Lepton Flavor Violating Higgs Boson Decay in G2HDM

Yu-Ming Chen

Advisor: Prof. Chia-Hung Chang

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Department of Physics

National Taiwan Normal University

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Abstract

We study lepton flavor violating Higgs boson decay process in the gauged two Higgs doublet model G2HDM. The model introduce an additional $SU(2)_H$ gauged symmetry and combine two Higgs doublet H_1 and H_2 into a $SU(2)_H$ doublet H. We examine the neutrino mass mechanisms, including Dirac Type or Majorana type. In the later neutrino mass is achieved by the type-II seesaw mechanism with an additional $SU(2)_H$ triplet Δ_N. Δ_N is included here in both the scalar potential and scalar mass spectrum. We present six Feynman diagrams of lepton flavor violating Higgs decays and calculate two classical diagrams among them. It is shown W' exchanging diagrams are the largest and we discuss how the amplitude will be determined by various parameters. In typical situation the branching ratio is in the order around 10^{-10} or smaller, probably beyond the ability of the Large Hadron Collider measurements to detect. However in some extreme case of parameters, especially large Yukawa couplings, the branching ratio could be as large as 10^{-6} and the H_2 exchanging diagrams could become more important.

Keywords : Higgs Physics, G2HDM, Neutrino Mass, Lepton Flavor Violation

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I Introduction

It has been several years since the historical discovery of the 125 GeV Higgs boson [1, 2]. While the observed Higgs boson fits its profile in Standard Model (SM) quite well, physicists are still contemplating whether it acts on its own or belongs to a larger structure which points to a theory beyond the SM. Among possible extensions of the SM, the two Higgs doublet model (2HDM) is one of the simplest options and also a very attractive one [3]. Here an additional Higgs doublet is added to pair up with the lonely Higgs boson. It is worth noting that this scenario arises naturally in the supersymmetric extension of SM, which could elegantly cure the hierarchy problem of SM. With the additional Higgs doublet, some of the long running problems, for example strong CP problem, baryogenesis, in particle phsyics could also be resolved. On the other hand, in order to eliminate tree level flavor changing neutral current, the vacuum expectation value (VEV) of the two Higgs doublets need to be carefully aligned. One elegant way to accomplish this is to impose a Z_2 symmetry and assign one of the Higgs doublet odd under Z_2 so as to render the VEV of this Higgs doublet zero. This model, sometimes called inert two Higgs doublet model (IHDM), has the further benefit that the Z_2 odd Higgs doublet could supply a Dark Matter (DM) candidate.

It was recently realised [4] that we can accomplish the same VEV assignment by introducing an additional gauged $SU(2)_H$ symmetry acting between the two Higgs doublets. In other words, the two Higgs doublets H_1 and H_2 are themselves $SU(2)_H$ doublet $H = \begin{pmatrix} H_1 & H_2 \end{pmatrix}^{\mathrm{T}}$. With the additional $SU(2)_H$ symmetry, it is just natural to rotate the VEV of the Higgs field into the first doublet and leave the VEV of the second doublet zero, in a configuration very similar to IHDM, but now a natural result of the additional gauged symmetry. The neutral component of H_2 could then be a DM candidate, its stability protected by $SU(2)_H$ gauge symmetry and Lorentz invariance. The authors in [4] named it as Gauged two Higgs Doublet Model (G2HDM) and they further embed the idea in a very elegant setting. The fermion Dirac masses are generated by extending the right handed SM fermions such as quarks and leptons into $SU(2)_H$ doublets. For example, right-handed electron will be part of a $SU(2)_H$ doublet ($SU(2)_L$ singlet): $E_R = \begin{pmatrix} e_R^H & e_R \end{pmatrix}^{\mathrm{T}}$. The heavy fermions such as e_R^H and the gauge bosons of $SU(2)_H$, called W', Z', in general need to be heavier than SM particles and hence a new scale is necessary. The easiest way is to introduce a $SU(2)_H$ doublet Φ_H, similar to the Higgs doublet in SM, with nonzero VEV to facilitate heavy fermions' Yukawa coupling. The VEV v_Φ then set the scale of heavy fermions and heavy gauge bosons. The authors in [4] employ a $SU(2)_H$ triplet Δ_H to induce $SU(2)_L$ and $SU(2)_H$ symmetry breaking simultaneously in a very crafted way.

From neutrino oscillation experiments, we now know that neutrinos have an extremely small mass $\sim 0.1\text{eV}$, about 10^6 times lighter than an electron. In G2HDM, with these novel fermions mentioned above, the possible mechanisms to generate neutrino mass are rich. It could be the Dirac mechanism which preserves lepton number. The $SU(2)_L$ doublet left-handed lepton $L_L = (\nu_L \;\; e_L)^{\mathrm{T}}$, $SU(2)_H$ doublet right-handed neutrino $N_R = \left(\nu_R \;\; \nu_R^H\right)^{\mathrm{T}}$ and the Higgs field H, doublet under both $SU(2)$s, could form Yukawa coupling. After electroweak symmetry breaking, the VEV of H gives rise to a Dirac mass between left-handed neutrino and right-handed neutrino, the same way quarks acquire masses. The feature of this mechanism is that the Yukawa coupling is extremely small compared with others due to the small neutrino mass splitting observed. On the other hand, If we are allowed to introduce an additional $SU(2)_H$ triplet Δ_N, it could give rise to a Majorana mass of right-handed neutrino. As a consequence, the seesaw mechanism could be arranged. The seesaw mechanism could naturally avoid extremely small coupling constants but would break lepton number conservation and a neutrino-less double beta decay signal is indicated.

While right-handed fermions such as ν_R pair up with heavy fermions ν_R^H to form doublet, it is clear that heavy femion flavor structure may not be consistent with SM particle flavor. That is, if we put in the flavor indices, the heavy fermion $\nu_R^{H,j}$ that pairs with a ν_R^i mass eigenstate may not be a mass eigenstate itself. Consequently a loop containing heavy fermion may be flavor-changing and generate charged lepton flavor violation (CLFV) processes in G2HDM. It has recently be pointed out that the experimental data of CLFV Higgs decay, especially $H \to \tau\mu$ may be as stringent as we expected. Hence we would like to study the process of CLFV Higgs decay in G2HDM in this paper. As we have mentioned, G2HDM introduces several scalars: a $SU(2)_H$ doublet Φ_H, a $SU(2)_H$ triplet Δ_H and a possible Δ_N if seesaw mechanism is preferred. These three scalars will have complicated mixing with H_1 and H_2. The observed 125 GeV Higgs boson is the mixture of the neutral component of H_1, Φ_H, Δ_H and Δ_N. The scalar Φ_H couple to heavy fermion and hence heavy fermion loop, on that account, the Higgs boson could potentially decay into flavor non-conserving lepton pairs. We construct CLFV processes in G2HDM and study their magnitudes.

The paper is organized as follows: in Sec. II, we briefly review the G2HDM setup with an additional Δ_N, including the matter content, the Higgs potential, the scalar spectrum, and the Yukawa interaction. In Sec. III, we scrutinize the neutrino mass generating mechanism through type-II seesaw. In Sec. IV, we examine two most significant diagrams of CLFV in G2HDM and calculate its amplitude. Finally, we conclude our findings in Sec. V.

II Gauged Two-Higgs-Doublet Model

G2HDM solved the vacuum alignment problem of two Higgs doublet model elegantly with the introduction of an additional gauged $SU(2)_H$ symmetry, acting horizontally between the two Higgs doublets. The $SU(2)_H$ can then rotate the nonzero VEV into one of the doublet. It has many intriguing features to explore. In this section, we review the content of this model, especially the particle assignment, and specify the slight difference in matter content, scalar potential and mass spectrum with an addtional Δ_N, if we introduce a seesaw mechanism for neutrino mass. Also, we write down the Yukawa interaction between scalar fields, SM fermions and heavy fermions. This will be relevant for our discussion of neutrino mass mechanism and lepton flavor violating Higgs decay.

A. Matter Content

The gauged symmetry group of G2HDM is $SU(2)_L \times U(1)_Y \times SU(2)_H \times U(1)_X$. Two $SU(2)_L$ doublets, H_1 and H_2, comprise a $SU(2)_H$ doublet $H = (H_1 \;\; H_2)^{\mathrm{T}}$.

The original $SU(2)_L$ doublets are singlet under $SU(2)_H$. But new heavy fermions can be introduced in various ways. The original paper of G2HDM [4] chooses the simplest extension. Heavy right-handed fermions are introduced to pair up with SM right-handed fermions to form $SU(2)_H$ doublets. To automatically cancel anomaly, heavy left-handed fermions are added, which are singlet under both $SU(2)$s. Extra scalar fields are introduced for various purposes. Only H carries quantum number in $SU(2)_L$, while other scalars are singlet under $SU(2)_L$ but transform under $SU(2)_H$. An $SU(2)_H$ doublet Φ_H is introduced to give masses to heavy gauged boson and heavy fermions. Hence its VEV v_Φ mostly specify the scale of the masses of extra particles. An $SU(2)_H$ triplet Δ_H develops a non-zero VEV and trigger through an elegant setting simultaneously the breaking of $SU(2)_L \times U(1)_Y \times SU(2)_H \times U(1)_X$ into $U(1)_EM$ by giving rise to nonzero VEV's to H_1 and Φ_H [4]. The detailed study of the scalar potential has been done in [5]. We will briefly review it later. In order to facilitate the Majorana mass term, we introduce yet another $SU(2)_H$ triplet Δ_N with X charge -2 as suggested in [4].

