The Higgs in the Standard Model and slightly beyond

Luís Sombreireiro - luis.sombreireiro@ist.utl.pt
Supervisor: Filipe Joaquim
Instituto Superior Técnico, Lisboa, Portugal
(Dated: November 2014)

In this thesis we study some theoretical aspects related with the Higgs field/boson in the context of the SM. We review the gauge structure of the SM and the role of the Higgs field in electroweak symmetry breaking to generate mass to gauge bosons, fermions and to the Higgs itself. Theoretical constrains to the Higgs boson mass are presented and its decays are analyzed in the context of the SM. Special attention is given to the decay $h^0 \to \gamma\gamma$. As an example of a scenario of new physics, we consider an extension fo the SM where scalar triplets are introduced. This model is particularly interesting because it can accommodate a mechanism to generate small neutrino masses (Type-II seesaw mehcanism). In the context of this model, the decays $h^0 \to \gamma\gamma$ and $h^0 \to \gamma Z$ are studied.
Keywords: Standard Model, Higgs field, Higgs decays, Seesaw mechanism

I. INTRODUCTION

Curiosity is the essence of Science. That need to know, to understand how Nature works, in the end, is what motivates scientists to do what they do. Physics is no different. Experimentalists around the world test well known theoretical predictions and search for evidences of new physics. Theorists try to either expand incomplete theories, or invent new ones all together. The Higgs "tale" is a rather interesting one; an experimental endeavor is highly motivated by a theoretical one. The Higgs boson had to be found so that a well-tested model, the Standard Model (SM), could be completed.

The SM of particle physics aims at describing all interactions, with the exception of gravity. Fermions are gathered into two groups of six quarks and six leptons. The photon is the carrier of the electromagnetic interaction, the $W^\pm$ and Z for the weak interaction, and the gluons for the strong interaction. There is also a scalar neutral particle, the Higgs boson, which is the quanta of the Higgs field. The Higgs is essential since without it particles are massless in the SM. The Higgs field gives particles its mass through the Higgs mechanism [1], which was adapted to the SM by Glashow, Weinberg and Salam [2, 3]. The Higgs field provides mass to particles for two reasons: first, it couples with all particles which later acquire mass and second, the structure of its potential is so that the Higgs field acquires a non-zero value at its lowest energy state. This means that, even in its ground state, the field does not vanish in all of space, contrary, for example, to the electromagnetic field, which is only nonzero after charged particles "create" it .

A good and known analogy for the mass-giving mechanism is to think of the Higgs field as water in a pool. When a swimmer tries to cross the pool, he will not reach the other side as fast as he could if the pool was empty. But what if there were swimmers that could just swim without interacting with the water? In the absence of interaction, they would cross the pool faster. We can say that the difference between swimmers is a characteristic called mass. A swimmer (or a particle) that interacts with the water (Higgs field) has mass and moves slower. Think of massless particles as the only ones that can travel at the speed of light in vacuum c. This is true because for massive particles, which always have a rest reference frame, there cannot be an inertial reference frame in which c is zero (according to the second postulate of Special Relativity). This means that there is no reference frame in which a massive particle and a photon are stationary. Therefore, they cannot have the same speed. We can think of the Higgs field as slowing down particles, like water slowing the swimmer, making so that they cannot reach c, giving them mass in the process.

In this work, we study various Higgs decays in the SM and in a model "slightly" modified. In sections II-VII, we review the structure of the SM and the role of the Higgs field in the mechanism for mass generation. Theoretical constraints to the Higgs boson mass are also presented. In section VIII, we present several decay widths: of Higgs decays into WW or ZZ, into fermions $f\bar{f}$, gluons gg, a photon pair $\gamma\gamma$ and γZ. In section IX, the Type-II seesaw model is introduced and its implications in the decays into $\gamma\gamma$ and γZ are studied. These decay modes, specially the one into two photons, were important for the Higgs search [4–6]. Their possible deviation from the SM predicted result is studied in terms of the parameters of the new theory, namely the masses of the new scalar particles which arise in the Type-II seesaw framework.

II. THE GAUGE STRUCTURE OF THE STANDARD MODEL

In aiming at finding a theory that explains fundamental particles and their interactions, particle physicists look at models described by Lagrangians which are invariant under certain symmetry transformations of the fundamental fields. A Lagrangian that is invariant to gauge transformations of its fields for a group G is said to be gauge invariant.

The SM Lagrangian is invariant under symmetries of the gauge group

$$SU(3)_C \times SU(2)_L \times U(1)_Y. \tag{1}$$

The gauge group $SU(2)_L \times U(1)_Y$ is related to electroweak interactions and $SU(3)_C$ is related to Quantum Chromodynamics (QCD). Left-handed leptons and quarks are grouped in doublets of $SU(2)$,

$$L_L = \begin{pmatrix} \nu_e \\ e \end{pmatrix}_L, \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L, \begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix}_L \tag{2}$$

and

$$Q_L = \begin{pmatrix} u \\ d \end{pmatrix}_L, \begin{pmatrix} c \\ s \end{pmatrix}_L, \begin{pmatrix} t \\ b \end{pmatrix}_L. \tag{3}$$

Hypercharge is defined as $Y = Q - T^3$ with Q as the charge of the particle and T^3 as the isospin. Right-handed particles transform as singlets of $SU(2)_L$, and are written as e_R, u_R, d_R (for all charged fermions). Neutrinos only have a left-handed component. There is no theoretical explanation for the observed number of families (why three and not more ?) nor for the quantum numbers of each particle. The connection between representations of gauge groups and how particles interact comes from imposing gauge invariance, leading to "non-cannonical" kinetic terms. In the case of leptons,

$$\mathcal{L} = \sum_{f=e,\mu,\tau} -i\overline{L_L^f}\gamma^\mu D_\mu L_L^f - i\overline{f_R}\gamma^\mu D_\mu f_R, \tag{4}$$

where D_μ, the covariant derivative, is defined in a general way as

$$D_\mu = \partial_\mu + igA_\mu^a T^a + ig'YB_\mu + ig_3.G_\mu^a\lambda^a. \tag{5}$$

Here λ^a are the generators of the $SU(3)$ group, $T^a = \frac{\tau^a}{2}$ the generators of the $SU(2)$ group, with τ^a being the Pauli matrices. The constants g, g' and g_3 are the gauge coupling constants for $SU(2)$, $U(1)$ and $SU(3)$ respectively. Their values determine the strength of the various interactions. In writing the covariant derivative, gauge bosons A_μ^a, B_μ and G_μ^a are introduced in a sum with the group generators. This is done to ensure gauge invariance of the kinetic terms when performing local gauge transformations. The kinetic term in Eq. (4) is invariant since $D_\mu L_L^e$ and $D_\mu e_R$ transform like the fields L_L^e and e_R themselves, for any transformation of the gauge group (1) [7],

$$U(1)_Y : \Psi \to e^{iq\alpha(x)}\Psi \text{ and } SU(2) : \Psi \to e^{iT^i\beta^i(x)}\Psi \tag{6}$$

After modifying the derivative to ensure gauge invariance, we can expand kinetic terms to see the field (and not the field doublets) interaction with the gauge bosons. We redefine the gauge bosons with definite charge as

$$W_\mu^\mp = \frac{1}{\sqrt{2}}\left(A_\mu^1 \pm iA_\mu^2\right), \tag{7}$$

for the charged gauge bosons $W^\pm$, and as

$$\begin{pmatrix} A_\mu \\ Z_\mu \end{pmatrix} = \begin{pmatrix} s_W & c_W \\ c_W & -s_W \end{pmatrix}\begin{pmatrix} A_\mu^3 \\ B_\mu \end{pmatrix}, \tag{8}$$

for the neutral gauge bosons. In (8), $s_W \equiv \sin\theta_W$ and $c_W \equiv \cos\theta_W$, with θ_W being the Weinberg angle [8].

