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MASTER THESIS

Spin-3/2 top quark excitations and Higgs physics

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To my parents

How wonderful that we have met with a paradox.
Now we have some hope of making progress.

N. Bohr

Contents

Chapter 1

Introduction

The Standard Model (SM) has successfully described many of the observed phenomena in high energy physics and is, up to the present, the most accurate theory we have in our understanding of the strong, weak and electromagnetic interactions. With the recent discovery at the Large Hadron Collider (LHC) of a particle with properties featuring the Higgs boson [1, 2], the last piece of the SM is thought to be found. Nevertheless, and since the SM leaves a number of questions unanswered (regarding *e.g.* the origin of dark matter, the masses of the neutrinos or the hierarchy problem), besides completing the SM this new state could also serve as a tool to explore physics beyond the Standard Model (BSM). It is, therefore, a great challenge for particle physicists nowadays to unravel all the New Physics (NP) hidden behind this newly observed particle.

Among the extensions of the SM, many of them contemplate the existence of exotic higher-spin particles. Examples of these are the *graviton* in Supergravity [3], a spin-3/2 partner of the top-quark in Randall-Sundrum models [4], or composite quarks and leptons in compositeness models [5, 6]. Furthermore, in some of these models, if these particles existed they would be expected to leave their footprint in the processes of Higgs production and decay. Thanks to the accuracy of the LHC, deviations from the SM could be detected experimentally providing indirect evidence of the existence of such particles and constituting a strong support to the underlying theory.

This work is, to the best of our knowledge, the first attempt to combine Higgs phenomenology with higher-spin fields. We speculate that the deviations from the SM of the $gg \to H \to \gamma\gamma$ partonic cross section observed at the LHC [7, 8] are caused by a spin-3/2 partner of the top-quark, similar to that presented in [4]. For the sake of generality, no concrete theory is chosen and the description of the spin-3/2 particle is carried out combining group theory formalism and an effective theory based approach, with an arbitrary high scale Λ. Consequently, our results can be easily generalized to any spin-3/2 field that couples to a massive spin-1/2 and a gauge field.

The script is organized as follows. In Chapter 2 we review the basics of the spontaneous symmetry breaking in the electroweak sector and derive the partial decay widths at leading order (LO) of the SM Higgs boson to photons and gluons. Chapter 3 is dedicated to the formalism of the propagation of spin-3/2 particles and their interaction with a spin-1/2 fermion and a gauge boson. In Chapter 4 the calculations and concepts from Chapters 2 and 3 are united to study the effect of a spin-3/2 excitation of the top-quark on Higgs production from gluon fusion and its decay to photons. In Chapter 5 we present the conclusions.

Chapter 2

The Higgs Boson in the Standard Model

A cornerstone of the SM is the mechanism of electroweak symmetry breaking (EWSB) which accounts for the weak bosons and the fermions to acquire their masses in a "clean" way, without violating the requirements of renormalizability and unitarity of the theory. This mechanism implies the existence of a new fundamental particle: the Higgs boson[1].

We start this chapter with a brief review on how the EWSB works and why it is crucial in the SM. Afterwards, we derive the Higgs decay into photons via a fermion loop and the decay into gluons, which will be useful in our future BSM calculations.

2.1 Review of the Electroweak Symmetry Breaking

This section is based on the reviews [9, 10, 11].

The Glashow-Weinberg-Salam electroweak theory is a Yang-Mills theory based on the symmetry group $SU(2)_L \times U(1)_Y$ which describes the weak and electromagnetic interactions between quarks and leptons mediated by gauge bosons. The coupling of these bosons to the matter fields is achieved mathematically through the covariant derivative (for the sake of simplicity, we do not consider here the strong interactions)

$$D_\mu \psi = \left(\partial_\mu - igT_a W^a_\mu - ig' \frac{Y}{2} B_\mu \right) \psi \,, \tag{2.1}$$

where ψ is some matter field, g and g' are respectively the $SU(2)$ and the $U(1)$ coupling constants, T_a with $a = 1, 2, 3$ are the generators of $SU(2)$,

[1] The reader is referred to some up-to-date reviews [9, 10, 11] and standard textbooks [12, 13, 14].

W_μ^a are the $SU(2)$ boson and B_μ the $U(1)$ boson fields in the interaction eigenbasis and Y is the hypercharge. The summation over repeated indices is implied.

In the Lagrangian, the kinetic and gauge interaction terms of the gauge bosons are

$$\mathcal{L}_K = -\frac{1}{4} W_{\mu\nu}^a W_a^{\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} \tag{2.2}$$

with

$$W_{\mu\nu}^a = \partial_\mu W_\nu^a - \partial_\nu W_\mu^a + g\, \epsilon_{bc}^a W_\mu^b W_\nu^c \,, \tag{2.3a}$$

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu \,, \tag{2.3b}$$

where ϵ_{bc}^a is an antisymmetric tensor.

These terms are invariant under local $SU(2)_L \times U(1)_Y$, as required. However, notice that there is no mass term for the bosons. While this is true for the case of the photons, which are massless, it is not for the weak bosons, with masses of the order of 100 GeV. Solving this problem is not trivial, for introducing by hand a mass term (of the form $M_W^2 W_{\mu\nu}^a W_a^{\mu\nu}$) does not preserve the symmetry and then the theory would not be renormalizable. The same thing happens with the fermions. Adding explicitly a mass term ($-m_f \bar{\psi}_f \psi_f$) breaks the invariance of the Lagrangian under the gauge transformation. The answer is what we know as Brout-Englert-Higgs mechanism (or Higgs mechanism, for short) [15, 16].

Consider a $SU(2)$ doublet of scalar fields

$$\Phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}, \quad Y_\Phi = +1 \tag{2.4}$$

which will couple to the gauge fields via $(D^\mu\Phi)^\dagger (D_\mu\Phi)$. Consider also a potential, function of this doublet

$$V(\Phi) = \mu^2 \Phi^\dagger \Phi + \lambda \left(\Phi^\dagger \Phi\right)^2 . \tag{2.5}$$

Let us call $\mathcal{L}_S$ the Lagrangian containing this scalar field:

$$\mathcal{L}_S = (D^\mu\Phi)^\dagger (D_\mu\Phi) - V(\Phi) \,. \tag{2.6}$$

We add this to the SM Lagrangian,

$$\mathcal{L}_{SM} = \mathcal{L}_K + (D^\mu\Phi)^\dagger (D_\mu\Phi) - V(\Phi) + \ldots \tag{2.7}$$

The parameter λ must be larger than zero in order for the potential to be bounded from below. From Fig. 2.1, for $\mu^2 > 0$ the potential has only the minimum $\Phi = 0$ and $\mathcal{L}_S$ describes a scalar particle with mass μ and

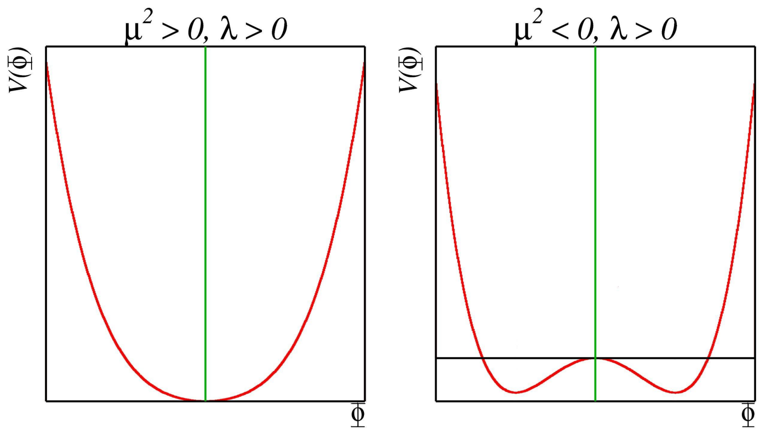


Figure 2.1: *Higgs potential in two dimensions before (left, $\mu^2 > 0$) and after spontaneous symmetry breaking (right, $\mu^2 < 0$).*

quartic coupling. If $\mu^2 < 0$, the potential $V(\Phi)$ has an infinite number of minima, given by the condition

$$|\Phi| = \sqrt{\frac{-\mu^2}{2\lambda}} \equiv \frac{v}{\sqrt{2}} \, . \tag{2.8}$$

The scalar field Φ is said to acquire a vacuum expectation value (vev)

$$\langle\Phi\rangle_0 = \langle 0|\Phi|0\rangle = \begin{pmatrix} 0 \\ \frac{v}{\sqrt{2}} \end{pmatrix} \, . \tag{2.9}$$

The vev must be developed in the neutral component, so that $U(1)_Q$ remains unbroken. Expanding the field around this vev

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} \theta_2 + i\theta_1 \\ v + h - i\theta_3 \end{pmatrix} \, . \tag{2.10}$$

Choosing a suitable gauge, called *unitary* gauge, one can get rid of the θ_i fields:

$$\Phi \longrightarrow \Phi'(x) = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h \end{pmatrix} \, . \tag{2.11}$$

Redefining the fields as

$$W^\pm_\mu = \frac{1}{\sqrt{2}}(W^1_\mu \mp iW^2_\mu) \, , \tag{2.12a}$$

$$Z_\mu = \frac{gW^3_\mu - g'B_\mu}{\sqrt{g^2 + g'^2}} \, , \tag{2.12b}$$

$$A_\mu = \frac{gW^3_\mu + g' B_\mu}{\sqrt{g^2 + g'^2}} \tag{2.12c}$$

and expanding the term $(D^\mu \Phi)^\dagger (D_\mu \Phi)$, one finds mass terms for the W and the Z bosons, while the photon remains massless.

$$M_W = \frac{1}{2} v g', \quad M_Z = \frac{1}{2} v \sqrt{g^2 + g'^2}, \quad M_A = 0 \,. \tag{2.13}$$

Finally, the mass of the fermions are generated through the Yukawa couplings

$$\mathcal{L}_F = -\lambda_e \bar{L} \Phi e_R - \lambda_d \bar{Q} \Phi d_R - \lambda_u \bar{Q} \tilde{\Phi} u_R + H.C. \,, \tag{2.14}$$

where $\tilde{\Phi}$ is the isodoublet $\tilde{\Phi} = i\tau_2 \Phi^*$ and we used the notation from [9] for the weak isodoublets and isosinglets:

$$\begin{aligned}
L_1 &= \begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L, e_{R_1} = e^-_R, \quad Q_1 = \begin{pmatrix} u \\ d \end{pmatrix}_L, u_{R_1} = u_R, \quad d_{R_1} = d_R \\
L_2 &= \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L, e_{R_1} = \mu^-_R, \quad Q_1 = \begin{pmatrix} c \\ s \end{pmatrix}_L, u_{R_2} = c_R, \quad d_{R_2} = s_R \\
L_3 &= \begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L, e_{R_3} = \tau^-_R, \quad Q_3 = \begin{pmatrix} t \\ b \end{pmatrix}_L, u_{R_3} = t_R, \quad d_{R_3} = b_R \,.
\end{aligned} \tag{2.15}$$

2.2 Higgs Decay in the Standard Model

2.2.1 Higgs decay into Photons

The decay of the Higgs boson into photons in the SM at LO involves fermion and W-boson loops. Due to its large mass, as the coupling to the Higgs is proportional to the mass, the contribution from the top (t) quark to the total amplitude is much larger than that from any other fermion. Consequently, to a good approximation we will neglect the contributions from the other fermions [9][2]. In this section we compute the matrix elements involving top-quark loops; the ones with W-boson loops will be taken from the literature [9, 17, 18].

There are two diagrams contributing to the decay into photons mediated by the t-quark:

[2] The procedure we follow here can be generalized to any other fermion just by making $m_t \to m_f$ and finally summing over all possible fermions [9].

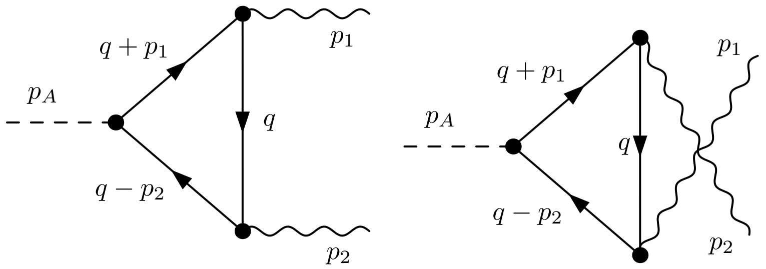


Diagram a) *Diagram b)*

Table 2.1: *Diagrams involving a t-quark loop at LO in the process of Higgs decay to photons. In our conventions p_A is incoming, p_1 and p_2 are outgoing. q is the loop momentum.*

Here, p_A is incoming and p_1, p_2 are outgoing. q is the loop momentum. We start computing the matrix element of *Diagram a)*, call it $\mathcal{M}_a$. Following the Feynman rules [12]:

$$\begin{aligned}\mathcal{M}_a &= -N_c \int \frac{d^4q}{(2\pi)^4} \left(-\frac{iem_t}{2\sin\theta_W M_W} \right) \left(i\frac{\not{q} - \not{p}_2 + m_t}{(q-p_2)^2 - m_t^2} \right) (-iee_q)\, \gamma_\nu \\ &\times \left(i\frac{\not{q} + m_t}{q^2 - m_t^2} \right) (-iee_q)\, \gamma_\mu \left(i\frac{\not{q} + \not{p}_1 + m_t}{(q+p_1)^2 - m_t^2} \right) \varepsilon^{\mu *}(p_1, \lambda_1) \varepsilon^{\nu *}(p_2, \lambda_2) \\ &= -\frac{N_c e^3 e_q^2 m_t}{2\sin\theta_W M_W} \varepsilon^{\mu *}(p_1, \lambda_1) \varepsilon^{\nu *}(p_2, \lambda_2) \\ &\times \int \frac{d^4q}{(2\pi)^4} \frac{\mathrm{Tr}\left[(\not{q} - \not{p}_2 + m_t) \gamma_\nu (\not{q} + m_t) \gamma_\mu (\not{q} + \not{p}_1 + m_t) \right]}{\left[(q-p_2)^2 - m_t^2 \right] \left[q^2 - m_t^2 \right] \left[(q+p_1)^2 - m_t^2 \right]} \,. \end{aligned} \tag{2.16}$$

In this expression θ_W is the weak mixing angle, M_W the mass of the W-boson, e_q is the charge of the t-quark in units of the electron charge e, λ_1 and λ_2 are the polarizations of the outgoing photons and γ_μ are the gamma matrices[3]. We have to integrate over the loop momentum, q, and include a global minus sign for we have a closed fermion line.

The factor N_c is the colour factor, which in this case is trivial. As there is no change in colour in any of the vertices, the colour factor is the number of colours of the quarks, which will be left as a parameter.