In order to get rid of some superfluous terms in scalar potential, all matter fields are also gauged in an additional $U(1)_X$ group. For example, a term like $\Phi_H^{\mathrm{T}} \Delta_H \Phi_H$ would be excluded. Therefore, the entire gauge group of G2HDM is composed of $SU(3)_C \times SU(2)_L \times U(1)_Y \times SU(2)_H \times U(1)_X$. There will be vector bosons W', Z', X' corresponding to these gauge groups, whose properties has been thoroughly studied in [6]. The complete matter content of G2HDM is listed in Table 1, we use the symbols denoted in [5].

Some features of the model are: (i) it is free of gauge anomaly; (ii) the additional Z_4 symmetry is not required to achieve an inert Higgs doublet, $SU(2)_H$ can naturally realize it; (iii) due to the gauge symmetries, the tree-level CLFV processes are forbidden for the SM sector; (iv) the triplet Δ_H induces SM Higgs VEV, breaking $SU(2)_L \times U(1)_Y$ down to $U(1)_Q$.

Matter Fields	$SU(3)_C$	$SU(2)_L$	$SU(2)_H$	$U(1)_Y$	$U(1)_X$
$Q_L = (u_L \;\; d_L)^{\mathrm{T}}$	3	2	1	1/6	0
$U_R = \left(u_R \;\; u_R^H\right)^{\mathrm{T}}$	3	1	2	2/3	1
$D_R = \left(d_R^H \;\; d_R\right)^{\mathrm{T}}$	3	1	2	$-1/3$	-1
u_L^H	3	1	1	2/3	0
d_L^H	3	1	1	$-1/3$	0
$L_L = (\nu_L \;\; e_L)^{\mathrm{T}}$	1	2	1	$-1/2$	0
$N_R = \left(\nu_R \;\; \nu_R^H\right)^{\mathrm{T}}$	1	1	2	0	1
$E_R = \left(e_R^H \;\; e_R\right)^{\mathrm{T}}$	1	1	2	-1	-1
ν_L^H	1	1	1	0	0
e_L^H	1	1	1	-1	0
$H = (H_1 \;\; H_2)^{\mathrm{T}}$	1	2	2	1/2	1
$\Phi_H = (\Phi_1 \;\; \Phi_2)^{\mathrm{T}}$	1	1	2	0	1
$\Delta_H = \begin{pmatrix} \Delta_3/2 & \Delta_p/\sqrt{2} \\ \Delta_m/\sqrt{2} & -\Delta_3/2 \end{pmatrix}$	1	1	3	0	0
$\Delta_N = \begin{pmatrix} \Delta_N^+/\sqrt{2} & \Delta_N^{++} \\ \Delta_N^0 & -\Delta_N^+/\sqrt{2} \end{pmatrix}$	1	1	3	0	-2

Table 1: Matter field contents and their quantum number assignments in G2HDM.

B. Scalar Potential

The invariance Higgs potential under all gauge groups can be separated into five components,

$$\mathrm{V_T} = \mathrm{V}(H) + \mathrm{V}(\Phi_H) + \mathrm{V}(\Delta_H) + \mathrm{V}(\Delta_N) + \mathrm{V_{mix}}(H, \Phi_H, \Delta_H, \Delta_N) \tag{1}$$

, with

$$
\begin{aligned}
\mathrm{V}(H) &= \mu_H^2(H^{\alpha i}H_{\alpha i}) + \lambda_H(H^{\alpha i}H_{\alpha i})^2 + \frac{1}{2}\lambda'_H\epsilon_{\alpha\beta}\epsilon^{\gamma\delta}(H^{\alpha i}H_{\gamma i})(H^{\beta j}H_{\delta j}) \\
&= \mu_H^2(H_1^\dagger H_1 + H_2^\dagger H_2) + \lambda_H(H_1^\dagger H_1 + H_2^\dagger H_2)^2 + \lambda'_H(-H_1^\dagger H_1 H_2^\dagger H_2 + H_1^\dagger H_2 H_2^\dagger H_1)
\end{aligned} \tag{2}
$$

, where $(\alpha,\beta,\gamma,\delta)$ and (i,j) refer to $SU(2)_H$ and $SU(2)_L$ indices respectively, and $H^{\alpha i} = H^*_{\alpha i}$;

$$
\begin{aligned}
\mathrm{V}(\Phi_H) &= \mu_\Phi^2\Phi_H^\dagger\Phi_H + \lambda_\Phi(\Phi_H^\dagger\Phi_H)^2 \\
&= \mu_\Phi^2(\Phi_1^*\Phi_1 + \Phi_2^*\Phi_2) + \lambda_\Phi(\Phi_1^*\Phi_1 + \Phi_2^*\Phi_2)^2
\end{aligned} \tag{3}
$$

, with $\Phi_H = (\Phi_1 \;\; \Phi_2)^{\mathrm{T}}$

$$
\begin{aligned}
\mathrm{V}(\Delta_H) &= -\mu_\Delta^2\mathrm{Tr}(\Delta_H^2) + \lambda_\Delta(\mathrm{Tr}(\Delta_H^2))^2 \\
&= -\mu_\Delta^2(\frac{1}{2}\Delta_3^2 + \Delta_p\Delta_m) + \lambda_\Delta(\frac{1}{2}\Delta_3^2 + \Delta_p\Delta_m)^2
\end{aligned} \tag{4}
$$

, where

$$
\Delta_H = \begin{pmatrix} \Delta_3/2 & \Delta_p/\sqrt{2} \\ \Delta_m/\sqrt{2} & -\Delta_3/2 \end{pmatrix} = \Delta_H^\dagger
$$

, with

$$
\Delta_m = (\Delta_p)^* \;\text{ and }\; \Delta_3 = (\Delta_3)^*
$$

; and

$$
\begin{aligned}
\mathrm{V}(\Delta_N) &= \mu_N^2\mathrm{Tr}(\Delta_N^2) + \lambda_N(\mathrm{Tr}(\Delta_N^2))^2 \\
&= \mu_N^2[\Delta_N^+\Delta_N^- + \Delta_N^{++}\Delta_N^{--} + (\Delta_N^0)^2] + \lambda_N[\Delta_N^+\Delta_N^- + \Delta_N^{++}\Delta_N^{--} + (\Delta_N^0)^2]^2
\end{aligned} \tag{5}
$$

, with

$$
\Delta_N = \begin{pmatrix} \frac{1}{\sqrt{2}}\Delta_N^+ & \Delta_N^{++} \\ \Delta_N^0 & -\frac{1}{\sqrt{2}}\Delta_N^+ \end{pmatrix}
$$

Finally, the mixed term

$$
\begin{aligned}
\mathrm{V}_{\mathrm{mix}} = &+ M_{H\Delta}(H^\dagger\Delta_H H) - M_{\Phi\Delta}(\Phi_H^\dagger\Delta_H\Phi_H) \\
&+ \lambda_{H\Phi}(H^\dagger H)(\Phi_H^\dagger\Phi_H) + \lambda'_{H\Phi}(H^\dagger\Phi_H)(\Phi_H^\dagger H) \\
&+ \lambda_{H\Delta}(H^\dagger H)\mathrm{Tr}(\Delta_H^2) + \lambda_{\Phi\Delta}(\Phi_H^\dagger\Phi_H)\mathrm{Tr}(\Delta_H^2) \\
&+ \lambda_{HN}(H^\dagger H)\mathrm{Tr}(\Delta_N^\dagger\Delta_N) + \lambda_{\Phi N}(\Phi_H^\dagger\Phi_H)\mathrm{Tr}(\Delta_N^\dagger\Delta_N) \\
&+ \lambda_{\Delta N}\mathrm{Tr}(\Delta_H^2)\mathrm{Tr}(\Delta_N^\dagger\Delta_N)
\end{aligned} \tag{6}
$$

Expanding the mixed potential terms with the component fields gives

$$
\begin{aligned}
\mathrm{V}_{\mathrm{mix}} = &+ M_{H\Delta}(\frac{1}{\sqrt{2}}H_1^\dagger H_2\Delta_p + \frac{1}{2}H_1^\dagger H_1\Delta_3 + \frac{1}{\sqrt{2}}H_2^\dagger H_1\Delta_m - \frac{1}{2}H_2^\dagger H_2\Delta_3) \\
&- M_{\Phi\Delta}(\frac{1}{\sqrt{2}}\Phi_1^*\Phi_2\Delta_p + \frac{1}{2}\Phi_1^*\Phi_1\Delta_3 + \frac{1}{\sqrt{2}}\Phi_2^*\Phi_1\Delta_m - \frac{1}{2}\Phi_2^*\Phi_2\Delta_3) \\
&+ \lambda_{H\Phi}(H_1^\dagger H_1 + H_2^\dagger H_2)(\Phi_1^*\Phi_1 + \Phi_2^*\Phi_2) \\
&+ \lambda'_{H\Phi}(H_1^\dagger H_1\Phi_1^*\Phi_1 + H_2^\dagger H_2\Phi_2^*\Phi_2 + H_1^\dagger H_2\Phi_2^*\Phi_1 + H_2^\dagger H_1\Phi_1^*\Phi_2) \\
&+ \lambda_{H\Delta}(H_1^\dagger H_1 + H_2^\dagger H_2)(\frac{1}{2}\Delta_3^2 + \Delta_p\Delta_m) \\
&+ \lambda_{\Phi\Delta}(\Phi_1^*\Phi_1 + \Phi_2^*\Phi_2)(\frac{1}{2}\Delta_3^2 + \Delta_p\Delta_m) \\
&+ \lambda_{HN}(H_1^\dagger H_1 + H_2^\dagger H_2)[\Delta_N^+\Delta_N^- + \Delta_N^{++}\Delta_N^{--} + (\Delta_N^0)^2] \\
&+ \lambda_{\Phi N}(\Phi_1^*\Phi_1 + \Phi_2^*\Phi_2)[\Delta_N^+\Delta_N^- + \Delta_N^{++}\Delta_N^{--} + (\Delta_N^0)^2] \\
&+ \lambda_{\Delta N}(\frac{1}{2}\Delta_3^2 + \Delta_p\Delta_m)[\Delta_N^+\Delta_N^- + \Delta_N^{++}\Delta_N^{--} + (\Delta_N^0)^2]
\end{aligned} \tag{7}
$$

From here, we would like to make some notes about the model characteristics:

(i) $U(1)_X$ gauges out some unwanted terms, for example $\Phi_H^{\mathrm{T}}\Delta_H\Phi$. To exclude a global symmetry, $U(1)_X$ is treated as a local symmetry.

