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LHC physics beyond the Standard Model

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The search for physics beyond the Standard Model at the LHC follows mainly two orientations. One consists in testing specific models in order to obtain precise constraints or evidence for discovery. The other is to find general parameterisations capable of testing classes of models. In the following I will discuss two such parametrisations relevant for LHC searches. One describes how vector-like fermions which couple mainly to the third generation quarks, via Yukawa interactions, can be relevant for the LHC searches in the short term. The other is a study of the $H \to \gamma\gamma$ decay process and the gluon fusion production of a light Higgs which provides a general framework for testing models of new physics beyond the Standard Model at the LHC.

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PoS(Kruger 2010)056

1. Introduction

The quest for a general parametrisation of physics beyond the standard model at the LHC is a very attractive idea, but a too general study may turn out to be difficult to implement. Seaching for specific particles or processes in a class of models can be a good compromise to describe physical effects expected at the LHC. Considering the integrated luminosity accumulated till now and the expected evolution in the short term, one of the most interesting signals is the possible discovery of heavy vector-like fermions. In many models of new physics, like for example extra dimensional models, Little Higgs models, dynamical models, there are heavy vector-like fermions which decay to Standard Model (SM) fermions plus a boson (W/Z and/or Higgs h). Moreover, the mixing of vector-like quarks with the third generation and in particular the top quark is a common feature in little Higgs models and composite Higgs models. Previous collider and precision data place limits on new heavy quarks and set the lowest mass scale for these resonances once some properties for these particles are assumed. Direct searches give mass constraints in the range around 300 GeV, typically assuming a charged current decay chain [1]. Mixing effects with the SM quarks give stringent bounds in the case of mixing with the first two generations but only mild bounds for the mixing with the third generation. In order to keep the discussion general I discuss an effective approach where the decays are induced by a new Yukawa coupling [2]. This coupling generates the mixing of the new heavy fermion with top and bottom. I shall ignore in the present discussion the possible mixing to the light generations, which is discussed elsewhere [3, 4].

The second example is a relatively general parametrisation of the Higgs searches into two photons, which is one of the main LHC goals and which may be within reach in the short term too. This mode is also a powerful probe of the electroweak symmetry breaking sector of the theory, because it is a loop-induced process, therefore it is sensitive to any particle with a large coupling to the Higgs. In the SM it depends primarily on the couplings of the Higgs boson with heavy quarks (the top) and gauge bosons (the W), whose masses are tightly related to the electroweak scale. In extensions of the SM, particles that do couple strongly to the Higgs, and therefore play a role in the breaking of the electroweak symmetry, will also contribute to this loop and modify the SM prediction. Many models in fact predict the existence of partners of the top and W: stops and gauginos in supersymmetry, heavy W's and tops in extra dimensional models and Little Higgs models, to cite only a few possible extensions of the SM.

2. Heavy vector-like quarks

In the following we shall assume that the new vector-like quarks interact with the SM fermions via Yukawa interactions. The quantum numbers of the new fermions with respect to the weak SU(2)$_L\times$ U(1)$_Y$ gauge group are therefore limited by the requirement of an interaction with the Higgs doublet and one of the SM quarks. The SM contains a doublet $q_L = \{u_L, d_L\} = (2, Y)$ and two singlets $u_R = (1, Y + \frac{1}{2})$ and $d_R = (1, Y - \frac{1}{2})$ where $Y = \frac{1}{6}$ for quarks, and the Higgs $H = (2, \frac{1}{2})$. The SM Yukawa couplings are:

$$\mathscr{L}_{\text{Yuk}} = -y_u\, \bar{q}_L H^c u_R - y_d\, \bar{q}_L H d_R + h.c.\,. \tag{2.1}$$

Taking into account the quantum numbers of the SM particles one can establish the possible quantum number assignments for the new fermions. One can add a new singlet fermion with the same

PoS(Kruger 2010)056

hypercharge assignments as in the SM, namely $Y \pm \frac{1}{2}$. There are 3 possible doublets: one with the SM hypercharge Y, and two others with $Y \pm 1$. Finally, one can add two triplets with hypercharge $Y \pm \frac{1}{2}$. U and D are in the following the heavy partners of the up and down SM particles, and they will mix with the SM fermions. X if present is the extra fermion that does not mix with the SM ones, because of a different electric charge.

The coupling λ connects the heavy fermions with the SM ones and generates a mixing between the two states. There are two types of mixing: the singlets and triplets couple to the left-handed doublet, while the doublets couple with the right-handed singlets. In the following we will study these two cases, adding two heavy states, U and D, and parametrising their mixing with the SM states. This formalism can be easily adapted to the different representations of the heavy fermions.

For singlets and triplets, after the Higgs doublet develops a vacuum expectation value

$$\langle H \rangle = \begin{pmatrix} 0 \\ \frac{v+h}{\sqrt{2}} \end{pmatrix} , \tag{2.2}$$

where $v \sim 246$ GeV and h is the physical Higgs boson, the mass is

$$\mathscr{L}_{\text{mass}} = -\frac{y_u v}{\sqrt{2}} \bar{u}_L u_R - x \bar{u}_L U_R - M \bar{U}_L U_R + h.c. , \tag{2.3}$$

where $x \sim \lambda v$ with the proportionality factor depending on the representation U belongs to (a similar expression holds for down-type fermions). In the singlet case, a mass term $\bar{U}_L u_R$ is also allowed, however one can find a combination of U_R and u_R to remove such a parameter and redefine the Yukawa couplings.

