A renormalizable left-right symmetric model with low scale seesaw mechanisms

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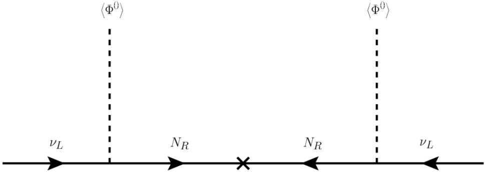
Introduction

Some reasons to extend the SM are:

1. SM charged fermion mass hierachy.
2. Tiny masses of light active neutrinos.
3. CKM and PMNS matrices very close and very different to the identity matrix, respectively.
4. Dark matter.
5. Number of SM fermion families.
6. Lepton asymmetry of the Universe.
7. $(g-2)_{\mu}$.
8. Tiny LFV signals several orders of magnitude below experimental sensitivity.



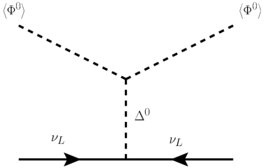
Type I seesaw

$LHN \quad \mathbf{2}\otimes\mathbf{2}\otimes\mathbf{1}$

Minkowski 1977, Gellman, Ramond, Slansky 1980

Glashow, Yanagida 1979, Mohapatra, Senjanovic 1980

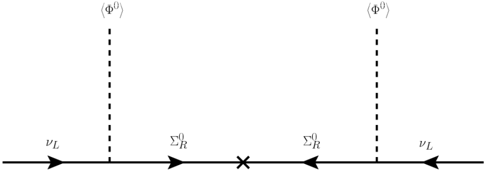
Lazarides Shafi Weterrich 1981, Schechter-Valle 1980 and 1982



Type II seesaw

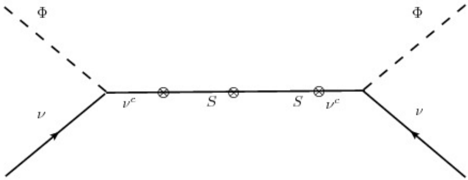
$L\Delta L \quad \mathbf{2}\otimes\mathbf{3}\otimes\mathbf{2}$

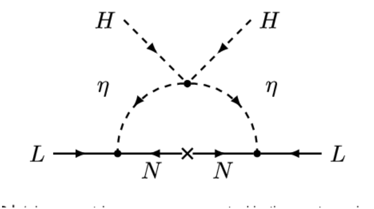
Schechter-Valle 1980 and 1982



Type III seesaw

$LH\Sigma \quad \mathbf{2}\otimes\mathbf{2}\otimes\mathbf{3}$





Inverse seesaw

$$-\mathcal{L}^{(\nu)}_{mass} = \frac{1}{2}\begin{pmatrix} \overline{\nu^C_L} & \overline{N}_R & \overline{S}_R \end{pmatrix} \mathbf{M}_\nu \begin{pmatrix} \nu_L \\ N^C_R \\ S^C_R \end{pmatrix} + H.c$$

$$\mathbf{M}_\nu = \begin{pmatrix} 0_{3\times3} & \mathbf{M}_1 & \mathbf{M_L} \\ \mathbf{M}_1^T & 0_{3\times3} & \mathbf{M}_2 \\ \mathbf{M_L}^T & \mathbf{M}_2^T & \mu \end{pmatrix}$$

$$\mathbf{M_L} = 0_{3\times3}$$

$$Q^{U(1)_L}_{\nu_L} = Q^{U(1)_L}_{S_R} = -Q^{U(1)_L}_{N_R} = 1$$

$$\widetilde{\mathbf{M}}_\nu = \mathbf{M}_1 \left(\mathbf{M}_2^T\right)^{-1} \mu \mathbf{M}_2^{-1} \mathbf{M}_1^T$$

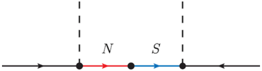
$$\mathbf{M}_\nu^{(1)} = -\frac{1}{2}\left(\mathbf{M}_2 + \mathbf{M}_2^T\right) + \frac{1}{2}\mu$$

$$\mathbf{M}_\nu^{(2)} = \frac{1}{2}\left(\mathbf{M}_2 + \mathbf{M}_2^T\right) + \frac{1}{2}\mu$$

One loop Ma radiative seesaw model

η and N are odd under a preserved Z_2

$$L\widetilde{\eta}N, \frac{\lambda_5}{2}\left(H^\dagger \cdot \eta\right)^2 + h.c$$



Linear seesaw:

$$\mu = 0_{3\times3}$$

$$\widetilde{\mathbf{M}}_\nu = -\mathbf{M}_L \mathbf{M}_2^{-1} \mathbf{M}_1^T - \mathbf{M}_1 \left(\mathbf{M}_2^T\right)^{-1} \mathbf{M}_L^T$$

An extended left-right symmetric model.

The main features of the extended LR symmetric model are:

1. The top quark gets its mass from a renormalizable Yukawa interaction.
2. The bottom, strange and charm quarks, as well as for the tau and muon leptons acquire their masses from a tree level Universal Seesaw mechanism.
3. The first generation SM charged fermions acquire one loop level masses.
4. The tiny neutrino masses are generated from an inverse seesaw mechanism at one loop level.

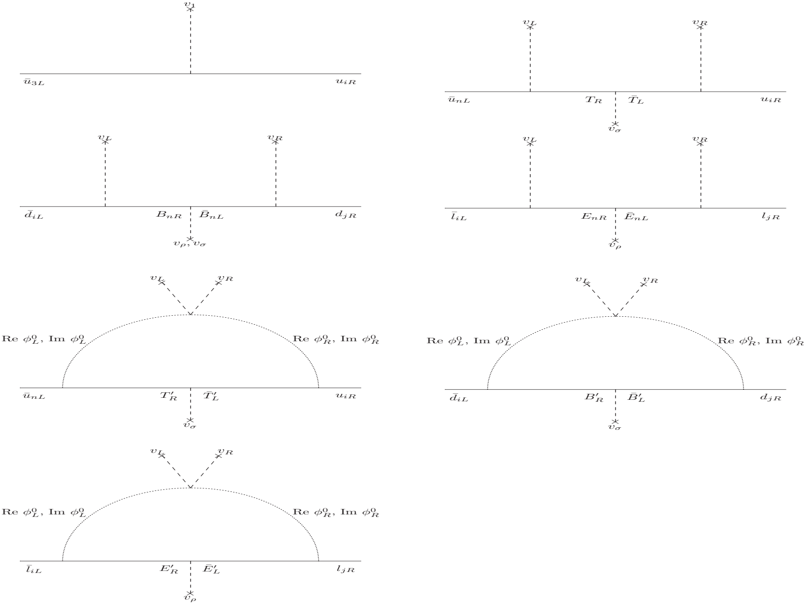


Figure: Feynman diagrams contributing to the entries of the SM charged fermion mass matrices. $n = 1, 2$, $i, j = 1, 2, 3$.

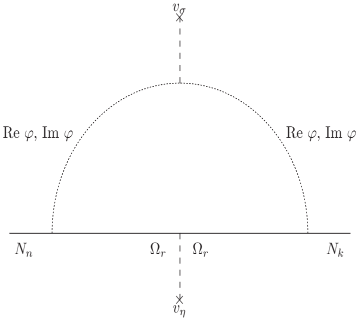


Figure: One-loop Feynman diagram contributing to the Majorana neutrino mass submatrix μ. Here, $n, k = 1, 2, 3$ and $r = 1, 2$.

$$
\begin{gathered}
\mathcal{G} = SU(3)_C \times SU\left(2\right)_L \times SU\left(2\right)_R \times U\left(1\right)_{B-L} \times Z_4^{(1)} \times Z_4^{(2)} \\
\Downarrow v_\sigma, v_\eta, v_\rho >> v_R \\
SU(3)_C \times SU\left(2\right)_L \times SU\left(2\right)_R \times U\left(1\right)_{B-L} \times Z_2 \\
\Downarrow v_R \sim \mathcal{O}(10)\, TeV \\
SU(3)_C \times SU\left(2\right)_L \times U\left(1\right)_Y \times Z_2 \\
\Downarrow v_1, v_L \\
SU(3)_C \otimes U\left(1\right)_Q \times Z_2
\end{gathered} \tag{1}
$$

The $Z_4^{(1)}$ symmetry is completely broken, whereas the $Z_4^{(2)}$ symmetry is broken down to the preserved Z_2 symmetry.