(ii) The quadratic terms for H_1 and H_2 have the following coefficients respectively,

$$\mu_H^2 - \frac{1}{2}M_{H\Delta}\cdot v_\Delta + \frac{1}{2}\lambda_{H\Delta}\cdot v_\Delta^2 + \frac{1}{2}\lambda_{H\Phi}\cdot v_\Phi^2 + \lambda_{HN}\cdot v_N^2 \tag{8}$$

$$\mu_H^2 + \frac{1}{2}M_{H\Delta}\cdot v_\Delta + \frac{1}{2}\lambda_{H\Delta}\cdot v_\Delta^2 + \frac{1}{2}(\lambda_{H\Phi} + \lambda'_{H\Phi})\cdot v_\Phi^2 + \lambda_{HN}\cdot v_N^2 \tag{9}$$

Notice that the five parameters $M_{H\Delta}$, $\lambda_{H\Delta}$, $\lambda_{H\Phi}$, $\lambda'_{H\Phi}$, λ_{HN} can take on either positive or negative value. Even with positive μ_H^2, one can still achieve $\langle H_1\rangle \neq 0$ and $\langle H_2\rangle = 0$ to break $SU(2)_L$.

(iii)The same situation for H can be applied to Φ_H, the coefficients of quadratic Φ_1, Φ_2 terms are

$$\mu_\Phi^2 - \frac{1}{2}M_{\Phi\Delta}\cdot v_\Delta + \frac{1}{2}\lambda_{\Phi\Delta}\cdot v_\delta^2 + \frac{1}{2}(\lambda_{H\Phi} + \lambda'_{H\Phi})\cdot v^2 + \lambda_{\Phi N}\cdot v_N^2 \tag{10}$$

$$\mu_\Phi^2 + \frac{1}{2}M_{\Phi\Delta}\cdot v_\Delta + \frac{1}{2}\lambda_{\Phi\Delta}\cdot v_\delta^2 + \frac{1}{2}\lambda_{H\Phi}\cdot v^2 + \lambda_{\Phi N}\cdot v_N^2 \tag{11}$$

respectively. Even with positive μ_Φ^2, we can still achieve $\langle\Phi_1\rangle = 0$ and $\langle\Phi_2\rangle \neq 0$.

(iv)In Eq. (4), if $-\mu_\Delta < 0$, $SU(2)_H$ is spontaneously breaking by $\langle\Delta_H\rangle = -v_\Delta \neq 0$ with $\langle\Delta_{p,m}\rangle = 0$. This also trigger the symmetry breaking of other gauge groups.

(v)The whole scalar potential is Hermitian. As a consequence, G2HDM is CP-conserving.

(vi)The triplet Δ_N will have the mixing effect with other scalars as expected. This will also appear in the following discussion.

C. Spontaneous Symmetry Breaking

The scalar fields can be shifted as follows to facilitate spontaneous symmetry breaking (SSB)

$$
\begin{aligned}
&H_1 = \begin{pmatrix} G^+ \\ \frac{v+h}{\sqrt{2}} + i\frac{G^0}{\sqrt{2}} \end{pmatrix}, H_2 = \begin{pmatrix} H^+ \\ H_2^0 \end{pmatrix} \\
&\Phi_H = \begin{pmatrix} G_H^p \\ \frac{v_\Phi+\phi_2}{\sqrt{2}} + i\frac{G_H^0}{\sqrt{2}} \end{pmatrix}, \Delta_H = \begin{pmatrix} \frac{-v_\Delta+\delta_3}{2} & \frac{1}{\sqrt{2}}\Delta_p \\ \frac{1}{\sqrt{2}}\Delta_m & \frac{v_\Delta-\delta_3}{2} \end{pmatrix}, \quad \Delta_N = \begin{pmatrix} \frac{1}{\sqrt{2}}G_N^+ & G_N^{++} \\ v_N+\delta_N & -\frac{1}{\sqrt{2}}G_N^+ \end{pmatrix}
\end{aligned} \tag{12}
$$

, where v, v_Φ, v_Δ, v_N are VEVs of H, Φ_H, Δ_H, Δ_N respectively, and G^+, G^0, G_H^p, G_H^0, G_N^+, G_N^{++} are goldstone bosons (GB).

Substituting the VEVs into the potential yields

$$
\begin{aligned}
\mathrm{V_T}(v, v_\Phi, v_\Delta, v_N) =& \frac{1}{4}[\lambda_H v^4 + \lambda_\Phi v_\Phi^4 + \lambda_\Delta v_\Delta^4 + 4\lambda_N V_N^4 \\
&+ 2(\mu_H^2 v^2 + \mu_\Phi^2 v_\Phi^2 - \mu_\Delta^2 v_\Delta^2 + 2\mu_N^2 v_N^2) \\
&- (M_{H\Delta} v^2 + M_{\Phi\Delta} v_\Phi^2) v_\Delta + \lambda_{H\Phi} v^2 v_\Phi^2 + \lambda_{H\Delta} v^2 v_\Delta^2 + \lambda_{\Phi\Delta} v_\Phi^2 v_\Delta^2 \\
&+ 2v_N^2(\lambda_{HN} v^2 + \lambda_{\Phi N} v_\Phi^2 + \lambda_{\Delta N} v_\Delta^2)]
\end{aligned} \tag{13}
$$

Notice that λ'_H and $\lambda'_{H\Phi}$ do not appear in Eq. (13), but there are additional terms correlated with Δ_N. The relation between VEVs can be obtained by performing partial derivative of the potential with respect to VEVs.

$$2\lambda_H v^2 + 2\mu_H^2 - M_{H\Delta} v_\Delta + \lambda_{H\Phi} v_\Phi^2 + \lambda_{H\Delta} v_\Delta^2 + 2\lambda_{HN} v_N^2 = 0 \tag{14}$$

$$2\lambda_\Phi v_\Phi^2 + 2\mu_\Phi^2 - M_{\Phi\Delta} v_\Delta + \lambda_{H\Phi} v^2 + \lambda_{\Phi\Delta} v_\Delta^2 + 2\lambda_{\Phi N} v_N^2 = 0 \tag{15}$$

$$4\lambda_\Delta v_\Delta^3 - 4\mu_\Delta^2 v_\Delta - M_{H\Delta} v^2 - M_{\Phi\Delta} v_\Phi^2 + 2v_\Delta(\lambda_{H\Delta} v^2 + \lambda_{\Phi\Delta} v_\Phi^2 + 2\lambda_{\Delta N} v_N^2) = 0 \tag{16}$$

$$4\lambda_N v_N^2 + 2\mu_N^2 + \lambda_{HN} v^2 + \lambda_{\Phi N} v_\Phi^2 + \lambda_{\Delta N} v_\Delta^2 = 0 \tag{17}$$

Despite there are some differences with [5], one can still solve for non-trivial v^2, v_Φ^2 and v_N^2 in terms of v_Δ and other parameters. Then, v_Δ is breaking the SM $SU(2)_L \times U(1)_Y$ and the $U(1)_X$, after breaking the $SU(2)_H$. Since Eq. (16) is cubic in terms of v_Δ, the solution of v_Δ could be either 1 real, 2 complex or 3 real. The complex soltion is unphysical, since G2HDM is CP-conseving. In the case of 3 real solutions, the one which consistent with SM Higgs VEV $v = 246$ GeV will be chosen.

