = \Gamma(\nu_i' \to \mu'^- + W_R^+) + \Gamma(\nu_i' \to e'^- + W_R^+) + \sum_{\alpha\beta\gamma} \Gamma(\nu_i' \to e_\alpha'^- + q_\beta'^{+\frac{1}{3}} + q_\gamma'^{+\frac{2}{3}})$$

$$\simeq \frac{3g^4}{256\pi^3} M_{\nu_i'} \left[1 + \frac{4\pi^2}{g^2} \left(|U_{ei}|^2 + |U_{\mu i}|^2 \right) \xi^2 \right] \quad \text{for} \quad M_{\mu'} < M_{\nu_i'} - M_{W_R} < M_{\tau'} .$$

$$\xi = \frac{M_{\nu_i'}^2 - M_{W_R}^2}{M_{\nu_i'}^2} < \frac{(M_{\nu_i'} + M_{W_R}) M_{\tau'}}{M_{\nu_i'}^2} \simeq \frac{2M_{\tau'}}{M_{W_R}} = \frac{2m_\tau \sin\beta'}{m_{W_L}} = 0.044 \sin\beta' \leq 0.031 \quad \text{for} \quad \sin\beta' \leq \frac{1}{\sqrt{2}} .$$



ν'_i
μ'^-
W_R^+
+
ν'_i
e'
W_R
ν'_j
μ'^-
W_R^+
+
ν'_i
W_R
μ'^-
ν'_j
e'
W_R^+
ν'_i
e'^-
W_R^+
+
ν'_i
μ'
W_R
ν'_j
e'^-
W_R^+
+
ν'_i
W_R
e'^-
ν'_j
μ'
W_R^+

$$\varepsilon^{\mu'}_{\nu'_i} = \frac{\Gamma(\mu'_i \to \mu'^- + W_R^+) - \Gamma(\nu'^c_i \to \mu'^+ + W_R^-)}{\Gamma_{\nu'_i}},$$

$$\varepsilon^{e'}_{\nu'_i} = \frac{\Gamma(\nu'_i \to e'^- + W_R^+) - \Gamma(\nu'^c_i \to e'^+ + W_R^-)}{\Gamma_{\nu'_i}}.$$

$$\varepsilon^{\mu'}_{\nu'_i} = -\varepsilon^{\mu'}_{e'_i} \simeq \begin{cases} \frac{3\pi}{4} \frac{\mathrm{Im}(U^*_{ei} U_{ej} U_{\mu i} U^*_{\mu j})}{1+\frac{4\pi^2}{g^2}\left(|U_{ei}|^2+|U_{\mu i}|^2\right)\xi^2} \frac{M^2_{\nu'_i}}{M^2_{\nu'_j}-M^2_{\nu'_i}} \xi^4 & \text{for } M_{\nu'_j} \simeq M_{\nu'_i}, \\ \frac{3\pi}{2} \frac{\mathrm{Im}(U^*_{ei} U_{ej} U_{\mu i} U^*_{\mu j})}{1+\frac{4\pi^2}{g^2}\left(|U_{ei}|^2+|U_{\mu i}|^2\right)\xi^2} \frac{M^2_{\nu'_i}}{M^2_{\nu'_j}} \xi^4 & \text{for } M^2_{\nu'_j} \gg M^2_{\nu'_i}. \end{cases}$$

If the ordinary neutrinos have a quasi-degenerate spectrum, i.e. $m_{\nu_1} \equiv m_1 \simeq m_{\nu_2} \equiv m_2 \simeq m_{\nu_3} \equiv m_3$, the decays of all of the three mirror neutrinos $\nu'_{1,2,3}$ will contribute to the mirror lepton asymmetry.