	Q_{nL}	Q_{3L}	Q_{iR}	L_{iL}	L_{iR}	T_L	T_R	T'_L	T'_R	B_{nL}	B_{1R}	B_{2R}	B'_L	B'_R	E_{nL}	E_{nR}	E'_L	E'_R	N_{iR}	Ω_{nR}
$SU(3)_C$	**3**	**3**	**3**	**1**	**1**	**3**	**3**	**3**	**3**	**3**	**3**	**3**	**3**	**3**	**1**	**1**	**1**	**1**	**1**	**1**
$SU(2)_L$	**2**	**2**	**1**	**2**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**
$SU(2)_R$	**1**	**1**	**2**	**1**	**2**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**
$U(1)_{B-L}$	$\frac{1}{3}$	$\frac{1}{3}$	$\frac{1}{3}$	-1	-1	$\frac{4}{3}$	$\frac{4}{3}$	$\frac{4}{3}$	$\frac{4}{3}$	$-\frac{2}{3}$	$-\frac{2}{3}$	$-\frac{2}{3}$	$-\frac{2}{3}$	$-\frac{2}{3}$	-2	-2	-2	-2	0	0
$Z_4^{(1)}$	-1	i	1	1	$-i$	1	1	1	1	1	$-i$	1	1	1	$-i$	-1	$-i$	-1	i	$-i$
$Z_4^{(2)}$	-1	-1	1	i	$-i$	1	-1	$-i$	i	1	-1	-1	i	$-i$	$-i$	i	1	-1	i	1

Table: Fermion assignments under $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times Z_4^{(1)} \times Z_4^{(2)}$. Here $i = 1, 2, 3$ and $n = 1, 2$

	Φ	χ_L	χ_R	ϕ_L	ϕ_R	φ	σ	η	ρ
$SU(3)_C$	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**
$SU(2)_L$	**2**	**2**	**1**	**2**	**1**	**1**	**1**	**1**	**1**
$SU(2)_R$	**2**	**1**	**2**	**1**	**2**	**1**	**1**	**1**	**1**
$U(1)_{B-L}$	0	1	1	1	1	0	0	0	0
$Z_4^{(1)}$	i	-1	1	-1	1	1	1	-1	i
$Z_4^{(2)}$	-1	1	1	$-i$	$-i$	i	-1	1	-1

Table: Scalar assignments under $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times Z_4^{(1)} \times Z_4^{(2)}$.

The scalar dark matter candidate is the lightest among the $Re\varphi$, $Im\varphi$, $Re\phi_L^0$, $Re\phi_R^0$, $Im\phi_L^0$ and $Im\phi_R^0$ fields. The fermionic dark matter candidate is the lightest among the right handed Majorana neutrinos N_{iR} $(i = 1, 2, 3)$.

The vacuum expectation values (VEVs) of the scalars Φ, χ_L and χ_R are:

$$\langle\Phi\rangle = \begin{pmatrix} v_1 & 0 \\ 0 & v_2 \end{pmatrix}, \quad \langle\chi_L\rangle = \begin{pmatrix} 0 \\ v_L \end{pmatrix}, \quad \langle\chi_R\rangle = \begin{pmatrix} 0 \\ v_R \end{pmatrix}, \tag{2}$$

We set $v_2 = 0$ to prevent a bottom quark mass arising from the Yukawa interaction involving the bidoublet scalar.

The bottom quark and tau lepton masses can be estimated as:

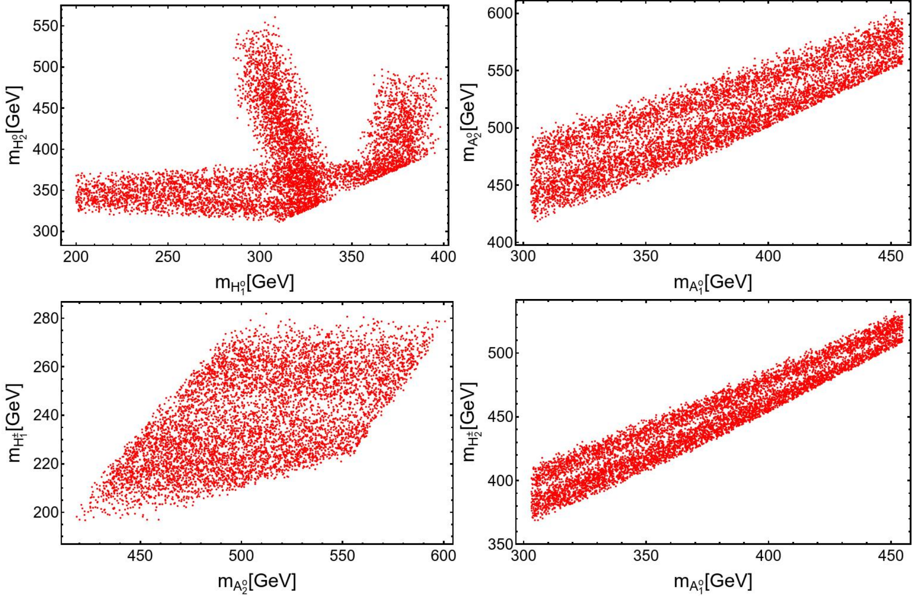
$$m_b \sim m_\tau \sim \frac{y^2 v_L v_R}{m_F} \tag{3}$$

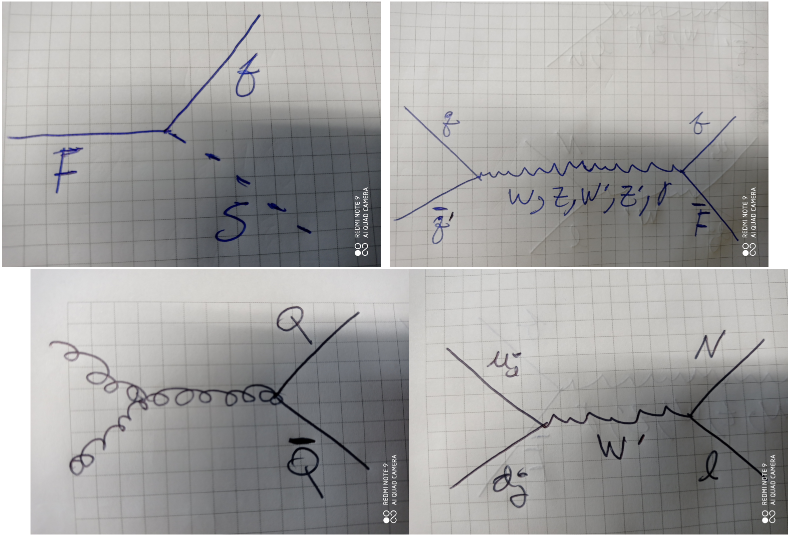
Taking $v_L \sim \mathcal{O}(100)$ GeV, $v_R \sim \mathcal{O}(10)$ TeV, $m_F \sim \mathcal{O}(100)$ TeV and $y \sim \mathcal{O}(0.4)$, it follows $m_b \sim m_\tau \sim \mathcal{O}(1)$ GeV.

The physical low energy scalar mass spectrum resulting from the scalar potential for the Φ, χ_L, ϕ_L, χ_R and ϕ_R scalars is composed of:

1. Physical electrically charged scalars $H_1^{\pm}$, $H_2^{\pm}$, $S_1^{\pm} = \phi_L^{\pm}$ and $S_2^{\pm} = \phi_R^{\pm}$.
2. Physical CP odd scalars A_1^0, A_2^0, P_1^0, P_2^0.
3. Physical CP even scalars h, H_1^0, H_2^0, H_3^0, $S_1^0 = \phi_{1I}^0$, $S_2^0 = \phi_{2I}^0$.