D. Mass Spectrum of Scalars

One can obtain the scalar mass spectrum by inspecting the quartic terms of fields. The scalar mass spectrum contains three blocks. Due to the extra scalar, the first block is now 4×4. In the flavor basis of $S=\{h,\phi_2,\delta_3,\delta_N\}$, the matrix is given as.

$$\mathcal{M}_0^2=\begin{pmatrix} \lambda_H v^2 & \lambda_{H\Phi}vv_\Phi & \frac{1}{2}M_{H\Delta}v-\lambda_{H\Delta}vv_\Delta & 2\lambda_{HN}vv_N \\ \lambda_{H\Phi}vv_\Phi & \lambda_\Phi v_\Phi^2 & \frac{1}{2}M_{\Phi\Delta}v_\Phi-\lambda_{\Phi\Delta}v_\Phi v_\Delta & 2\lambda_{\Phi N}v_\Phi v_N \\ \frac{1}{2}M_{H\Delta}v-\lambda_{H\Delta}vv_\Delta & \frac{1}{2}M_{\Phi\Delta}v_\Phi-\lambda_{\Phi\Delta}v_\Phi v_\Delta & \lambda_\Delta v_\Delta^2-\frac{1}{8v_\Delta}(M_{H\Delta}v^2+M_{\Phi\Delta}v_\Phi^2) & -2\lambda_{\Delta N}v_\Delta v_N \\ 2\lambda_{HN}vv_N & 2\lambda_{\Phi N}v_\Phi v_N & -2\lambda_{\Delta N}v_\Delta v_N & 4\lambda_N v_N^2 \end{pmatrix} \tag{18}$$

The matrix can be diagonalized by a similarity transformation with an orthogonal matrix. O, defining as $|f\rangle_\alpha\equiv O_{\alpha\beta}|m\rangle_\beta$ with $|f\rangle$ and $|m\rangle$ referring to the flavor and mass eigenstates respectively,

$$O^T\cdot\mathcal{M}_0^2\cdot O=\mathrm{Diag}(m_{h_1}^2,m_{h_2}^2,m_{h_3}^2,m_{h_4}^2) \tag{19}$$

The lightest eigenstate will be the observed 125 GeV Higgs boson. The physical Higgs h_α is a linear combination of the four components of S: $h_\alpha=O_{\beta\alpha}h_\beta$. For instance, the observed 125 GeV Higgs boson could be a composite of the neutral components of H_1 and $SU(2)_H$ doublet Φ_H, together with the real components of the $SU(2)_H$ triplets Δ_H and Δ_N.

The second block is in the basis of $G=\{G_H^p,H_2^{0*},\Delta_p\}$

$$\mathcal{M}_0'^2=\begin{pmatrix} M_{\Phi\Delta}v_\Delta+\frac{1}{2}\lambda'_{H\Phi}v^2 & \frac{1}{2}\lambda'_{H\Phi}vv_\Phi & -\frac{1}{2}M_{\Phi\Delta}v_\Phi \\ \frac{1}{2}\lambda'_{H\Phi}vv_\Phi & M_{H\Delta}v_\Delta+\frac{1}{2}\lambda'_{H\Phi}v_\Phi^2 & \frac{1}{2}M_{H\Delta}v \\ -\frac{1}{2}M_{\Phi\Delta}v_\Phi & \frac{1}{2}M_{H\Delta}v & \frac{1}{4v_\Delta}(M_{H\Delta}v^2+M_{\Phi\Delta}v_\Phi^2) \end{pmatrix} \tag{20}$$

When we diagonalize to mass eigenstate, the GB acquires zero mass, while the other two are physical fields D and $\tilde{\Delta}$. The field D could be the DM candidate in G2HDM as well as ν_R, ν_R^H, ν_L^H and $W'^{(p,m)}$, determined by their indivitual mass.

The final block is diagonlaized, including physical charged Higgs $H^\pm$, where

$$m_{H^\pm}^2=M_{H\Delta}v_\Delta-\frac{1}{2}\lambda'_H v^2+\frac{1}{2}\lambda'_{H\Phi}v_\Phi^2 \tag{21}$$

, and five GBs $G^\pm$, G^0, G_H^0, G_N^+, G_N^{++}. Note that the minimization conditions Eqs. (14)-(17) have been used to simplify various matrix elements. Despite a little difference in the first block and extra Goldstone bosons in the third block, the rest are the same as [5].

E. Yukawa Coupling

As described above, the quark $SU(2)_L$ doublet Q_L is singlet under $SU(2)_H$, while $SU(2)_L$ singlet u_R and d_R pair up with u_R^H and d_R^H respectively to form a $SU(2)_H$ doublet, for example, $U_R^{\mathrm{T}} = \begin{pmatrix} u_R^H & u_R \end{pmatrix}^{\mathrm{T}}$. Therefore, the Lagrangian for the Yukawa coupling is

$$
\begin{aligned}
\mathcal{L}_{\text{Yuk}} &\supset y_d \bar{Q}_L (D_R \cdot H) + y_u \bar{Q}_L (U_R \cdot \tilde{\tilde{H}}) + \text{H.c.} \\
&= y_d \bar{Q}_L (d_R^H H_2 - d_R H_1) - y_u \bar{Q}_L (u_R \tilde{H}_1 + u_R^H \tilde{H}_2) + \text{H.c.},
\end{aligned}
\tag{22}
$$

where $\tilde{\tilde{H}} \equiv \begin{pmatrix} \tilde{H}_2 & -\tilde{H}_1 \end{pmatrix}^{\mathrm{T}}$ with $\tilde{H}_{1,2} = i\tau_2 H_{1,2}^*$. $A \cdot B$ is defined as $\epsilon_{ij} A^i B^j$, where A and B are two-dimensional spinor representations of $SU(2)_H$. After SSB, $\langle H_1 \rangle \neq 0$, u and d acquire their masses, but $\langle H_2 \rangle = 0$, meaning u_R^H and d_R^H are still massless.

To give a mass to the heavy fermions, D2HDM uses the $SU(2)_H$ scalar doublet $\Phi_H = \begin{pmatrix} \Phi_1 & \Phi_2 \end{pmatrix}^{\mathrm{T}}$ together with $SU(2)_{L,H}$ singlets u_L^H and d_L^H to write additional Yukawa couplings:

$$
\begin{aligned}
\mathcal{L}_{\text{Yuk}} &\supset -y_d' \bar{d}_L^H (D_R \cdot \Phi_H) + y_u' \bar{u}_L^H (U_R \cdot \tilde{\Phi}_H) + \text{H.c.} \\
&= -y_d' \bar{d}_L^H (d_R^H \Phi_2 - d_R \Phi_1) - y_u' \bar{u}_L^H (u_R \Phi_1^* + u_R^H \Phi_2^*) + \text{H.c.}
\end{aligned}
\tag{23}
$$

, where $\tilde{\Phi}_H = \begin{pmatrix} \Phi_2^* & -\Phi_1^* \end{pmatrix}^{\mathrm{T}}$. With $\langle \Phi_2 \rangle = v_\Phi/\sqrt{2}$, u_R^H and u_L^H obtain a mass $y_u' v_\Phi/\sqrt{2}$; also, d_R^H and d_L^H obtain a mass $y_d' v_\Phi/\sqrt{2}$.

The Yukawa couplings in the lepton sector is similar to the quark sector:

$$
\begin{aligned}
\mathcal{L}_{\text{Yuk}} &\supset y_e \bar{L}_L (E_R \cdot H) + y_\nu \bar{L}_L (N_R \cdot \tilde{\tilde{H}}) - y_e' e_L^H (E_R \cdot \Phi_H) + y_\nu' \nu_L^H (N_R \cdot \tilde{\Phi}_H) + \text{H.c.} \\
&= y_e \bar{L}_L (e_R^H H_2 - e_R H_1) - y_\nu \bar{L}_L (\nu_R \tilde{H}_1 + \nu_R^H \tilde{H}_2) \\
&\quad - y_e' e_L^H (e_R^H \Phi_2 - e_R \Phi_1) - y_\nu' \nu_L^H (\nu_R \Phi_1^* + \nu_R^H \Phi_2^*) + \text{H.c.}
\end{aligned}
\tag{24}
$$

The charged leptons receive their masses the same way as in SM. The neutrinoes get Dirac masses in this setup: $m_\nu^D = y_\nu v/\sqrt{2}$. Heavy lepton e^H and ν^H receive masses characterised by v_Φ: $y_e' v_\Phi/\sqrt{2}$ and $y_\nu' v_\Phi/\sqrt{2}$ respectively.

For the benefit of discussions in Sec. IV, we substitute the shifted field into the lepton sector while ignoring the GBs to arrive at the interaction between SM lepton, heavy leton and heavy scalars:

$$
\begin{aligned}
\mathcal{L}_{\text{Yuk}} = &+ y_e (\bar{\nu}_L e_R^H H^+ - \bar{e}_L e_R^H H_2^0 - \bar{e}_L e_R \frac{v+h}{\sqrt{2}} \\
&- y_\nu (\bar{\nu}_L \nu_R \frac{v+h}{\sqrt{2}} + \bar{\nu}_L \nu_R^H H_2^0 - \bar{e}_L \nu_R^H H^-) \\
&- y_e' e_L^H e_R^H \frac{v_\Phi + \phi_2}{\sqrt{2}} - y_\nu' \nu_L^H \nu_R^H \frac{v_\Phi + \phi_2}{\sqrt{2}}
\end{aligned}
\tag{25}
$$

It is very useful to observe that H_1 only couples to SM fermions, H_2 couples to one SM fermion and one heavy fermion, and Φ_H couples to two heavy fermions.

III Neutrino Mass Generation

In this original setup in G2HDM, neutrinos are purely Dirac. By the Yukawa coupling, SM neutrino receives mass $m_\nu^D = y_\nu v/\sqrt{2}$. Since the neutrino masses are most likely to be very small, in the order around 10^{-5}eV. In such setup, the Yukawa coupling constant y_ν of the neutrino is down to $\mathcal{O}(10^{-11})$.