The mass matrix is diagonalised by the two mixing matrices

$$V_u^{L,R} = \begin{pmatrix} \cos\theta_u^{L,R} & \sin\theta_u^{L,R} \\ -\sin\theta_u^{L,R} & \cos\theta_u^{L,R} \end{pmatrix} , \tag{2.4}$$

defined as

$$\begin{pmatrix} \cos\theta_u^L & -\sin\theta_u^L \\ \sin\theta_u^L & \cos\theta_u^L \end{pmatrix} \begin{pmatrix} \frac{y_u v}{\sqrt{2}} & x \\ 0 & M \end{pmatrix} \begin{pmatrix} \cos\theta_u^R & \sin\theta_u^R \\ -\sin\theta_u^R & \cos\theta_u^R \end{pmatrix} = \begin{pmatrix} m_t & 0 \\ 0 & m_{t'} \end{pmatrix} , \tag{2.5}$$

where $m_{t'} \geq M \geq m_t$. The relations between the three input parameters, the mixing angles and the masses is

$$\frac{y_u^2 v^2}{2} = m_t^2 \left(1 + \frac{x^2}{M^2 - m_t^2} \right) , \tag{2.6}$$

$$m_{t'}^2 = M^2 \left(1 + \frac{x^2}{M^2 - m_t^2} \right) , \tag{2.7}$$

$$\sin\theta_u^L = \frac{Mx}{\sqrt{(M^2 - m_t^2)^2 + M^2 x^2}} , \tag{2.8}$$

$$\sin\theta_u^R = \frac{m_t}{M} \sin\theta_u^L . \tag{2.9}$$

For $M \gg m_t$, the right-handed mixing angle is much smaller than the left-handed one.

PoS(Kruger 2010)056

For doublets:

$$\mathscr{L}_{\text{mass}} = -\frac{y_u v}{\sqrt{2}}\,\bar{u}_L u_R - x\bar{U}_L u_R - M\,\bar{U}_L U_R + h.c. \tag{2.10}$$

where :

$$\begin{pmatrix} \cos\theta_u^L & -\sin\theta_u^L \\ \sin\theta_u^L & \cos\theta_u^L \end{pmatrix} \begin{pmatrix} \frac{y_u v}{\sqrt{2}} & 0 \\ x & M \end{pmatrix} \begin{pmatrix} \cos\theta_u^R & \sin\theta_u^R \\ -\sin\theta_u^R & \cos\theta_u^R \end{pmatrix} = \begin{pmatrix} m_t & 0 \\ 0 & m_{t'} \end{pmatrix}; \tag{2.11}$$

note that now the formulas for the left- and right-handed mixing angles are exchanged:

$$\sin\theta_u^R = \frac{Mx}{\sqrt{(M^2-m_t^2)^2+M^2x^2}}\,, \tag{2.12}$$

$$\sin\theta_u^L = \frac{m_t}{M}\sin\theta_u^R\,. \tag{2.13}$$

In this case, therefore, it is the left-handed angle which should be small for large M. The couplings to Z, W and h can be written as two by two matrices in the mass eigenstate basis (the couplings with the photon and gluon stay diagonal due to gauge invariance). If we denote by g_W^{sm} and g_W^{ψ} the couplings of the W with the SM doublet and the new fermion respectively, the left-handed couplings can be written as:

$$g_{WL} = \left(V_d^L\right)^\dagger \cdot \begin{pmatrix} g_W^{sm} & 0 \\ 0 & g_W^{\psi} \end{pmatrix} \cdot V_u^L = \begin{pmatrix} g_W^{sm} c_d^L c_u^L + g_W^{\psi} s_d^L s_u^L & g_W^{sm} c_d^L s_u^L - g_W^{\psi} s_d^L c_u^L \\ g_W^{sm} s_d^L c_u^L - g_W^{\psi} c_d^L s_u^L & g_W^{sm} s_d^L s_u^L + g_W^{\psi} c_d^L c_u^L \end{pmatrix}, \tag{2.14}$$

where s and c stand for the sin and cos of the mixing angles. The same formula applies for the right-handed couplings, with $g_W^{sm}=0$:

$$g_{WR} = \left(V_d^R\right)^\dagger \cdot \begin{pmatrix} 0 & 0 \\ 0 & g_W^{\psi} \end{pmatrix} \cdot V_u^R = \begin{pmatrix} g_W^{\psi} s_d^R s_u^R & -g_W^{\psi} s_d^R c_u^R \\ -g_W^{\psi} c_d^R s_u^R & g_W^{\psi} c_d^R c_u^R \end{pmatrix}. \tag{2.15}$$