At first sight, this integral would be divergent: there will be some term in the numerator proportional to q^2 (no term proportional to q^3 will arise, as the trace of an odd number of γ's is zero) that will yield

$$\int \frac{d^4q\, q^2}{q^6} \sim \int \frac{dq}{q} \,,$$

[3]In our conventions $\eta_{\mu\nu} = diag(1, -1, -1, -1)$ and $\{\gamma^\mu, \gamma^\nu\} = 2\eta^{\mu\nu}$.

which is logarithmically divergent. So we need to introduce a dimensional regularization: 4→ n.

$$\mathcal{M}_a = -\frac{N_c e^3 e_q^2 m_t}{2\sin\theta_W M_W}\varepsilon^{\mu *}(p_1,\lambda_1)\varepsilon^{\nu *}(p_2,\lambda_2)\mu^{4-n}\int \frac{d^n q}{(2\pi)^n}\frac{Tr\,[\ldots]}{\mathcal{D}_1\mathcal{D}_2\mathcal{D}_3}\ , \quad (2.17)$$

where

$$\mathcal{D}_1 = q^2 - m_t^2,\quad \mathcal{D}_2 = (q-p_2)^2 - m_t^2,\quad \mathcal{D}_3 = (q+p_1)^2 - m_t^2. \quad (2.18)$$

This matrix element is computed step by step in Appendix A. After some calculations one arrives at

$$\mathcal{M}_a = \frac{N_c i e^3 e_q^2}{32\pi^2 \sin\theta_W M_W} A_{1/2}^H(\tau_t)\varepsilon^{\mu *}(p_1,\lambda_1)\varepsilon^{\nu *}(p_2,\lambda_2)\left(p_{1\nu}p_{2\mu} - \eta_{\mu\nu}\frac{m_H^2}{2}\right) \quad (2.19)$$

being $A_{1/2}^H(\tau_t)$ defined in Appendix A as

$$A_{1/2}^H(\tau_t) = 2\left[\tau_t + (\tau_t - 1)\arcsin^2\sqrt{\tau_t}\right]\tau_t^{-2} \quad (2.20)$$

with $\tau_t = m_H^2/4m_t^2$.

Before squaring and averaging Eq.(2.19), recall that this is not the matrix element of the process, but that of the first diagram, Fig.2.1. At this point we have to compute the matrix element of the second diagram, $\mathcal{M}_b$, but we will see that there is no need no repeat the whole calculation. Notice that both diagrams are the same if we exchange the outgoing photons, therefore, the matrix element, $\mathcal{M}_b$ is equal to $\mathcal{M}_a$ switching $(p_1,\lambda_1)\leftrightarrow(p_2,\lambda_2)$. However, this yields exactly the same result, it can be easily seen by switching also the indices $\mu\leftrightarrow\nu$ in Eq.(2.19), which does not modify the result. In conclusion, we can assert that $\mathcal{M}_a = \mathcal{M}_b$ and if we call $\mathcal{M}_{H\to\gamma\gamma}^{(t)}$ the matrix element of the Higgs decay via t-quarks loops,

$$\begin{aligned}\mathcal{M}_{H\to\gamma\gamma}^{(t)} &= \mathcal{M}_a + \mathcal{M}_b = 2\mathcal{M}_a \\ &= \frac{N_c i e^3 e_q^2}{16\pi^2\sin\theta_W M_W}A_{1/2}^H(\tau_t) \\ &\times \varepsilon^{\mu *}(p_1,\lambda_1)\varepsilon^{\nu *}(p_2,\lambda_2)\left(p_{1\nu}p_{2\mu} - \eta_{\mu\nu}\frac{m_H^2}{2}\right)\ . \end{aligned} \quad (2.21)$$

The total matrix element of the process is the sum of this plus the contributions from the W-bosons: $\mathcal{M}_{H\to\gamma\gamma} = \mathcal{M}_{H\to\gamma\gamma}^{(t)} + \mathcal{M}_{H\to\gamma\gamma}^{(W)}$. The latter is far from trivial and a detailed calculation of it is beyond the scope of this work. We take it from [18], with the notation from [9]:

$$\mathcal{M}^{(W)}_{H\to\gamma\gamma} = \frac{ie^3}{16\pi^2 \sin\theta_W M_W}\ A_1^H(\tau_W) \\ \times \left(p_{1\nu}p_{2\mu} - \eta_{\mu\nu}\frac{m_H^2}{2}\right)\varepsilon^{\mu *}(p_1,\lambda_1)\varepsilon^{\nu *}(p_2,\lambda_2)\ , \tag{2.22}$$

being $A_1^H(\tau_W)$ defined as [9]

$$A_1^H(\tau_W) = -\left[2\tau_W^2 + 3\tau_W + 3(2\tau_W - 1)\arcsin^2\sqrt{\tau_W}\right]\tau_W^{-2} \tag{2.23}$$

with the parameter, $\tau_W = m_H^2/4M_W^2$, again smaller than one.

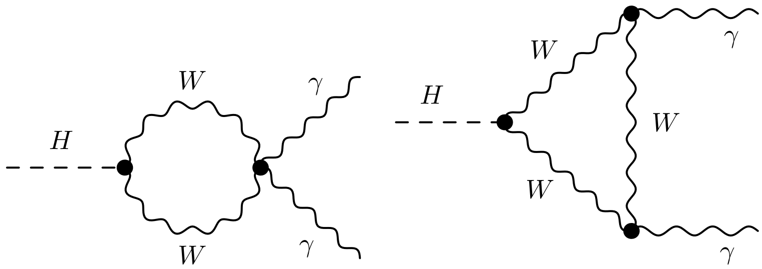


Table 2.2: *Some diagrams with W-loops that contribute to the process of Higgs decays to photons.*

The total matrix element of Higgs decay into photons is then

$$\mathcal{M}_{H\to\gamma\gamma} = \frac{ie^3}{16\pi^2 \sin\theta_W M_W}\left(N_c e_q^2 A_{1/2}^H(\tau_t) + A_1^H(\tau_W)\right) \\ \times \left(p_{1\nu}p_{2\mu} - \eta_{\mu\nu}\frac{m_H^2}{2}\right)\varepsilon^{\mu *}(p_1,\lambda_1)\varepsilon^{\nu *}(p_2,\lambda_2)\ , \tag{2.24}$$

which is the quantity that has to be squared and averaged:

$$\overline{|\mathcal{M}|^2} = \frac{e^6}{256\pi^4 \sin^2\theta_W M_W^2}\left|N_c e_q^2 A_{1/2}^H(\tau_t) + A_1^H(\tau_W)\right|^2 \\ \times \sum_{\lambda_1\lambda_2} \varepsilon^{\mu *}(p_1,\lambda_1)\varepsilon^{\nu *}(p_2,\lambda_2)\varepsilon^{\rho}(p_1,\lambda_1)\varepsilon^{\sigma}(p_2,\lambda_2) \\ \times \left(p_{1\nu}p_{2\mu} - \eta_{\mu\nu}\frac{m_H^2}{2}\right)\left(p_{1\sigma}p_{2\rho} - \eta_{\sigma\rho}\frac{m_H^2}{2}\right)\ . \tag{2.25}$$

In the Feynman gauge, according to the Ward identity, the sum of the ε's over λ_1, λ_2 can be replaced by $\eta^{\mu\rho}\eta^{\nu\sigma}$, so the tensor contraction is equal to

$$\left(p_{1\nu}p_{2\mu} - \eta_{\mu\nu}\frac{m_H^2}{2}\right)\left(p_1^\nu p_2^\mu - \eta^{\mu\nu}\frac{m_H^2}{2}\right) \\ = -2\eta_{\mu\nu}p_1^\nu p_2^\mu\frac{m_H^2}{2} + \eta_{\mu\nu}\eta^{\mu\nu}\frac{m_H^4}{4} = \frac{m_H^4}{2}\ . \tag{2.26}$$

And the averaged squared matrix element:

$$\overline{|\mathcal{M}|^2} = \frac{e^6 m_H^4}{512\pi^4 \sin^2\theta_W M_W^2} \left| N_c e_q^2 A_{1/2}^H(\tau_t) + A_1^H(\tau_W) \right|^2 \tag{2.27}$$

$$= \frac{N_c^2 \alpha^3 m_H^4}{8\pi \sin^2\theta_W M_W^2} \left| N_c e_q^2 A_{1/2}^H(\tau_t) + A_1^H(\tau_W) \right|^2 . \tag{2.28}$$

This is related to the partial decay width via the formula

$$\Gamma = \frac{1}{2} \frac{|\vec{p}|}{8\pi m_H^2} \overline{|\mathcal{M}|^2} , \tag{2.29}$$

where $|\vec{p}|$ is the modulus of the momentum of any of the final particles in the reference system in which the initial particle is at rest. A trivial calculation shows that $|\vec{p}| = m_H/2$. The symmetry factor 1/2 at the beginning has to be included because the two final particles are identical. Therefore

$$\Gamma_{H\to\gamma\gamma} = \frac{\alpha^3 m_H^3}{256\pi^2 \sin^2\theta_W M_W^2} \left| N_c e_q^2 A_{1/2}^H(\tau_t) + A_1^H(\tau_W) \right|^2 . \tag{2.30}$$

This result is often expressed in terms of G_F. Performing the change

$$\frac{\alpha}{M_W^2 \sin^2\theta} = \frac{2G_F}{\pi\sqrt{2}} , \tag{2.31}$$

we get

$$\Gamma_{H\to\gamma\gamma} = \frac{G_F m_H^3}{128\sqrt{2}\pi} \left(\frac{\alpha}{\pi}\right)^2 \left| N_c e_q^2 A_{1/2}^H(\tau_t) + A_1^H(\tau_W) \right|^2 , \tag{2.32}$$

which coincides with [9, 17, 18].

2.2.2 Higgs decay into Gluons

The partial decay width of the process $H \to gg$ can be easily derived from the previous one. First of all, since the weak bosons have no colour, they do not couple to the gluons, so there is no contribution coming from W-loops. Looking at Eq.(2.32), we need to change $e_q e \to g_s$ or, equivalently $\alpha e_q^2 \to \alpha_s$; make $A_1^H = 0$; divide by the previous colour factor and multiply by the new one. The new colour factor is

$$\sum_{a,b} \left(Tr\left[t^a t^b\right] \right)^2 = \frac{1}{4} \sum_{a,b} \delta^{ab} = \frac{N_c^2 - 1}{4} . \tag{2.33}$$

So the decay rate is

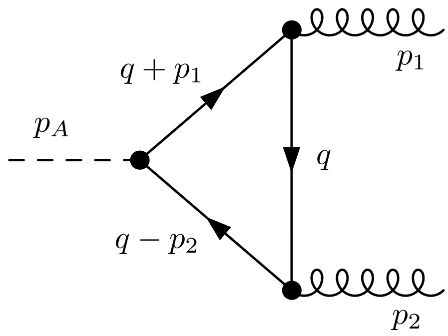


Figure 2.2: *Diagram of the process $H \to gg$ (plus gluons crossed). The momenta convention is the same as in the decay to photons.*

$$\Gamma_{H\to gg} = \frac{(N_c^2 - 1)G_F m_H^3}{512\sqrt{2}\pi} \left(\frac{\alpha_s}{\pi}\right)^2 \left|A_{1/2}^H(\tau_t)\right|^2 . \tag{2.34}$$

Substituting $N_c = 3$ we get the same result as [9, 13, 17].

Finally, we show in Fig. 2.3 the SM prediction for the Higgs branching ratios plotted against the Higgs mass M_H. Quantum Chromodynamics (QCD) next-to-leading order (NLO) corrections have been taken into account. As explained in [9], these can be quite important: they can be of the order of the LO contributions. At a mass of around 125 GeV, at which the new particle has been found [1, 2], the branching ratio of the decay to gluons is of the order of 100 times larger than the branching ratio to photons. Comparing Eq.(2.32) and Eq.(2.34) this difference is due to the fact that the electromagnetic constant, α, is around 10 times smaller than the strong coupling, α_s, so that when squared gives a difference of order 100. Even though, for experimental purposes, the decay to photons is more convenient, since it produces a clearer signal. The reason for this is that the outgoing photons can be directly measured, whereas the gluons have to be detected in the form of gluon-initiated jets of hadrons. Since gluons are coloured particles, according to QCD confinement, they can not exist in free form. Therefore, the outgoing gluons will interact with quarks, antiquarks and other gluons to produce a narrow cone of hadrons known as jet. These jets are what will be measured in the detectors and their signal will have to be examined to determine the properties of the original gluons.

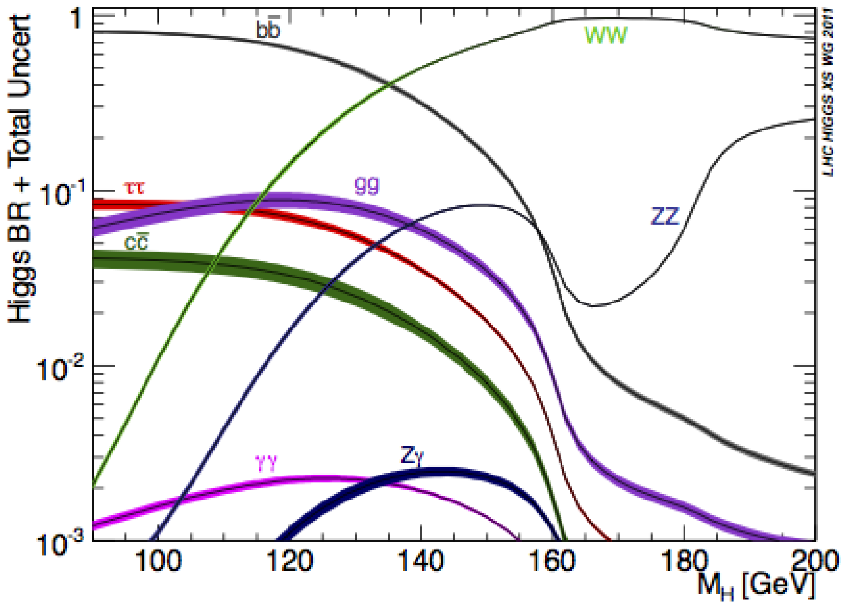


Figure 2.3: *SM predictions for the Higgs branching ratios are plotted against the mass. For $H \to \gamma\gamma$ and $H \to gg$ NLO QCD corrections are included. At a mass of M_H = 125 GeV, the branching ratio to gluons is around 100 times larger than that to photons. In spite of that, the decay to photons is more convenient for experimental purposes. Image taken from [19].*

Chapter 3

Models for Spin-3/2 Top Excitations

A possible spin-3/2 excitation has been widely studied in the literature and is expected to provide insights into New Physics [3, 4, 20, 21]. The first to derive a mathematical formalism for spin-3/2 fields were Rarita and Schwinger [23]. In their pioneer work, these fields are described by a Dirac spinor with a Lorentz index, that we will denote as ψ^ν.

In group theory, the so-called Rarita-Schwinger fields come from the direct product of two representations of the Lorentz group: a vector $\left(\frac{1}{2},\frac{1}{2}\right)$ and a Dirac spinor $\left(\frac{1}{2},0\right)\oplus\left(0,\frac{1}{2}\right)$ [20][1]. Performing this direct product yields

$$\left(\frac{1}{2},\frac{1}{2}\right)\otimes\left[\left(\frac{1}{2},0\right)\oplus\left(0,\frac{1}{2}\right)\right]=\left(\frac{1}{2},1\right)\oplus\left(\frac{1}{2},0\right)\oplus\left(1,\frac{1}{2}\right)\oplus\left(0,\frac{1}{2}\right). \tag{3.1}$$

Notice that in this product a Dirac spinor $\left(\frac{1}{2},0\right)\oplus\left(0,\frac{1}{2}\right)$ shows up. Moreover, another pair of Dirac spin-1/2 fermions is present in the representation $\left(\frac{1}{2},1\right)\oplus\left(1,\frac{1}{2}\right)$. Hence, in order to isolate the spin-3/2 particle, we need to impose the conditions

$$\gamma^\nu\psi_\nu = 0\ , \tag{3.2a}$$

$$\partial^\nu\psi_\nu = 0\ . \tag{3.2b}$$

With these constraints, the relation

$$\left(-i\not{\partial}+M\right)\psi_\nu = 0\ , \tag{3.3}$$

where M is the mass of the spin-3/2 particle, can be interpreted as the equation of motion [26].