Furthermore, the massless scalar eigenstates $Im\chi_L^0$, $Im\chi_R^0$, $\chi_L^{\pm}$ and $\chi_R^{\pm}$ are associated with the Goldstone bosons corresponding to the longitudinal components of the Z, Z', $W^{\pm}$ and $W'^{\pm}$ gauge bosons, respectively.


m_{H_2^0}[GeV]
m_{H_1^0}[GeV]
m_{A_2^0}[GeV]
m_{A_1^0}[GeV]
m_{H_1^±}[GeV]
m_{A_2^0}[GeV]
m_{H_2^±}[GeV]
m_{A_1^0}[GeV]



The sterile neutrinos have the following decay modes $N_a^{\pm} \to l_i^{\pm} W^{\mp}$, $N_a^{\pm} \to \nu_i Z$ and $N_a^{\pm} \to \nu_i S$, $N_a^{\pm} \to l_i^{+} l_j^{-} \nu_k$, $N_a^{\pm} \to l_i^{-} u_j \bar{d}_k$, $N_a^{\pm} \to b\bar{b}\nu_k$

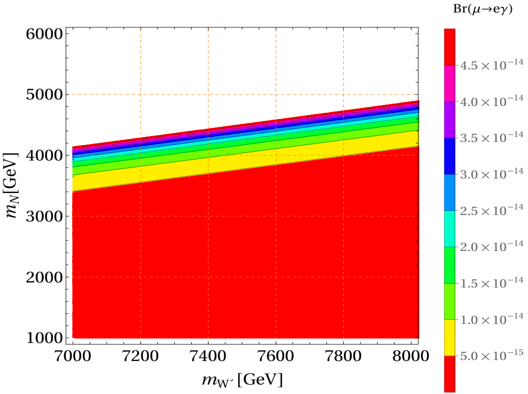


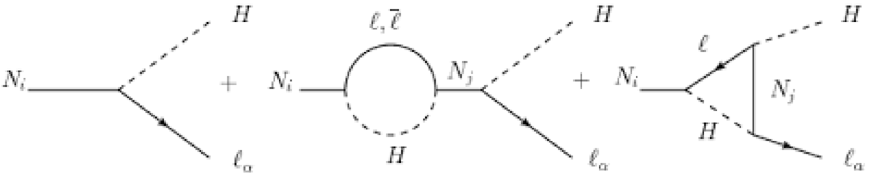
Figure: Allowed parameter space in the $m_{W'} - m_N$ plane consistent with the LFV constraints.

In the region of parameter space consistent with $\mu \to e\gamma$ decay rate constraints, the maximum obtained branching ratios for the $\tau \to \mu\gamma$ and $\tau \to e\gamma$ decays can reach values below their corresponding upper experimental bounds of 4.4×10^{-8} and 3.3×10^{-8}, respectively.

In our analysis of Leptogenesis we assume:

1. The first generation of pseudo-Dirac fermions $N_a^{\pm}$, i.e, $N_1^{\pm}$ will be much lighter than the second and third generation ones.
2. The off-diagonal Dirac and Majorana neutrino mass submatrices are diagonal with 11 entries much smaller than the remaining ones.
3. The initial temperature is taken to be much larger than the mass $m_{N^{\pm}}$ of the lightest pair of pseudo-Dirac fermions $N_1^{\pm} = N^{\pm}$.
4. The exotic leptonic fields E_{nR}, E' and Ω_{nR} $(n = 1, 2)$ are heavier than the lightest pseudo-Dirac fermions $N_1^{\pm} = N^{\pm}$.

The lepton asymmetry parameter, is induced by decay process of $N^{\pm}$.



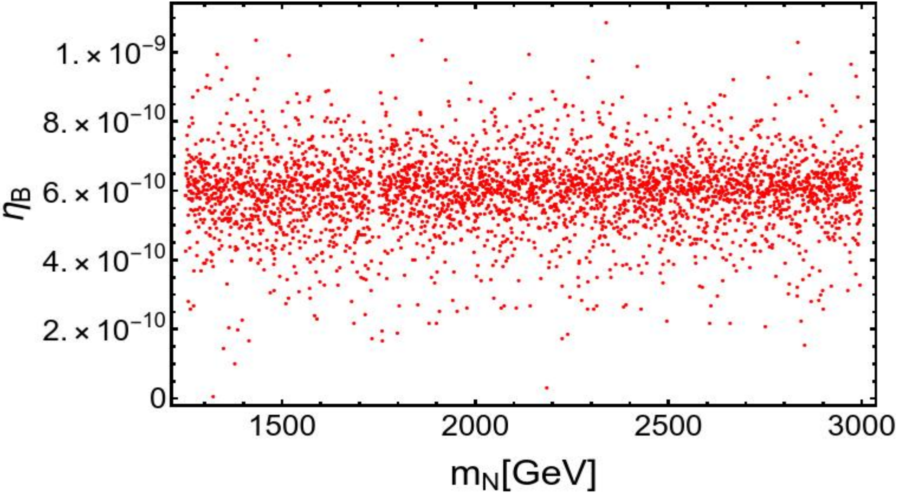
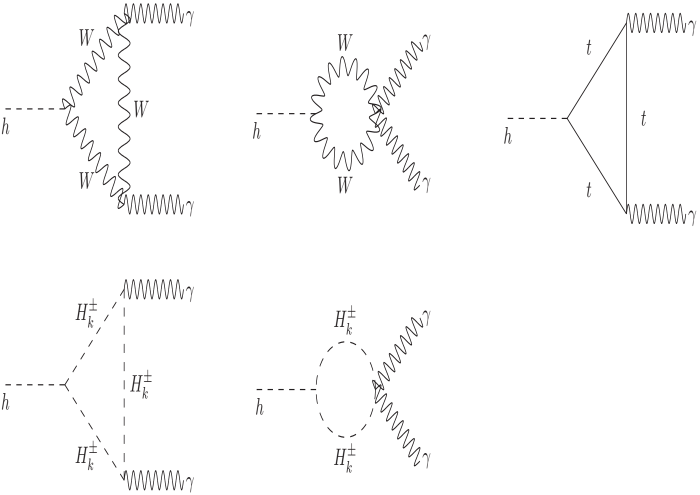


Figure: Correlation of the baryon asymmetry and the mass m_N of the lightest pair of pseudo-Dirac fermions $N_1^{\pm} = N^{\pm}$. Here we have set $v_R = 14$ TeV, $m_{W'} = 7$ TeV, $m_{Z'} = 7.2$ TeV, $m_{N_2^{\pm}} = 14$ TeV, $m_{N_3^{\pm}} = 28$ TeV and $\left|y_{22}^{(L)}\right| = \left|y_{33}^{(L)}\right| = \left|y_2^{(L)}\right|$


h
W
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$H_k^{\pm}$
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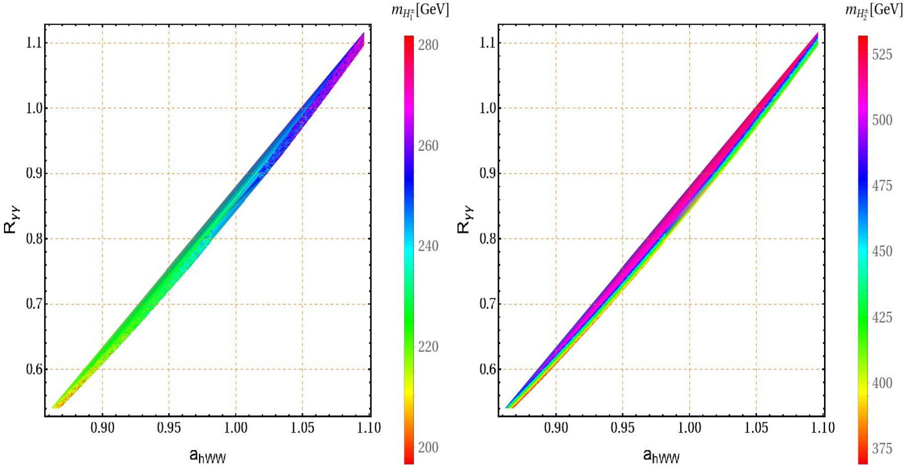
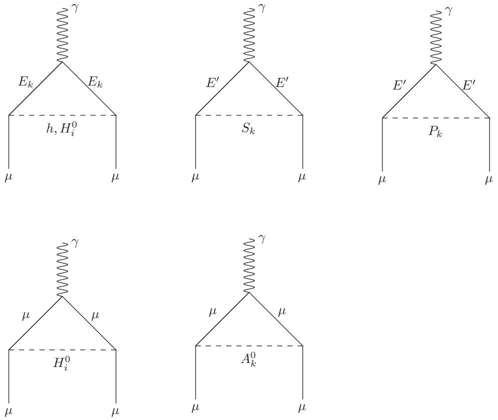
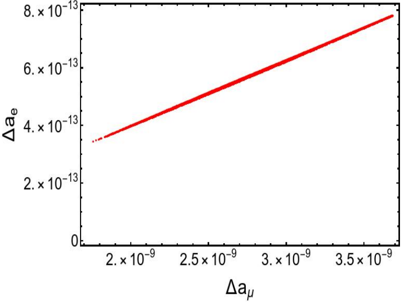


Figure: Correlation of the Higgs diphoton signal strength with the a_{hWW} deviation factor from the SM Higgs-W gauge boson coupling.