However, It is also suggested in [4] there is another option. By adding a new $SU(2)_H$ triplet Δ_N, one can generate Majorana mass terms for neutrino and achieve the type-II see-saw mechanism for Neutrino mass. In this way, the small parameters in Dirac Mass setting could avoided and y_ν could be elevated to the magnitude of around y_e. The Lorentz invariant Majorana Lagrangian for Δ_N is

$$\begin{aligned}\mathcal{L}_{\mathrm{M}} &\supset g_\nu N_R^{\mathrm{T}} i\sigma^2 \Delta_N N_R \\ &= g_\nu \begin{pmatrix} \nu_R & \nu_R^H \end{pmatrix} i\sigma^2 \begin{pmatrix} \frac{1}{\sqrt{2}}\Delta_N^+ & \Delta_N^{++} \\ \Delta_N^0 & -\frac{1}{\sqrt{2}}\Delta_N^+ \end{pmatrix} \begin{pmatrix} \nu_R \\ \nu_R^H \end{pmatrix}\end{aligned} \tag{26}$$

with $\langle \Delta_N^0 \rangle = v_N$. Therefore, the right-handed SM neutrino acquires a Majorana mass $m_R^M = g_\nu v_N$. Notice that the heavy right-handed neutrino does not receive a Majorana mass, so the heavy neutrino remains as Dirac fermion. This setup is very similar to that in [7], besides the different gauge groups and matter content.

The tree-level Majorana mass matrix in the basis of $\{\nu_L, \nu_R\}$ is

$$\mathcal{M}_\nu = \begin{pmatrix} 0 & y_\nu \frac{v}{\sqrt{2}} \\ y_\nu \frac{v}{\sqrt{2}} & g_\nu v_N \end{pmatrix} \tag{27}$$

Diagonalizing the mass matrix gives two eigenvalues corresponding to the mass of one heavy neutrino (notice that it is different from ν^H) and one light neutrino

$$\begin{aligned} M_R &= g_\nu v_N + \frac{y_\nu^2 v^2}{2 g_\nu v_N} \\ m_\nu &= -\frac{y_\nu^2 v^2}{2 g_\nu v_N} \end{aligned} \tag{28}$$

If we let $m_R^M \gg m_\nu^D$, M_R is proportional to m_R^M. To achieve light neutrino mass in the order of sub-eV, the Yukawa coupling constant is about

$$\frac{y_\nu}{\sqrt{2}} \sim 4.06 \times 10^{-8} \sqrt{\frac{M_R}{\mathrm{GeV}}} \tag{29}$$

which is in the same magnitude with that of the electron if we choose $M_R \sim 1\text{TeV}$. Of course, the Yukawa coupling constant can be further increased as M_R becomes larger.

In the discussion up to this point, we haven't explicitly shown the flavor indices of the extra fermions. In order to cancel the anomaly, presumably they, right-handed neutrino and all the heavy fermions, should all come in three flavors. Our concerns here are the leptons and hence we shall focus on the three flavors of $\nu_R, \nu_R^H, \nu_L^H, e_L^H, e_R^H$. The mass of heavy fermions are Dirac and hence the mass matrix can be diagonalized by bi-unitary matrices:

$$M_\nu^H = U_\nu^H \text{Diag}(m_{\nu 1}^H, m_{\nu 2}^H, m_{\nu 3}^H) V_\nu^H, \;\; M_e^H = U_e^H \text{Diag}(m_{e1}^H, m_{e2}^H, m_{e3}^H) V_e^H \tag{30}$$

We have the freedom to rotate ν_L^H and e_L^H to thier respective mass eigenstates $\nu_L'^H, e_L'^H$:

$$\nu_L'^H = U_\nu^H \nu_L^H, \;\; e_L'^H = U_e^H e_L^H \tag{31}$$

But ν_R^H and e_R^H has been locked into doublets with ν_R and e_R to form $N_R = \begin{pmatrix}\nu_R & \nu_R^H\end{pmatrix}$ and $E_R = \begin{pmatrix}e_R^H & e_R\end{pmatrix}^{\text{T}}$. So the extra matrix matrices V_ν^H and V_e^H needed to rotate ν_R^H and e_R^H to diagonalize mass matrix will show up as a mixing between the mass eigenstates and gauge eigenstates of ν^H and e^H. Let's call the mixing V_ν^H and V_e^H respectively. These two mixing matrices will appear in the interaction vertices as shown below. Please observe that the exception is the case of degenerate heavy fermions. In the degenerate case, V_ν^H and V_e^H can be absorbed into U_ν^H and U_e^H since $\text{Diag}(m_{\nu,e1}^H, m_{\nu,e2}^H, m_{\nu,e3}^H) \propto I$. Only one rotation matrix is needed to make the mass matrix diagonal to identity matrix. We can always rotate just ν_L^H and e_L^H to diagonalize ν^H and e^H mass matrices and hence there is no mixing.

IV Lepton Flavor Violating Higgs Decays

In G2HDM, the appearance of extra heavy fermions and their coupling with SM fermions through exotic scalars and gauge bosons implies that the rotation defining the heavy fermion mass eigenstates does not diagonalize their Yukawa coupling nor the gauge coupling as described above. At the loop level these flavor-changing interactions will induce CLFV processes involving the charged leptons. Here we will study the CLFV Higgs decays $h \to l^i l^j$ in G2HDM. These decay channels are currently searched at the LHC; at 95% C.L., at ATLAS and CMS:

$$\mathcal{B}(h \to \mu e) < 6.1 \times 10^{-5}(\text{ATLAS}); 3.5 \times 10^{-4}(\text{CMS})$$
$$\mathcal{B}(h \to \tau e) < 2.8 \times 10^{-3}(\text{ATLAS}); 6.1 \times 10^{-3}(\text{CMS})$$
$$\mathcal{B}(h \to \tau \mu) < 4.7 \times 10^{-3}(\text{ATLAS}); 2.5 \times 10^{-3}(\text{CMS})$$

, where $h \to l^i l^j$ stands for $h \to l^i \bar{l}^j, l^j \bar{l}^i$. Recall that the CLFV Higgs couplings, which may lead to the above decay, are absent at tree level in SM. On the other hand, at one loop this process can only be induced by the non-vanishing minuscule neutrino masses as implied by various neutrino oscillation experiments and thus the contribution is expected to be quite small. A positive measurement of this branching ratio(BR) in the near future at the percent level would be a clear indication of new physics beyond the SM.

Feynman diagrams of CLFV processes are shown in Fig 1. We'll present our calculation in three group. It will be shown that the diagrams with W' exchange is the largest after various constraints are included.

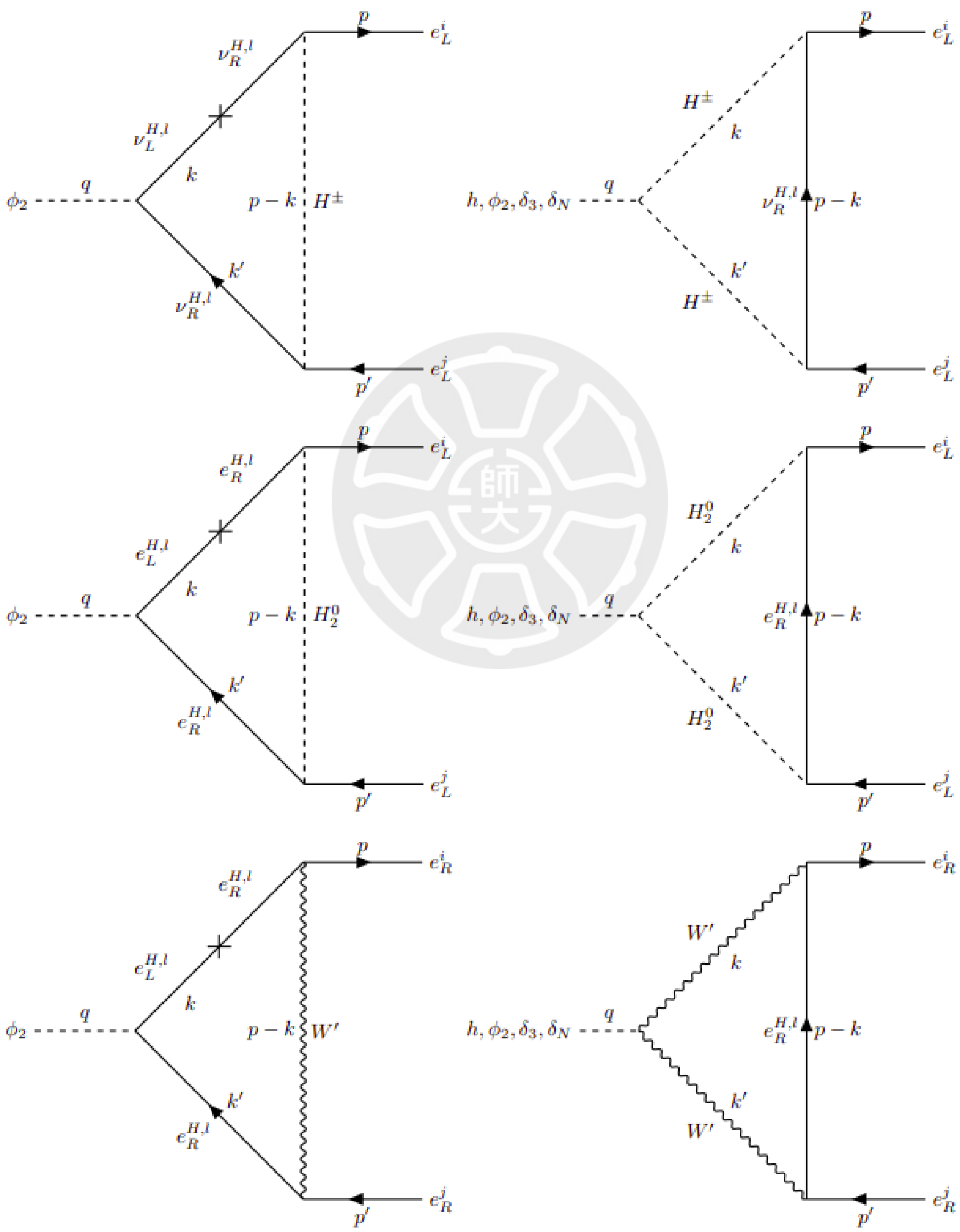

Figure 1: Feynman diagrams of CLFV processes in G2HDM

For all the diagrams, in the loop we need to sum over the three flavors of the heavy fermions e_R^H and ν_R^H. It is interesting to observe that CLFV processes will cancel if the heavy fermions are degenerate in masses, as described above. Therefore, we present our numerical result and formula at the other extreme, as two of the heavy fermion masses are sent to infinity and hence the two decouple from the theory, leaving only one flavor of heavy fermions in the picture. This is by no means the preferred scenario but the one where the formula is easier to write down. The reality would be in between the degeneracy case with zero CLFV and the extreme case in which we present results.