Note that $g_W^{sm} = \frac{g}{\sqrt{2}}$, and g_W^{ψ}, the same for left- and right-handed components, depends on the representation: it is equal to the SM one for a doublet and equal to $\pm g$ for a triplet. Note also that in the case where either U or D are absent, the same formulas can be used by setting $g_W^{\psi}=0$ and setting to zero the absent mixing angle. Similarly, a general matrix formula can be written for the Z couplings of both left- and right-handed ups and downs:

$$g_{Zf} = (V_f)^\dagger \cdot \begin{pmatrix} g_Z^{sm} & 0 \\ 0 & g_Z^{\psi} \end{pmatrix} \cdot V_f = \begin{pmatrix} g_Z^{sm}c^2 + g_Z^{\psi}s^2 & (g_Z^{sm}-g_Z^{\psi})sc \\ (g_Z^{sm}-g_Z^{\psi})sc & g_Z^{sm}s^2 + g_Z^{\psi}c^2 \end{pmatrix}. \tag{2.16}$$

The Z couplings can be always expressed as functions of the weak isospin and charge of the fermion:

$$g_Z(T_3,Y) = \frac{g}{\cos\theta_W}\left(T_3 - \sin^2\theta_W Q\right). \tag{2.17}$$

Finally, the Higgs couplings in the two cases are:

$$\lambda_h = \left(V^L\right)^\dagger \cdot \begin{pmatrix} \frac{y}{\sqrt{2}} & \frac{x}{v} \\ 0 & 0 \end{pmatrix} \cdot V^R = \begin{pmatrix} c^L\left(c^R\frac{y}{\sqrt{2}} - s^R\frac{x}{v}\right) & c^L\left(s^R\frac{y}{\sqrt{2}} + c^R\frac{x}{v}\right) \\ s^L\left(c^R\frac{y}{\sqrt{2}} - s^R\frac{x}{v}\right) & s^L\left(s^R\frac{y}{\sqrt{2}} + c^R\frac{x}{v}\right) \end{pmatrix}, \tag{2.18}$$

$$\lambda_h = \left(V^L\right)^\dagger \cdot \begin{pmatrix} \frac{y}{\sqrt{2}} & 0 \\ \frac{x}{v} & 0 \end{pmatrix} \cdot V^R = \begin{pmatrix} c^R\left(c^L\frac{y}{\sqrt{2}} - s^L\frac{x}{v}\right) & s^R\left(c^L\frac{y}{\sqrt{2}} - s^L\frac{x}{v}\right) \\ c^R\left(s^L\frac{y}{\sqrt{2}} + c^L\frac{x}{v}\right) & s^R\left(s^L\frac{y}{\sqrt{2}} + c^L\frac{x}{v}\right) \end{pmatrix}. \tag{2.19}$$

PoS(Kruger 2010)056

2.1 Constraints from existing data

The bound to the coupling of the W to top and bottom arises from the observation of single top production at the TeVatron: we allow a variation of $\pm 20\%$ [5]. A tighter constraint originates from the unitarity of the CKM mixing matrix and flavour physics: however such a bound is not applicable here because it does not take into account the effect of the heavy fermions, and it is also sensitive to the contributions of other particles in the model (see [6] for a detailed study of bounds for singlet quarks and [7] for a vector-like up-type quark or a fourth generation). The couplings of the Z to the bottom are very constrained[8]: in the left-handed coupling a $+1\%$ and -0.2% deviations are allowed; in the right-handed one, $+20\%$ and -5%. In the cases under study, only one of the two is affected, so that those limits are sufficient even though the bounds are correlated; in the case where they are both present, a more detailed fit is required.

For the oblique corrections, we calculated the contribution to the T parameter [9]. A detailed study is given in [10, 11]. We allow for a deviation of $+0.4$ and -0.2: we consider a tighter bound on negative values because it is generically more difficult to accommodate for a negative shift in T. For instance, increasing the Higgs mass with respect to the reference value will generate an effective negative contribution. This is a very conservative bound and we use it just to underline the power of oblique constraints with respect to the tree level ones. As mentioned above, model dependent contributions from other heavier particles may be relevant and give rise to cancellations, therefore significantly modifying the allowed parameter space.

Direct searches at the TeVatron are another important bound : 335 GeV for a t' state [12], and 385 GeV for a b'[13], however, these bounds assume 100% branching ratios $t' \to Wq$ and $b' \to Wt$. This is true for a fourth generation, but in our case decays involving neutral bosons will play an important role.

2.2 Non-standard doublets

Among the possible cases under study, the less constrained by these data is the one of non-standard doublets, on which we focus in the following. In the case $\psi = (2, \frac{7}{6}) = \{X, U\}^T$, the vector-like fermion contains a top partner together with a new fermion X with charge $\frac{5}{3}$. The Yukawa couplings involve the left-handed component of ψ:

$$\begin{aligned} \mathscr{L}_{\text{Yuk}} &= -y_u\, \bar{q}_L H^c u_R - \lambda\, \bar{\psi}_L H u_R - M\, \bar{\psi}_L \psi_R + h.c. \\ &= -\frac{y_u v}{\sqrt{2}}\, \bar{u}_L u_R - \frac{\lambda v}{\sqrt{2}}\, \bar{U}_L u_R - M\,(\bar{U}_L U_R + \bar{X}_L X_R) + h.c.\,. \end{aligned} \tag{2.20}$$

In this case $x = \frac{\lambda v}{\sqrt{2}}$: only the up mixing is present, and the left-handed angle is smaller. The only tree level bound comes from the left-handed Wtb coupling:

$$\frac{\delta g_W}{g_W^{sm}} = \cos\theta_u^L - 1 \sim -\frac{1}{2}\frac{x^2 m_t^2}{M^4}\,. \tag{2.21}$$

Due to the extra m_t^2 suppression, the tree level bounds are negligible. The T parameter can receive both a positive and a negative contribution. For positive T we fix the bound at 0.4, and the curve does not depend much on the precise value (solid blue line in Figure 1).