γ
E_k E_k
h, H_i^0
μ μ
γ
E′ E′
S_k
μ μ
γ
E′ E′
P_k
μ μ
γ
μ μ
H_i^0
μ μ
γ
μ μ
A_k^0
μ μ



8. × 10^−13
6. × 10^−13
4. × 10^−13
2. × 10^−13
0
Δa_e
2. × 10^−9
2.5 × 10^−9
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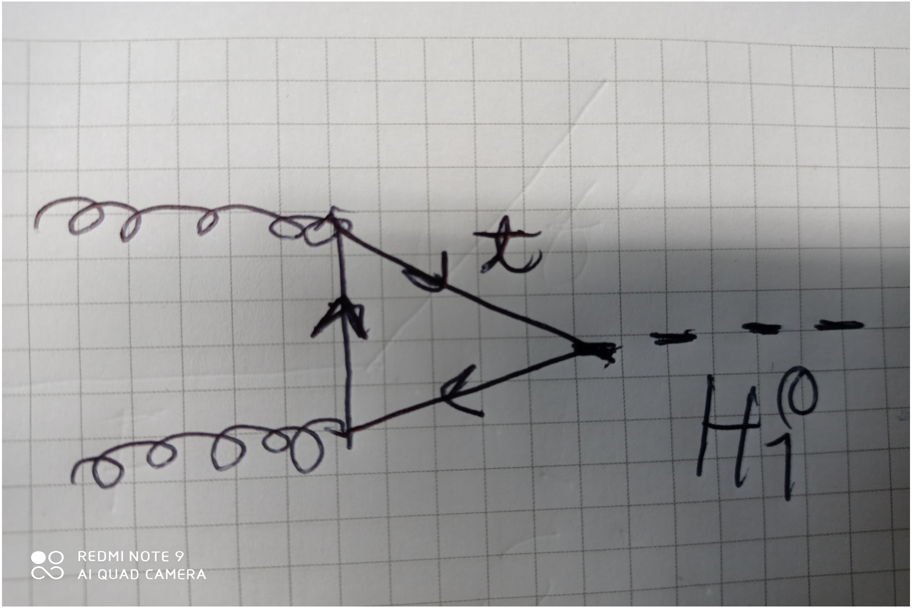
3.5 × 10^−9
Δa_μ


t
H_1^0
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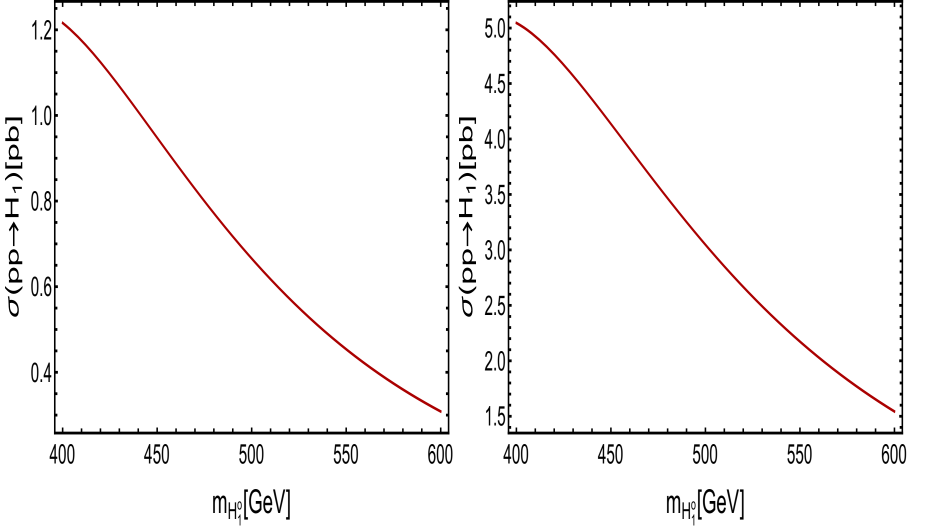
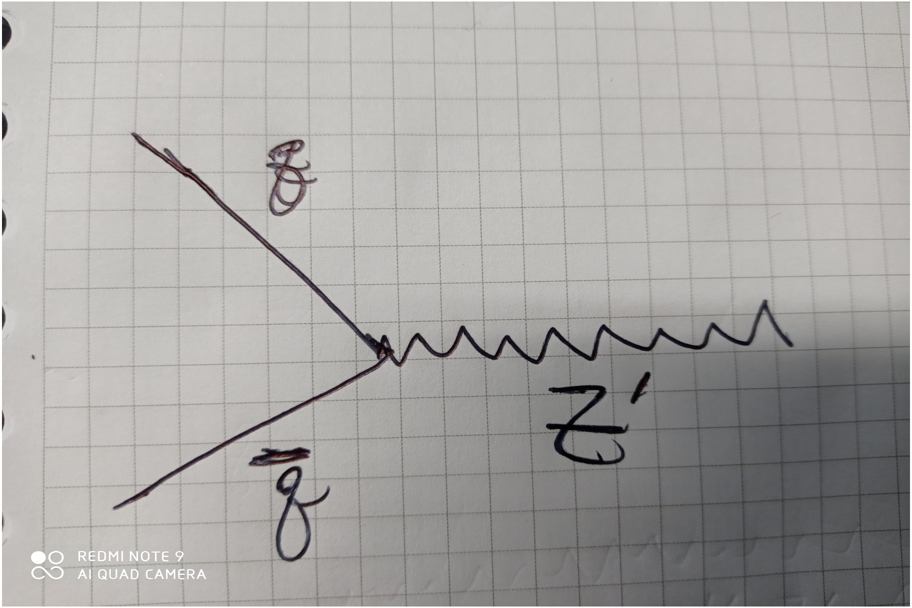


Figure: Total cross section for the H_1^0 production via gluon fusion mechanism at the LHC for $\sqrt{s} = 14$ TeV (left-panel) and $\sqrt{S} = 28$ (right-panel) TeV as a function of the heavy scalar mass $m_{H_1^0}$.


q
$\bar{q}$
Z'
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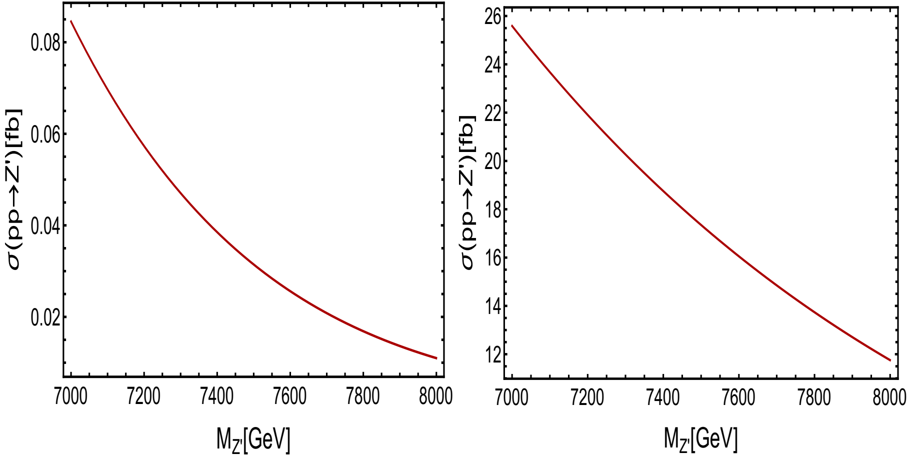

Figure: Total cross section for the Z' production via Drell-Yan mechanism at a proton-proton collider for $\sqrt{S} = 14$ TeV (left-panel) and $\sqrt{S} = 28$ (right-panel) TeV as a function of the Z' mass.