A. W' Exchanging Process

The first group of the diagrams is the W' exchanging diagrams. We use unitary gauge in our calculation. The amplitude of the bottom-left diagram in Fig. 1 is

$$
\begin{aligned}
i\mathcal{M} = \int \frac{d^4k}{(2\pi)^4} \frac{-i}{(p-k)^2 - m_{W'}^2 + i\varepsilon} (g^{\mu\nu} - \frac{k^\mu k^\nu}{m_{W'}^2}) \bar{u}(p) P_L \frac{g_H}{\sqrt{2}} \gamma_\mu \left(V_e^H\right)^{jl} \frac{i m_{e^H}}{k^2 - m_{e_l^H}^2 + i\varepsilon} \\
\times (-i y_e') \frac{i(\not{k}' + m_{e^H})}{k'^2 - m_{e^H}^2 + i\varepsilon} \frac{g_H}{\sqrt{2}} \gamma_\nu \left(V_e^{H\dagger}\right)^{li} P_R v(p'),
\end{aligned}
\tag{32}
$$

, where O_{21} is the element of the orthogonal matrix O that diagonalizes the mass mastrix $\mathcal{M}_0^2$ displayed in Eq. 18. To calculate, we need to sum over internal heavy e^H flavor l while i, j are external lepton flavor, with V_e^H the mixing matrix of e^H as described above. It is obvious that in the case of degenerate heavy e^H, the amplitude is proportional to $V_e^H V_e^{H\dagger} = I$. As expected the flavor changing part vanishes.

The above amplitude is ostensibly infinite and we arrives at a pole in ϵ when we use dimensional regularization to regularize the infinity. The pole arises at large loop momentum. When the loop momentum is large, heavy fermion mass is unimportant and could be ignored and it results in a degenerate case. Hence it is no surprise that the pole is independent of the internal heavy fermion flavor l. When we concentrate on flavor changing process, the pole will be canceled out due to $V_e^H V_e^{H\dagger} = I$. This absence of infinity justifies our use of unitary gauge, which, as a limit of R_ξ gauge, could have problems in loop calculation if the result is infinite.

Hence in the following we only include amplitude that is dependent on the heavy fermion mass. In reality the heavy e^H and ν^H are unlikely to be exactly degenerate unless there is a symmetry or coincidence. As noted above, to consistently present our result and give a rough idea of the magnitude of flavor changing processes, we study the opposite extreme situation that two of the heavy fermions are very heavy and decouple from the theory.

The amplitudes in this group are basically determined by the mass $M_{W'}$ of W' and the gauge coupling constant g_H as well as the heavy lepton Yukawa coupling y'_e. The mass $M_{W'}$ is in turn determined by g_H, v, v_Φ, v_Δ and v_N

$$M_{W'}^2 = \frac{1}{4} g_H^2 (v^2 + v_\Phi^2 + 4v_\Delta^2 + 4v_N^2), \tag{33}$$

with $v \ll v_\Phi \cong v_\Delta \cong v_N$. Additionally, according to [5], v_Φ is constrained by LEP Z observables to be roughly larger than at least 10 TeV, while v_Δ depends on other unrelated parameters, we can safely ignore the role played by the triplet Δ_H at the moment. The maximum BR is calculated with $v_\Phi = v_N = 10$ TeV, $g_H = y'_e = 1$ to be

$$\mathcal{B}(h \to \mu\tau) = 8.47 \times 10^{-10} \tag{34}$$

, which would be outside the LHC detecting width. We will fixed $v_N = 10$ TeV, $y'_e = 0.5$ and vary the scale of v_Φ at some values starting with 10 TeV. In Fig. 2, we plot $v_\Phi = 10, 30, 80$ TeV with $g_H \sim (0.01, 1)$. We also plot Fig. 3 in a more extreme situation that $y'_e = 3$. Although it may not be a favored choice, the one-loop calculation allows us to have Yukawa coupling merely larger than 1.

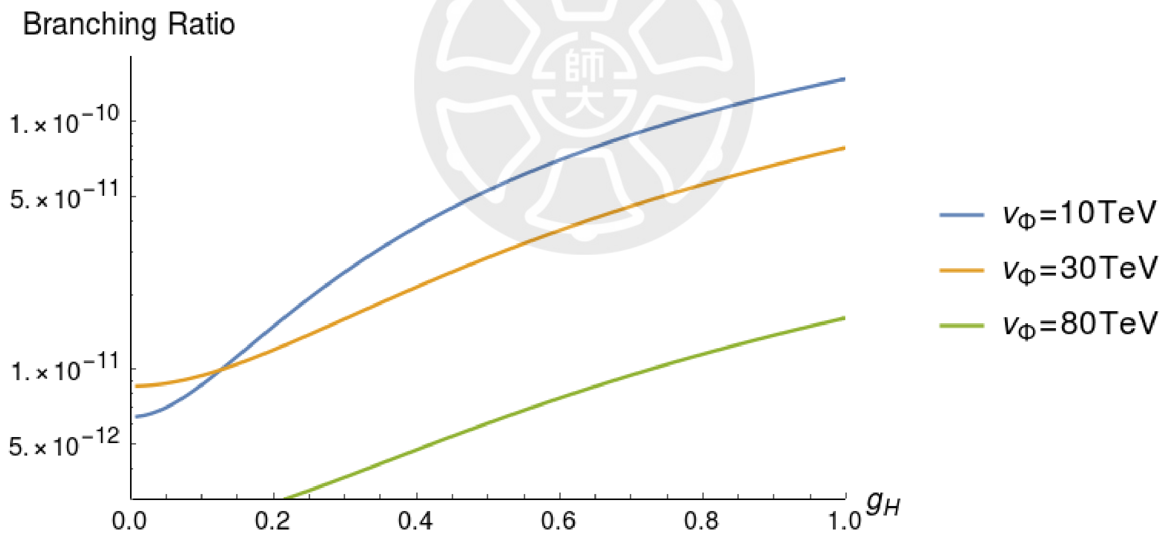

Figure 2: The $h \to \mu\tau$ branching ratio dependece on g_H with $y'_e = 0.5$ in W' exchanging diagram

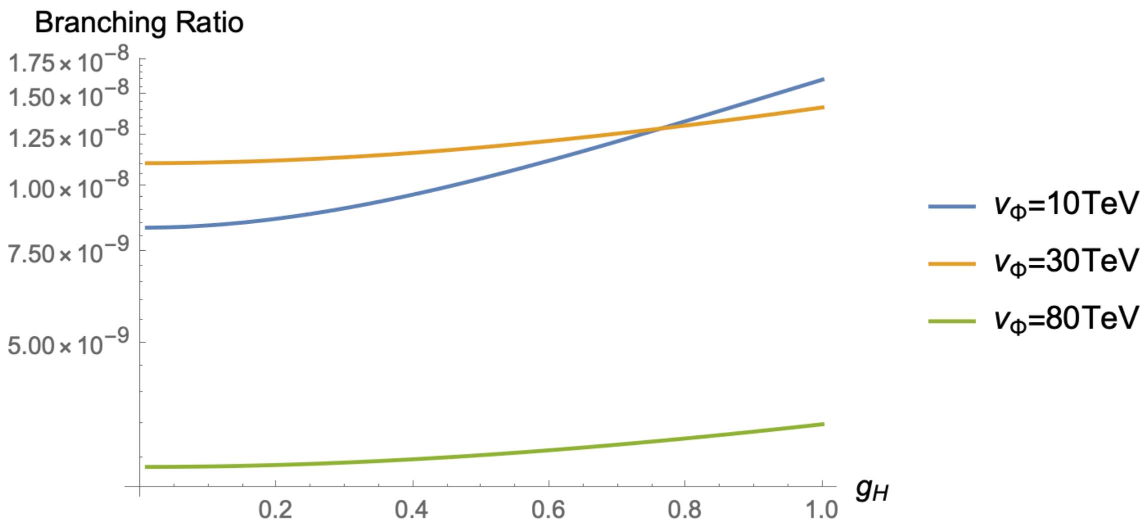

Figure 3: The $h \to \mu\tau$ branching ratio dependece on g_H with $y'_e = 3$ in W' exchanging diagram

In this case $M_{W'}$ is intrinsically determined by g_H. It is not hard to make a guess that the amplitude will increase as $M_{W'}$ and g_H decrease. However, there is a factor of g_H^2 in the overall amplitude, which will reduce the amplitude as g_H decreases. Two trends are roughly canceled each other and as a result the amplitude is pretty stable as g_H varied. Finally, there are overlaps between $(y'_e, v_\Phi) = (0.5, 10)$ and $(y'_e, v_\Phi) = (0.5, 30)$ in both graphs which indicate the BR is proportional to v_Φ in small g_H region.