$$\varepsilon^{\mu'}_{\nu'_1} = -\varepsilon^{e'}_{\nu'_1} \simeq \frac{3\pi}{4} J_{CP} \left(\frac{m_1^2}{\Delta m_{21}^2} - \frac{m_1^2}{\Delta m_{31}^2} \right) \xi^4 ,$$

$$\varepsilon^{\mu'}_{\nu'_2} = -\varepsilon^{e'}_{\nu'_2} \simeq \frac{3\pi}{4} J_{CP} \left(\frac{m_2^2}{\Delta m_{21}^2} + \frac{m_2^2}{\Delta m_{32}^2} \right) \xi^4 ,$$

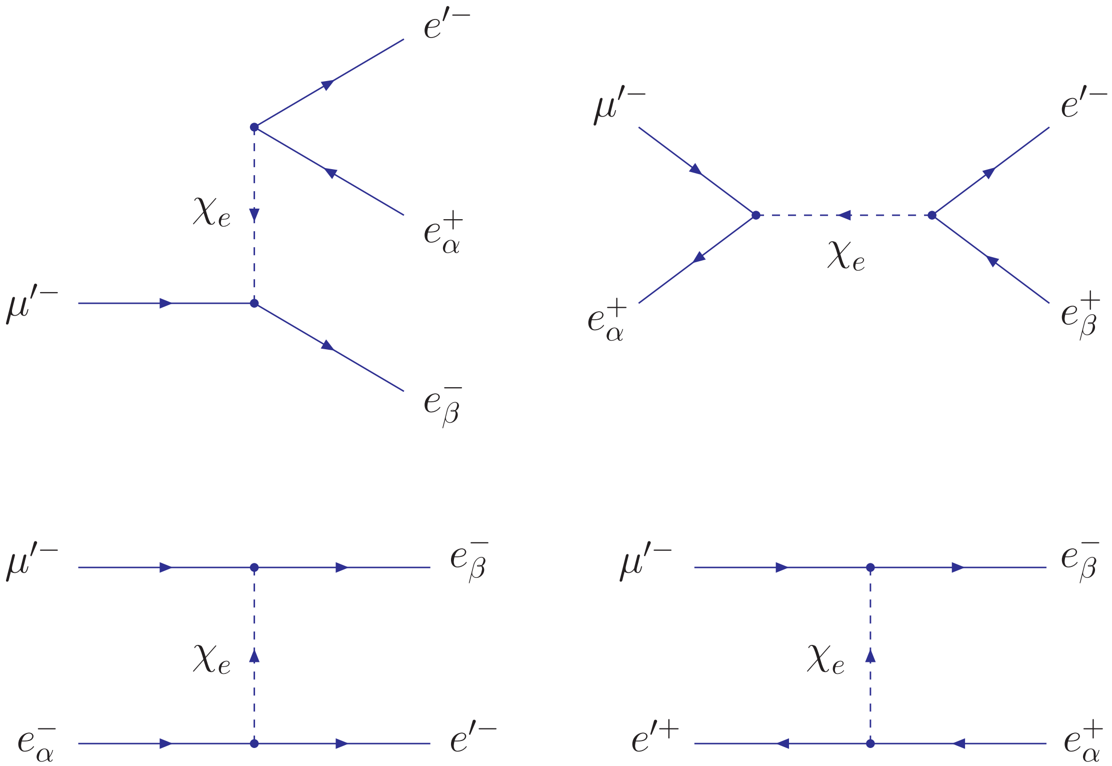
$$\varepsilon^{\mu'}_{\nu'_3} = -\varepsilon^{e'}_{\nu'_3} \simeq \frac{3\pi}{4} J_{CP} \left(\frac{m_3^2}{\Delta m_{32}^2} - \frac{m_3^2}{\Delta m_{31}^2} \right) \xi^4 .$$

If the ordinary neutrinos have a hierarchical spectrum, only the decays of the lightest mirror neutrinos ν'_1 or ν'_3 will dominate the mirror lepton asymmetry.