PoS(Kruger 2010)056

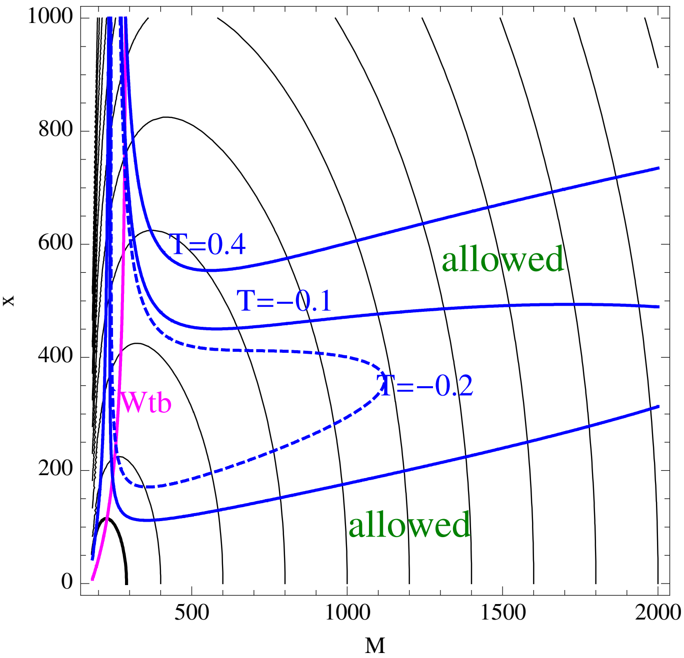

Figure 1: *non-SM doublet : in magenta the bound from Wtb, in blue from the T parameter, in black the direct exclusion limit from the TeVatron. The grey lines mark constant values of the $m_{t'}$ mass (the value can be read from the intersection with the $x = 0$ axis).*

For negative contributions, we impose a tighter bound at -0.2, as it is generically more difficult to accommodate for a negative shift in T. In the latter case, the curve is very sensitive to the precise value (blue dashed line in Figure 1) and two fine tuned regions on the small M branch are allowed.

The physical spectrum contains a top partner t' and a lighter new fermion X with charge $\frac{5}{3}$ and mass $m_X = M$. The only decay channel for X is into W^+t, where the W is virtual if $M < m_t + m_W$. If $m_{t'} - m_X > m_W$, then $t' \to W^+X$ mostly, with a sub-leading channel in $t' \to (Z,h)t$. The $t' \to W^+b$ channel is suppressed by an extra power of v/M in the coupling. If $m_{t'} - m_X < m_W$, then $t' \to (Z,h)t$ is the main channel with a small contribution to $t' \to W^+b$.

In the case $\psi = (2, -\frac{5}{6}) = \{D, X\}^T$, a bottom partner and a fermion X with charge $-\frac{4}{3}$ are present, and a tighter constraint comes from deviation of the Z couplings with the bottom. The strongest tree level bound comes from corrections to the $Z\bar{b}b$ coupling:

$$\frac{\delta g_{ZbL}}{g^{sm}_{ZbL}} = -\frac{2}{1-\frac{2}{3}\sin^2\theta_W}\sin^2\theta^L_d \sim -\frac{2}{1-\frac{2}{3}\sin^2\theta_W}\frac{x^2 m_b^2}{M^4}\,, \tag{2.22}$$

$$\frac{\delta g_{ZbR}}{g^{sm}_{ZbR}} = \frac{3\sin^2\theta^R_d}{2\sin^2\theta_W} \sim \frac{3}{2\sin^2\theta_W}\frac{x^2}{M^2}\,. \tag{2.23}$$

The tree level decays are similar to the other doublet, replacing top with bottom.

For small values of x, the tree level BR saturates rapidly to $BR(ht) \sim BR(Zt) \sim 50\%$ while the decay in Wb is very suppressed. Intermediate values of x are excluded by a negative contribution to T, while large values of x are again allowed: however, $x > 500$ GeV is at the edge of the non