In the major part of Fig. 4 and Fig. 5, BR decreases as v_Φ becomes larger as we expect. However, there is always a region that the amplitude is propotional to the overall v_Φ no matter what value g_H takes. To put it another way, the maximum BR does not locate at the minimum of v_Φ. We can see that clearly especially in Fig. 5.

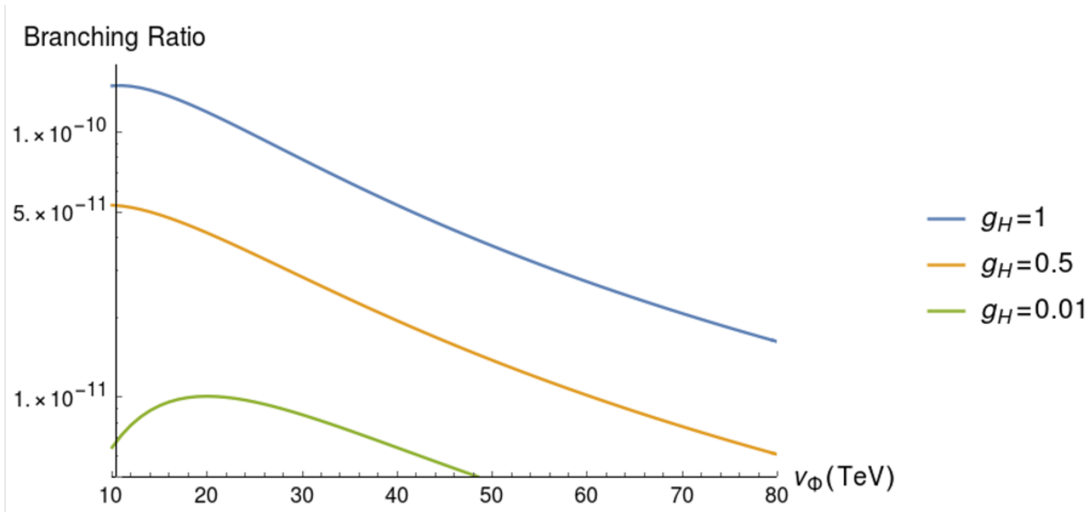

Figure 4: The $h \to \mu\tau$ branching ratio dependece on v_Φ with $y'_e = 0.5$ in W' exchanging diagram

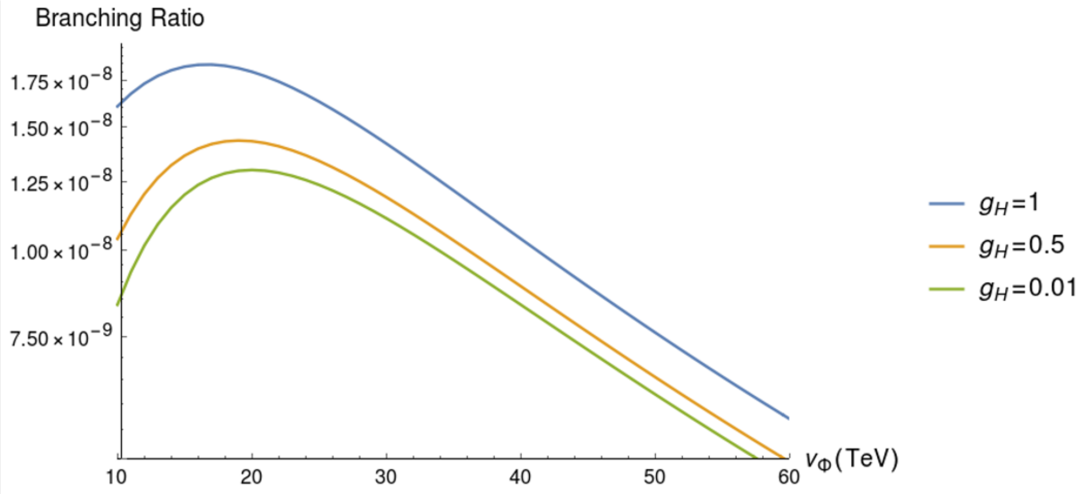

Figure 5: The $h \to \mu\tau$ branching ratio dependece on v_Φ with $y'_e = 3$ in W' exchanging diagram

B. $H^\pm$ and H_2^0 Exchanging Process

The second group is the $H^\pm$ exchanging diagrams. The amplitude of the top-left diagram in Fig. 1 is

$$i\mathcal{M} = \int \frac{d^4k}{(2\pi)^4} \frac{i}{(p-k)^2 - m_{H^\pm}^2 + i\varepsilon} \bar{u}(p) P_R(-iy_\nu U_\nu^{il}) \frac{i m_{\nu^H}}{k^2 - m_{\nu^H}^2 + i\varepsilon} \times (-iy'_\nu) \frac{i(\not{k}' + m_{\nu^H})}{k'^2 - m_{\nu^H}^2 + i\varepsilon} (-iy_\nu U_\nu^{jl}) P_L v(p') \tag{35}$$

The amplitude is set by the neutrino Yukawa coupling constant y_ν, the heavy neutrino Yukawa coupling constant y'_ν and the $H^\pm$ mass $m_{H^\pm}$. Similar to the W' diagrams, the amplitude increases as $H^\pm$ becomes lighter, i.e. $m_{H^\pm}$ decreases. However, there is no connection between the factors in the numerator: y_ν, y'_ν and v_Φ. Hence the three parameters can be changed more or less independently. The central value of $m_{H^\pm}$ is put around 15 TeV with a 1σ range between 9 TeV to 20 TeV in [5]. Therefore, it is possible for $m_{H^\pm}$ to be between 3 TeV to 30 TeV. It is obvious that the amplitude will become small if the Yukawa couplings are small. However, heavy neutrino Yukawa coupling y'_ν could not be too small, or the heavy fermions will be light. The neutrino Yukawa coupling constant y_ν is set by Eq. (29) to be around $10^{-8}\sqrt{M_R/\text{GeV}}$.

We argue if the scale M_R is large, y_ν could be large too. Hence, we would fix $v_\Phi = 10$ TeV and survey a range of $y'_\nu \sim (0.01, 1)$ in Fig. 6. When y'_ν is small, $m_{H^\pm} \gg m_{\nu^H}$ in the denominator of the amplitude. Thus, the BR depends mostly on y'_ν in the overall amplitude. However, as y'_ν getting larger, the m_{ν^H} in the denominator cancels the trends and the BR dependece on y'_ν becomes lower. It is not hard to argue that the rate will only become large in a corner of the parameter space: with large y_ν, y'_ν and small $m_{H^\pm}$. It is calculated that the maximum BR for

$y_\nu \cong y'_\nu \sim 1$ and $m_{H^\pm} = 3$TeV is

$$\mathcal{B}(h \to \mu\tau) = 5.18 \times 10^{-10} \tag{36}$$

This also would be invisible in the LHC current level of detection. In general, besides the maximum, the BR of $H^\pm$ diagrams will be much smaller than that of the W' diagrams.

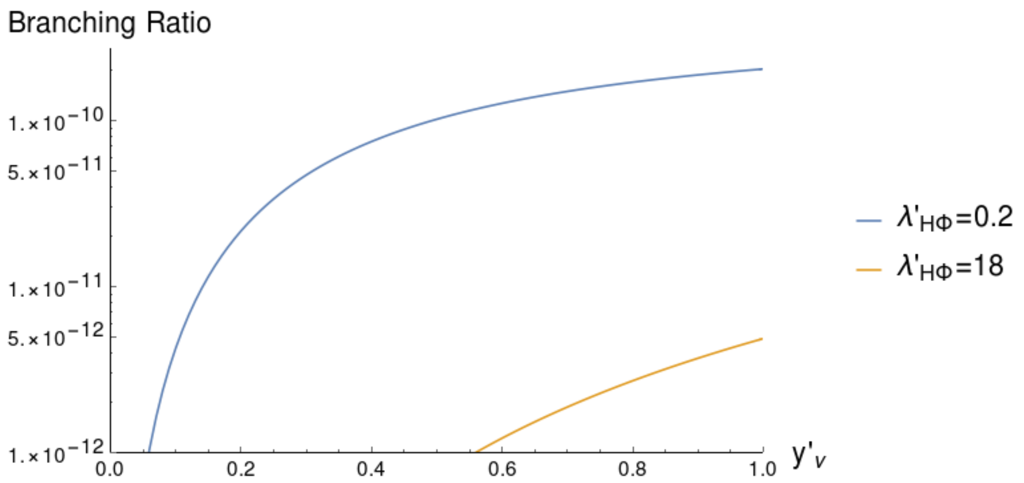

Figure 6: The $h \to \mu\tau$ branching ratio dependece on y'_ν in $H^\pm$ exchanging diagram