$$\varepsilon^{\mu'}_{\nu'_1} = -\varepsilon^{e'}_{\nu'_1} \simeq \frac{3\pi}{2} J_{CP} \left(\frac{m_1^2}{m_2^2} - \frac{m_1^2}{m_3^2} \right) \xi^4 ,$$

$$\varepsilon^{\mu'}_{\nu'_3} = -\varepsilon^{e'}_{\nu'_3} \simeq \frac{3\pi}{2} J_{CP} \left(\frac{m_3^2}{m_1^2} - \frac{m_3^2}{m_2^2} \right) \xi^4 .$$

$$J_{CP} = s_{12} c_{12} s_{23} c_{23} s_{13} c_{13}^2 \sin\delta .$$



Note in the presence of a fast conversion between the mirror muons and the mirror electrons, the mirror muon asymmetry and the mirror electron asymmetry will be cancelled each other once they are produced.

In order to avoid the cancellation between the mirror muon asymmetry and the mirror electron asymmetry, we can require

$$\Gamma_s^{\mu'^- \to e'^-} < H(T_{\rm sph}) \,,$$

$$\Gamma_t^{\mu'^- \to e'^-} < H(T_{\rm sph}) \,, \qquad H(T) = \left(\frac{8\pi^3 g_*}{90}\right)^{\frac{1}{2}} \frac{T^2}{M_{\rm Pl}} \,.$$

$$\Gamma_t^{\mu'^- + e'^+} < H(T_{\rm sph}) \,,$$

$$\Gamma_D^{\mu'^- \to e'^-} < H(T_{\rm sph}) \,.$$

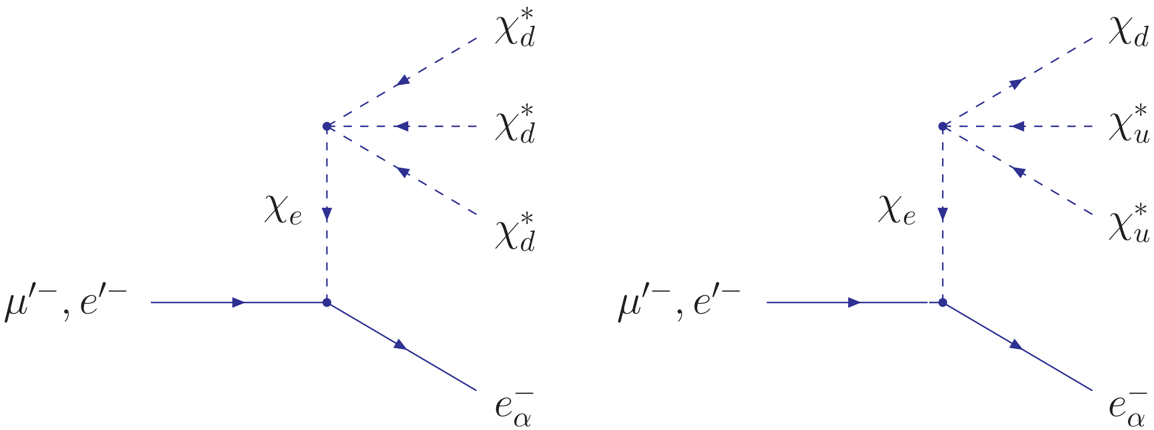
$$\Gamma_s^{\mu'^- \to e'^-} = \sum_{\alpha\beta} \Gamma(\mu'^- + e_\alpha^+ \to e'^- + e_\beta^+) \simeq \frac{3\left(f_e^\dagger f_e\right)_{\mu'\mu'}\left(f_e^\dagger f_e\right)_{e'e'}}{4\pi^3} \frac{T^5}{M_{\chi_e}^4} \quad \text{for} \quad M_{\mu'} < T < M_{\chi_e} \,,$$