PoS(Kruger 2010)056

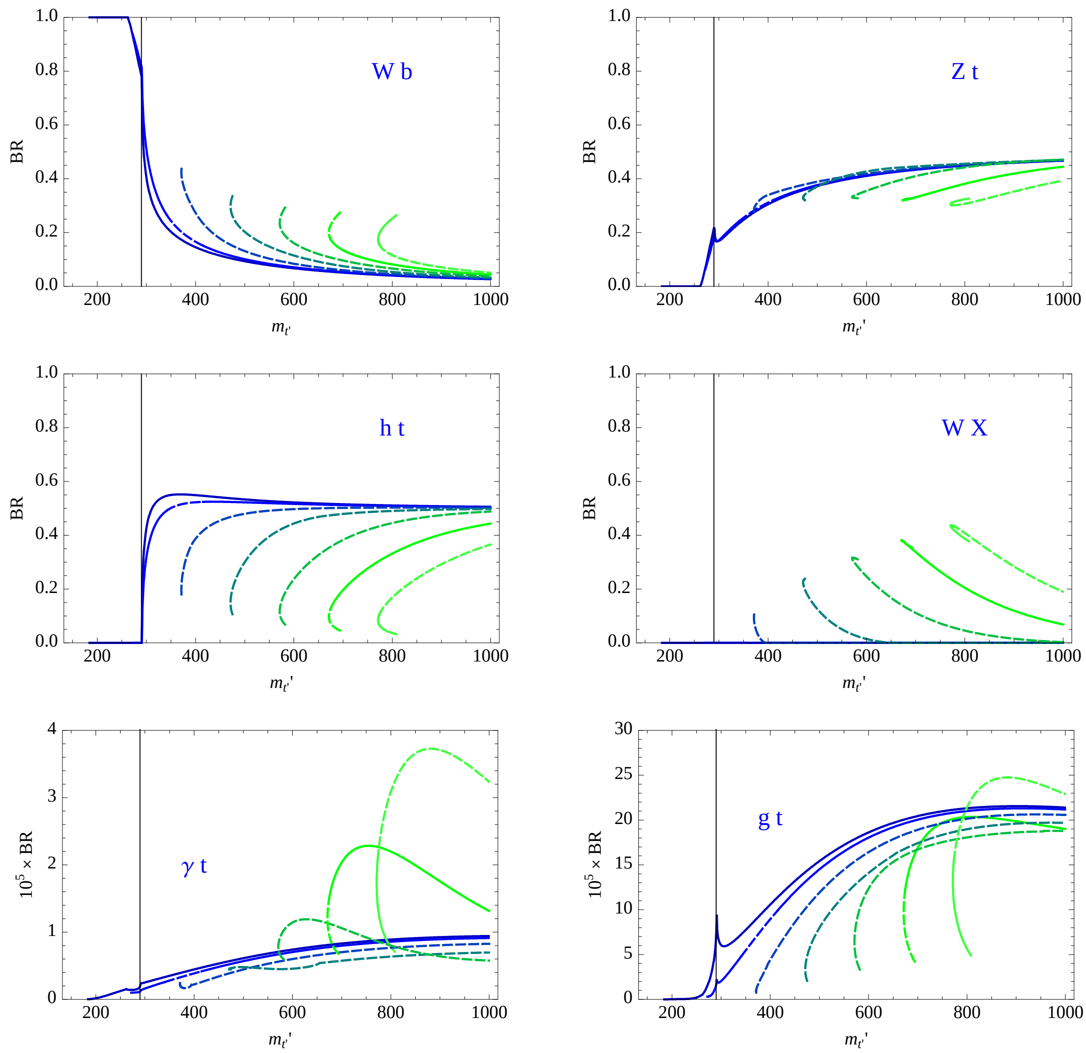

Figure 2: *non-SM doublet: the lines correspond to $x = 10, 100, 200, 300, 400, 500, 600$ GeV from darker (blue) to lighter grey (green). The dotted portions are excluded by the T parameter. The vertical line marks the direct exclusion by the TeVatron.*

PoS(Kruger 2010)056

perturbative regime and the tree level and one-loop results cannot be trusted. In this case, we also observe larger values of the loop BR compared to other cases.

The new states will be abundantly produced both in the 7 TeV and in the 14 TeV phases of the LHC for fairly low masses. At 7 TeV, the strong pair production cross section is above 1 pb for $m_{t'} < 600$ GeV, while at 14 TeV we have cross sections larger than 100 fb for masses below a TeV. Another even more interesting channel at the LHC is the single production of the heavy fermion, see for example [14, 15]. Even though the process is mediated by electroweak bosons and Higgs, the cross sections can be large and depend crucially on the value of the new Yukawa coupling.

3. A general parametrisation of $H \to \gamma\gamma$

In the SM the contribution of heavy particles to $H \to \gamma\gamma$ and $H \to gg$ processes does not decouple for particle masses much larger than the Higgs boson one. The reason is that these SM masses are uniquely generated by the coupling to the Higgs boson and the mass dependence of their coupling cancels the mass dependence in the loop integral. In general extensions of the SM this is not in general the case, as the masses may receive other contributions. In the following I will review a general parametrisation introduced in [16]. The effect on the decay can therefore be sensitive to the mass scale of the new physics. Studying this channel in detail can give some hints about the model of new physics, and this information will be complementary to the direct discovery of new states at the LHC.