The kinetic terms for leptons can be expanded and written as (for the electron)

$$\begin{aligned}\mathcal{L} = &-\frac{g}{\sqrt{2}}\overline{e_L}\gamma^\mu\nu_{eL}W_\mu^- + -\frac{g}{\sqrt{2}}\overline{\nu_{eL}}\gamma^\mu e_L W_\mu^+ \\ &+q\bar{e}\gamma^\mu e A_\mu \\ &-\sum_{f=e,\nu_e}\overline{f}\gamma^\mu(g_V^f - g_A^f\gamma_5)fZ_\mu,\end{aligned} \tag{9}$$

with $g_V^f = \frac{1}{2}T_3^f - Q^f(\sin\theta_W)^2$, $g_A^f = \frac{1}{2}T_3^f$ and $q = |e|$. The charged currents only exist for left-handed particles. The kinetic terms for quarks are

$$\begin{aligned}\mathcal{L} = &-i\overline{Q_L^{ud}}\gamma^\mu D_\mu Q_L^{ud} - i\overline{u_R}\gamma^\mu D_\mu u_R - i\overline{d_R}\gamma^\mu D_\mu d_R \\ &+ (\text{with } c, s \text{ and } t, b),\end{aligned} \tag{10}$$

which lead to the electroweak currents:

$$\begin{aligned}\mathcal{L} = &-\frac{g}{2\sqrt{2}}\bar{d}\gamma^\mu(1-\gamma_5)uW_\mu^- + \text{H.c.} \\ &-q\left(\frac{2}{3}\bar{u}\gamma^\mu u - \frac{1}{3}\bar{d}\gamma^\mu d\right)A_\mu \\ &-\sum_{f=u,d}\overline{f}\gamma^\mu(g_V^f - g_A^f\gamma_5)fZ_\mu.\end{aligned} \tag{11}$$

Apart from the kinetic terms for fermions, we also have the ones for gauge bosons, namely

$$\mathcal{L} = -\frac{1}{4}F_{\mu\nu}^a F^{a\mu\nu} - \frac{1}{4}G_{\mu\nu}G^{\mu\nu}, \tag{12}$$

where

$$F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a - g\epsilon^{abc}A_\mu^b A_\nu^c, \tag{13}$$

and

$$G_{\mu\nu}^a = \partial_\mu B_\nu^a - \partial_\nu B_\mu^a. \tag{14}$$

The final problem to address is the one of how to give mass for all these fermions and bosons. In the SM when asking, "how does a particle gain mass?", one should ask, "what terms do I need to add to the Lagrangian in order to give particles mass?". These terms, called mass terms, contribute to equations of motion for massive particles through the Euler-Lagrange equations. The problem is that these mass terms cannot be simply written in the theory without introducing something else, as explained in the following sections.

III. WHY WE NEED THE HIGGS?

Experimentally, we know that the W and the Z bosons, the carriers of the weak interactions, are massive [9, 10]. Then, in some way, a mass term should be present in the Lagrangian of our theory. However, writing bare mass terms for bosons is problematic, since they break gauge invariance, i.e bare mass terms for W and Z are not invariant under $SU(2)$ transformations. Non-abelian gauge fields usually transform as $B_\mu \to B_\mu + \delta B_\mu$ with

$$\delta B^c_\mu(x) T^c = -f^{abc}\epsilon^a(x) B^b_\mu(x) T^c - \frac{1}{g}\partial_\mu \epsilon^b(x) T^c. \tag{15}$$

The mass terms for the W and Z boson, $m_W^2 W^+_\mu W^-_\mu$ and $m_Z^2 Z_\mu Z_\mu$ are not invariant in the SM because the last term in (15) is present.

For quarks and leptons, their mass terms $m_f \overline{\Psi}\Psi = m_f \overline{\Psi_L}\Psi_R$ are not invariant under both $U(1)_Y$ and $SU(2)_L$ transformations. In particular, the mass term will be invariant only if its total hypercharge is zero. Since Ψ_L and Ψ_R have different hypercharge, (because left-handed leptons and quarks are grouped in a doublet and their isospin is not zero) their mass terms are not allowed in the theory's Lagrangian. A mechanism is needed to provide mass to both gauge fields and fermions in a non trivial way.

IV. THE HIGGS MECHANISM

The spontaneous symmetry breaking mechanism responsible for generating mass in the SM was proposed by Higgs [1], Brout, Englert [11], Guralnik, Hagen and Kibble [12]. The Higgs mechanism mixes gauge invariance (which brings to the theory massless bosons) with spontaneous symmetry breaking (SSB).One of the goals is to give mass to the $W^\pm$ and Z boson, while keeping the photon massless and electrodynamics an exact symmetry. For three masses to be generated, at least three new degrees of freedom are needed. The simplest choice to achieve this task is to consider an $SU(2)$ scalar doublet

$$\Phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} \qquad \text{with} \qquad Y(\phi) = +1.$$

The most general invariant and renormalizable Lagrangian for this field is

$$\mathcal{L} = (D^\mu \Phi)^\dagger (D_\mu \Phi) - V(\Phi) \tag{16}$$

with

$$V(\Phi) = \mu^2 \Phi^\dagger \Phi + \frac{\lambda}{4}\left(\Phi^\dagger \Phi\right)^2. \tag{17}$$

If both terms are positive, then $V(\Phi)$ is positive and the minimum of the potential is at $\langle \phi^0 \rangle = \langle \phi^+ \rangle = 0$. This would be the Lagrangian of a spin zero particle with mass μ. If, for example, $\mu^2 < 0$, the neutral part of the doublet Φ can acquire a non-vanishing minimum energy value,

$$\langle \phi \rangle = \langle 0|\phi|0 \rangle = \begin{pmatrix} 0 \\ \frac{v_d}{\sqrt{2}} \end{pmatrix} \tag{18}$$

where v_d is the vacuum expectation value (VEV).It is convenient to redefine the Higgs doublet in a way that the unphysical fields are eliminated through a gauge transformation (unitary gauge), i.e.

$$\Phi = \begin{pmatrix} 0 \\ \frac{1}{\sqrt{2}}\left(v_d + H(x)\right) \end{pmatrix}. \tag{19}$$

The Higgs field has been shifted so that the minimum of the potential is zero. We can rewrite Eq. (16) and (17) as a function of the redefined Higgs doublet and expand the kinetic term to extract the following boson masses: $m_{h^0} = \sqrt{\frac{-\mu^2}{2}} = \sqrt{\frac{\lambda v_d^2}{4}}$, $m_W = \frac{g v_d}{2}$, $m_A = 0$ and $m_Z = \frac{v_d\sqrt{g^2+g'^2}}{2}$. This shows that the existence of the Higgs VEV combined with the covariant derivatives lead to mass terms for the gauge bosons, except for the photon. The above masses are not free parameters, as they depend on v_d and on the gauge couplings. After SSB, the Lagrangian is not invariant to (1). The electroweak symmetry group is broken spontaneously into the electromagnetic subgroup $U(1)_Q$,

$$SU(3)_C \times SU(2)_L \times U(1)_Y \to SU(3)_C \times U(1)_Q, \tag{20}$$

Even after SSB, a $U(1)$ transformation of the form $\Psi \to e^{iQ\eta(x)}\Psi$, with Q defined as $Q = T^3 + Y$, leaves the Lagrangian invariant. Three of the four Higgs doublet degrees of freedom are absorbed by the longitudinal polarization of the $W^\pm$ and Z bosons "making" them acquire mass. The remaining degree of freedom corresponds to a scalar particle, the Higgs boson.