We take into account the constraint $\frac{M_{Z'}}{g_R} > 7$ TeV arising from LEP I and II measurements of $e^+e^- \to l^+l^-$, which is also consistent with LHC bounds.

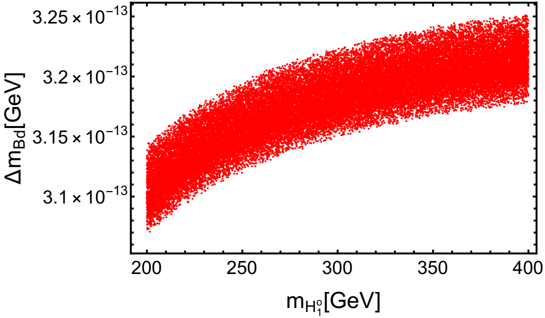

Figure: Correlation between the Δm_{B_d} mass splitting and the heavy CP even scalar mass $m_{H_1^0}$. The couplings of the flavor changing neutral Yukawa interactions that produce the $B_d^0 - \bar{B}_d^0$ oscillations have been set to be equal to 10^{-4}. We have also fixed $m_{H_3^0} = 10$ TeV and considered the ranges 200 GeV$\leqslant m_{H_1^0} \leqslant$ 400 GeV, 350 GeV$\leqslant m_{H_2^0} \leqslant$ 550 GeV and 300 GeV$\leqslant m_{A_1^0} \leqslant$ 450 GeV, whereas we have also set $m_{A_2^0} = m_{A_1^0} + 150 GeV$.

$B_s^0 - \bar{B}_s^0$ and $K^0 - \bar{K}^0$ are consistent with the experimental data for flavor violating Yukawa couplings equal to 2.5×10^{-4} and 10^{-6}, respectively.

Conclusions

- In the extended LR model, the top and exotic fermions get tree level masses whereas the masses of the bottom, charm and strange quarks, tau and muon leptons are generated from a tree level. The masses of the first generation of SM charged fermions arise at one loop level.
- Neutrino masses can be generated via one-loop inverse seesaw.
- Fermion masses and mixings, DM, meson oscillations, CLFV, $(g-2)_{e,\mu}$ anomalies, lepton and baryon asymmetries can be accounted for.
- Dark matter stability can arise from a residual parity symmetry.

Acknowledgements

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Extra Slides

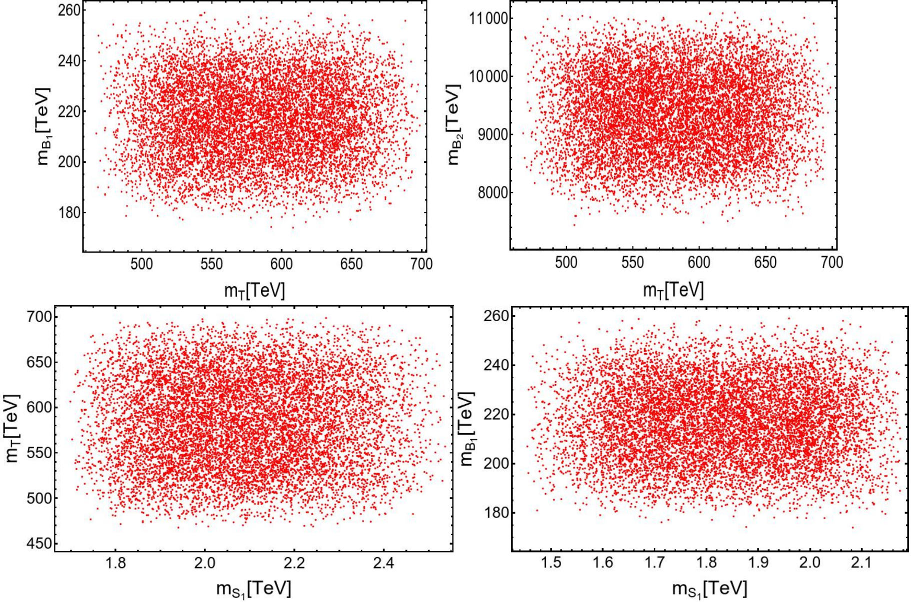


Figure: Correlations between the heavy exotic quark masses.

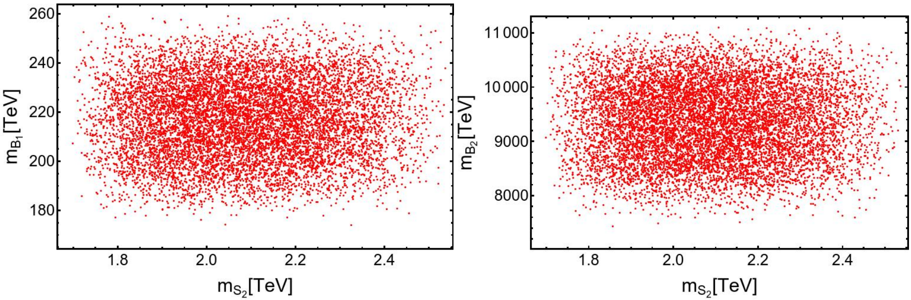


Figure: Correlations between the exotic quark masses and the masses m_{S_1} and m_{S_2} of the inert scalars S_1 and S_2, respectively.

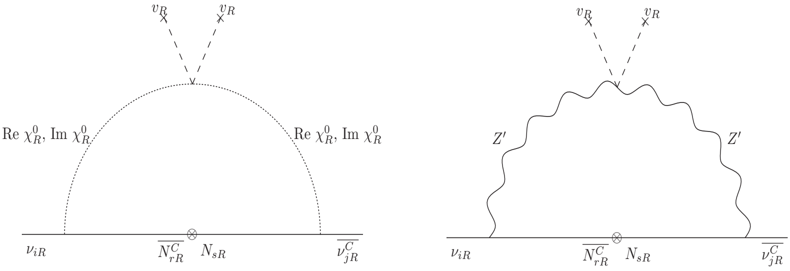
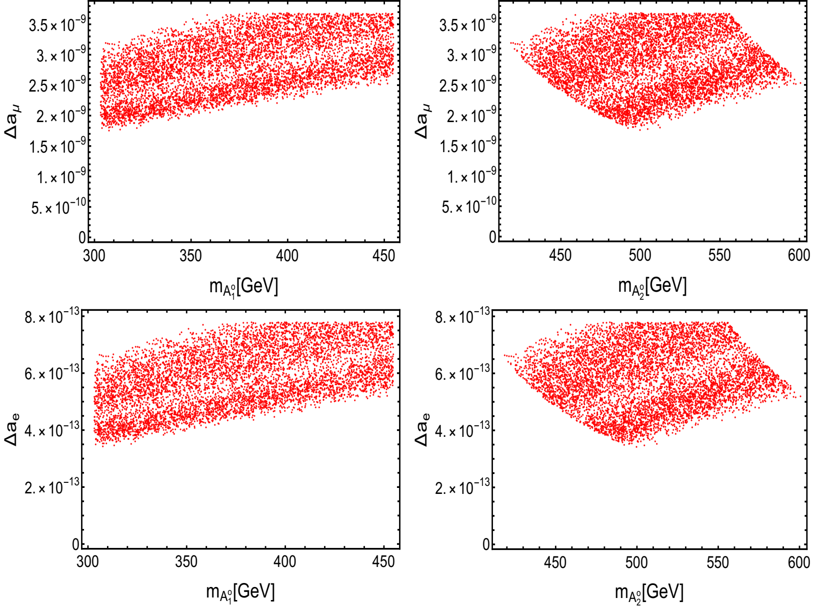


Figure: Feynman diagram contributing to the Majorana neutrino mass submatrix $\tilde{\mu}$. Here, $i, j, r, s = 1, 2, 3$ and the cross mark $\otimes$ in the internal lines corresponds to the one loop level induced Majorana mass term.