[1] Some nice introductions to the Lorentz group representations can be found in [24] and in [25].

Recall that these equations only apply on shell, *i.e.*, the two pairs of spin-1/2 fermions are only removed as long as the Rarita-Schwinger field is on-shell. Otherwise, the spin-1/2 fermions can be present [20].

There are many different models that predict the existence of this new state[2]. Here we review models based on extra dimensions and on the compositeness of quarks and leptons.

3.1 Extra-Dimensions

In the framework of warped extra dimensions, spin-3/2 excitations have been successfully introduced via a 5D extension of the Rarita-Schwinger field [27]. More concretely, it is suggested in [4] a Randall-Sundrum scenario in which SM fields are confined in the infrared brane, but fermions are allowed to propagate in the bulk. There, it is argued that under these conditions a hypothetical spin-3/2 resonance of the right-handed top quark would be the lightest of the Regge excitations, and that it could be even lighter than the Kaluza-Klein modes [4].

The Randall-Sundrum (RS) model is a five-dimensional theory in which the additional spatial dimension is compactified on an interval: $0 \leq y \leq L$, where y denotes the coordinate of the extra dimension [28]. This interval is called *bulk* and its boundaries are the *branes*, corresponding $y = 0$ to the ultraviolet brane and $y = L$ to the infrared one. In the RS-model, the geometry of the bulk is an Anti de-Sitter (AdS_5) space, with metric

$$ds^2 = e^{-2ky}\eta_{\mu\nu}dx^\mu dx^\nu + dy^2 \equiv g_{MN}dx^M dx^N \ , \tag{3.4}$$

where $\eta_{\mu\nu}$ is the 4-dimensional Minkowski metric[3], $\eta_{\mu\nu} = diag(-1, 1, 1, 1)$, capital letters refer to five dimensional indices $M = (\mu, 5)$ and k is the curvature of the space. The "warp" factor, e^{-2ky}, was proposed as a solution to the hierarchy problem: the exponential function enables large differences in scales without requiring an extremely large L [29].

The description of a five-dimensional vector-spinor Ψ_M was carried out in [27]; we will follow the notation from [4]. The starting point is the 5D-Rarita Schwinger action,

$$S = \int d^4x \int dy \sqrt{-G}\bar{\Psi}_M \Gamma^{MNP} \left(D_N + \frac{1}{3} M \Gamma_N \right) \Psi_P \ , \tag{3.5}$$

where the covariant derivative is

$$D_M \Psi_N = \partial_M \Psi_N - \Gamma^P_{MN} \Psi_P + \frac{1}{2} \omega^{AB}_M \gamma_{AB} \ . \tag{3.6}$$

[2]Probably, the most well known is Supergravity, where the spin-3/2 gravitino arises as the superpartner of the graviton [3]. The author not being very familiar with supersymmetries, we will focus on other kinds of models.

[3]Only in this section, in order to reproduce the results from [4], the mostly plus convention will be adopted

Here Γ^P_{MN} are the five-dimensional Christoffel symbols and ω^{AB}_M the associated spin-connection. $\Gamma^M = e^a_M \gamma_a$, being the γ_a's the 5D flat space gamma matrices and e^a_M the fünfbeins with $e^a_\mu = e^{-ky}\ \delta^a_\mu$ for $a, \mu = 0, \ldots, 3$, $e^5_5 = 1$ and all other components equal to zero. The symbols γ_{AB} and Γ^{MNP} are defined as $\gamma_{AB} = \frac{1}{4}[\gamma_A, \gamma_B]$ and $\Gamma^{MNP} = \Gamma^{[M}\Gamma^N\Gamma^{P]}$ [4, 27].

It is stated in [27] that, making use of the gauge freedom, a gauge can be fixed setting $\Psi_5 = 0$. Having done this, we are left with $\Psi_M = (\Psi_\mu, 0)$, so we only have to concern ourselves with the four dimensional wavefunction Ψ_μ. This can be splitted into its left- and right-handed components $\Psi^{L,R}_\mu = \pm\gamma_5 \Psi^{L,R}_\mu$, being $\Psi_\mu = \Psi^L_\mu + \Psi^R_\mu$.

At this point, it is convenient to perform the Kaluza-Klein decomposition of the $\Psi^{L,R}$. It consists on separating variables (the 4D coordinates x from the fifth dimensional y) using a series expansion of the form

$$\Psi^{L,R}_\mu(x, y) = \frac{1}{\sqrt{L}} \sum_n \psi^{L,R}_{\mu n}(x) e^{2ky} f^{L,R}_n(y) \tag{3.7}$$

and plugging this expression into the equations of motion. The functions $f^{L,R}_n(y)$, which only depend on the y-coordinate, are the Kaluza-Klein wavefunctions and are normalized as [27]

$$\frac{1}{L} \int_0^L dy\ e^{ky} f^{L,R*}_m f^{L,R}_n = \delta_{mn}\ . \tag{3.8}$$

After performing some simplifications on the action [27] one can derive the equations of motion, according to the principle of stationary action, *i.e.* requiring $\delta S = 0$. It is found that the KK-wavefunctions satisfy the coupled differential equations [4]

$$(\partial_5 + M) f^R_n = m_n e^{ky} f^L_n\ , \tag{3.9a}$$

$$(-\partial_5 + M) f^L_n = m_n e^{ky} f^R_n\ . \tag{3.9b}$$

In addition, when working with extra dimensions one needs to impose boundary conditions for the KK-wavefunctions in the branes, namely Neumann (which will be denoted by a +) or Dirichlet (−) conditions in the UV and in the IR brane.

As explained in [4], imposing the same b.c. in both (UV, IR) branes, *i.e.* $(+,+)$ or $(-,-)$, would be inconsistent with the phenomenology since it would lead the spin-3/2 particles to have masses of the order of those of their SM spin-1/2 counterparts. Therefore, only the b.c. $(+,-)$ and $(-,+)$ are studied for different spin-3/2 excitations of the SM quarks. It is found that the singlet top excitation is significantly lighter and could be of the order of a few TeV.

3.2 Compositeness

It has been shown in the past that as one goes to higher and higher energies new degrees of freedom appear. In this sense, at the beginning of the last century the atomic nucleus was thought to be indivisible. Also nucleons were regarded as fundamental until the experiments at Stanford Linear Accelerator (SLAC) in the 1960's proved this wrong [30]. Therefore, it seems plausible that quarks might not be the last link of the chain.

Compositeness models suggest that spin-1/2 quarks and leptons are made up of more fundamental constituents, often called "preons". There should then be a new interaction that binds preons together, and which would be characterized by a scale Λ much larger than the scale of the strong interaction Λ_{QCD} [5].

In some of these models a spin-3/2 excited state shows up naturally [5, 6]. For instance, if quarks are formed by three spin-1/2 preons, a change in a preon spin wavefunction would lead to a spin-3/2 lowest excitation, as proposed in [6]. It is hoped that the mass of this excited state is much smaller than the compositeness-scale, so that searching for a signal of this excitation could provide evidence of the underlying preon structure.

3.3 Effective Field Theory Approach

It is a fact that there are many scales in nature: from elementary particles to galaxies one finds nuclei, atoms, molecules,... each of them with their own interaction laws. Luckily, one does not need to understand the physics at all scales in order to describe phenomena that happen at a given scale. This is the idea behind effective field theories (EFT's)[4].

EFT's are a very useful theoretical tool in describing the relevant physics involved at low-energies, being *low* defined with respect to some scale Λ. There are basically two kind of EFT approaches: "top-down" and "bottom-up". In the former, the high-energy theory is known, but when it comes to describing the physics at much lower energies, $E << \Lambda$, the heavy degrees of freedom can be appropriately eliminated, simplifying the calculations. On the other hand, if the underlying high-energy theory is not known, one can follow a bottom-up approach by imposing symmetry and naturalness constraints in the Lagrangian [32].

As we have seen, the spin-3/2 interactions often involve large scales (*e.g.* the compositeness scale) and there are various candidates for a complete high-energy theory. Hence, it is natural to adopt a bottom-up EFT approach in their description.

[4] For general, pedagogical introductions to EFT's the reader is referred to [31, 32, 33, 34].

The most general form of the free Lagrangian for a spin-3/2 field is given by [35, 36]

$$\mathcal{L} = \bar{\psi}_\mu \Lambda^{\mu\nu} \psi_\nu \ , \tag{3.10}$$

with

$$\begin{aligned}\Lambda^{\mu\nu} = &\ (i\not{\partial} - M)\eta^{\mu\nu} + iA(\gamma^\mu \partial^\nu + \gamma^\nu \partial^\mu) \\ &+ \frac{i}{2}(3A^2 + 2A + 1)\gamma^\mu \not{\partial} \gamma^\nu + (3A^2 + 3A + 1)M\gamma^\mu\gamma^\nu \ .\end{aligned} \tag{3.11}$$

This is obtained imposing invariance under the so-called *contact* (or *point*) transformation:

$$\psi^\mu \longrightarrow \psi'^\mu = (\eta^{\mu\nu} + a\gamma^\mu\gamma^\nu)\,\psi_\nu \equiv \mathcal{O}^{\mu\nu}(a)\psi_\nu \ , \tag{3.12a}$$

$$A \longrightarrow A' = \frac{A - 2a}{1 + 4a} \ . \tag{3.12b}$$

A and a are arbitrary real parameters. From Eq.(3.12b) we see that a can not take the value $a = -1/4$, since in that case A' becomes infinite. Similarly, when the propagator is derived from Eq.(3.10), it is found that it has a singularity at $A = -1/2$ (see *e.g.* [37]).

Probably the most popular choice is assuming $A = -1$, which leads to the well-known Rarita-Schwinger propagator [4, 20, 21, 23]

$$G^{\mu\nu}_{(RS)} = \frac{\not{p} - M}{p^2 + M^2}\left[\eta^{\mu\nu} - \frac{1}{3}\gamma^\mu\gamma^\nu - \frac{2p^\mu p^\nu}{3M^2} - \frac{\gamma^\mu p^\nu - \gamma^\nu p^\mu}{3M}\right] \ . \tag{3.13}$$

However, in this work we will also follow an approach suggested by Haberzettl [35], in which the "forbidden" values of the parameters, $a = -1/4$ and $A = -1/2$, are chosen.

Recall from the group representation the presence of two Dirac spinors in Eq.(3.1): one represented by $\left(\frac{1}{2}, 0\right) \oplus \left(0, \frac{1}{2}\right)$ and the other contained in $\left(\frac{1}{2}, 1\right) \oplus \left(1, \frac{1}{2}\right)$. Let us identify the former with [35]

$$\left(\frac{1}{2}, 0\right) \oplus \left(0, \frac{1}{2}\right) \longrightarrow \gamma^\nu\psi_\nu \ , \tag{3.14}$$

which can be seen that transforms as a spinor under the Lorentz group. Imposing that this component of the field is zero, $\gamma^\nu\psi_\nu = 0$, the complementary component is obtained [35]:

$$\left(\frac{1}{2}, 1\right) \oplus \left(1, \frac{1}{2}\right) \longrightarrow \left(\eta^{\mu\nu} - \frac{1}{4}\gamma^\mu\gamma^\nu\right)\psi_\nu \equiv D^{\mu\nu}\psi_\nu \ , \tag{3.15}$$

being the contraction of $D^{\mu\nu}$ with γ_ν equal to zero. In other words, the operator $D^{\mu\nu}$ projects our vector spinor ψ_ν onto the $\left(\frac{1}{2}, 1\right) \oplus \left(1, \frac{1}{2}\right)$ subspace, also called Rarita-Schwinger (RS) subspace [35]. In this way one of the Dirac spinors is eliminated.

Notice that this projection is achieved by the point transformation Eq.(3.12a) for $a = -1/4$. Since for this value of a the transformation Eq.(3.12b) is singular, we are forced to simultaneously choose $A = -1/2$, leading to an undefined $\frac{0}{0}$ situation. This limit is resolved in [35], resulting in the "new" point transformation

$$\psi^\mu \longrightarrow \psi'^\mu = \left(\eta^{\mu\nu} - \frac{1}{4}\gamma^\mu\gamma^\nu\right)\psi_\nu = D^{\mu\nu}\psi_\nu \; , \tag{3.16a}$$

$$A \longrightarrow A' = -\frac{1}{2} \; . \tag{3.16b}$$

The corresponding tensor is then

$$\begin{aligned}\Lambda^{\mu\nu} = \; & (i\not\partial - M)\eta^{\mu\nu} - \frac{i}{2}(\gamma^\mu\partial^\nu + \gamma^\nu\partial^\mu) \\ & + \frac{3i}{8}\gamma^\mu\not\partial\gamma^\nu + \frac{1}{4}M\gamma^\mu\gamma^\nu \; .\end{aligned} \tag{3.17}$$

Solving the equation

$$D_{\mu\alpha}\Lambda^{\alpha\beta}D_{\beta\rho}G^{\rho\nu} = D_\mu^\nu \; , \tag{3.18}$$

the propagator for the singular case can be obtained [35]:

$$G^{\mu\nu}_{(SC)} = D^{\mu\rho}\left[\frac{(\not p + M)\eta_{\rho\sigma}}{p^2 - M^2} + \frac{2p_\rho(\not p + 2M)p_\sigma}{(p^2 - M^2)(p^2 - 4M^2)}\right]D^{\sigma\nu} \; . \tag{3.19}$$

As one should expect, this propagator only acts in the RS subspace, that is [35]

$$\gamma_\mu G^{\mu\nu} = G^{\mu\nu}\gamma_\nu = 0 \; . \tag{3.20}$$

It is suggested in [35] that the pole of Eq.(3.19) at $p^2 = 4M^2$ represents the spin-1/2 piece contained in the $\left(\frac{1}{2}, 1\right) \oplus \left(1, \frac{1}{2}\right)$ RS field. However, this does not satisfy the usual requirements to be considered a particle.

Having reviewed the propagation of the Rarita-Schwinger field, it is now time to turn to its interaction. The interaction of the spin-3/2 and spin-1/2 states are described by dimension-five operators of the form [21]

$$\mathcal{L}_{int} = i\,\frac{g}{\Lambda}\,\bar\psi_\sigma(\eta^{\sigma\mu} + z\gamma^\sigma\gamma^\mu)\gamma^\nu T_a\xi G^a_{\mu\nu} + H.C. \tag{3.21}$$

Here Λ sets the strength of the interaction. ξ is the spin-1/2 fermion spinor and $\bar\psi$ the spin-3/2 vector spinor, $G^a_{\mu\nu}$ is the field strength tensor of the gauge field (not to be confused with the propagator), being g the coupling constant and T_a the generators of the group. The parameter z measures the "off-shellness" of the interaction: see that if the field is on-shell, $\gamma^\mu\psi_\mu = 0$ and the term on z vanishes, which is equivalent to choosing $z = 0$.