V. FERMION MASSES AND MIXING

Gauge-invariant mass terms for fermions cannot be written since they involve a coupling between left and right-handed fields, which are doublets and singlets of $SU(2)$, respectively. The product $\overline{L_L} l_R$ has hypercharge $Y = -1$.Since the Higgs field is a doublet with hypercharge $Y = 1$, the coupling of both is invariant to $U(1)_Y$ and $SU(2)_L$. We can write the Yukawa term for leptons

$$\mathcal{L} = -\sum_{\alpha,\beta=e,\mu,\tau} Y^l_{\alpha\beta}\overline{L_{\alpha L}}\phi l_{\beta R} + \text{ H.c}, \tag{21}$$

where the matrix Y^l is a general complex matrix, usually known as Yukawa coupling matrix. Upon

SSB, the Yukawa term becomes, in the unitary gauge,

$$\mathcal{L} = -\left(\frac{v_d + H}{\sqrt{2}}\right) \sum_{\alpha,\beta=e,\mu,\tau} Y^l_{\alpha\beta} \overline{l_{\alpha L}} l_{\beta R} + \text{ H.c..} \tag{22}$$

The term proportional to v_d is a mass term for charged-leptons. In order to obtain the fermion masses, we need to diagonalize Y^l as $Y^l \to V_L^{l\dagger} Y^l V_R^l$. In doing so, we redefine the lepton flavour states into mass states with $l_L \to V_L^{l\dagger} l_L$ and $l_R \to V_R^{l\dagger} l_R$ and obtain

$$\mathcal{L} = -\frac{v_d}{\sqrt{2}} \sum_{\alpha,\beta=e,\mu,\tau} y^l_{\alpha\alpha} \overline{l_{\alpha L}} l_{\alpha R} + \text{ H.c.}, \tag{23}$$

a mass term with $m_\alpha = \frac{y^l_\alpha v_d}{\sqrt{2}}$. It is straightforward to show that under these rotations the neutral and electromagnetic currents remain diagonal. There can be, however, a mixing between neutrino and chargedlepton mass eigenstates if neutrinos are not massless due to the appearance of $V_L^{l\dagger} V_L^{\nu_l}$ in charged currents. In the SM, neutrinos are strictly massless since there are no right-handed neutrinos. The rotation of neutrino flavour states is $V_L^{\nu_l} = V_L^l$ so that there is no mixing.

Yukawa terms can also be written for quarks by introducing $\tilde{\phi} = i\tau_2 \phi^*$ with $Y(\tilde{\phi}) = -1$. With this, the quark Yukawa Lagrangian is:

$$-\mathcal{L} = \sum_{\alpha,\beta=d,s,b} Y^D_{\alpha\beta} \overline{Q_{\alpha L}} \phi q^D_{\beta R} + \sum_{\alpha,\beta=u,c,t} Y^U_{\alpha\beta} \overline{Q_{\alpha L}} \tilde{\phi} q^U_{\beta R} \tag{24}$$

where q^D can be any of the d, s or b quarks, q^U is either u, c or t. After SSB, mass terms for quarks can be obtained via the diagonalization of the Yukawa matrices Y^D e Y^U. Thus, there is a rotation of the quarks flavour states, similarly to the lepton case. This implies rewriting the whole Lagrangian in function of the quark mass eigenstates. A mixing matrix $V^{U\dagger} V^D$ will arise in the charged currents, combining rotations from the quarks with different isospin. The matrix $V_{\text{CKM}} = V^{U\dagger} V^D$ is known as the Cabibbo-Kobayashi-Maskawa matrix, which contains new parameters that are unrelated to electroweak symmetry breaking and v_d.

VI. CONSTRAINTS FROM UNITARITY

Setting unitarity constraints is equivalent to imposing that a certain interaction cross-section does not lead to probabilities that sum up to more than one. These unitarity conditions were first used to seek problems with weak interactions at high energies and later to constrain the Higgs boson mass [13]. This leads to an unitary condition of the type:

$$|\text{Re}(a_l)| \le \begin{cases} 1 & \text{final state identical particles} \\ \frac{1}{2} & \text{distinguishable particles} \end{cases} . \tag{25}$$

This unitary condition, when applied, for example, to the partial wave expansion of the scattering process $W^+W^- \to W^+W^-$, leads to the upper bound

$$m_{h^0} < 870\text{GeV}. \tag{26}$$

By analyzing various scattering processes (as done in [6]), one can obtain a slightly more constraining bound, namely

$$m_{h^0} < 710\text{GeV}. \tag{27}$$

This bound was obtained by requiring that the Higgs boson mass does not exceed values obtained from tree-level contributions to the first partial wave decomposition. For high Higgs masses, unitarity is violated at tree level for high energies. This bound might change due to high order corrections or to the existence of new physics.

VII. TRIVIALITY AND STABILITY BOUNDS

Before the Higgs discovery at the LHC [4, 5], experiments at LEP [14] have set upper and lower bonds for the Higgs mass. These where crucial to know in which energy range should the Higgs be looked for in experiments. Apart from experimental limits, there were also theoretical limits coming from triviality and vacuum stability requirements. Such bounds stem from the fact that the parameters of any renormalizable theory change with the energy scale.

The triviality bound for the Higgs boson mass comes from observing that λ (defined in (17)), goes to infinity at a certain scale, for certain initial values. Large values of λ may make the theory non-perturbative. So, we set that the coupling does not exceed a certain an arbitrary value 4π. There is another bound called the stability bound, wich comes from imposing that λ does not become negative. Therefore, the lowest value λ can take is $\lambda = 0$.