Δa_μ
$m_{A_1^0}$[GeV]
Δa_μ
$m_{A_2^0}$[GeV]
Δa_e
$m_{A_1^0}$[GeV]
Δa_e
$m_{A_2^0}$[GeV]

For the case of two heavy seesaw mediators, one has:

$$M = \begin{pmatrix} F_1 G_1 + X_1 Y_1 & F_1 G_2 + X_1 Y_2 & F_1 G_3 + X_1 Y_3 \\ F_2 G_1 + X_2 Y_1 & F_2 G_2 + X_2 Y_2 & F_2 G_3 + X_2 Y_3 \\ F_3 G_1 + X_3 Y_1 & F_3 G_2 + X_3 Y_2 & F_3 G_3 + X_3 Y_3 \end{pmatrix}, \tag{4}$$

$$M_{ij} = M^i_j = F_i G_j + X_i Y_j = F^i G_j + X^i Y_j \tag{5}$$

$$F^i = F_i, \qquad X^i = X_i \tag{6}$$

$$\det\left(M^i_j\right) = \frac{1}{3!}\varepsilon_{j_1 j_2 j_3}\varepsilon^{i_1 i_2 i_3} M^{j_1}_{i_1} M^{j_2}_{i_2} M^{j_3}_{i_3}, \qquad M^i_j = M_{ij} \tag{7}$$

$$\begin{aligned} \det\left(M^i_j\right) &= \frac{1}{3!}\varepsilon_{j_1 j_2 j_3}\varepsilon^{i_1 i_2 i_3}\left(F^{j_1} G_{i_1} + X^{j_1} Y_{i_1}\right) \\ &\times \left(F^{j_2} G_{i_2} + X^{j_2} Y_{i_2}\right)\left(F^{j_3} G_{i_3} + X^{j_3} Y_{i_3}\right) = 0 \end{aligned} \tag{8}$$

For the case of three fermionic seesaw mediators, one has:

$$M_{ij} = M^i_j = F_i G_j + X_i Y_j + R_i S_j = F^i G_j + X^i Y_j + R^i S_j \tag{9}$$

Then, it follows that:

$$\det\left(M^i_j\right) = \varepsilon_{j_1 j_2 j_3} \varepsilon^{i_1 i_2 i_3} F^{j_1} G_{i_1} X^{j_2} Y_{i_2} R^{j_3} S_{i_3} \neq 0 \tag{10}$$

provided that:

$$G_{i_1} \neq Y_{i_2} \neq S_{i_3}, \qquad F^{j_1} \neq X^{j_2} \neq R^{j_3} \tag{11}$$

Therefore, we have shown that in order to generate the masses of three fermion families via a seesaw mechanism, there should be at least three fermionic seesaw mediators. Furthermore, the number of the massless states obtained in a mass matrix resulting from a seesaw mechanism is $3 - n$, where n is the number of fermionic seesaw mediators.

In the analysis of Leptogenesis, the following processes are considered:

1. Decays $N_1^{\pm} \to l_i H_r^{+}$ $(r = 1, 2)$, $N_1^{\pm} \to \nu_i h$ ($i = 1, 2, 3$ and $N_1^{\pm} \to l_i^{-} u_j \bar{d}_k$ $(i, j, k = 1, 2, 3)$.
2. Z' mediated scattering processes $N_1^{\pm} N_1^{\pm} \longleftrightarrow l_i \bar{l}_j$ $(i, j = 1, 2, 3)$, $N_1^{\pm} N_1^{\pm} \longleftrightarrow u_i \overline{u}_j$ and $N_1^{\pm} N_1^{\pm} \longleftrightarrow d_i \overline{d}_j$.
3. W' mediated processes $N_1^{\pm} l_{iR} \longleftrightarrow \overline{u}_{jR} d_{kR}$, $N_1^{\pm} \overline{u}_{iR} \longleftrightarrow l_{jR} \overline{d}_{kR}$, $N_1^{\pm} d_{iR} \longleftrightarrow l_{jR} u_{kR}$ $(i, j, k = 1, 2, 3)$.
4. Inverse two and three body decays of $N_1^{\pm}$.

The Boltzmann equations take the form:

$$\begin{aligned}\frac{dN_{N_1^\pm}(z)}{dz} &= -\left[D(z)+S(z)\right]\left[N_{N_1^\pm}(z)-N_{N_1^\pm}^{eq}(z)\right],\\ \frac{dN_{N_{B-L}}(z)}{dz} &= -\varepsilon_\pm D(z)\left[N_{N_1^\pm}(z)-N_{N_1^\pm}^{eq}(z)\right]-W(z)\,N_{N_{B-L}}(z),\end{aligned}$$

where $z=\frac{m_{N_1^\pm}}{T}$, whereas $N_{N_1^\pm}$ and $N_{N_{B-L}}$ are the number density and the amount of $B-L$ asymmetry, respectively. Here $\varepsilon_\pm$ reads:

$$\begin{aligned}\varepsilon_\pm &= \sum_{i=1}^{3}\sum_{r=1}^{2}\frac{\left[\Gamma\left(N_\pm\to l_iH_r^+\right)-\Gamma\left(N_\pm\to\bar{l}_iH_r^-\right)\right]}{\left[\Gamma\left(N_\pm\to l_iH_r^+\right)+\Gamma\left(N_\pm\to\bar{l}_iH_r^-\right)\right]}\\ &+ \sum_{i=1}^{3}\frac{\left[\Gamma\left(N_\pm\to h\nu_i\right)-\Gamma\left(N_\pm\to h\nu_i\right)\right]}{\left[\Gamma\left(N_\pm\to h\nu_i\right)+\Gamma\left(N_\pm\to h\nu_i\right)\right]}\\ &\simeq \frac{\Im\left\{\left(\left[\left(y_{N_+}\right)^\dagger\left(y_{N_-}\right)\right]^2\right)_{11}\right\}}{8\pi A_\pm}\frac{r}{r^2+\frac{\Gamma_\pm^2}{m_{N_+}^2}},\end{aligned} \tag{12}$$

with:

$$r = \frac{m_{N_+}^2 - m_{N_-}^2}{m_{N_+} m_{N_-}}, \quad A_\pm = \left[\left(y_{N_\pm} \right)^\dagger y_{N_\pm} \right]_{11}, \quad \Gamma_\pm = \frac{A_\pm m_{N_\pm}}{8\pi},$$
$$y_{N_\pm} = \frac{y^{(L)}}{\sqrt{2}} (1 \mp S) = \frac{y^{(L)}}{\sqrt{2}} \left[1 \pm \frac{1}{4} M^{-1} \left(\mu + \widetilde{\mu} \right) \right], \quad (13)$$

The baryon to photon ratio is:

$$\eta_B = \frac{n_B}{n_\gamma} = \frac{3}{4} a_{sph} N_{B-L}, \qquad a_{sph} = \frac{8n_f + 4n_H}{22n_f + 13n_H}, \quad (14)$$

where a_{sph} is the L to B sphaleron conversion rate. Furthermore, n_f is the number of fermion families and n_H is the number of Higgs doublets.

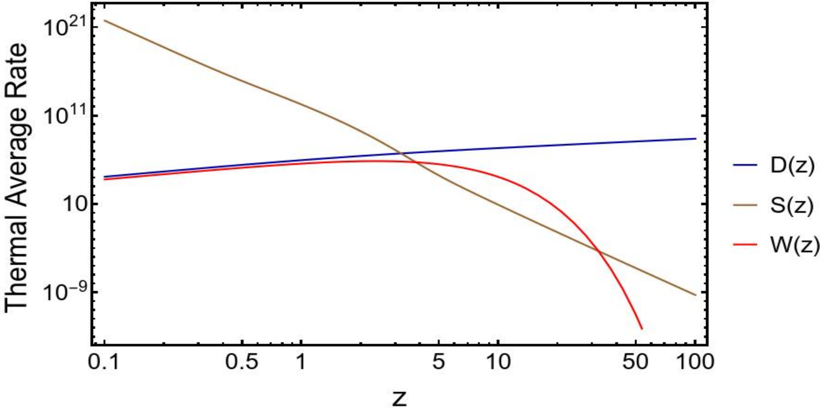


Figure: Thermally averaged scattering rates $D(z)$, $S(z)$ and $W(z)$ as functions of $z = \frac{m_N}{T}$, with m_N the mass of the lightest pair of pseudo-Dirac fermions $N_1^{\pm} = N^{\pm}$ and T the temperature. Here we have set $v_R = 14$ TeV, $m_{W'} = 7$ TeV and $m_{Z'} = 7.2$ TeV.