Chapter 4

Higgs Production and Decay via Spin-3/2 Excitations

In this section we study how a spin-3/2 partner of the top quark, presented in [4], affects the process of Higgs production and decay at the LHC. More concretely, we consider the Higgs production from gluon-gluon fusion, which takes place via fermion loops analogous to those from Section 2.2.2; and the decay to photons, whose signal is easier to detect than that of the gluons, since the latter can only be detected as gluon-initiated jets of hadrons. By the arguments given in Section 2.2.2, we restrict ourselves once again to the contribution from the top-quark, the most massive of the SM fermions. We will compute the matrix element of the process of gluon-gluon fusion; that of the decay to photons can be easily derived from this.

Let us consider the possibility that some of the top-quarks from Fig.2.2 are replaced by their spin-3/2 partners, henceforth t^*. From Eq.(3.21), each vertex involving an excited top includes a factor $1/\Lambda$, where Λ is some large scale. Hence, staying at order $1/\Lambda^2$ only one t^* will be present (one propagator and two vertices). Furthermore, assuming that this particle is a pure spin-3/2 state, its coupling to a scalar and a spin-1/2 spinor is not allowed by the group representation, so it must be the particle in the loop that couples to both gluons.

We changed the conventions with respect to Section 2.2.2 for convenience: we want the denominators to have the suitable form from Appendix C to directly use those formulae when reducing the loop integrals. Now, although the momenta of the external particles are labelled in the same way, they are all *incoming*.

In our new conventions, the matrix element reads

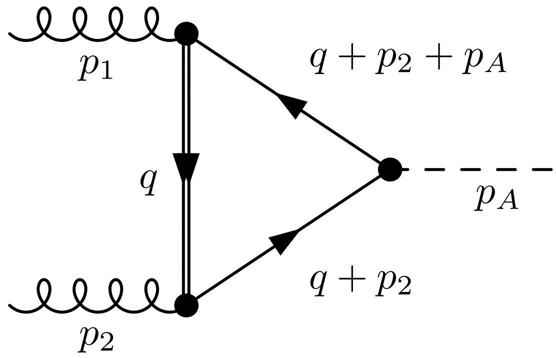


Figure 4.1: *Higgs production at the LHC from gluon fusion coming from proton-proton collision. The double line in the triangle represents the t^*-quark and the other two lines are t-quarks. All external momenta are incoming.*

$$
\begin{aligned}
\mathcal{M} = &-2\int \frac{d^4q}{(2\pi)^4}\left(-\frac{iem_t}{2\sin\theta_W M_W}\right)\left(i\,\frac{\not{q}+\not{p}_2+m_t}{(q+p_2)^2-m_t^2}\right)(i\,\Gamma_2^{\nu\rho}(p_2))\\
&\times (i\,G_{\rho\sigma}(q))\,(i\,\Gamma_1^{\sigma\mu}(p_1))\left(i\frac{\not{q}+\not{p}_2+\not{p}_A+m_t}{(q+p_2+p_A)^2-m_t^2}\right)\\
&\times \varepsilon_\mu^*(p_1,\lambda_1)\varepsilon_\nu^*(p_2,\lambda_2)\ .
\end{aligned}
\tag{4.1}
$$

The 2 in front of the integral accounts for the two contributing diagrams (the one shown in Fig. 4.1 and the same with the gluons crossed: $p_1 \leftrightarrow p_2$) which, as seen is Section 2.2 yield the same result. The tt^*g vertices $\Gamma_2^{\mu\nu}(p)$ and $\Gamma_1^{\mu\nu}(p)$ and the propagator $G_{\mu\nu}(q)$ can be found in Appendix B.

4.1 Using the Rarita-Schwinger Propagator

In the study of the t^* the Rarita-Schwinger case has been the most common choice [4, 21, 22]. Accordingly, we will start this section computing the diagram in Fig. 4.1, Eq.(4.1), using the Rarita Schwinger propagator. When doing this, we will find that the matrix element yields infinity and we will try to understand the nature of this divergence.

4.1.1 Higgs Production via Fermion Triangle

In the computation of the matrix element Eq.(4.1) some sort of regularization is required. As in Section 2.2, here we use dimensional regularization and we will work in $n = 4 - 2\epsilon$ dimensions. To that end, the n-dimensional form of the t^*-propagator and the vertices are needed. These are presented in Appendix B. Substituting them in the matrix element we get

$$\mathcal{M} = -\frac{m_t g_s^2 \mu^{n-4}}{2\sin\theta_W M_W} \frac{1}{\Lambda^2}\delta_{c_1 c_2}\, \varepsilon^*_\mu(p_1,\lambda_1)\, \varepsilon^*_\nu(p_2,\lambda_2) \int \frac{d^n q}{(2\pi)^n} \frac{\mathrm{Tr}\,\mathcal{A}^{\nu\mu}}{\mathcal{D}_1\mathcal{D}_2\mathcal{D}_3}\ , \quad (4.2)$$

where c_1 and c_2 are the colours of the gluons. The denominators are $\mathcal{D}_1 = q^2 - M_{t^*}^2$, $\mathcal{D}_2 = (q+p_2)^2 - m_t^2$ and $\mathcal{D}_3 = (q+p_2+p_A)^2 - m_t^2$. The quantity

$$\begin{aligned}\mathcal{A}^{\nu\mu} &= (\not{q} + \not{p}_2 + m_t)\left(-\gamma^\nu p_2^\rho + \not{p}_2\eta^{\nu\rho} + z\not{p}_2\gamma^\nu\gamma^\rho - z\gamma^\nu\not{p}_2\gamma^\rho\right)\\ &\quad\times(\not{q} + M_{t^*})\left[\eta_{\rho\sigma} - \frac{(n-2)q_\rho q_\sigma}{(n-1)M_{t^*}^2} - \frac{\gamma_\rho\gamma_\sigma}{n-1} + \frac{(q_\rho\gamma_\sigma - q_\sigma\gamma_\rho)}{(n-1)M_{t^*}}\right]\\ &\quad\times\left(p_1^\sigma\gamma^\mu - \eta^{\sigma\mu}p_1 + z\gamma^\sigma\not{p}_1\gamma^\mu - z\gamma^\sigma\gamma^\mu\not{p}_1\right)(\not{q} + \not{p}_2 + \not{p}_A + m_t) \quad (4.3)\end{aligned}$$

will be traced with `FORM` [38]. Again, the mostly minus convention for $\eta_{\mu\nu}$ is adopted.

When tracing these quantities, some helpful simplifications can be made. First of all, outgoing photons are transverse and massless, so that all terms containing p_1^μ, p_2^ν, p_1^2 or p_2^2 can be discarded. Furthermore, do not forget that the matrix element Eq.(4.2) will have to be squared and averaged. Summing over the polarizations of the gluon p_1 will yield

$$\sum_{\lambda_1=1,2} \varepsilon_\mu(p_1,\lambda_1)\varepsilon^*_\nu(p_1,\lambda_1) = -\eta_{\mu\nu} + \frac{p_{1\mu}n_{1\nu} + p_{1\nu}n_{1\mu}}{p_1\cdot n_1}\ , \quad (4.4)$$

where $n_{1\mu}$ is some vector that together with $\varepsilon_\mu(p_1,1)$, $\varepsilon_\mu(p_1,2)$ and $p_{1\mu}$ spans the space. We are allowed by the gauge freedom to choose $n_1 = p_2$ and then, from Eq.(4.4), the sum:

$$\sum_{\lambda_1=1,2} \varepsilon_\mu(p_1,\lambda_1)p_2^\mu\varepsilon^*_\nu(p_1,\lambda_1) = 0\ . \quad (4.5)$$

That is, already at the level of the matrix element, terms proportional to p_2^μ can be set to zero, simplifying a lot the intermediate calculations. By the same token, the choice of $n_2 = p_1$ when summing over λ_2, allows us to set terms in p_1^ν equal to zero. By Lorentz covariance, this means that in our gauge choice the traces of $M_1^{\nu\mu}$ and $M_2^{\nu\mu}$ can only be proportional to the metric tensor $\eta^{\mu\nu}$.

The full result for the matrix element will not be exhibited here due to its length. Instead, it is shown in Appendix D. Somewhat surprisingly, we encounter terms proportional to $1/\epsilon$, which, when taking the limit $\epsilon \to 0$ will diverge. This divergent part of the matrix element is

$$\begin{aligned}
\mathcal{M}^{div} = & -\frac{i m_t m_H^2 g_s^2 \mu^{n-4}}{48\pi^2 \sin\theta_W M_W}\, \delta_{c_1 c_2}\, \frac{1}{\Lambda^2}\, \varepsilon^*_\mu(p_1,\lambda_1)\, \varepsilon^{*\,\mu}(p_2,\lambda_2) \\
& \times \frac{1}{\bar\epsilon}\left[-\frac{2m_H^2}{M_{t^*}}(1+8z+16z^2)+8M_{t^*}\right. \\
& \left. -\frac{2m_t m_H^2}{M_{t^*}^2}(1+4z)+\frac{96 m_t^2 z}{M_{t^*}}(1+2z)+\frac{16 m_t^3}{M_{t^*}^2}(1+4z)\right]
\end{aligned} \tag{4.6}$$

with $\bar\epsilon$ defined as

$$\frac{1}{\bar\epsilon} \equiv \frac{1}{\epsilon} - \gamma_E + \log(4\pi)\ . \tag{4.7}$$

From here, the matrix element for the decay to photons can be obtained by simply exchanging $g_s \to e_q e$ and replacing the colour factor $1/2\ \delta_{c_1c_2} \to N_c$.

In order to get a finite, consistent result to be compared with the experiments, we need to renormalize this matrix element.

4.1.2 Renormalization

Renormalization is an important issue when dealing with ETF's [31, 33, 34, 39]. In "top-down" approaches renormalization is achieved by imposing *matching conditions* that relate the EFT with the full theory. In our case, as we ignore the underlying high-energy theory, this is not possible [31].

A first guess would be a cut-off regularization: since the theory is expected to be valid only up to some scale Λ, it seems reasonable to cut the integrals at that scale. However, besides violating Poincaré invariance, the cut-off regularization in EFT's presents the disadvantage of breaking the expansion in $1/\Lambda$. Suppose that the integral in Eq.(4.2) diverges as Λ^2, then after multiplying by the $1/\Lambda^2$ coming from the two vertices, the matrix element will not be suppressed by any power of Λ: it would be of $O(1)$. Then, we should also consider, for consistency, operators of dimension 6 (that is, of $O(1/\Lambda^2)$) with integrals that diverge as Λ^4, and in general higher dimension operators with integrals that diverge as Λ^n [31].

To address this problem, it is suggested to adopt a mass-independent renormalization scheme, as dimensional regularization with minimal subtraction. In this renormalization scheme, the bare couplings and masses of the Lagrangian are redefined into renormalized ones, giving rise to counterterms that cancel the divergences[1]. In an attempt to apply this to our process, we need to identify first, which are the parameters to be renormalized. To that end, we will look at our problem from another perspective.

[1]An example of this kind of calculation at one loop in the context of EFT's is shown in [34].

In the EFT-based approach we have adopted, higher order operators that parametrize BSM effects are added to the SM Lagrangian. In a basis of independent operators $\mathcal{O}_i$ [40],

$$\mathcal{L} = \mathcal{L}_{SM} + \sum_i \bar{c}_i \mathcal{O}_i \ , \tag{4.8}$$

where $\bar{c}_i$ are called Wilson coefficients. Our process under study is contained in the operator

$$\mathcal{O}_g = \frac{g_s^2 \bar{c}_g}{M_W^2} \Phi^\dagger \Phi G^a_{\mu\nu} G^{\mu\nu}_a \ . \tag{4.9}$$

$G^{\mu\nu}_a$ is the gluon field tensor and $\bar{c}_g$ the Wilson coefficient. After EWSB, *i.e.*, when the Higgs acquires its vev, and in the unitary gauge,

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h \end{pmatrix} \tag{4.10}$$

and the $\mathcal{O}_g$ operator takes the form

$$\mathcal{O}_g = \frac{g_s^2 \bar{c}_g}{2M_W^2} \left(v^2 G^a_{\mu\nu} G^{\mu\nu}_a + 2vh G^a_{\mu\nu} G^{\mu\nu}_a + hh G^a_{\mu\nu} G^{\mu\nu}_a \right) \ . \tag{4.11}$$

Thus, after spontaneous symmetry breaking $\mathcal{O}_g$ gives rise to three kinds of diagrams:

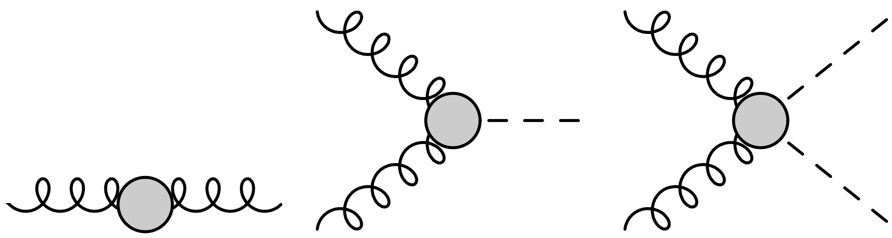


Diagram a) *Diagram b)* *Diagram c)*

Table 4.1: *Diagrams that arise from the operator $\mathcal{O}_g$ after spontaneous symmetry breaking. Diagram a) corresponds to the gluon self-energy, diagram b) to Higgs production and diagram c) to Higgs pair production.*

The parameter $\bar{c}_g$ should then absorb the divergences coming from these three diagrams. The process we studied in Section 4.1.1 corresponds to a diagram of the *b-type*. If we want to renormalize this process, we should renormalize in a general way the contributions from the spin-3/2 states to the $\mathcal{O}_g$ operator. It is therefore necessary to study the divergences coming from *a-* and *c-type* diagrams as well, so that when our coupling is redefined: $\bar{c}_g^R = \bar{c}_g + \delta\bar{c}_g$, we ensure that the counter-terms originated by $\delta\bar{c}_g$ cancel the divergences from the three processes: gluon-self energy (*a-type* diagram), Higgs production from gluon fusion (*b-type*) and Higgs pair production (*c-type*). In this section we compute the divergent part of the matrix element of such processes and discuss the possibility of renormalizing them altogether.