In order to determine the triviality and stability bounds, we have to solve the RGEs (taken from [15]) for all the parameters in the theory. Apart from the Higgs quartic coupling and the gauge coupling constants, we take into account only the effect of the top Yukawa coupling. We solved the two-loop RGEs and imposed the triviality ($\lambda = 4\pi$) and the stability ($\lambda = 0$) constraints. The results of our numerical simulations are shown in Fig. 1. From them, we conclude that, for a 125 GeV Higgs,

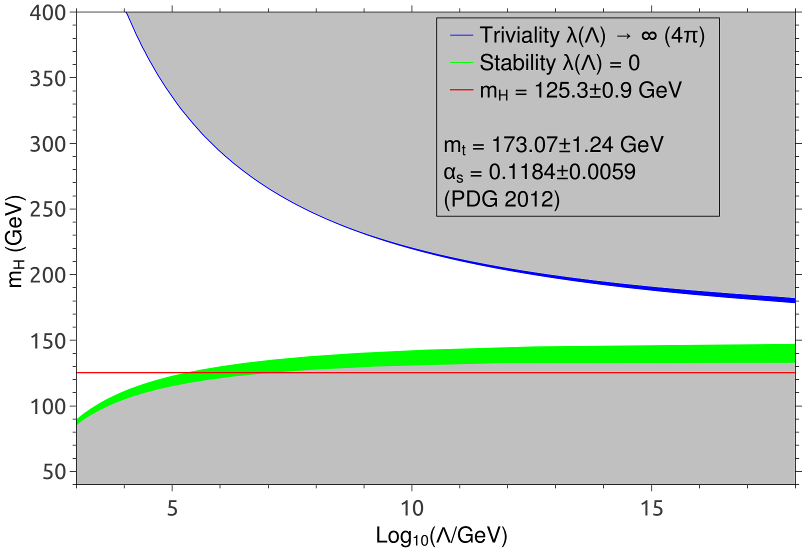

FIG. 1. Triviality and stability bounds for m_{h^0} in the SM, as a function of Λ. The calculations start at the scale $m_Z = 91.2$ GeV. The triviality (stability) bound is shown is blue (green). The grey regions are excluded.

λ will become negative, at an energy scale ranging from 10^5 to 10^7 GeV, if we assume than there is no new physics. This energy interval changes considerably depending on the values and the errors of the top mass ($y_t = m_t\sqrt{2}/v_d$) and the strong gauge coupling ($\alpha_s = g_3^2/4\pi$), as shown by the green band in Fig. 1. Taking the uncertainties into account, our results are compatible with those of [16, 17]. We should, however, remember that this analysis is very sensitive to experimental errors and to the presence of new physics. This can be used, however, to test the consistency of beyond the SM contributions.

VIII. HIGGS DECAYS IN THE STANDARD MODEL

We present the dominant Higgs decays in the SM. The decay widths into massive gauge bosons and fermions are proportional to the tree-level $h^0WW(ZZ)$ and $h^0f\overline{f}$ couplings, respectively. Other decays do not occur at the classical level, but with virtual exchange of particles in a loop; usually fermions and the W. For fermion loops, we only take into account the loop with the top quark, which is dominant. Our calculations will be performed in the unitary gauge. Most of the results were obtained analytically and with the help of the *FeynCalc* package [18]. The decay width formulas obtained are:

$$\begin{aligned}\Gamma(h^0 \to VV^*) = \\ \frac{G_F m_{h^0}^3}{16\sqrt{\pi}}\delta_V\sqrt{1-\tau_V^{-1}}\left[1-\tau_V^{-1}+3\tau_V^{-2}\right] \\ +\frac{3G_F^2 m_V^4}{16\pi^3}m_{h^0}\delta'_V F(\epsilon);\end{aligned} \tag{28}$$

with $\delta_V = 2$ if $V = W$ and 1 if $V = Z$, $\delta'_W = 1$, $\delta'_Z = \frac{7}{12} - \frac{10}{9}\sin^2\theta_W + \frac{40}{9}\sin^4\theta_W$, $\epsilon = \frac{m_V}{m_{h^0}}$ and

$$\begin{aligned}F(\epsilon) = {}& \frac{3(1-8\epsilon^2+20\epsilon^4)}{(4\epsilon^2-1)^{1/2}}\arccos\left(\frac{3\epsilon^2-1}{2\epsilon^3}\right) \\ & -\frac{1-\epsilon^2}{2\epsilon^2}(2-13\epsilon^2+47\epsilon^4) \\ & -3(1-6\epsilon^2+4\epsilon^4)\log\epsilon,\end{aligned} \tag{29}$$

$$\Gamma(h^0 \to ff) = \frac{G_F N_c}{16\sqrt{2}\pi}m_{h^0}m_f^2\left(1-\frac{4m_f^2}{m_{h^0}^2}\right)^{\frac{3}{2}}, \tag{30}$$

with $N_c = 1(3)$ if the fermion in the loop is a lepton(quark);

$$\Gamma(h^0 \to gg) = \frac{G_F m_{h^0}^3\alpha_s^2}{36\sqrt{2}\pi^3}\left|\frac{3}{4}\sum_q A_{1/2}^H(\tau_q)\right|^2, \tag{31}$$

with $\tau_f = \frac{m_{h^0}^2}{4m_f^2}$ and

$$A_{1/2}^H(\tau_f) = 2\tau_f^{-1}\left[1+\left(1-\tau_f^{-1}\right)f(\tau_f)\right], \tag{32}$$

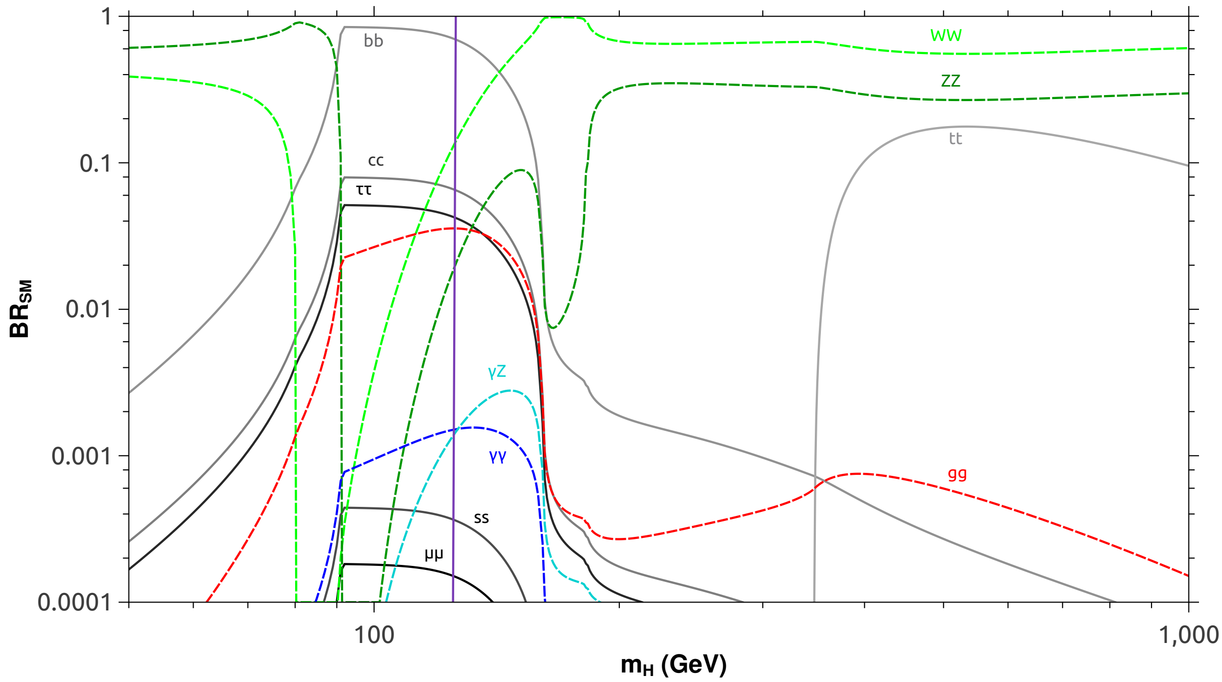

FIG. 2. Dominant decays in the Standard Model. The vertical line, in Purple, is the latest Higgs mass measurement [19]. Full lines are decays into fermions while dashed lines are decays into gauge bosons. For this plot, the following values were used: $m_\mu = 0.10565$ GeV, $m_\tau = 1.77682$ GeV, $m_s = 0.095$ GeV, $m_c = 1.275$ GeV, $m_b = 4.18$ GeV, $m_t = 173.34$ GeV, $m_W = 80.385$ GeV, $m_Z = 91.1876$ GeV, $s_W = 0.23126$, $G_F = 1.166378 \times 10^{-5}$ GeV^{-2}, $\alpha = 1/(137.03599)$ and $\alpha_s = 0.1185$.