The contributions arising from the scattering processes, as well as from the inverse decays, are subdominant for temperatures sufficiently lower than the mass m_N, i.e, $z >> 1$.

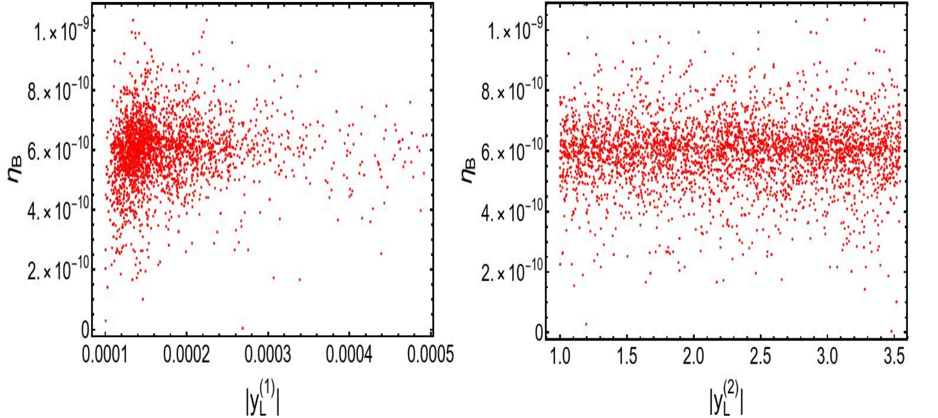
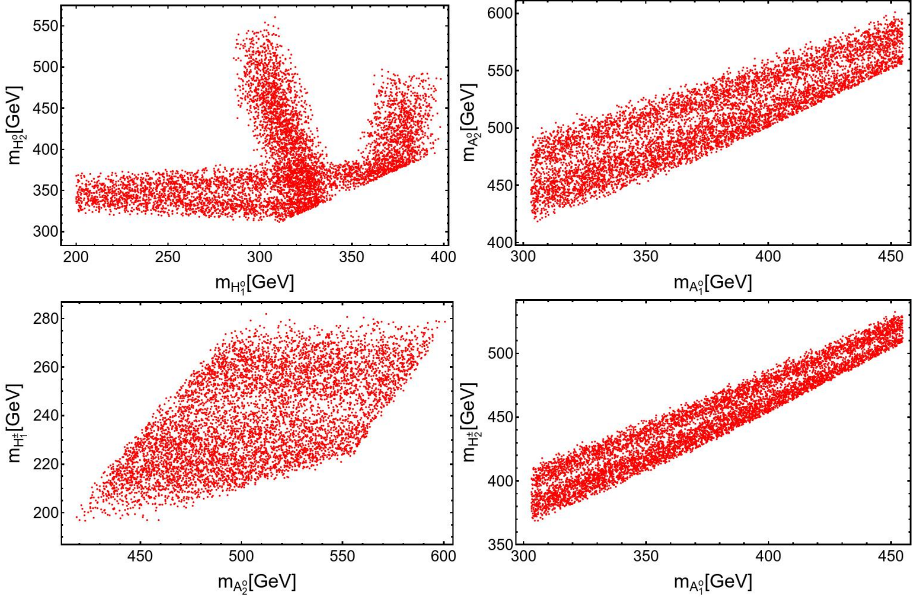


Figure: Correlation of the baryon asymmetry and the magnitude of the Dirac neutrino Yukawa couplings $y_{11}^{(L)}$ and $y_{22}^{(L)}$. Here we have set $v_R = 14$ TeV, $m_{W'} = 7$ TeV, $m_{Z'} = 7.2$ TeV, $m_{N_2^{\pm}} = 14$ TeV, $m_{N_3^{\pm}} = 28$ TeV and $\left|y_{22}^{(L)}\right| = \left|y_{33}^{(L)}\right| = \left|y_2^{(L)}\right|$.



550
500
450
400
350
300
$m_{H_2^0}$[GeV]
200
250
300
350
400
$m_{H_1^0}$[GeV]
600
550
500
450
400
$m_{A_2^0}$[GeV]
300
350
400
450
$m_{A_1^0}$[GeV]
280
260
240
220
200
$m_{H_1^\pm}$[GeV]
450
500
550
600
$m_{A_2^0}$[GeV]
500
450
400
350
$m_{H_2^\pm}$[GeV]
300
350
400
450
$m_{A_1^0}$[GeV]

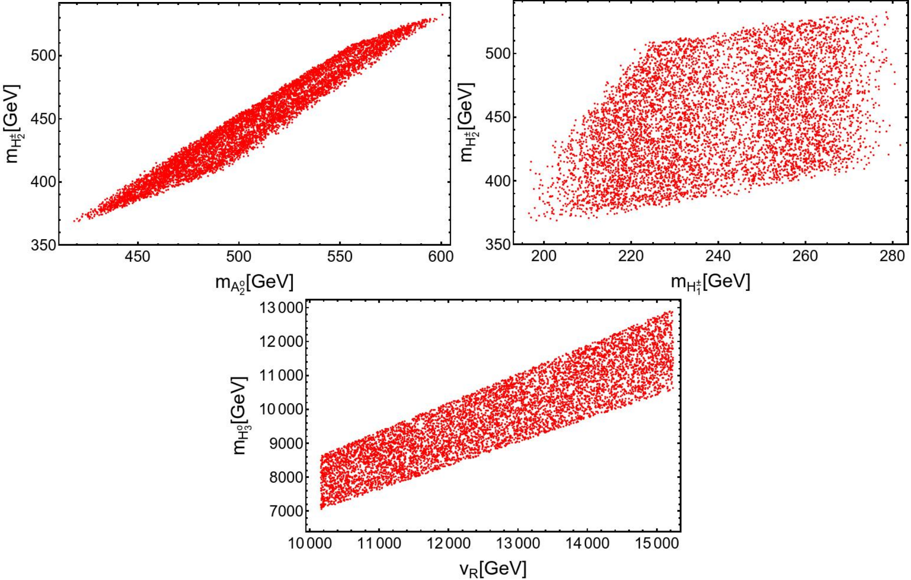


Figure: Correlations between the non SM scalar masses (top plots). Correlation between the mass of the CP even neutral scalar H_3^0 and the scale v_R of breaking of the left-right symmetry.

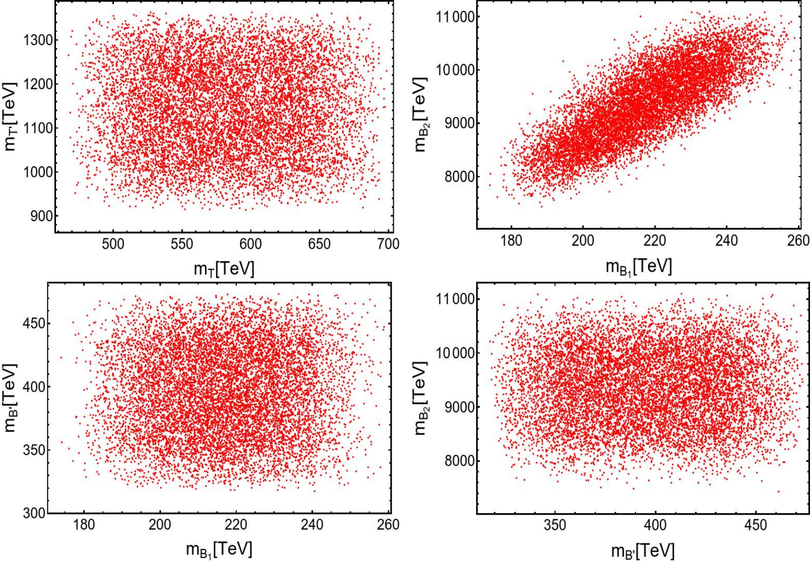


Figure: Correlations between the heavy exotic quark masses.