Gluon Self-Energy

In this case there is only one diagram:

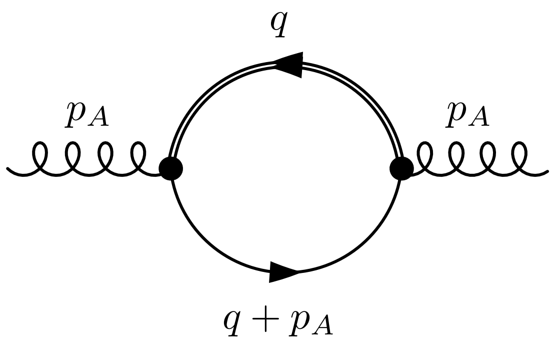


Figure 4.2: *Diagram of the process of gluon self-energy via one t^*. The t^* is represented by the double line and the other fermion line is a t-quark.*

The matrix element is

$$\begin{aligned}\mathcal{M} = &- \int \frac{d^n q}{(2\pi)^n} \left(i\, \Gamma_2^{\mu\alpha}(p_A)\right) \left(i\, G^{RS}_{\alpha\beta}(q)\right) \left(i\, \Gamma_1^{\beta\nu}(-p_A)\right) \\ &\times \left(i\, \frac{\not{q} + \not{p}_A + m_t}{(q+p_A)^2 - m_t^2}\right) \varepsilon^*_\mu(p_A, \lambda_A)\, \varepsilon_\nu(p_A, \lambda_A) \\ = &- \frac{g_s^2}{2\Lambda^2} \delta_{c_1 c_1}\, \varepsilon^*_\mu(p_A, \lambda_A)\, \varepsilon_\nu(p_A, \lambda_A) \int \frac{d^n q}{(2\pi)^n} \frac{\operatorname{Tr} A^{\mu\nu}}{\mathcal{D}_1 \mathcal{D}_2}\ . \end{aligned} \tag{4.12}$$

The denominators are $\mathcal{D}_1 = q^2 - M_{t^*}^2$ and $\mathcal{D}_2 = (q+p_A)^2 - m_t^2$, and the $A^{\mu\nu}$ tensor

$$\begin{aligned} A^{\mu\nu} = &\left(-\gamma^\mu p_A^\alpha + \not{p}_A \eta^{\mu\alpha} + z \not{p}_A \gamma^\mu \gamma^\alpha - z \gamma^\mu \not{p}_A \gamma^\alpha\right) \left(\not{q} + M_{t^*}\right) \\ &\times \left(\eta_{\alpha\beta} - \frac{1}{3}\gamma_\alpha \gamma_\beta - \frac{2}{3M_{t^*}^2} q_\alpha q_\beta + \frac{1}{3M_{t^*}} \left(q_\alpha \gamma_\beta - q_\beta \gamma_\alpha\right)\right) \\ &\times \left(-p_A^\beta \gamma^\nu + \eta^{\beta\nu} \not{p}_A - z \gamma^\beta \not{p}_A \gamma^\nu + z \gamma^\beta \gamma^\nu \not{p}_A\right) \left(\not{q} + \not{p}_A + m_t\right)\ . \end{aligned} \tag{4.13}$$

As we are only interested in the divergent part, we can use the the 4-dimensional form of the propagator. In this case the gluon is off-shell, that is, we can not make $p_A^2 = 0$, and it can have longitudinal polarization, so p_A^μ, p_A^ν can not be set to zero either. We get for the divergent part

$$\begin{aligned} \mathcal{M}^{div} = &- \frac{-i g_s^2 M_{t^*} m_t}{6\pi^2 \Lambda^2} \varepsilon^*_\mu(p_A, \lambda_A)\, \varepsilon_\nu(p_A, \lambda_A) \\ &\times \frac{1}{\epsilon} \left(p_A^2 \eta^{\mu\nu} - p_A^\mu p_A^\nu\right) \left(1 - 4z - 8z^2\right)\ . \end{aligned} \tag{4.14}$$

Higgs Pair Production

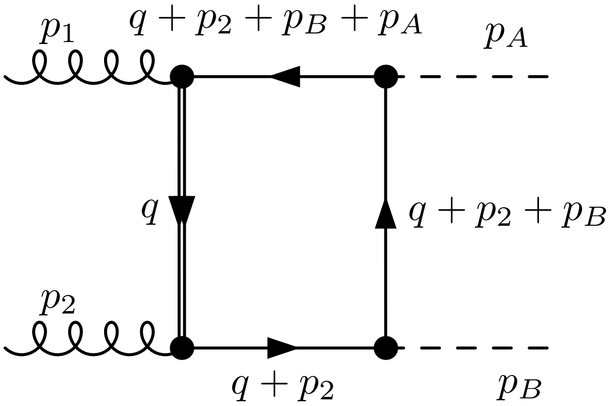


Figure 4.3: *One of the diagrams of the process of di-Higgs production from gluon fusion via a loop with one excited top quark. As usual, the double line represents the spin-3/2 t^* and the other fermion lines are the SM t-quarks. The other three diagrams are obtained crossing the gluons: $p_1 \leftrightarrow p_2$; the final Higgs bosons: $p_A \leftrightarrow p_B$ and both: $p_1 \leftrightarrow p_2$ and $p_A \leftrightarrow p_B$.*

There are four diagrams contributing to Higgs pair production involving a t^*: the one shown in Fig. (4.3); the same interchanging gluons, $p_1 \leftrightarrow p_2$; interchanging Higgs bosons, $p_A \leftrightarrow p_B$; and interchanging both, $p_1 \leftrightarrow p_2$ and $p_A \leftrightarrow p_B$. Since the four diagrams yield the same result, we can just compute the matrix element of one of them and multiply it by four. Then the total matrix element is simply

$$\begin{aligned}
\mathcal{M} =& -4\,\mu^{4-n} \int \frac{d^n q}{(2\pi)^n} \left(-\frac{iem_t}{2\sin\theta_W M_W} \right) \left(i\, \frac{\not{q} + \not{p}_2 + m_t}{(q+p_2)^2 - m_t^2} \right) \\
& \times \left(i\, \Gamma_2^{\nu\rho}(p_2)\right) \left(i\, G_{\rho\sigma}^{RS}(q)\right) \left(i\, \Gamma_1^{\sigma\mu}(p_1)\right) \left(i\, \frac{\not{q} + \not{p}_2 + \not{p}_B + \not{p}_A + m_t}{(q+p_2+p_B+p_A)^2 - m_t^2} \right) \\
& \times \left(-\frac{iem_t}{2\sin\theta_W M_W} \right) \left(i\, \frac{\not{q} + \not{p}_2 + \not{p}_B + m_t}{(q+p_2+p_B)^2 - m_t^2} \right) \varepsilon_\mu^*(p_1,\lambda_1)\, \varepsilon_\nu^*(p_2,\lambda_2) \\
=& -\frac{e^2 m_t^2 g_s^2}{2\sin^2\theta M_W^2} \frac{1}{\Lambda^2} \delta_{c_1 c_2} \varepsilon_\mu^*(p_1,\lambda_1)\, \varepsilon_\nu^*(p_2,\lambda_2) \mu^{4-n} \int \frac{d^n q}{(2\pi)^n} \frac{\mathrm{Tr}\, A^{\nu\mu}}{\mathcal{D}_1\mathcal{D}_2\mathcal{D}_3\mathcal{D}_4}\,.
\end{aligned} \tag{4.15}$$

The denominators are $\mathcal{D}_1 = q^2 - M_{t^*}^2$, $\mathcal{D}_2 = (q+p_2)^2 - m_t^2$, $\mathcal{D}_3 = (q + p_2 +$

$p_B)^2 - m_t^2$ and $\mathcal{D}_4 = (q + p_2 + p_B + p_A)^2 - m_t^2$. The quantity to be traced is

$$
\begin{aligned}
A^{\nu\mu} = & \left(\not{q} + \not{p}_2 + m_t\right)\left(-\gamma^\nu p_2^\rho + \not{p}_2 \eta^{\nu\rho} + z\not{p}_2\gamma^\nu\gamma^\rho - z\gamma^\nu\not{p}_2\gamma^\rho\right) \\
& \times \left(\not{q} + M_{t^*}\right)\left[\eta_{\rho\sigma} - \frac{1}{3}\gamma_\rho\gamma_\sigma - \frac{2}{3M_{t^*}^2} q_\rho q_\sigma + \frac{1}{3M_{t^*}}\left(q_\rho\gamma_\sigma - q_\sigma\gamma_\rho\right)\right] \\
& \times \left(p_1^\sigma\gamma^\mu - \eta^{\sigma\mu}\not{p}_1 + z\gamma^\sigma\not{p}_1\gamma^\mu - z\gamma^\sigma\gamma^\mu\not{p}_1\right) \\
& \times \left(\not{q} + \not{p}_2 + \not{p}_B + \not{p}_A + m_t\right)\left(\not{q} + \not{p}_2 + \not{p}_B + m_t\right) \ . \qquad (4.16)
\end{aligned}
$$

Once again, since we are only computing the divergent part, there is no need to use the n-dimensional propagator. The divergent part reads

$$
\begin{aligned}
\mathcal{M}^{div} = & -\frac{ie^2 m_t^2 g_s^2}{32\pi^2 \sin^2\theta M_W^2}\frac{1}{\Lambda^2}\delta_{c_1c_2}\varepsilon_\mu^*(p_1,\lambda_1)\ \varepsilon_\nu^*(p_2,\lambda_2) \\
& \times \frac{1}{\epsilon}\ \left(p_B^\mu p_B^\nu \mathcal{B}_1 + \eta^{\mu\nu}\mathcal{B}_2\right) \ , \qquad (4.17)
\end{aligned}
$$

with

$$
\mathcal{B}_1 = \frac{-4m_H^2}{3M_{t^*}^2}\left(\frac{1}{3} + 4z + 8z^2\right) + \frac{2(u+t)}{3M_{t^*}^2}\left(\frac{1}{3} + 4z + 8z^2\right) \ , \qquad (4.18)
$$

$$
\begin{aligned}
\mathcal{B}_2 = & (u+t)\left[\frac{-m_H^2}{3M_{t^*}^2}(1+4z) + \frac{16m_t}{M_{t^*}}(z+2z^2) + \frac{4m_t^2}{M_{t^*}^2}(1+4z)\right] \\
& - \frac{tu}{3M_{t^*}^2}\left(\frac{1}{3} + 4z + 8z^2\right) + \frac{m_H^4}{M_{t^*}^2}\left(\frac{7}{9} + 4z\frac{8}{3}z^2\right) \\
& - \frac{32 m_t m_H^2}{M_{t^*}}(z+2z^2) - \frac{8m_t^2 m_H^2}{M_{t^*}^2}(1+4z) \ . \qquad (4.19)
\end{aligned}
$$

Here u and t stand for the Mandelstam variables. This result is, as one should expect, symmetric under the change $u \leftrightarrow t$, which corresponds to changing $p_1 \leftrightarrow p_2$ or $p_A \leftrightarrow p_B$.

Having collected all the needed pieces, we could now search for a renormalization procedure that would lead us to a finite result. Unfortunately, we anticipate that this will not be possible. To start with, notice that the diagram at hand, from Fig. 4.1, is at LO: in a theory that contemplates the existence of the t^*-quark this diagram would be, together with that from Fig. 2.2, the LO contributions to the process $gg \to H$. In the perturbative expansion they are both of $O(g_s^2)$ and there is no ggH vertex, *i.e.* no $O(g_s^0)$, since the gluons are massless. If the LO contribution is already divergent, this means that the theory is *non-renormalizable*. But even non-renormalizable theories can be renormalized in some way! [41, 24] This can be achieved adding new terms to the Lagrangian that will absorb the divergences, *as long as* they are all respect the symmetry of the theory. The reason why we can not renormalize the expression Eq.(4.2) is *not only* that the theory is non-renormalizable, *but also* that all the terms consistent with the symmetry are already in the Lagrangian, since we are dealing with a bottom-up EFT. The Lagrangian of the EFT was built-up as the most general gauge-invariant Lagrangian [40]: by definition, there are no more terms to be added that are consistent with the symmetry. Therefore, this way of proceeding, proposed *e.g.* in [41], is of no use in our case.

The only thing that could that could be done is to redefine the Wilson coefficient: $\bar{c}_g^R = \bar{c}_g + \delta\bar{c}_g$, but from the form of the divergences: Eq.(4.6), Eq.(4.14) and Eq.(4.17) one sees that there is no way that the $\delta\bar{c}_g$ will absorb the divergences coming from the three diagrams.

4.2 Using the Singular Case Propagator

The problems with the infinities that we encountered in the last section clearly indicate that a better procedure could be found that yields a finite matrix element. In this section we will try a different approach and comment on the results.

Recall from the beginning of Chapter 4, that we are assuming the t^* to be a "pure" spin-3/2 state. We suggest that when using the RS-propagator, since it acts on the entire space of the vector spinor ψ_μ, it also describes the propagation of the spin-1/2 pieces we found in Chapter 3 (keep in mind that the conditions $\partial^\mu \psi_\mu = 0$, $\gamma^\mu \psi_\mu = 0$ only apply on-shell). It might happen that these spin-1/2 degrees of freedom are the origin of the divergences. Therefore, a possible solution could be restricting ourselves to the RS subspace. As we saw in Chapter 3, this is achieved in the approach introduced by Haberzettl, going through the "singular" case $a = -1/4$, $A = -1/2$ [35].

In this section we compute the matrix element of the process of Higgs production involving a t^* in the loop, Eq.(4.1), but replacing the RS-propagator by the one proposed by Haberzettl [35]. The n-dimensional form of this propagator was computed and is shown in Appendix B. Plugging it into Eq.(4.1), we get

$$\mathcal{M} = \frac{-m_t g_s^2}{2\sin\theta_W M_W}\, \delta_{c_1 c_2}\, \frac{1}{\Lambda^2}\, \varepsilon^*_\mu(p_1,\lambda_1)\, \varepsilon^*_\nu(p_2,\lambda_2) \times \int \frac{d^n q}{(2\pi)^4} \left(\frac{\mathrm{Tr}\, M_1^{\nu\mu}}{\mathcal{D}_1\mathcal{D}_2\mathcal{D}_3} + \frac{\mathrm{Tr}\, M_2^{\nu\mu}}{\mathcal{D}_1\tilde{\mathcal{D}}_1\mathcal{D}_2\mathcal{D}_3} \right), \tag{4.20}$$

where $\mathcal{D}_1 = q^2 - M_{t^*}^2$, $\tilde{\mathcal{D}}_1 = q^2 - (n/(n-2))^2 M_{t^*}^2$, $\mathcal{D}_2 = (q+p_2)^2 - m_t^2$ and