$$f\left(\frac{1}{x}\right) = \begin{cases} \arcsin^2\sqrt{1/x} & x \geq 1 \\ -\dfrac{1}{4}\left[\log\dfrac{1+\sqrt{1-x}}{1-\sqrt{1-x}} - i\pi\right]^2 & x < 1 \end{cases} ; \tag{33}$$

$$\Gamma(h^0 \to \gamma\gamma) = \frac{G_F \alpha^2 m_{h^0}^3}{128\sqrt{2}\pi^3} \times \left|\sum_f N_c Q^2 A_{1/2}^H(\tau_f) + A_1^H(\tau_W)\right|^2 , \tag{34}$$

with $\alpha = e^2/4\pi$ and

$$A_1^H(\tau_W) = -2 + 3\tau_W^{-1} + 3(2\tau_W^{-1} - \tau_W^{-2}) f(\tau_W); \tag{35}$$

$$\Gamma(h^0 \to \gamma Z) = \frac{G_F^2 m_W^2 \alpha m_{h^0}^3}{64\pi^4}\left(1 - \frac{m_Z^2}{m_{h^0}^2}\right)^3 \times \left|\sum_f N_f \frac{Q_f \hat{v}_f}{c_W} A_{1/2}^H(\tau_f, \kappa_f) + A_1^H(\tau_W, \kappa_W)\right|^2 \tag{36}$$

with $\hat{v}_f = 4g_V^f$, $\kappa_f = m_Z^2/4m_f^2$

$$A_1^H(\tau_W, \kappa_W) = c_W\Bigg\{\left(3 - \frac{s_W^2}{c_W^2}\right) 4m_W^2 C_0(m_W^2) - \left[(1 + 2\tau_W)\frac{s_W^2}{c_W^2} - (5 + 2\tau_W)\right] 4m_W^2 C_2(m_W^2)\Bigg\}, \tag{37}$$

$$A_{1/2}^H(\tau_f, \kappa_f) = -m_f^2\left[C_0(m_f^2) + 4C_2(m_f^2)\right] . \tag{38}$$

These SM dominant Higgs decay widths are represented in Fig. 2.

IX. $h^0 \to \gamma\gamma$ AND $h^0 \to \gamma Z$ IN THE TYPE II SEESAW MODEL

On the 4th of July 2014, the ATLAS and CMS collaboration announced the discovery of a new particle with mass 125.6±0.3 GeV [19]. Until now, all data indicates that this particle matches the SM Higgs Boson. It is interesting to explore scenarios beyond the SM that could contribute to the signal strength of the Higgs decay modes, and compare them to experiment, manely $h^0 \to \gamma\gamma$. One of these models, the one we are going to study, is the so-called Type II seesaw model.

In the SM neutrinos are strictly massless. This is so because there are no right-handed neutrinos and, thus, no Driac mass term for neutrinos can be written. The problem at hand is: how do we extend the SM in order to account for small but non-zero neutrino masses.

It is possible to generate a Majorana mass term

$$\begin{aligned}\mathcal{L}_{Majorana} &= -\frac{1}{2}\overline{(\nu_L)^c} M_{\nu_L}\nu_L + H.c \\ &= -\frac{1}{2}\nu_L^{\intercal} C^{-1} M_{\nu_L}\nu_L + H.c,\end{aligned} \tag{39}$$

after SSB. Here, in which M_{ν_L} is a matrix in the flavour space and C is the charge conjugation matrix defined as $-\gamma_\mu^\intercal = C^{-1}\gamma_\mu C$ (in Dirac space). A bare Majorana neutrino mass term is not allowed in the SM because it is not gauge invariant.

In the Type II seesaw model, a mass term of the type (39) stems from SSB, meaning that we must introduce a scalar just like we did for the remaining fermions. The new scalar must couple with $\nu_L^\intercal C^{-1}\nu_L$ in gauge invariant way .

We introduce the triplet of $SU(2)$

$$\vec{\Delta} = \left(\Delta^1 \ \ \Delta^2 \ \ \Delta^3\right)^\intercal . \tag{40}$$

Since we are dealing with lepton doublets (which lead us to 2×2 $SU(2)$ transformations), it is convenient to write this scalar in the form

$$\begin{aligned}\Delta \equiv \sqrt{2}T^i\Delta^i &\equiv \begin{pmatrix} \Delta^3 & \Delta^1 - i\Delta^2 \\ \Delta^1 + i\Delta^2 & -\Delta^3 \end{pmatrix} \\ &\equiv \begin{pmatrix} \Delta^+/\sqrt{2} & \Delta^{++} \\ \Delta^0 & -\Delta^+/\sqrt{2} \end{pmatrix}. \end{aligned} \tag{41}$$

With well-defined charge, these scalars have $T^3(\Delta^{++}) = 1$, $T^3(\Delta^+) = 0$ and $T^3(\Delta^0) = -1$, and Δ transforms under an $SU(2)$ transformation as

$$\Delta \to U\Delta U^\dagger, \tag{42}$$

allowing us to write the invariant Yukawa term

$$\mathcal{L}_{Yukawa} = -\frac{1}{2}\overline{(\nu_{\alpha L})^c} Y^\Delta_{\alpha\beta}\varepsilon\Delta\nu_{\beta L} + H.c \tag{43}$$

where $Y^\Delta_{\alpha\beta}$ are the Yukawa couplings and

$$\varepsilon = i\tau_2 = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}, \tag{44}$$

The Lagrangian of the scalar sector is

$$\begin{aligned}\mathcal{L} = &(D_\mu\phi)^\dagger(D^\mu\phi) + Tr\left[(D_\mu\Delta)^\dagger(D^\mu\Delta)\right] \\ &-V(\phi,\Delta) + \mathcal{L}_{Yukawa}\end{aligned} \tag{45}$$

where $\mathcal{L}_{Yukawa}$ contains all of the Yukawa sector of the SM. $V(\phi,\Delta)$ is the most general renormalizable and gauge invariant scalar potential defined by

$$\begin{aligned}V(\phi,\Delta) = &-m_H^2\phi^\dagger\phi + \frac{\lambda}{4}(\phi^\dagger\phi)^2 + M_\Delta^2 Tr\left[\Delta^\dagger\Delta\right] \\ &+\lambda_1(\phi^\dagger\phi)Tr\left[\Delta^\dagger\Delta\right] + \lambda_2\left(Tr\left[\Delta^\dagger\Delta\right]\right)^2 \\ &+\lambda_3 Tr\left[\Delta^\dagger\Delta\Delta^\dagger\Delta\right] + \lambda_4\phi^\dagger\Delta\Delta^\dagger\phi \\ &+\left[\mu\left(\phi^\intercal\varepsilon\Delta^\dagger\phi\right) + h.c\right].\end{aligned} \tag{46}$$