In order to establish the notations, we will briefly review the decay of the Higgs into photons and gluons (the decay width in gluons is directly related to the gluon-fusion production cross section at hadronic colliders). The decay widths can be written as:

$$\Gamma_{\gamma\gamma} = \frac{G_F \alpha^2 m_H^3}{128\sqrt{2}\pi^3} \left| A_W(\tau_W) + \sum_{\text{fermions}} N_{c,f} Q_f^2 A_F(\tau_f) + \sum_{NP} N_{c,NP} Q_{NP}^2 A_{NP}(\tau_{NP}) \right|^2 , \tag{3.1}$$

$$\Gamma_{gg} = \frac{G_F \alpha_s^2 m_H^3}{16\sqrt{2}\pi^3} \left| \frac{1}{2} \sum_{\text{quarks}} A_F(\tau_f) + \sum_{NP} C(r_{NP}) A_{NP}(\tau_{NP}) \right|^2 , \tag{3.2}$$

where $\tau_x = \frac{m_H^2}{4m_x^2}$, $N_{c,x}$ is the number of colour states in the colour representation (3 for quarks, 1 for leptons), the constant $C(r)$ is an SU(3) colour factor (defined as $\mathrm{Tr}[t_r^a t_r^b] = C(r)\delta^{ab}$ where t_r^a are the SU(3) generators in the representation r; it is equal to $1/2$ for the quarks and 3 for an adjoint), Q_x is the electric charge of the particle in the loop, and the functions $A(\tau)$ depends on the spin and couplings to the Higgs of the particle running in the loop.

In the SM, all masses are proportional to the Higgs vacuum expectation value (VEV) v, therefore the couplings to the Higgs can be written as

$$y^{SM}_{h\bar{f}f} = \frac{m_f}{v} \quad \text{for fermions}\,, \tag{3.3}$$

$$g^{SM}_{h\phi\phi} = 2\frac{m_\phi^2}{v} \quad \text{for bosons}\,. \tag{3.4}$$

Under this assumption, the amplitudes are given by (F stands for spin-$1/2$ fermions, W for vector bosons and S for scalar bosons) [17]

$$A_F(\tau) = \frac{2}{\tau^2}\left(\tau + (\tau-1) f(\tau)\right) , \tag{3.5}$$

$$A_W(\tau) = -\frac{1}{\tau^2}\left(2\tau^2 + 3\tau + 3(2\tau-1) f(\tau)\right) , \tag{3.6}$$

$$A_S(\tau) = -\frac{1}{\tau^2}\left(\tau - f(\tau)\right) ; \tag{3.7}$$

where

$$f(\tau) = \begin{cases} \arcsin^2\sqrt{\tau} & \tau \leq 1 \\ -\frac{1}{4}\left[\log\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi\right]^2 & \tau > 1 \end{cases} . \tag{3.8}$$

PoS(Kruger 2010)056

For our study we are particularly interested in the limit of such functions for large mass of the particle in the loop with respect to the Higgs mass, $\tau \ll 1$:

$$A_F(0) = \frac{4}{3}\,, \quad A_W(0) = -7\,, \quad A_S(0) = \frac{1}{3}\,. \tag{3.9}$$

Note that the particle in the loop does not decouple for large mass because the (SM) coupling to the Higgs is also proportional to the mass of the particle. As we are interested in Higgs masses below the W threshold and above the LEP limit (where the $\gamma\gamma$ signal is non negligible), the light Higgs approximation is useful for the top and the new physics. For the W this approximation is not valid, and the function $A_W(\tau_W)$ ranges from -8 for $m_H = 115$ GeV to -9.7 for $m_H = 150$ GeV.

The new physics can be parametrised by two independent parameters describing the contribution of the new particles to the two decay widths, however using the actual amplitude is not a convenient way of treating the new contributions. One convenient possibility is to normalise the new contribution to the top one. The main reason is that the top gives the main contribution to the amplitudes in the SM, and any new physics (which addresses the problem of the Higgs mass naturalness). Moreover, as it will soon be clear, these two parameters can give some intuitive information about what kind of new physics runs into the loop. The widths can be rewritten as

$$\Gamma_{\gamma\gamma} = \frac{G_F \alpha^2 m_H^3}{128\sqrt{2}\pi^3} \left| A_W(\tau_W) + 3\left(\frac{2}{3}\right)^2 A_t(\tau_t)\,[1+\kappa_{\gamma\gamma}] + \dots \right|^2, \tag{3.10}$$

$$\Gamma_{gg} = \frac{G_F \alpha_s^2 m_H^3}{16\sqrt{2}\pi^3} \left| \frac{1}{2} A_t(\tau_t)\,[1+\kappa_{gg}] + \dots \right|^2, \tag{3.11}$$

where the dots stand for the negligible contribution of the light quarks and leptons, and the coefficients κ can be written as:

$$\kappa_{\gamma\gamma} = \sum_{NP} \frac{3}{4} N_{c,NP} Q_{NP}^2\, \frac{v}{m_{NP}} \frac{\partial m_{NP}}{\partial v}\, \frac{A_{F,W,S}(m_{NP})}{A_t}\,, \tag{3.12}$$

$$\kappa_{gg} = \sum_{NP} 2C(r_{NP})\, \frac{v}{m_{NP}} \frac{\partial m_{NP}}{\partial v}\, \frac{A_{F,W,S}(m_{NP})}{A_t}\,, \tag{3.13}$$

where the ratio of A functions depends on the spin and masses of the new particles (and top). In the light Higgs approximation, however, the ratio only depends on the spin of the new particle:

$$\frac{A_{NP}}{A_t} = \begin{cases} 1 & \text{for fermions} \\ -\frac{21}{4} & \text{for vectors} \\ \frac{1}{4} & \text{for scalars} \end{cases} \tag{3.14}$$

An interesting feature of this parameterisation is that a particle with the same quantum numbers of the top will give $\kappa_{\gamma\gamma} = \kappa_{gg}$, and a single particle will give a contribution to the two coefficients with the same sign. In this way, if the experimental data allows us to point to a specific quadrant in the $\kappa_{\gamma\gamma}$–κ_{gg} parameter space, we can have a hint of the underlying new physics model. Note that the modification of the SM couplings will affect in general the other production channels, and the branching ratio into heavy gauge bosons. These effects will, however, have a minor impact on our analysis, and their inclusion will be necessary in a more detailed model-dependent analysis, if a model is preferred by data.