$$\Gamma_t^{\mu'^- \to e'^-} = \sum_{\alpha\beta} \Gamma(\mu'^- + e_\alpha^- \to e'^- + e_\beta^-) \simeq \frac{\left(f_e^\dagger f_e\right)_{\mu'\mu'}\left(f_e^\dagger f_e\right)_{e'e'}}{4\pi^3} \frac{T^5}{M_{\chi_e}^4} \quad \text{for} \quad M_{\mu'} < T < M_{\chi_e} \,,$$

$$\Gamma_t^{\mu'^- + e'^+} = \sum_{\alpha\beta} \Gamma(\mu'^- + e_\alpha^+ \to e'^- + e_\beta^+) \simeq \frac{\left(f_e^\dagger f_e\right)_{\mu'\mu'}\left(f_e^\dagger f_e\right)_{e'e'}}{4\pi^3} \frac{T^5}{M_{\chi_e}^4} \quad \text{for} \quad M_{\mu'} < T < M_{\chi_e} \,,$$

$$\Gamma_D^{\mu'^- \to e'^-} = \sum_{\alpha\beta} \Gamma(\mu'^- \to e'^- + e_\alpha^- + e_\beta^+) \simeq \frac{\left(f_e^\dagger f_e\right)_{\mu'\mu'}\left(f_e^\dagger f_e\right)_{e'e'}}{6144\pi^3} \frac{M_{\mu'}^5}{M_{\chi_e}^4} \quad \text{for} \quad M_{\mu'} < M_{\chi_e} \,.$$

Baryon asymmetry



$$\Gamma_{\mu'} = \sum_\alpha \left[\Gamma(\mu'^- \to e_\alpha^- + \chi_d^* + \chi_d^* + \chi_d^*) + \Gamma(\mu'^- \to e_\alpha^- + \chi_d + \chi_u^* + \chi_u^*)\right] \simeq \frac{\left(f_e^\dagger f_e\right)_{\mu'\mu'}}{32768\pi^5}\left(|\kappa_3|^2 + \frac{1}{3}|\kappa_2|^2\right)\frac{M_{\mu'}^5}{M_{\chi_e}^4},$$

$$\Gamma_{e'} = \sum_\alpha \left[\Gamma(e'^- \to e_\alpha^- + \chi_d^* + \chi_d^* + \chi_d^*) + \Gamma(e'^- \to e_\alpha^- + \chi_d + \chi_u^* + \chi_u^*)\right] \simeq \frac{\left(f_e^\dagger f_e\right)_{e'e'}}{32768\pi^5}\left(|\kappa_3|^2 + \frac{1}{3}|\kappa_2|^2\right)\frac{M_{e'}^5}{M_{\chi_e}^4}.$$

The four-body decays of a mirror muon or electron into an ordinary lepton and three scalars can become efficient at different temperatures. Only the mirror muon asymmetry can be partially converted to the final baryon asymmetry.

$$\left.\begin{array}{l}\Gamma_{\mu'} > H(T_{\rm sph}) \\ \Gamma_{e'} > H(T_{\rm sph})\end{array}\right\} \Longrightarrow \frac{\eta_B}{7.04} \simeq -\frac{28}{79} \times \left\{\begin{array}{lll}(\varepsilon^{\mu}_{\nu'_1} + \varepsilon^{\mu}_{\nu'_2} + \varepsilon^{\mu}_{\nu'_3})/g_* & \text{for} & m_1 \simeq m_2 \simeq m_3\,, \\ \varepsilon^{\mu}_{\nu'_1}/g_* & \text{for} & m_1 \ll m_2 \ll m_3\,, \\ \varepsilon^{\mu}_{\nu'_3}/g_* & \text{for} & m_3 \ll m_1 \ll m_2\,.\end{array}\right.$$

The observed baryon asymmetry then can predict a correlation between the neutrinos' CP violation and mass scale.

$$\sin\delta = -3.2 \times 10^{-4} \quad \text{for} \quad m_1 \simeq m_2 \simeq m_3 = 0.2\,\text{eV} \Longrightarrow \sum m_\nu = 0.6\,\text{eV}\,;$$

$$m_1 = 0.00393\,\text{eV} < m_2 = 0.00953\,\text{eV} < m_3 = 0.0494\,\text{eV} \Longrightarrow \sum m_\nu = 0.0620\,\text{eV} \quad \text{for} \quad \sin\delta = -1\,.$$

(e.g. Capozzi *et al.*, 1312.2878.)