After SSB, we redefine the neutral fields as

$$\phi^0 \to \frac{1}{\sqrt{2}}(v_d + H_d + iZ_1), \tag{47}$$

$$\Delta^0 \to \frac{1}{\sqrt{2}}(v_t + H_t + +iZ_2). \tag{48}$$

Through the minimization of the potential we find the parameters M_Δ and m_H to be

$$\begin{aligned}M_\Delta^2 &= \frac{2\mu v_d^2 - \sqrt{2}(\lambda_1+\lambda_4)v_d^2 v_t - 2\sqrt{2}(\lambda_2+\lambda_3)v_t^3}{2\sqrt{2}v_t}, \\ m_H^2 &= \frac{\lambda v_d^2}{2} - \sqrt{2}\mu v_t + \frac{(\lambda_1+\lambda_4)}{2}v_t^2.\end{aligned} \tag{49}$$

The parameter M_Δ will be one of the inputs in the computation of the decay rates. We find the following masses for the scalar particles in this model:

$$m^2_{H^{\pm\pm}} = \frac{\sqrt{2}\mu v_d^2 - \lambda_4 v_d^2 v_t - 2\lambda_3 v_t^3}{v_t}, \tag{50}$$

$$m^2_{H^\pm} = \frac{(v_d^2 + 2v_t^2)(2\sqrt{2}\mu - \lambda_4 v_t)}{2v_t}, \tag{51}$$

$$m^2_{G^\pm} = 0, \tag{52}$$

$$m^2_{h^0} = \frac{1}{2}\left[A + C - \sqrt{(A-C)^2 + 4B^2}\right], \tag{53}$$

$$m^2_{H^0} = \frac{1}{2}\left[A + C + \sqrt{(A-C)^2 + 4B^2}\right], \tag{54}$$

with $A = \lambda v_d^2/4$, $B = 2v_d\left[-\sqrt{2}\mu + (\lambda_1+\lambda_4)v_t\right]$ and $C = \sqrt{2}\mu v_d^2 + 4(\lambda_2+\lambda_3)v_t^3/v_t$ and $m_{H^0} > m_{h^0}$ for most of the parameter space. The scalars $H^\pm$ and $G^\pm$ are mass eigenstates that come from diagonalizing the mass matrix obtained by collecting terms with combinations of $\phi^\pm$ and $\Delta^\pm$. Similarly, the mass eigenstates h^0 and H^0 come from diagonalizing the CP-even mass matrix for the singly-charged scalars. Note that there is also a CP-odd neutral scalar mass matrix that we do not consider since these states will not be relevant for our future analysis.

Motivated by the fact that neutrino masses are very small and by electroweak precision measurements for $\delta\rho$, which provide us an upper bound on v_t, we will henceforth consider the limit $v_t << v_d$. In this limit, the decay rates for $h^0 \to \gamma\gamma$ and $h^0 \to \gamma Z$ depend solely on the parameters λ_1 and λ_4. The relation of these parameters with the masses in our model is given by

$$\left(\lambda_1 + \frac{\lambda_4}{2}\right) = \frac{2}{v_d^2}(m^2_{H^\pm} - M_\Delta^2), \tag{55}$$

$$\lambda_1 = \frac{2}{v_d^2}(m^2_{H^{\pm\pm}} - M_\Delta^2), \tag{56}$$

from which we can see that for $M_\Delta >> v_d$, we have $m_{H^\pm} \simeq m_{H^\pm} \simeq M_\Delta$.

The decays $h^0 \to \gamma\gamma$ and $h^0 \to \gamma Z$ in the Type-II seesaw model involve extra diagrams when compared to the SM, shown in Fig. 3. Our results for $\Gamma(h^0 \to \gamma\gamma)$ and $\Gamma(h^0 \to \gamma Z)$ are Eqs. (58) and (59), respectively, which use the definitions

$$A_0^H(\tau_S) = -\left[\tau_S - f(\tau_S)\right]\tau_S^{-2} \tag{57}$$

and $Q_S^Z = T_S^3 - Q_S s_W^2$. These decay widths agree with those of [20] and [21].

$$\Gamma(h^0 \to \gamma\gamma) = \frac{G_F \alpha^2 m_{h^0}^3}{128\sqrt{2}\pi^3} \left| \sum_f N_c Q^2 A_{1/2}^{h^0}(\tau_f) + A_1^{h^0}(\tau_W) + \frac{(m_{H^\pm}^2 - M_\Delta^2)}{m_{H^\pm}^2} A_0^{h^0}(\tau_{H^\pm}) + 4\frac{(m_{H^{\pm\pm}}^2 - M_\Delta^2)}{m_{H^{\pm\pm}}^2} A_0^{h^0}(\tau_{H^{\pm\pm}}) \right|^2, \tag{58}$$

$$\Gamma(h^0 \to \gamma Z) = \frac{G_F^2 m_W^2 \alpha h^{0^3}}{64\pi^4} \left(1 - \frac{m_Z^2}{m_{h^0}^2}\right)^3 \left| \sum_f N_f \frac{Q_f \hat{v}_f}{c_W} A_{1/2}^{h^0}(\tau_f, \kappa_f) + A_1^{h^0}(\tau_W, \kappa_W) - \frac{2Q_{H^\pm}^Z Q_{H^\pm}}{c_W} \frac{(m_{H^\pm}^2 - M_\Delta^2)}{m_{H^\pm}^2} A_0^{h^0}(\tau_{H^\pm}, \kappa_{H^\pm}) - \frac{2Q_{H^{\pm\pm}}^Z Q_{H^{\pm\pm}}}{c_W} \frac{(m_{H^{\pm\pm}}^2 - M_\Delta^2)}{m_{H^{\pm\pm}}^2} A_0^{h^0}(\tau_{H^{\pm\pm}}, \kappa_{H^{\pm\pm}}) \right|^2, \tag{59}$$

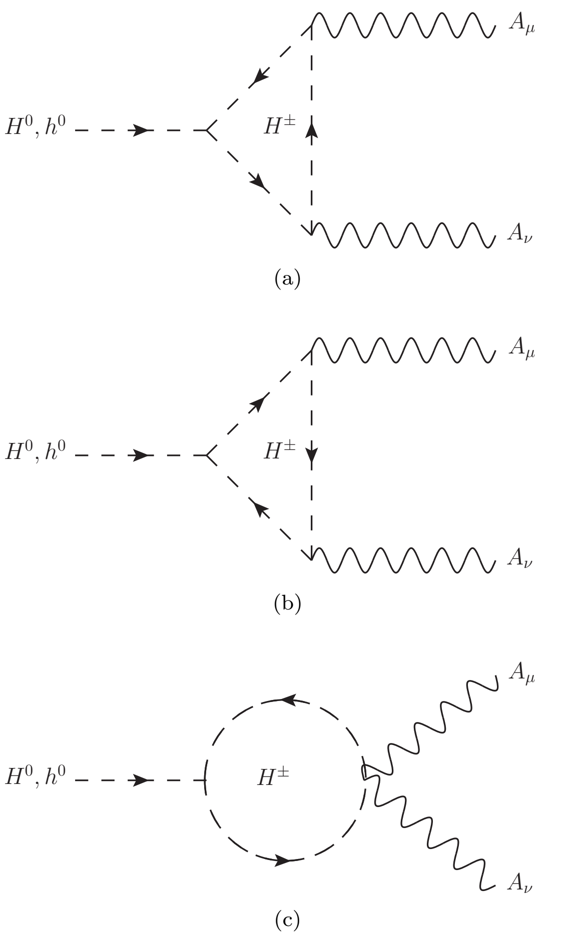

FIG. 3. Extra diagrams involving $H^\pm$ scalar loops for the decay $h^0 \to \gamma\gamma$ in the Type II seesaw model. In the SM, only fermions and the charged bosons run in the loop.