$\mathcal{D}_3 = (q + p_2 + p_A)^2 - m_t^2$. The quantities:

$$\begin{aligned} M_1^{\nu\mu} = & (\not{q} + \not{p}_2 + m_t) \left(-\gamma^\nu p_2^\rho + \not{p}_2 \eta^{\nu\rho} + z\not{p}_2 \gamma^\nu \gamma^\rho - z\gamma^\nu \not{p}_2 \gamma^\rho \right) \\ & \times \left(\eta_{\rho\alpha} - \frac{1}{n}\gamma_\rho \gamma_\alpha \right) (\not{q} + M_{t^*}) \eta^{\alpha\beta} \left(\eta_{\beta\sigma} - \frac{1}{n}\gamma_\beta \gamma_\sigma \right) \\ & \times \left(p_1^\sigma \gamma^\mu - \eta^{\sigma\mu} p_1 + z\gamma^\sigma \not{p}_1 \gamma^\mu - z\gamma^\sigma \gamma^\mu \not{p}_1 \right) (\not{q} + \not{p}_2 + \not{p}_A + m_t) \,, \end{aligned} \tag{4.21}$$

$$\begin{aligned} M_2^{\nu\mu} = & (\not{q} + \not{p}_2 + m_t) \left(-\gamma^\nu p_2^\rho + \not{p}_2 \eta^{\nu\rho} + z\not{p}_2 \gamma^\nu \gamma^\rho - z\gamma^\nu \not{p}_2 \gamma^\rho \right) \\ & \times \left(\eta_{\rho\alpha} - \frac{1}{n}\gamma_\rho \gamma_\alpha \right) \frac{n}{n-2} q^\alpha \left(\not{q} + \frac{n}{n-2} M_{t^*} \right) q^\beta \left(\eta_{\beta\sigma} - \frac{1}{n}\gamma_\beta \gamma_\sigma \right) \\ & \times \left(p_1^\sigma \gamma^\mu - \eta^{\sigma\mu} p_1 + z\gamma^\sigma \not{p}_1 \gamma^\mu - z\gamma^\sigma \gamma^\mu \not{p}_1 \right) (\not{q} + \not{p}_2 + \not{p}_A + m_t) \end{aligned} \tag{4.22}$$

were traced with `FORM` [38]. Following the instructions from [35] we set the off-shell parameter at $z = -1/4$. The matrix element is

$$\begin{aligned} \mathcal{M} = & \frac{-i m_t m_H^2 g_s^2}{32\pi^2 \sin\theta_W M_W} \frac{1}{\Lambda^2} \, \delta_{c_1 c_2} \, \varepsilon_\mu^*(p_1, \lambda_1) \, \varepsilon^{*\,\mu}(p_2, \lambda_2) \\ & \times \Bigg\{ \frac{1}{\bar{\epsilon}}(-4M_{t^*}) - \frac{8}{3}M_{t^*} - 4m_t + B_0^{fin}(0, M_{t^*}^2, m_t^2) \left[-\frac{8}{9}M_{t^*} + \frac{2}{3}\frac{m_t^2}{M_{t^*}} \right] \\ & + B_0^{fin}(0, 4M_{t^*}^2, m_t^2) \left[-\frac{40}{9}M_{t^*} - \frac{2}{3}\frac{m_t^2}{M_{t^*}} \right] + B_1^{fin}(0, M_{t^*}^2, m_t^2) \left[\frac{8}{9}M_{t^*} - \frac{8}{9}\frac{m_t^2}{M_{t^*}} \right] \\ & + B_1^{fin}(0, 4M_{t^*}^2, m_t^2) \left[-\frac{32}{9}M_{t^*} + \frac{8}{9}\frac{m_t^2}{M_{t^*}} \right] + C_0(0, m_H^2, 0, M_{t^*}^2, m_t^2, m_t^2) \Big[M_{t^*} m_H^2 \\ & - \frac{10}{3} m_t^2 M_{t^*} + \frac{2}{3}\frac{m_t^4}{M_{t^*}} \Big] + C_0(0, m_H^2, 0, 4M_{t^*}^2, m_t^2, m_t^2) \left[-\frac{8}{3} m_t^2 M_{t^*} - \frac{2}{3}\frac{m_t^4}{M_{t^*}} \right] \Bigg\} . \end{aligned} \tag{4.23}$$

B_0^{fin}, B_1^{fin} denote the finite part of the integrals B_0 and B_1 that can be computed numerically, *e.g.* using *LoopTools* [42]. These integrals, together with C_0, are defined in Appendix C. Once again, the di-photon decay can be derived by exchanging $g_s \to e_q e$ and $1/2 \; \delta_{c_1 c_2} \to N_c$.

This matrix element still has a $\frac{1}{\bar{\epsilon}}$-dependent part: the singular case approach, although reasonable from a theoretical point of view, has not solved our problem with the infinities.

4.3 Results

Seeking a renormalizable theory of the propagation of the t^*-quark and its interaction with the gluon field is, apparently, much more complicated than we expected and has turned out to be beyond the scope of this work. At this point, in order to get an idea of how the presence of a t^* influences Higgs production and decay, probably the best one can do is to introduce the until now avoided cut-off. This is, indeed, the most common choice in the literature when evaluating processes involving one or more t^*'s. [4, 21, 22].

In this section we regularize with a cut-off the matrix element for Higgs production and decay when the singular case propagator is used, Eq.(4.20). We set one cut-off at $\Lambda = 7M_{t^*}$, guided by the estimation of unitarity-violation from [4], and other two cut-offs at $\Lambda = 8M_{t^*}$ and $\Lambda = 6M_{t^*}$ to see how the cut-off choice modifies the final result.

We started computing the ratio of the partonic cross section for $gg \to H$ involving the t^*-quark, $\hat{\sigma}^{BSM}_{gg\to H}$, divided by the SM partonic cross section $\hat{\sigma}^{SM}_{gg\to H}$. This was done by simply dividing the total BSM squared averaged matrix element, which is the sum of the t- and the t^*-quarks contributions, by the SM one:

$$\frac{\hat{\sigma}^{BSM}_{gg\to H}}{\hat{\sigma}^{SM}_{gg\to H}} = \frac{\overline{\left|\mathcal{M}^{(t)}_{gg\to H} + \mathcal{M}^{(t^*)}_{gg\to H}\right|^2}}{\overline{\left|\mathcal{M}^{(t)}_{gg\to H}\right|^2}} \,. \tag{4.24}$$

In this expression we are assuming that the QCD corrections at NLO, which can be quite significant [9], alter in the same way the decays mediated by the t^* and by the t-quarks, so that they factorize in the numerator and cancel the SM corrections of the denominator.

The Higgs mass was taken at $m_H = 125$ GeV and the t-quark mass at $m_t = 173$ GeV. The ratio σ^{BSM}/σ^{SM} using the three different cut-offs is represented against the t^* mass in Fig. 4.4. The vertical line drawn at $M_{t^*} = 803$ GeV represents the lower bound for the t^* mass recently found by CMS [43].

Following these same steps, the ratio

$$\frac{\Gamma^{BSM}_{H\to\gamma\gamma}}{\Gamma^{SM}_{H\to\gamma\gamma}} = \frac{\overline{\left|\mathcal{M}^{(W)}_{H\to\gamma\gamma} + \mathcal{M}^{(t)}_{H\to\gamma\gamma} + \mathcal{M}^{(t^*)}_{H\to\gamma\gamma}\right|^2}}{\overline{\left|\mathcal{M}^{(W)}_{H\to\gamma\gamma} + \mathcal{M}^{(t)}_{H\to\gamma\gamma}\right|^2}} \tag{4.25}$$

was computed and is represented against the t^* mass in Fig. 4.5. Here we assumed once again that the QCD corrections at NLO factorize.

Finally, since the Higgs mass is much smaller than its total decay width, $\Gamma = 4,04$ MeV [44], the so-called narrow-width approximation (NWA) can be used to obtain the partonic cross section of the process $gg \to H \to \gamma\gamma$ as

$$\hat{\sigma}_{gg\to H\to\gamma\gamma} = \hat{\sigma}_{gg\to H} \times \Gamma_{H\to\gamma\gamma} \,. \tag{4.26}$$

We made use of this NWA to compute the signal strength, $\hat{\mu}$, defined as [7]

$$\hat{\mu}_{ggF} = \frac{\hat{\sigma}^{BSM}_{gg\to H\to\gamma\gamma}}{\hat{\sigma}^{SM}_{gg\to H\to\gamma\gamma}} \ , \tag{4.27}$$

by multiplying the ratios we had computed in Eq.(4.24) and Eq.(4.25):

$$\hat{\mu}_{ggF} = \frac{\hat{\sigma}^{BSM}_{gg\to H}}{\hat{\sigma}^{SM}_{gg\to H}} \times \frac{\Gamma^{BSM}_{H\to\gamma\gamma}}{\Gamma^{SM}_{H\to\gamma\gamma}} \ . \tag{4.28}$$

When writing this expression, another approximation is being made. In this product we are at some point multiplying both the production and the decay mediated by the t^*, which is suppressed by $1/\Lambda^4$ and, hence, goes two orders further in the $1/\Lambda$ expansion we are considering. However, since the moment when we introduced the cut-off, this perturbative expansion can no longer be regarded as such. Furthermore, the contributions coming from $1/\Lambda^4$ should be very small, specially at large masses. If one wanted to do things more properly one should compute the matrix element of the total process $gg \to H \to \gamma\gamma$ summing all possible diagrams and only then square and average it. We will assume that the approximation is good enough and analyze the results.

This signal strength is plotted in Fig. 4.6 for the three cut-offs together with the results from ATLAS [7] and CMS [8]. We see that our results for masses above the lower bound found by CMS [43] are within the errors of both ATLAS and CMS experiments. A mass for the t^* can be read off at around 800 GeV and 3700 GeV by comparison with ATLAS and CMS respectively. This discrepancy can be due to the difference in the $\hat{\mu}$'s measured in the two experiments ($\sim$ 0,2) and their wide errors ($\sim \pm$ 0,4).

Looking at three three plots Fig. 4.4, Fig. 4.5 and Fig. 4.6 together, the curves corresponding to the three cut-offs always converge as one goes to larger masses. Nevertheless, this convergence is not fast enough and the cut-off dependence is quite significant over a large region of the parameter space. This indicates that a more accurate approach should be sought to describe the effects of the t^*-quark in Higgs production and decay.

Concretely, there might exist a renormalizable theory for these spin-3/2 excitations that produces finite results, with no need to introduce a cut-off at all. A possible solution to our problem could be the procedure proposed by Pascalutsa in [45]. We ensured that the propagator did not describe additional spin-1/2 degrees of freedom by projecting to the RS subspace, but we did not check this for the vertices. The presence of redundant lower spin fields in the interactions imply that they are *inconsistent* and violate degree of freedom counting. This is solved in [45] by redefining the spin-3/2 field ψ_μ, resulting in a *consistent* coupling and additional "counter terms". It is hoped that applying this method to the present problem might give rise to a finite expression for the matrix element.

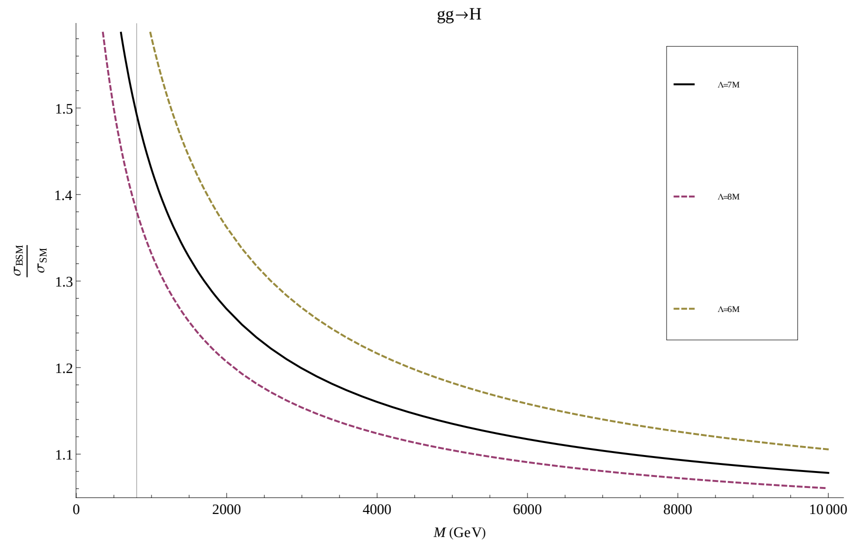


Figure 4.4: *Partonic cross section for Higgs production from gluon fusion in extensions of the SM involving a* t^**-quark,* σ_{BSM}*, divided by the SM one,* σ_{SM}*. The results are shown for three different cut-offs: at* $\Lambda = 7M_{t^*}$ *(as suggested in [4]),* $\Lambda = 8M_{t^*}$ *and* $\Lambda = 6M_{t^*}$*. At* $M_{t^*} = 803\,GeV$*, a vertical line represents the lower bound for the* M_{t^*} *found at 95% confidence by CMS in [43].*

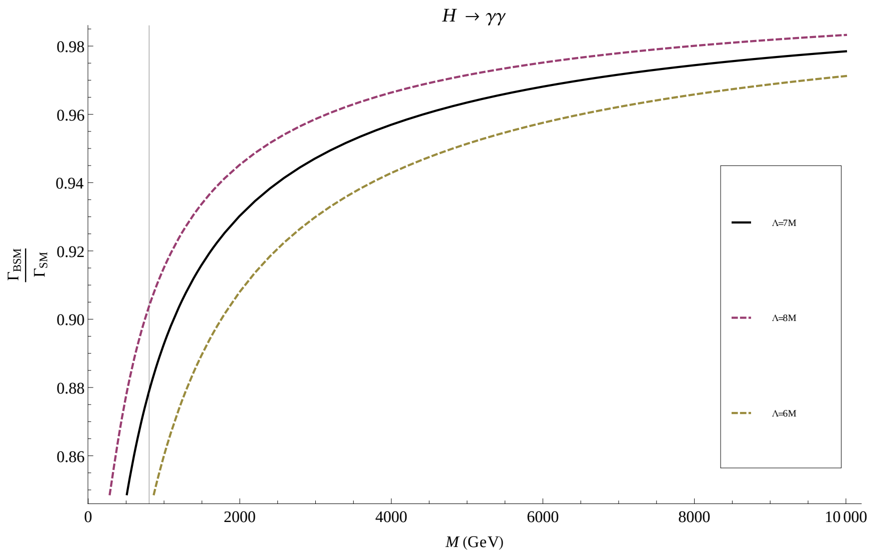


Figure 4.5: *Partial decay width of the Higgs boson into two photons in extensions of the SM involving a t^*-quark, Γ_{BSM}, divided by the SM one, Γ_{SM}. The results are shown for three different cut-offs: at $\Lambda = 7M_{t^*}$ [4], $\Lambda = 8M_{t^*}$ and $\Lambda = 6M_{t^*}$. At $M_{t^*} = 803 GeV$, a vertical line represents the lower bound for the M_{t^*} found at 95% confidence by CMS in [43].*

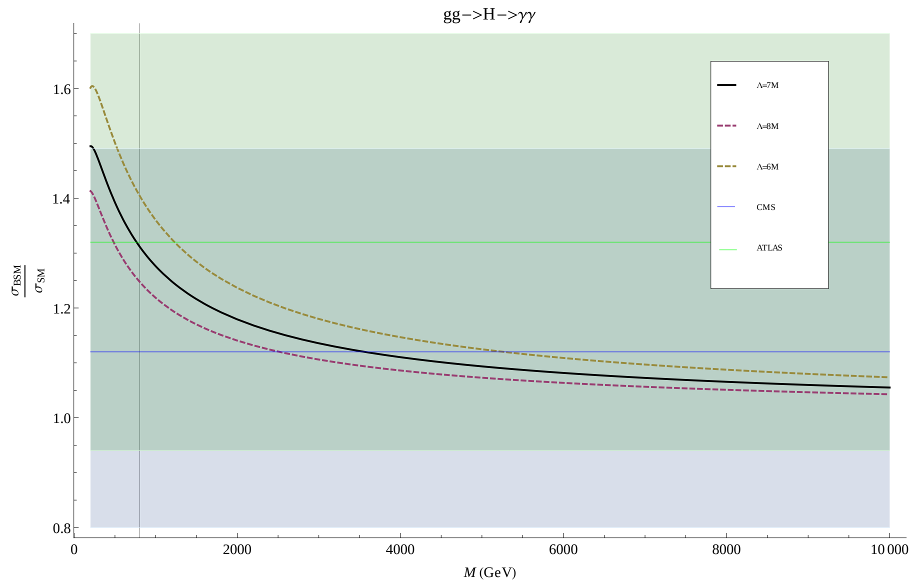


Figure 4.6: *Partonic cross section of Higgs production from gluon fusion combined with further decay into di-photons in extensions of the SM involving a* t^**-quark,* σ_{BSM}*, divided by the SM one,* σ_{SM}*. The results are shown for three different cut-offs: at* $\Lambda = 7M_{t^*}$ *[4],* $\Lambda = 8M_{t^*}$ *and* $\Lambda = 6M_{t^*}$*. The green line shows the result from ATLAS with the error band in light green [7]:* $\hat{\mu}^{ATLAS} = 1.32 \pm 0.38$*. The blue line with the light blue coloured region represents the result from CMS [8] at* $\hat{\mu}^{CMS} = 1.12^{+0.37}_{-0.32}$*. At* $M_{t^*} = 803\,GeV$*, a vertical line represents the lower bound for the* M_{t^*} *found at 95% confidence by CMS in [43].*

Chapter 5

Conclusions

In this work we have reviewed the basic concepts of spontaneous symmetry breaking in the electroweak sector and the Higgs decays to photons and gluons in the SM. Afterwards, a theoretical description of a spin-3/2 partner of the t-quark was carried out in the framework of EFT's. Putting this together allowed us to study the effect of the t^* quark in the processes of Higgs production and decay mediated by fermion loops. Firstly the Rarita-Schwinger case was chosen, leading to divergences in the matrix element that could not be removed by usual renormalization schemes. Identifying the spin-1/2 arising from the group representation as a possible origin of the divergences, the approach suggested by Habberzettl was adopted. However, even when using the singular case propagator the obtained expression for the matrix element is divergent. The author believes that the answer relies on eliminating the lower spin fields from the vertices, which could be achieved by the method proposed by Pascalutsa [45]. This would lead to a renormalizable theory in the description of spin-3/2 particles that would yield a finite result for the matrix element. It is likely that by the time when this is achieved, the statistics from the experiments at LHC will be better shrinking the error bands and allowing for a better comparison with our results. Eventually, the explanation, or not, of possible deviations from the SM in the Higgs production and decay cross section, could strongly support, or rule out, some models predicting the existence of these particles, such as those based on extra-dimensions or compositeness.