NH: $\Delta m^2_{21} = 7.54 \times 10^{-5}\,\text{eV}^2\,,\ \Delta m^2_{31} = 2.43 \times 10^{-3}\,\text{eV}^2\,,\ s^2_{12} = 0.308\,,\ s^2_{23} = 0.437\,,\ s^2_{13} = 0.0234\,;$

IH: $\Delta m^2_{21} = 7.54 \times 10^{-5}\,\text{eV}^2\,,\ \Delta m^2_{31} = 2.38 \times 10^{-3}\,\text{eV}^2\,,\ s^2_{12} = 0.308\,,\ s^2_{23} = 0.455\,,\ s^2_{13} = 0.0240\,.$

Dark matter

The scalars $\chi_{d,u}$ can keep stable as a result of the unbroken $Z_2 \times Z_2'$ symmetry.

The stable scalars $\chi_{d,u}$ can annihilate into the standard model particles through their couplings with the standard model Higgs scalar and then can obtain a relic density to serve as the dark matter particles. This dark matter scenario through Higgs portal has been studied by many people.

The dark matter scalars $\chi_{d,u}$ may efficiently annihilate into the ordinary quarks through the mediation of the first-generation mirror quarks even if the Higgs portal is not strong enough.

Summary

- Currently, no experimental results require the neutrinos to be the Majorana particles rather than the Dirac particles.
- On the theoretical hand, the Dirac neutrinos can nicely acquire their tiny masses in some Dirac seesaw models.
- We have proposed a class of $G_{\text{SM}} \times G_{\text{SM}'}$ models, the interactions for simultaneously generating the ordinary and dark matter asymmetries can also account for the Dirac neutrino masses. The asymmetric dark matter can have a predictive mass at about $6\,\text{GeV}$. If the dark proton is the dark matter particle, it can have a testable scattering off the ordinary nucleons in the presence of a small $U(1)_Y \times U(1)'_Y$ kinetic mixing input by hand. By imposing a proper discrete mirror symmetry, our Dirac seesaw models will not require more unknown parameters than the conventional Majorana seesaw models.

- We have presented an $SO(10) \times SO(10)'$ model with a $(\mathbf{16} \times \overline{\mathbf{16}'})_H$ scalar to give a common origin for the Dirac neutrino masses, the ordinary and dark matter asymmetries. The lightest dark nucleon as the dark matter particle should have a predictive mass at about $15\,\text{GeV}$. The dark proton can scatter off the ordinary nucleons at a testable level if it is the dark matter particle. Unlike the ordinary proton decay, the dark proton/neutron decay can be fast enough to open a window for the dark matter indirect detection experiments. The $U(1)_{em} \times U(1)'_{em}$ kinetic mixing which is the key for dark matter detection can be induced by the $(\mathbf{16} \times \overline{\mathbf{16}'})_H$ scalar.
- We have demonstrated a new leptogenesis scenario where the ordinary neutrinos can have a seesaw-suppressed Dirac mass term, while the mirror Dirac neutrinos can decay to generate a mirror muon asymmetry and an opposite mirror electron asymmetry although the net lepton number is zero. The four-body decays of a mirror muon into an ordinary lepton and three dark matter scalars can convert the mirror muon asymmetry to an ordinary lepton asymmetry before the sphaleron processes do not actively work. Benefited from the parity symmetry for solving the strong CP problem, this leptogenesis can give a correlation between the neutrinos' CP violation and mass scale.

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