In order to compare the Type-II seesaw model predictions with the ones from the SM, we define the ratios

$$R_{\gamma\gamma}^{\rm exp} = \frac{[\sigma(pp \to H \to \gamma\gamma)]_{\rm exp}}{[\sigma(pp \to H \to \gamma\gamma)]_{\rm SM}} = \frac{[\sigma(pp \to H){\rm Br}(H \to \gamma\gamma)]_{\rm exp}}{[\sigma(pp \to H){\rm Br}(H \to \gamma\gamma)]_{\rm SM}} \tag{60}$$

and

$$R_{\gamma\gamma}^{\rm Type-II} = \frac{[\sigma(pp \to H){\rm Br}(H \to \gamma\gamma)]_{\rm Type-II}}{[\sigma(pp \to H){\rm Br}(H \to \gamma\gamma)]_{\rm SM}} \approx \frac{\Gamma(H \to \gamma\gamma)_{\rm Type-II}}{\Gamma(H \to \gamma\gamma)_{\rm SM}}. \tag{61}$$

The last approximation is due to the fact that the dominant production channels for $pp \to H$ are substantially the same for both models (see [22]).

In the type-II seesaw model, the loops with the new Higgs charged states bring additional contributions to $H \to \gamma\gamma$ as seen in Eq. (58). Due to the factor of four in (58), we expect the contribution from $H^{\pm\pm}$ to dominate over the one from $H^\pm$ for a given spectrum of the charged Higgs masses. The value of M_Δ was chosen to be 200 GeVs so that our results can be compared directly with the ones from [20]. We found agreement between our analysis and theirs.

For $h^0 \to \gamma\gamma$, most of the parameter space allowed by $R_{\gamma\gamma}^{\rm exp}$ for m_{H^+} and $m_{H^{++}}$ happens when M_Δ is larger than the charged-scalar masses, as seen in Fig. 4. If we only account for the $H^\pm$ ($H^{\pm\pm}$) contribution, for $M_\Delta = 200$ GeV, we get a minimal value for the charged-scalar of $m_{H^\pm} \approx 110$ GeV($m_{H^{\pm\pm}} \approx 150$ GeV). The behavior of $R_{\gamma\gamma}^{\rm Type-II}$ for a larger seesaw scale M_Δ is a translation (in the scalar mass axis) of the graph for $M_\Delta = 200$ GeV. In that case, the contributions are finite for $m_{H^+}, m_{H^{++}} \to 0$. Lowering the value of M_Δ leads mostly to large destructive contributions for both singly and doubly charged scalar in both $h^0 \to \gamma\gamma$ and $h^0 \to \gamma Z$ decays. The positive contributions happen for low values of m_{H^+} and $m_{H^{++}}$. Unlike

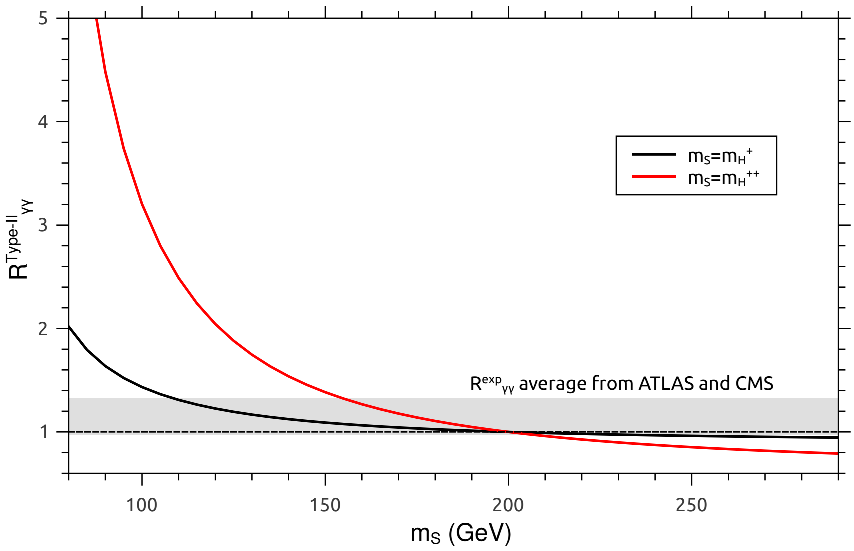

FIG. 4. Comparison between $R_{\gamma\gamma}^{\text{Type-II}}$ with the experimental value $R_{\gamma\gamma}^{\text{exp}} = 1.58^{+0.27}_{-0.23}$ We use $M_\Delta = 200$ GeV, $m_{\text{top}} = 173.34$ GeV, $m_W = 80.385$ GeV, $m_Z = 91.1876$ GeV, $s_w = 0.23126$, $G_F = 1.166378 \times 10^{-5}$ GeV^{-2} and $\alpha = 1/(137.03599)$.

what happens for large M_Δ values, we obtain null Type-II contributions for $m_{H^+}, m_{H^{++}} \to 0$.

We can also see what happens if both charged-scalar contributions are present. In the case in which one of the charged scalars mass is smaller and the other is larger than M_Δ, Ref [23] gives us a constraint for the charged-scalar mass difference $\Delta M \equiv |m_{H^{\pm\pm}} - m_{H^\pm}| \leq 40$ GeV from fits for the electroweak precision observables. Another analysis that can be made is to look at vacuum stability, unitary, and potential stability bounds for this model. In [24] we see that if we want $m_{H^{\pm\pm}}, m_{H^\pm} < M_\Delta$ or $\lambda_1 + \frac{\lambda_4}{2}, \lambda_1 < 0$ to be true, then this to leads an unstable potential when we set the remaining potential parameters to zero ($\lambda = \lambda_2 = \lambda_3 = 0$). When taking into account both contributions, there exists a region in the parameter space were $R_{\gamma\gamma}^{\text{Type-II}}$ is within the experimental bounds, as shown in Fig. 5. There is a large range of possible charged-scalar masses were both contribution cancel each other, giving the SM result, as shown by the dashed line in Fig. 5.

X. CONCLUSIONS

In this work, we reviewed the Standard Model of electroweak interactions and studied the dominant Higgs decays (Fig. 2). Of these, special attention was given to the decay $h^0 \to \gamma\gamma$ in the SM. Its decay width was calculated analytically. This calculation was motivated by its need for the studies in section IX and by recent discussions on how to use dimensional regularization in this decay, specifically whether we should set $d \to 4$ before or after the integration in the loop's momentum. The conclusion to take from this work is that when a divergent integral is present, one should find a method to regularize the integral properly in a general case, before setting approximations, i.e. always work in d dimensions and only take the limit $d \to 4$ in the end.