PoS(Kruger 2010)056

3.1 Observables at the LHC and model testing

The LHC will measure the inclusive $\gamma\gamma$ Higgs decays and new physics will modify both the total production cross section and the branching fraction in photons. For large masses, close to the W threshold, the decay into two heavy gauge bosons (one is virtual) becomes relevant and will also yield a relatively early measurement. At large luminosities one may also measure the $\gamma\gamma$ decays in a specific production channel, for instance the vector boson fusion one that can be isolated using two forward jet tagging: in this case one may probe directly the branching ratios.

In the Higgs mass range of interest, between 115 and 150 GeV, the main production channel is gluon fusion with a SM cross section of $40-25$ pb, followed by vector boson fusion ($5-4$ pb) and by other channels (WH, ZH, $\bar{t}tH$) which sum up to $4-2$ pb. Here we will assume that the new physics significantly contributes only to the loop in the gluon fusion channel, while the other cross sections are unaffected. The total production cross section normalised with the SM one, that we denote as $\bar{\sigma}$, can be written as:

$$\bar{\sigma}(H) = \left(\frac{\sigma_{gg}^{NP} + \sigma_{VBF}^{SM} + \sigma_{VH,\bar{t}tH}^{SM}}{\sigma_{gg}^{SM} + \sigma_{VBF}^{SM} + \sigma_{VH,\bar{t}tH}^{SM}} \right) \simeq \left(\frac{(1+\kappa_{gg})^2 \sigma_{gg}^{SM} + \sigma_{VBF}^{SM} + \sigma_{VH,\bar{t}tH}^{SM}}{\sigma_{gg}^{SM} + \sigma_{VBF}^{SM} + \sigma_{VH,\bar{t}tH}^{SM}} \right) . \tag{3.15}$$

In the SM the Higgs branching fraction in photons amounts to 2×10^{-3}. In the presence of new physics, the branching fraction will also be sensitive to the gluon loop via the total width, as the gluon channel is significant: it amounts to 7% of the total for $m_H = 115$ GeV, decreasing to 3% for $m_H = 150$ GeV. We therefore define a branching ratio normalised to the SM value, $\overline{BR}$

$$\begin{aligned} \overline{BR}(H\to\gamma\gamma) &= \frac{\Gamma_{\gamma\gamma}^{NP}}{\Gamma_{\gamma\gamma}^{SM}} \frac{\Gamma_{\text{tot}}^{SM}}{\Gamma_{gg}^{NP} + \Gamma_{\gamma\gamma}^{NP} + \Gamma_{\text{others}}^{SM}} \\ &\simeq \left(1 + \frac{\kappa_{\gamma\gamma}}{\frac{9}{16} A_W(\tau_W) + 1} \right)^2 \frac{\Gamma_{\text{tot}}^{SM}}{(1+\kappa_{gg})^2 \Gamma_{gg}^{SM} + (\Gamma_{\text{tot}}^{SM} - \Gamma_{gg}^{SM})} . \end{aligned} \tag{3.16}$$

The branching ratio in heavy vectors will depend on κ_{gg} via the total width of the Higgs, therefore the normalised $\overline{BR}$ is

$$\overline{BR}(H\to VV^*) = \frac{\Gamma_{\text{tot}}^{SM}}{\Gamma_{gg}^{NP} + \Gamma_{\gamma\gamma}^{NP} + \Gamma_{\text{others}}^{SM}} \simeq \frac{\Gamma_{\text{tot}}^{SM}}{(1+\kappa_{gg})^2 \Gamma_{gg}^{SM} + (\Gamma_{\text{tot}}^{SM} - \Gamma_{gg}^{SM})} . \tag{3.17}$$

We considered as an example the following models:

- [♦] a fourth generation (the result is independent on the masses and Yukawa couplings);

- [♣] supersymmetry in the MSSM golden region: we only included the contribution of the stops with the spectrum given by the benchmark point in [18]. In this case the result is very sensitive to the parameters in the superpotential and in the SUSY breaking terms, therefore the general MSSM will cover a region of the parameter space;

- [▲] Simplest Little Higgs, the result scales with the W' mass (in the plots, $m_{W'} = 2$ TeV);

- [$*$] Littlest Higgs, the result scales with the symmetry breaking scale f and has a mild dependence on the triplet VEV x (we set $x = 0$): for a model with T-parity we use $f = 500$ GeV, without T parity $f = 5$ TeV;

PoS(Kruger 2010)056

- [■] colour octet model, the result depends on 2 free parameters: for illustration we use in the plots $X_1 = 1/9$ and $X_2 = 1/36$;