Appendix A

SM Higgs decay calculation

Here the matrix element of the diagram shown in Fig. A.1 will be computed explicitly[1].

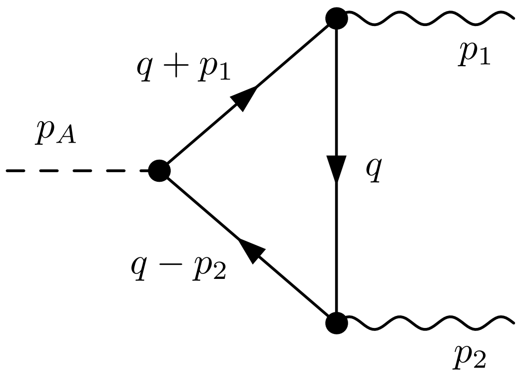


Figure A.1: *One of the two diagrams contributing to Higgs decay into photons in the SM via a t-quark loop.*

The starting point is Eq.(2.17),

$$\begin{aligned}\mathcal{M}_a = &- \frac{N_c e^3 e_q^2 m_t}{2\sin\theta_W M_W} \varepsilon^{\mu *}(p_1,\lambda_1)\varepsilon^{\nu *}(p_2,\lambda_2)\mu^{4-n} \\ &\times \int \frac{d^n q}{(2\pi)^n} \frac{\mathrm{Tr}\left[(\not{q} - \not{p}_2 + m_t)\gamma_\nu(\not{q} + m_t)\gamma_\mu(\not{q} + \not{p}_1 + m_t)\right]}{\mathcal{D}_1\mathcal{D}_2\mathcal{D}_3} \;, \end{aligned} \tag{A.1}$$

with the denominators defined in Section 2.2.1:

$$\mathcal{D}_1 = q^2 - m_t^2, \quad \mathcal{D}_2 = (q - p_2)^2 - m_t^2, \quad \mathcal{D}_3 = (q + p_1)^2 - m_t^2 \;. \tag{A.2}$$

[1]This calculation can be found *e.g.* in [46], but there some 1/2 factor is missing.

Firstly, use Feynman parametrization to rewrite the denominator:

$$\frac{1}{\mathcal{D}_1\mathcal{D}_2\mathcal{D}_3} = 2\int_0^1 dx \int_0^{1-x} dy \left\{\mathcal{D}_1(1-x-y) + \mathcal{D}_2 x + \mathcal{D}_3 y\right\}^{-3} . \tag{A.3}$$

Expanding the squares and summing we arrive at

$$2\int_0^1 dx \int_0^{1-x} dy \left(q^2 - m_t^2 - 2qp_2x + 2qp_1y\right)^{-3} . \tag{A.4}$$

In order to symmetrize the denominator, perform the change of variables $q' = q - xp_2 + yp_1$:

$$\begin{aligned} q'^2 = (q - xp_2 + yp_1)^2 = q^2 - 2qp_2x + 2qp_1y - 2xyp_1p_2 \Longrightarrow \\ \Longrightarrow q^2 - 2qp_2x + 2qp_1y = q'^2 + 2xyp_1p_2 . \end{aligned} \tag{A.5}$$

This can be plugged in the expression above and the denominator reads

$$2\int_0^1 dx \int_0^{1-x} dy \left(q'^2 + 2xyp_1p_2 - m_t^2\right)^{-3} = 2\int_0^1 dx \int_0^{1-x} dy \left(q'^2 - A\right)^{-3} \tag{A.6}$$

with $A = m_t^2 - 2xyp_1p_2$.

At this point, the trace of the numerator has to be computed. This was done with the help of `FORM`.

$$\begin{aligned} &\mathrm{Tr}\left[(\not{q} - \not{p}_2 + m_t)\gamma_\nu(\not{q} + m_t)\gamma_\mu(\not{q} + \not{p}_1 + m_t)\right] \\ &= 16q_\mu q_\nu m_t - 8q_\mu p_{2\nu} m_t + 8q_\nu p_{1\mu} m_t - 4p_{1\mu}p_{2\nu}m_t + 4p_{1\nu}p_{2\mu}m_t + 4\eta_{\mu\nu}m_t^3 \\ &- 4\eta_{\mu\nu}q^2 m_t - 4\eta_{\mu\nu}p_1p_2m_t \end{aligned} , \tag{A.7}$$

where $\eta_{\mu\nu}$ stands for the metric of the Minkowski space in the mostly minus convention: $\eta_{\mu\nu} = diag(1, -1, -1, -1)$. Take into account that the outgoing photons are on-shell and have no longitudinal polarization, that is: $p_{1\mu}\varepsilon^{\mu *}(p_1) = p_{2\nu}\varepsilon^{\nu *}(p_2) = 0$. We are left with

$$4m_t\left[4q_\mu q_\nu + p_{1\nu}p_{2\mu} - \eta_{\mu\nu}\left(q^2 - m_t^2 + p_1p_2\right)\right] . \tag{A.8}$$

The next step is to write q in terms of q'. This gives

$$\begin{aligned} &4m_t[4\left(q'_\mu + xp_{2\mu} - yp_{1\mu}\right)\left(q'_\nu + xp_{2\nu} - yp_{1\nu}\right) + p_{1\nu}p_{2\mu} \\ &\qquad - \eta_{\mu\nu}\left(q'^2 + 2q'p_2x - 2q'p_1y - 2xyp_1p_2 - m_t^2 + p_1p_2\right)] \\ &= 4m_t[4q'_\mu q'_\nu - 4q'_\mu yp_{1\nu} + 4q'_\nu xp_{2\mu} - 4xyp_{1\nu}p_{2\mu} + p_{1\nu}p_{2\mu} \\ &\qquad - \eta_{\mu\nu}\left(q'^2 + 2q'p_2x - 2q'p_1y - 2xyp_1p_2 - m_t^2 + p_1p_2\right)]. \end{aligned} \tag{A.9}$$

Bringing all pieces together, we get for the matrix element

$$\begin{aligned}
\mathcal{M}_a = &-\frac{4N_c e^3 e_q^2 m_t^2}{\sin\theta_W M_W}\varepsilon^{\mu *}(p_1,\lambda_1)\varepsilon^{\nu *}(p_2,\lambda_2)\mu^{4-n} \\
&\times \int \frac{d^n q'}{(2\pi)^n}\int_0^1 dx \int_0^{1-x} dy \left(q'^2 - A\right)^{-3} \\
&\times [4q'_\mu q'_\nu - 4q'_\mu y p_{1\nu} + 4q'_\nu x p_{2\mu} - 4xyp_{1\nu}p_{2\mu} + p_{1\nu}p_{2\mu} \\
&\quad - \eta_{\mu\nu}\left(q'^2 + 2q'p_2 x - 2q'p_1 y - 2xyp_1p_2 - m_t^2 + p_1p_2\right)] \; .
\end{aligned} \tag{A.10}$$

Terms proportional to q will cancel after integration and can be dropped. Renaming q' as q:

$$\begin{aligned}
\mathcal{M}_a = &-\frac{4N_c e^3 e_q^2 m_t^2}{\sin\theta_W M_W}\varepsilon^{\mu *}(p_1,\lambda_1)\varepsilon^{\nu *}(p_2,\lambda_2)\mu^{4-n} \\
&\times \int_0^1 dx \int_0^{1-x} dy \int \frac{d^n q}{(2\pi)^n} \left(q^2 - A\right)^{-3} \\
&\times [4q_\mu q_\nu - 4xyp_{1\nu}p_{2\mu} + p_{1\nu}p_{2\mu} \\
&\quad - \eta_{\mu\nu}\left(q^2 - 2xyp_1p_2 - m_t^2 + p_1p_2\right)] \; .
\end{aligned} \tag{A.11}$$

As discussed before, terms proportional to q^2 are expected to lead to a divergence. Let's compute the integral of these terms (this can be found in [13]).

$$\begin{aligned}
&\int \frac{d^n q}{(2\pi)^n}\,\frac{4q_\mu q_\nu}{(q^2-A)^3} - \int \frac{d^n q}{(2\pi)^n}\,\frac{\eta_{\mu\nu} q^2}{(q^2-A)^3} \\
&= 4\frac{(-1)^2 i}{(4\pi)^{n/2}}\frac{\eta_{\mu\nu}}{2}\frac{\Gamma(2-n/2)}{\Gamma(3)}\left(\frac{1}{A}\right)^{2-n/2} - \eta_{\mu\nu}\frac{(-1)^2 i}{(4\pi)^{n/2}}\frac{n}{2}\frac{\Gamma(2-n/2)}{\Gamma(3)}\left(\frac{1}{A}\right)^{2-n/2} \\
&= \frac{i\eta_{\mu\nu}}{2(4\pi)^{n/2}}\frac{\Gamma(2-n/2)}{\Gamma(3)}\left(\frac{1}{A}\right)^{2-n/2}(4-n) = \frac{i\eta_{\mu\nu}}{2(4\pi)^{2-\epsilon}}\epsilon\Gamma(\epsilon)\left(\frac{1}{A}\right)^{\epsilon} \; .
\end{aligned} \tag{A.12}$$

We have made the change $n = 4 - 2\epsilon$. In the limit $n \to 4$, that is, $\epsilon \to 0$, $\Gamma(\epsilon)$ can be expanded as $(1/\epsilon - \gamma_E + O(\epsilon))$ and we get

$$\frac{i\eta_{\mu\nu}}{2(4\pi)^2}\epsilon\left(\frac{1}{\epsilon} - \gamma_E + O(\epsilon)\right) = \frac{i\eta_{\mu\nu}}{2(4\pi)^2} + O(\epsilon) \; . \tag{A.13}$$

We see that there is indeed no divergence, as the ϵ coming from the trace in the numerator cancels the pole of the $\Gamma(\epsilon)$.

We calculate now the term q^0 from Eq.(A.11).

$$\int \frac{d^n q}{(2\pi)^n} \frac{-4xyp_{1\nu}p_{2\mu} + p_{1\nu}p_{2\mu} - \eta_{\mu\nu}\left(-2xyp_1p_2 - m_t^2 + p_1p_2\right)}{(q^2 - A)^3}$$

$$= b_{\mu\nu} \int \frac{d^n q}{(2\pi)^n}(q^2 - A)^{-3} = \frac{(-1)^3 i}{(4\pi)^{n/2}} \frac{\Gamma(3 - n/2)}{\Gamma(3)} \left(\frac{1}{A}\right)^{3-n/2} b_{\mu\nu} \; . \quad \text{(A.14)}$$

Taking the limit $n \to 4$ is straightforward:

$$-\frac{i}{2(4\pi)^2 A} b_{\mu\nu} \; . \quad \text{(A.15)}$$

Summing these two results gives

$$\begin{aligned} i\frac{\eta_{\mu\nu}A - b_{\mu\nu}}{2(4\pi)^2 A} &= \frac{i}{2(4\pi)^2 A}[\eta_{\mu\nu}(m_t^2 - 2xyp_1p_2) + 4xyp_{1\nu}p_{2\mu} - p_{1\nu}p_{2\mu} \\ &+ \eta_{\mu\nu}\left(-2xyp_1p_2 - m_t^2 + p_1p_2\right)] \\ &= \frac{i}{2(4\pi)^2 A}\left[4xyp_{1\nu}p_{2\mu} - p_{1\nu}p_{2\mu} + \eta_{\mu\nu}\left(-4xyp_1p_2 + p_1p_2\right)\right] \\ &= \frac{i}{2(4\pi)^2 A}\left(p_{1\nu}p_{2\mu} - \eta_{\mu\nu}p_1p_2\right)(4xy - 1) \; . \end{aligned} \quad \text{(A.16)}$$

The matrix element is

$$\begin{aligned} \mathcal{M}_a = &-\frac{4N_c e^3 e_q^2 m_t^2}{\sin\theta_W M_W} \varepsilon^{\mu *}(p_1, \lambda_1)\varepsilon^{\nu *}(p_2, \lambda_2) \\ &\times \int_0^1 dx \int_0^{1-x} dy \; \frac{i}{2(4\pi)^2 A}\left(p_{1\nu}p_{2\mu} - \eta_{\mu\nu}p_1p_2\right)(4xy - 1) \; . \end{aligned} \quad \text{(A.17)}$$

If we impose that the final photons are on shell, according to four momentum conservation, $m_H^2 = p_A^2 = (p_1 + p_2)^2 = 2p_1p_2$ and A can be written as

$$A = m_t^2 - 2xyp_1p_2 = m_t^2 - xym_H^2 = m_t^2\left(1 - xy\frac{m_H^2}{m_t^2}\right) \; . \quad \text{(A.18)}$$

Substituting this in the expression for the matrix element,

$$\begin{aligned} \mathcal{M}_a = &\frac{2N_c i e^3 e_q^2}{(4\pi)^2 \sin\theta_W M_W} \varepsilon^{\mu *}(p_1, \lambda_1)\varepsilon^{\nu *}(p_2, \lambda_2) \\ &\times \left(p_{1\nu}p_{2\mu} - \eta_{\mu\nu}\frac{m_H^2}{2}\right) I\left(\frac{m_H^2}{m_t^2}\right) \; , \end{aligned} \quad \text{(A.19)}$$

where $I(m_H^2/m_t^2)$ was defined as

$$I\left(\frac{m_H^2}{m_t^2}\right) = \int_0^1 dx \int_0^{1-x} dy \frac{1-4xy}{1-xy\frac{m_H^2}{m_t^2}} \ . \tag{A.20}$$

This integral can be found *e.g.* in [17] and reads (using the notation from [9])

$$I\left(\frac{m_H^2}{m_t^2}\right) = \frac{1}{4} A_{1/2}^H(\tau_t) \ , \tag{A.21}$$

with

$$A_{1/2}^H(\tau_t) = 2\left[\tau_t + (\tau_t - 1)\arcsin^2\sqrt{\tau_t}\right]\tau_t^{-2} \ . \tag{A.22}$$

The parameter $\tau_t = m_H^2/4m_t^2$ was taken to be smaller than one, according to the recent discoveries at LHC [1, 2].