$$
\begin{aligned}
m_u^{\exp}(M_Z) &= 1.24 \pm 0.22 \text{ MeV} , \\
m_c^{\exp}(M_Z) &= 0.626 \pm 0.020 \text{ GeV} , \\
m_t^{\exp}(M_Z) &= 172.9 \pm 0.04 \text{ GeV} , \\
m_d^{\exp}(M_Z) &= 2.69 \pm 0.19 \text{ MeV} , \\
m_s^{\exp}(M_Z) &= 53.5 \pm 4.6 \text{ MeV} , \\
m_b^{\exp}(M_Z) &= 2.86 \pm 0.03 \text{ GeV} ,
\end{aligned}
$$

and the CKM parameters are [19]

$$
\begin{aligned}
|\mathbf{V}_{12}^{\exp}| &= 0.22452 \pm 0.00044 , \\
|\mathbf{V}_{23}^{\exp}| &= 0.04214 \pm 0.00076 , \\
|\mathbf{V}_{13}^{\exp}| &= 0.00365 \pm 0.00012 , \\
J_q^{\exp} &= (3.18 \pm 0.15) \times 10^{-5} .
\end{aligned}
$$

The magnitudes of the quark Yukawa couplings are randomly varied in the range [0.1, 1.5], whereas their complex phases are ranged between 0 and 2π. Furthermore, we have fixed $v_L = 100$ GeV and $v_R = 10$ TeV and randomly varied $\theta = \theta_S = -\theta_P$ in a small range around $\frac{\pi}{3}$. The masses of the vector like quarks and inert scalar mediators are varied in the ranges:

$$
\begin{aligned}
&0.5 \text{ TeV} \le m_{S_1} = m_{P_1} \le 10 \text{ TeV}, \qquad 1.01 m_{S_1} \le m_{S_2} = m_{P_2} \le 1.03 m_{S_1}, \\
&1 \text{ TeV} \le m_{T'}, m_{B'} \le 10^3 \text{ TeV}, \qquad 10^2 \text{ TeV} \le m_{B_1} \le 2 \times 10^2 \text{ TeV}, \\
&10^2 \frac{m_b}{m_c} \text{ TeV} \le m_T \le 2 \times 10^2 \frac{m_b}{m_c} \text{ TeV}, \qquad 10^2 \frac{m_b}{m_s} \text{ TeV} \le m_{B_2} \le 2 \times 10^2 \frac{m_b}{m_s} \text{ TeV},
\end{aligned}
$$

In the above described range of parameters, we find that the minimization of the χ^2 function yields the following benchmark point, consistent with the experimental values of the SM quark masses and CKM parameters:

$$
\begin{aligned}
&\theta \simeq 85.9^\circ, \qquad m_{S_1} = m_{P_1} \simeq 1.9 \text{ TeV}, \qquad m_{S_2} = m_{P_2} \simeq 2.1 \text{ TeV}, \\
&v_L \simeq 100 \text{ GeV}, \qquad v_R \simeq 10 \text{ TeV}, \\
&m_T \simeq 583 \text{ TeV}, \qquad m_{T'} \simeq 1.1 \times 10^3 \text{ TeV}, \qquad m_{B_1} \simeq 216 \text{ TeV}, \\
&m_{B_2} \simeq 9.3 \times 10^3 \text{ TeV}, \qquad m_{B'} \simeq 396 \text{ TeV}, \\
&x_1^{(T)} \simeq 0.24 - 0.02i, \qquad x_2^{(T)} \simeq 0.96 - 0.06i, \qquad z_1^{(T)} = z_2^{(T)} \simeq -0.16 + 0.08i, \\
&x_{12}^{(B)} \simeq -0.05 - 0.03i, \qquad x_{22}^{(B)} \simeq -0.62 - 0.06i, \qquad x_3^{(B)} \simeq 0.07 - 0.62i, \\
&z_{11}^{(B)} \simeq 0.25, \qquad z_{12}^{(B)} \simeq 0.49, \qquad z_{13}^{(B)} \simeq -0.44, \\
&z_{21}^{(B)} \simeq -1.17i, \qquad z_{22}^{(B)} \simeq 0.95i, \qquad z_{23}^{(B)} \simeq 0.80i, \\
&w_1^{(T')} \simeq -0.39 + 0.197i, \qquad w_2^{(T')} \simeq -0.58 + 0.29i, \\
&r_1^{(T')} \simeq -0.154 + 0.084i, \qquad r_2^{(T')} \simeq -0.875 + 0.48i, \qquad r_3^{(T')} \simeq 0.30 - 0.16i, \\
&w_1^{(B')} \simeq 0.12 - 0.087i, \qquad w_2^{(B')} \simeq 0.44 + 0.74i, \\
&r_1^{(B')} \simeq -0.17 - 0.98i, \qquad r_2^{(B')} \simeq -1.22, \qquad r_3^{(B')} \simeq 1.299 + 0.698i . \qquad (26)
\end{aligned}
$$

With the above particle content, the following relevant Yukawa terms arise:

$$
\begin{aligned}
-\mathcal{L}_Y =& \sum_{i=1}^{3} \alpha_i \overline{Q}_{3L} \Phi Q_{iR} + \sum_{n=1}^{2} x_n^{(T)} \overline{Q}_{nL} \widetilde{\chi}_L T_R + \sum_{i=1}^{3} z_i^{(T)} \overline{T}_L \widetilde{\chi}_R^{\dagger} Q_{iR} \\
&+ \sum_{n=1}^{2} w_n^{(T')} \overline{Q}_{nL} \widetilde{\phi}_L T'_R + \sum_{i=1}^{3} r_i^{(T')} \bar{T}'_L \widetilde{\phi}_R^{\dagger} Q_{iR} \\
&+ x_3^{(B)} \overline{Q}_{3L} \chi_L B_{1R} + \sum_{n=1}^{2} x_{n2}^{(B)} \overline{Q}_{nL} \chi_L B_{2R} + \sum_{n=1}^{2} \sum_{i=1}^{3} z_{ni}^{(B)} \overline{B}_{nL} \chi_R^{\dagger} Q_{iR} \\
&+ \sum_{n=1}^{2} w_n^{(B')} \overline{Q}_{nL} \phi_L B'_R + \sum_{i=1}^{3} r_i^{(B')} \bar{B}'_L \phi_R^{\dagger} Q_{iR} \\
&+ y_T \overline{T}_L \sigma T_R + y_{T'} \bar{T}'_L \sigma T'_R + y_{B_1} \overline{B}_{1L} \rho B_{1R} + y_{B_2} \overline{B}_{2L} \sigma B_{2R} + y_{B'} \bar{B}'_L \sigma B'_R \\
&+ \sum_{n=1}^{2} y_{E_n} \overline{E}_{nL} \rho E_{nR} + y_{E'} \bar{E}'_L \rho E'_R + \sum_{i=1}^{3} \sum_{n=1}^{2} x_{in}^{(E)} \overline{L}_{iL} \chi_L E_{nR} \\
&+ \sum_{n=1}^{2} \sum_{i=1}^{3} z_{nj}^{(E)} \overline{E}_{nL} \chi_R^{\dagger} L_{jR} + \sum_{i=1}^{3} w_i^{(E')} \overline{L}_{iL} \phi_L E'_R + \sum_{i=1}^{3} r_i^{(E')} \bar{E}'_L \phi_R^{\dagger} L_{iR} \\
&+ \sum_{i=1}^{3} \sum_{j=1}^{3} y_{ij}^{(L)} \overline{L}_{iL} \Phi L_{jR} + \sum_{i=1}^{3} \sum_{j=1}^{3} x_{ij}^{(N)} \overline{N_{iR}^{C}} \widetilde{\chi}_R^{\dagger} L_{jR} + \sum_{n=1}^{2} (y_{\Omega})_n \overline{\Omega}_{nR} \Omega_{nR}^{C} \eta \\
&+ \sum_{i=1}^{3} \sum_{k=1}^{2} x_{ik}^{(S)} \overline{N}_{iR} \Omega_{kR}^{C} \varphi + H.c. \qquad (12)
\end{aligned}
$$