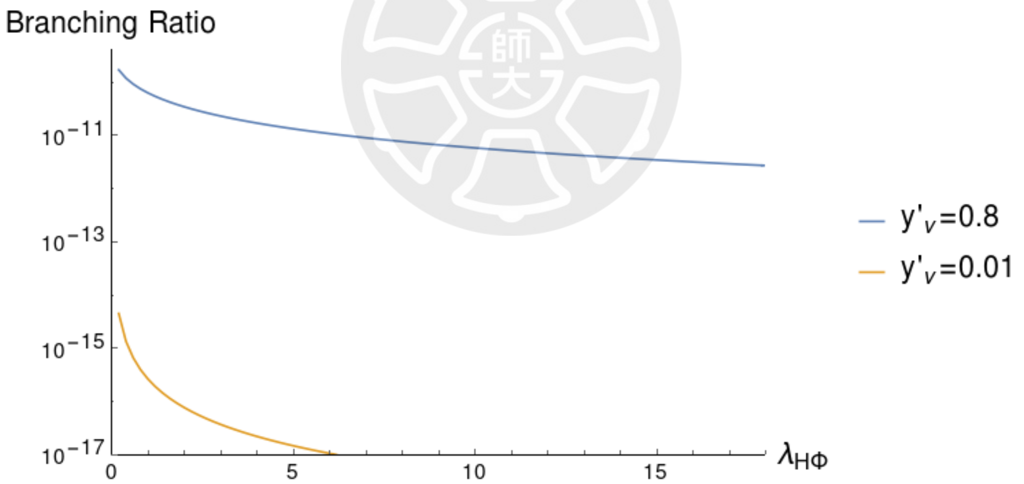

Figure 7: The $h \to \mu\tau$ branching ratio dependece on $\lambda'_{H\Phi}$ in $H^\pm$ exchanging diagram

Also, in Fig. 7, we vary $\lambda'_{H\Phi} \sim (0.2,\ 18)$, i.e. varying $m_{H^\pm}$ from around 3 TeV to 30 TeV. We can see the BR is not sensitive to $m_{H^\pm}$, besides a small region with $y'_\nu = 0.01$. The reason comes from the fact that the amplitude only depends on $1/\lambda'_{H\Phi}$, therefore the trend on $\lambda'_{H\Phi}$ will be much flatter than the trends on other parameters.

The third group are H_2^0 exchanging diagrams. The amplitude of the middle-left diagram in Fig. 1 is

$$i\mathcal{M} = \int \frac{d^4k}{(2\pi)^4} \frac{i}{(p-k)^2 - m_{H_2^0}^2 + i\varepsilon} \bar{u}(p) P_R (-iy_e U_e^{il}) \frac{i m_{e^H}}{k^2 - m_{e^H}^2 + i\varepsilon} \times (-iy'_e) \frac{i(\not{k}' + m_{e^H})}{k'^2 - m_{e^H}^2 + i\varepsilon} (-iy_e U_\nu^{jl}) P_L v(p') \tag{37}$$

These diagrams are set by charged lepton Yukawa coupling y_e, heavy lepton Yukawa coupling y'_e and H_2^0 mass $m_{H_2^0}$. Although the mass of H_2^0 centers around a smaller value than $m_{H^\pm}$ at 1 TeV, potentially putting the rate to the 10^{-10} range, charged lepton Yukawa coupling y_e is fixed by charged lepton mass to be quite small, The rate is estimated to be as small as 10^{-32}. We can safely ignore this contribution in our discussion.

C. Features in W' and $H^\pm$ Exchanging Process

Notice that in order to make the BR for $H^\pm$ diagram reaching the maximum, we need v_N to be around 10^{12} TeV, which will eliminate the possibility of W' diagram. On the contrary, the maximum BR for W' diagrams occur when the magnitude of v_N is in TeV range, which diminishes the $H^\pm$ diagram's BR to $\mathcal{O}(10^{-34})$. To put it briefly, the W' diagrams and $H^\pm$ diagrams will not be visible at the same time.

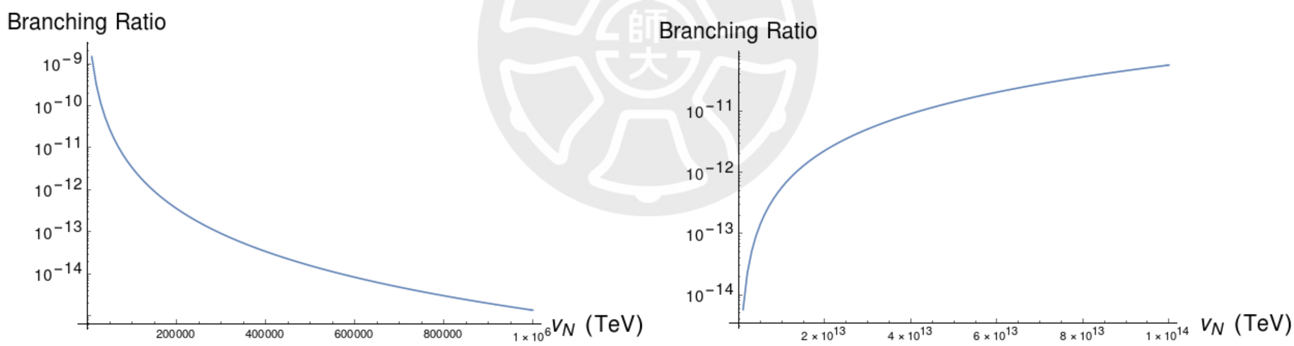

Figure 8: The $h \to \mu\tau$ branching ratio dependece on v_N in W' (left) and $H^\pm$ (right) exchanging diagram

W' and $H^\pm$ diagrams are predominantly depending on Yukawa coupling y'_e and y_ν respectively. In general, both couplings should not exceed 1 for perturbation reason. However, there is no further constraint on y'_e and y_ν. Therefore in the scenerio that the couplings are raised to around 10, the BRs for both diagram can reach to $\mathcal{O}(10^{-6})$, within the margin of LHC detection range. It is also interested to observe that in large Yukawa coupling case, the BR of $H^\pm$ is larger than the BR of W' diagram.

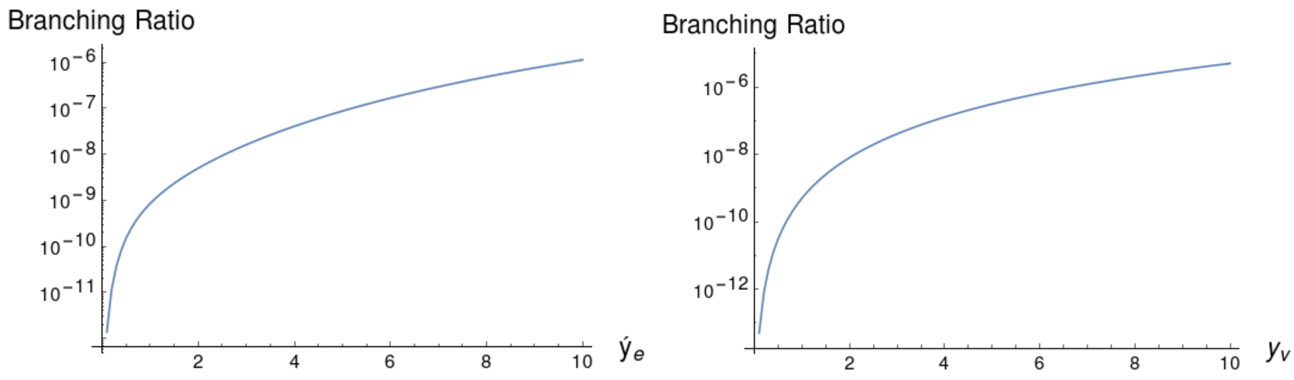

Figure 9: The $h \to \mu\tau$ branching ratio dependece on y'_e in W' (left) and y_ν in $H^\pm$ (right) exchanging diagram

V Conclusion

G2HDM provides a rich matter content and elaborate coupling to construct various circumstances. In this work, we add another $SU(2)_H$ triplet Δ_N into the matter content, which alter the scalar potential and mass spectrum. The first block of mass spectrum is now 4×4, indicates that Δ_N's VEV v_N also contribues in 125 GeV Higgs mass. While the second block and the charged Higgs mass $m_{H^\pm}$ is the same because of the minimum conditions, Eq. (14)-(17). Due to the same reason, the scenario that $SU(2)_H$ triplet Δ_N triggers symmetry breaking is unchanged.

By adding Δ_N and using type-II seesaw mechanism, we can achieve Majorana-type nuetrino and elevate neutrino Yukawa coupling y_ν to the magnitude comparable to electron Yukawa coupling y_e with heavy neutrino mass $M_R \sim 1$TeV. We can further increase y_ν by a heavier M_R to facilitate CLFV process mediated by ν_R^H.

Finally, we scrutinize the CLFV processes in G2HDM. The Yukawa coupling between SM fermions and heavy fermions allows us to construct various CLFV processes. The maximum BRs of the W' diagrams and $H^\pm$ diagrams fall within the magnitude of 10^{-10}, hardly visible in LHC. But if there is a scenerio that Yukawa coupling can exceed 1, the BRs of the W' diagrams and $H^\pm$ diagrams can reach to $\mathcal{O}(10^{-6})$, marginally visible in LHC.

Due to the gauge kinetic term, Δ_N's VEV v_N also contribute in the W' mass. Since W' diagrams require small v_N while $H^\pm$ diagrams need large v_N to achieve the seesaw mechanism, the two types of diagrams will not exist at the same time. In contrast, the BR of H_2^0 diagrams is constrainted by the charged lepton Yukawa coupling to be no hope in any detection.

The various constraints in the theory still need to be examined, especially the contribution of v_N in 125 GeV Higgs, charged Higgs and gauge bosons. A more concrete scheme of the heavy fermions flavor effect in CLFV process should also be included in the future study.

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