Appendix B

Feynman Rules for Spin-3/2

We present the Feynman rules for the propagation and interaction of spin-3/2 particles suitable for dimensional regularization.

B.1 Propagators

The Lagrangian for the propagation of a spin-3/2 particle in n-dimensions was found in [47]:

$$\begin{aligned}\Lambda_A^{\mu\nu} &= \left(i\partial\!\!\!/ - M\right)\eta^{\mu\nu} + iA\left(\gamma^\mu\partial^\nu + \gamma^\nu\partial^\mu\right) \\ &+ \frac{i}{n-2}\left[(n-1)A^2 + 2A + 1\right]\gamma^\mu\partial\!\!\!/\gamma^\nu \\ &+ \frac{M\gamma^\mu\gamma^\nu}{(n-2)^2}\left[n(n-1)A^2 + 4(n-1)A + n\right] \ . \end{aligned} \tag{B.1}$$

B.1.1 Rarita-Schwinger Propagator

The Rarita-Schwinger propagator is obtained making $A = -1$ in Eq.(B.1) and inverting the Lagrangian

$$\Lambda_{-1}^{\rho\mu} G_{\mu\nu}^{(R.S.)} = \eta_\nu^\rho \tag{B.2}$$

one gets [47]

$$iG_{\mu\nu}^{(R.S.)}(p) = \frac{p\!\!\!/ + M}{(p^2 - M^2)}\left[\eta_{\mu\nu} - \frac{(n-2)p_\mu p_\nu}{(n-1)M^2} - \frac{\gamma_\mu\gamma_\nu}{n-1} + \frac{(p_\mu\gamma_\nu - p_\nu\gamma_\mu)}{(n-1)M}\right] . \tag{B.3}$$

B.1.2 Singular Case Propagator

The Singular Case propagator was derived following the procedure from [35]. Firstly, we computed the n-dimensional projector,

$$D_{\mu\nu} = \eta_{\mu\nu} - \frac{1}{n}\gamma_\mu\gamma_\nu \ . \tag{B.4}$$

Then we make $A = -1/2$ in Eq.(B.1) and solve the equation

$$D_{\mu\alpha}\Lambda^{\alpha\beta}_{-1/2}D_{\beta\rho}G^{\rho\nu}_{(S.C.)} = D^{\nu}_{\mu} \,. \tag{B.5}$$

With the help of the package `Tracer` [48] for *Mathematica* one can find[1]:

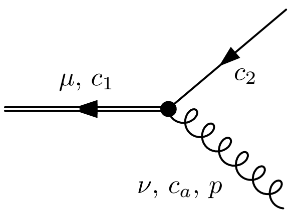
$$G^{(S.C.)}_{\mu\nu} = D_{\mu\rho}\left[\frac{\not{p}+M}{p^2-M^2}\eta^{\rho\sigma} + \frac{n}{n-2}p^{\rho}\frac{\not{p}+\frac{n}{n-2}M}{(p^2-M^2)\left(p^2-(n/(n-2))^2M^2\right)}p^{\sigma}\right]D_{\sigma\nu} \,. \tag{B.6}$$

B.2 Vertices

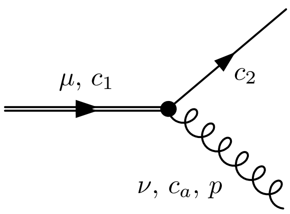
The vertices were derived from the interaction Lagrangian for the case when the gauge field is a gluon:

$$\mathcal{L} = i\,\frac{g}{\Lambda}\,\bar{\psi}_{\sigma}(\eta^{\sigma\mu}+z\gamma^{\sigma}\gamma^{\mu})\gamma^{\nu}T_a\xi G^a_{\mu\nu} - i\,\frac{g}{\Lambda}\bar{\xi}\gamma^{\nu}(\eta^{\sigma\mu}+z\gamma^{\mu}\gamma^{\sigma})\psi_{\sigma}T_aG^a_{\mu\nu} \,. \tag{B.7}$$

Notice that this Lagrangian and, thus the vertices, do not depend on the dimension n. From the first term one obtains the vertex $\Gamma_1^{\mu\nu}(p)$ and from the second term the $\Gamma_2^{\nu\mu}(p)$. The vertices read



$$i\,\Gamma_1^{\mu\nu}(p) = i\,\frac{g_s}{\Lambda}T^{c_a}_{c_1c_2}\left[p^{\mu}\gamma^{\nu} - \eta^{\mu\nu}\not{p} + z\gamma^{\mu}\not{p}\gamma^{\nu} - z\gamma^{\mu}\gamma^{\nu}\not{p}\right] \,,$$



$$i\,\Gamma_2^{\nu\mu}(p) = i\,\frac{g_s}{\Lambda}T^{c_a}_{c_1c_2}\left[-\gamma^{\nu}p^{\mu} + \not{p}\eta^{\mu\nu} + z\not{p}\gamma^{\nu}\gamma^{\mu} - z\gamma^{\nu}\not{p}\gamma^{\mu}\right] \,,$$

with all momenta *incoming*.

[1] The Author thanks B. Fuks for this contribution.

Appendix C

Passarino-Veltman Integral Reduction

The reduction of the tensor integrals were carried out according to the method proposed by Passarino and Veltman in [49]. The expressions that we present here were taken from [50]. The integrals are defined as

$$\begin{aligned} A_0(m_1) &\equiv (2\pi\mu)^{4-n} \int \frac{d^n q}{i\pi^2} \frac{1}{\mathcal{D}_1} \ , \\ B_{\{0,\,\mu,\,\mu\nu\}}(p_1, m_1, m_2) &\equiv (2\pi\mu)^{4-n} \int \frac{d^n q}{i\pi^2} \frac{\{1, q_\mu, q_\mu q_\nu\}}{\mathcal{D}_1 \mathcal{D}_2} \ , \\ C_{\{0,\,\mu,\,\mu\nu,\,\mu\nu\rho\}}(p_1, p_2, m_1, m_2, m_3) &\equiv (2\pi\mu)^{4-n} \int \frac{d^n q}{i\pi^2} \frac{\{1, q_\mu, q_\mu q_\nu, q_\mu q_\nu q_\rho\}}{\mathcal{D}_1 \mathcal{D}_2 \mathcal{D}_3} \ , \end{aligned} \tag{C.1}$$

with the denominators

$$\begin{aligned} \mathcal{D}_1 &= q^2 - m_1^2 + i\varepsilon \ , \\ \mathcal{D}_2 &= (q + p_1)^2 - m_2^2 + i\varepsilon \ , \\ \mathcal{D}_3 &= (q + p_1 + p_2)^2 - m_3^2 + i\varepsilon \ . \end{aligned} \tag{C.2}$$

These integrals can be decomposed applying Lorentz covariance:

$$\begin{aligned} B_\mu &= p_{1\mu} B_1 \ , \\ B_{\mu\nu} &= p_{1\mu} p_{1\nu} B_{21} + \eta_{\mu\nu} B_{22} \ , \\ C_\mu &= p_{1\mu} C_{11} + p_{2\mu} C_{12} \ , \\ C_{\mu\nu} &= p_{1\mu} p_{1\nu} C_{21} + p_{2\mu} p_{2\nu} C_{22} + \{p_1\ p_2\}_{\mu\nu} C_{23} + \eta_{\mu\nu} C_{24} \ , \\ C_{\mu\nu\rho} &= p_{1\mu} p_{1\nu} p_{1\rho} C_{31} + p_{2\mu} p_{2\nu} p_{2\rho} C_{32} + \{p_1\ p_1\ p_2\}_{\mu\nu\rho} C_{33} \\ &\quad + \{p_1\ p_2\ p_2\}_{\mu\nu\rho} C_{34} + \{p_1\ \eta\}_{\mu\nu\rho} C_{35} + \{p_2\ \eta\}_{\mu\nu\rho} C_{36} \ . \end{aligned} \tag{C.3}$$

Where we used the notation [50]

$$\{p_i\ p_j\ p_k\}_{\mu\nu\rho} \equiv \sum_{\sigma(i,j,k)} p_{\sigma(i)\mu} p_{\sigma(j)\nu} p_{\sigma(k)\rho} \ , \tag{C.4}$$

$\sigma(i,j,k)$ denoting all different permutations of (i,j,k), and

$$\{p_i\ \eta\}_{\mu\nu\rho} \equiv p_{i\mu}\eta_{\nu\rho} + p_{i\nu}\eta_{\mu\rho} + p_{i\rho}\eta_{\mu\nu} \ . \tag{C.5}$$

All the C_{ij}, B_{ij} and A_{ij} integrals can be further reduced to different C_0's, B_0's and A_0's. The relations needed for the B-integrals read

$$\begin{aligned}
B_1 =& \ \frac{1}{2p_1^2}\left[f_1 B_0 + A_0(M_1) - A_0(M_2)\right] \ , \\
B_{22} =& \ \frac{1}{n-1}\left\{M_1^2 B_0 - \frac{1}{2}\left[f_1 B_1 - A_0(M_2)\right]\right\} \ , \\
B_{21} =& \ \frac{1}{p_1^2}\left\{\frac{1}{2}\left[f_1 B_1 + A_0(M_2)\right] - B_{22}\right\} \ ,
\end{aligned} \tag{C.6}$$

with

$$f_1 = M_2^2 - M_1^2 - p_1^2 \ . \tag{C.7}$$

The reduction of the C-integrals is achieved by recursive matrix relations, much more cumbersome. It will not be shown here: it can be found, *e.g.* in [50].

For the B-integrals in which $p_1^2 = 0$ an alternative expression can be found for the B_1 and the B_{22}. Here, in that case we will leave them as B_1 and B_{22} and compute these numerically with *LoopTools* [42].

Furthermore, the divergent part of the integrals will be useful to isolate the divergences from the finite part of the matrix element. These are

$$\begin{aligned}
A_0^{div} =& \ \frac{M_1^2}{\bar\epsilon} \ , \\
B_0^{div} =& \ \frac{1}{\bar\epsilon} \ , \\
B_1^{div} =& \ -\frac{1}{2\bar\epsilon} \ , \\
B_{21}^{div} =& \ \frac{1}{3\bar\epsilon} \ , \\
B_{22}^{div} =& \ \frac{3M_1^2 + 3M_2^2 - p_1^2}{12\bar\epsilon} \ , \\
C_{24}^{div} =& \ \frac{1}{4\bar\epsilon} \ , \\
C_{35}^{div} =& \ -\frac{1}{6\bar\epsilon} \ , \\
C_{36}^{div} =& \ -\frac{1}{12\bar\epsilon} \ .
\end{aligned} \tag{C.8}$$

Appendix D

Triangle Diagram in the Rarita Schwinger approach

We show the result for the matrix element, Eq.(4.2), of the contributions from the t^*-quark to Higgs production from gluon fusion, Fig. 4.1, using the Rarita Schwinger propagator.

$$\mathcal{M} = -\frac{i m_t m_H^2 g_s^2 \mu^{4-n}}{48\pi^2 \sin\theta_W M_W}\, \delta_{c_1 c_2}\, \frac{1}{\Lambda^2}\, \varepsilon^*_\mu(p_1,\lambda_1)\, \varepsilon^{*\,\mu}(p_2,\lambda_2)$$
$$\times \Bigg[\frac{1}{\bar{\epsilon}}\mathcal{A}^{(1)} + A_0^{fin}(m_t^2)\mathcal{A}^{(2)} + B_0^{fin}(m_H^2, m_t^2, m_t^2)\mathcal{A}^{(3)} + B_0^{fin}(0, M_{t^*}^2, m_t^2)\mathcal{A}^{(4)}$$
$$+ B_1^{fin}(0, M_{t^*}^2, m_t^2)\mathcal{A}^{(5)} + C_0(0, m_H^2, 0, M_{t^*}^2, m_t^2, m_t^2)\mathcal{A}^{(6)} + \mathcal{A}^{(7)}\Bigg]\ , \tag{D.1}$$

with

$$\mathcal{A}^{(1)} = -\frac{2m_H^2}{M_{t^*}}(1+8z+16z^2)+8M_{t^*} \\ -\frac{2m_t m_H^2}{M_{t^*}^2}(1+4z)+\frac{96m_t^2 z}{M_{t^*}}(1+2z)+\frac{16m_t^3}{M_{t^*}^2}(1+4z)\ , \tag{D.2}$$

$$\mathcal{A}^{(2)} = \frac{8}{M_{t^*}}\left(\frac{1}{3}+4z+8z^2\right)+\frac{20m_t}{3M_{t^*}^2}(1+4z)\ , \tag{D.3}$$

$$\mathcal{A}^{(3)} = -\frac{2m_H^2}{M_{t^*}}\left(1+8z+16z^2\right)+\frac{2m_t m_H^2}{M_{t^*}^2}(1-4z) \\ +\frac{4m_t^2}{M_{t^*}}(1+16z+32z^2)+\frac{8m_t^3}{M_{t^*}^2}(1+4z)\ , \tag{D.4}$$

$$\mathcal{A}^{(4)} = \frac{16M_{t^*}}{3}+\frac{4m_t}{3}(1+4z)-\frac{4m_t^2}{M_{t^*}}\ , \tag{D.5}$$

$$\mathcal{A}^{(5)} = -\frac{16M_{t^*}}{3}+\frac{8m_t}{3}(1+4z)+\frac{16m_t^2}{3M_{t^*}}-\frac{8m_t^3}{3M_{t^*}^2}(1+4z)\ , \tag{D.6}$$

$$\mathcal{A}^{(6)} = -6M_{t^*}m_H^2+20m_t^2 M_{t^*}-\frac{4m_t^4}{M_{t^*}}\ , \tag{D.7}$$

$$\mathcal{A}^{(7)} = \frac{m_H^2}{3M_{t^*}}\left(-1+16z+80z^2\right)-16M_{t^*}+\frac{m_H^2 m_t}{M_{t^*}^2}\left(2+\frac{16}{3}z-\frac{32}{3}z^2\right) \\ +\frac{m_t}{3}(50+8z)+\frac{32m_t^2}{M_{t^*}}\left(\frac{1}{3}-z-5z^2\right)+\frac{m_t^3}{M_{t^*}^2}\left(-\frac{38}{3}-\frac{152}{3}z+64z^2\right)\ . \tag{D.8}$$

A_0^{fin}, B_0^{fin} and B_1^{fin} stand for the finite part of the integrals presented in Appendix C. These, together with C_0, will be computed with the *LoopTools* package for *Mathematica* [42]. From this expression the matrix element for the di-photon decay is obtained by exchanging $g_s \to e_q e$ and replacing the colour factor $1/2\ \delta_{c_1c_2} \to N_c$.

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Declaration of Authoship

I Alberto Marín González hereby declare that this thesis and the work presented in it is entirely my own, except where otherwise indicated. The use of all material from other sources has been properly and fully acknowledged

Granada, 8th of September, 2014

Signed:

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