A new scalar triplet was introduced in section IX, which leads to new contributions to the decays $h^0 \to \gamma\gamma$ and $h^0 \to \gamma Z$ due to the presence of additional charged scalars which couple to h^0, γ and Z. For $h^0 \to \gamma\gamma$, the experimental results for $R_{\gamma\gamma}^{\text{exp}}$ (defined in Eq. 60) can accommodate the inclusion of the new scalar triplet, depending on the masses of the charged scalars and on the remaining parameters of the theory. For an approximation where $v_t << v_d$ and the values of the parameters in the potential (46) are neglected when compared to the magnitude of $\frac{v_d}{v_t}$, the new contributions to the decay width of $h^0 \to \gamma\gamma$ depend only on the masses of the charged scalars. For this approximation, the results are shown in Fig. 4. The inclusion of a new triplet scalar, apart from giving an possible explanation to the smallness of the neutrino masses through a Type-II seesaw mechanism, can be accommodated by the recent experimental value of $h^0 \to \gamma\gamma$ for a large set of both charged-scalars masses, as seen in Fig. 5 for $v_d << v_t$.

The study presented in this thesis illustrates how analysis of Higgs decays can be used to constrain new physics scenarios. As the LHC continues its operation, new measurements will improve our knowledge of the SM electroweak sector and, hopefully, find evidences for physics beyond the SM. Who knows...

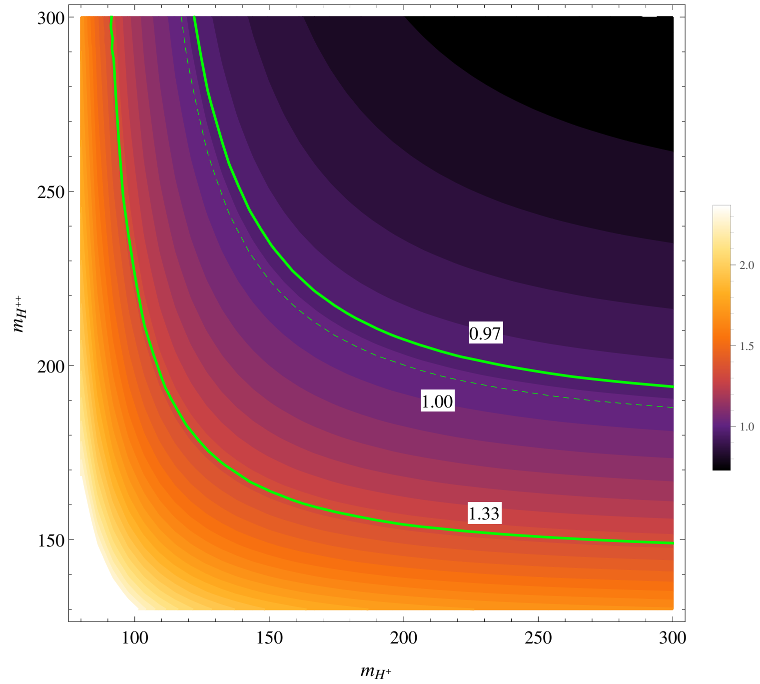

FIG. 5. Contour plot of $R_{\gamma\gamma}^{\text{Type-II}}/R_{\gamma\gamma}^{\text{SM}}$ for contributions of both $H^{\pm}$ and $H^{\pm\pm}$ in the approximation $v_d >> v_t$. The parameter space allowed by the experimental value $R_{\gamma\gamma}^{\text{exp}} = 1.15 \pm 0.18$ is the area between the two green lines.

[1] P. W. Higgs, Phys.Rev. **145**, 1156 (1966).
[2] S. Weinberg, Phys.Rev.Lett. **27**, 1688 (1971).
[3] S. Glashow, Nucl.Phys. **22**, 579 (1961).
[4] G. Aad *et al.* (ATLAS Collaboration), Phys.Lett. **B716**, 1 (2012), arXiv:1207.7214 [hep-ex].
[5] S. Chatrchyan *et al.* (CMS Collaboration), Phys.Lett. **B716**, 30 (2012), arXiv:1207.7235 [hep-ex].
[6] A. Djouadi, Physics Reports **457**, 1 (2008), ISSN 0370-1573, `http://www.sciencedirect.com/science/article/pii/S0370157307004334`.
[7] K. Brading and E. Castellani, Symmetries in Physics, by Katherine Brading, Elena Castellani, Cambridge, UK: Cambridge University Press, 2010 **1** (2010).
[8] S. Weinberg, Phys.Rev.Lett. **19**, 1264 (1967).
[9] G. Arnison *et al.* (UA1 Collaboration), Phys.Lett. **B122**, 103 (1983).
[10] G. Arnison *et al.* (UA1 Collaboration), Phys.Lett. **B126**, 398 (1983).
[11] F. Englert and R. Brout, Phys.Rev.Lett. **13**, 321 (1964).
[12] G. Guralnik, C. Hagen, and T. Kibble, Phys.Rev.Lett. **13**, 585 (1964).
[13] B. W. Lee, C. Quigg, and H. B. Thacker, Phys. Rev. D **16**, 1519 (Sep 1977), `http://link.aps.org/doi/10.1103/PhysRevD.16.1519`.
[14] R. Barate *et al.* (LEP Working Group for Higgs boson searches, ALEPH Collaboration, DELPHI Collaboration, L3 Collaboration, OPAL Collaboration), Phys.Lett. **B565**, 61 (2003), arXiv:hep-ex/0306033 [hep-ex].
[15] H. Arason, D. J. Castaño, B. Kesthelyi, S. Mikaelian, E. J. Piard, P. Ramond, and B. D. Wright, Phys. Rev. D **46**, 3945 (Nov 1992), `http://link.aps.org/doi/10.1103/PhysRevD.46.3945`.
[16] T. Hambye and K. Riesselmann, Phys.Rev. **D55**, 7255 (1997), arXiv:hep-ph/9610272 [hep-ph].
[17] J. Ellis, J. Espinosa, G. Giudice, A. Hoecker, and A. Riotto, Phys.Lett. **B679**, 369 (2009), arXiv:0906.0954 [hep-ph].
[18] R. Mertig, M. Bohm, and A. Denner, Comput.Phys.Commun. **64**, 345 (1991), see http://feyncalc.org/.
[19] J. Beringer et al, Phys. Rev. D **86**, 010001 ((2012) and 2013 partial update for the 2014 edition).
[20] C.-S. Chen, C.-Q. Geng, D. Huang, and L.-H. Tsai, Physics Letters B **723**, 156 (2013).
[21] A. Arhrib, R. Benbrik, M. Chabab, G. Moultaka, and L. Rahili, Journal of High Energy Physics **2012**, 1 (2012).
[22] P. Bhupal Dev, D. K. Ghosh, N. Okada, and I. Saha, JHEP **1303**, 150 (2013), arXiv:1301.3453.
[23] E. J. Chun, H. M. Lee, and P. Sharma, JHEP **1211**, 106 (2012), arXiv:1209.1303 [hep-ph].
[24] A. Arhrib, R. Benbrik, M. Chabab, G. Moultaka, M. Peyranere, L. Rahili, and J. Ramadan, Physical Review D **84**, 095005 (2011).