- [▶] Lee-Wick Standard Model, the result scales with the LW Higgs mass: in the plots we set it to 1 TeV for illustration;

- [$\otimes$] Universal Extra Dimension model [19], where only the top and W resonances contribute and the result scales with the size of the extra dimension: here we set $m_{KK} = 500$ GeV close to the experimental bound;

- [★] the model of Gauge Higgs unification in flat space in [20], where only the W and top towers contribute ($\beta = m_t L$), with the first W resonance at 2 TeV;

- [•] the Minimal Composite Higgs [21] (Gauge Higgs unification in warped space) with the IR brane at $1/R' = 1$ TeV: only W and top towers contribute significantly. The point only depends on the overall scale of the KK masses, as the other parameters are fixed by the W and top masses;

- [▼] a flat (W' at 2 TeV) and [♠] warped ($1/R'$ at 1 TeV) version of brane Higgs models, in both cases the hierarchy in the fermionic spectrum is explained by the localisation, and all light fermion towers contribute. The result only depends on the overall scale of the KK masses.

In the numerical results the value of the mass of the new particles is at or around the lower bound given by precision electroweak tests; for larger masses, the contribution scales like the inverse squared mass (with the exception of the fourth generation). Note that in many cases, the result only depends on one mass scale, and is insensitive to other free parameters present in the model: for example, in extra dimensional models with flavour, the final result does not depend on the precise localisation pattern of the bulk fields.

The LHC will surely be able to measure the inclusive cross section $\sigma(pp \to H \to \gamma\gamma)$, as this is one of the golden channels for the discovery of a light Higgs. For an integrated luminosity of 10 fb^{-1} we can expect a 10% accuracy with respect to the Standard Model one [22]. We plotted the inclusive cross section normalised by the SM value in the $\kappa_{\gamma\gamma}$–κ_{gg} parameter space for a light Higgs ($m_H = 120$ GeV) in Figure 3 and for a Higgs near the VV-threshold ($m_H = 150$ GeV) in Figure 4: many models lie very far from such a line, and a 10% measurement would allow us to probe new physics masses up to a few TeV in some cases. Note that many of the models we studied predict a reduction of the inclusive signal: the measurement of an enhancement at the LHC may be a sign of unexpected new physics. Note also that some very different models can give the same prediction, like the fourth generation case where a suppression in the $\gamma\gamma$ decay is accidentally compensated by an enhancement in the gluon fusion cross section. Therefore, we need to measure another observable at the LHC in order to distinguish such models. For the light Higgs case, in Figure 3 we plotted the vector boson fusion channel, which is sensitive to the $\gamma\gamma$ branching fraction directly. This channel is orthogonal to the inclusive one, and therefore offers the best discrimination power.

Such a simple parametrisation assumes that the tree level physics is Standard model-like. This hypothesis applies to many models of new physics, in particular to those discussed in the text. An

PoS(Kruger 2010)056

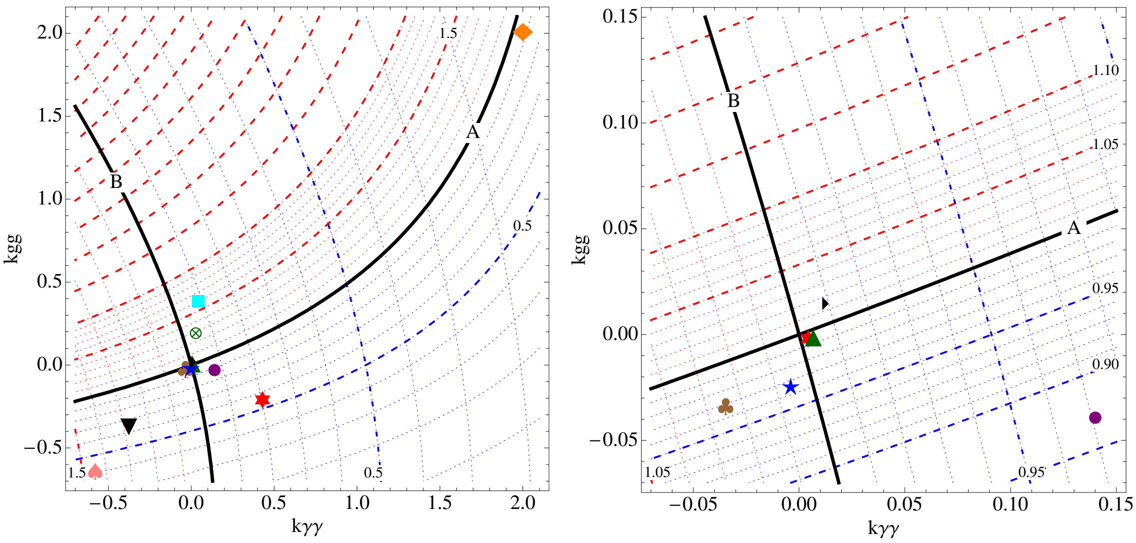

Figure 3: $\kappa_{\gamma\gamma}$ and κ_{gg} at the LHC for a light Higgs ($m_H = 120$ GeV). The two solid lines correspond to the SM values of the inclusive $\gamma\gamma$ channel (**A**), and the vector boson fusion production channel (**B**). On the left panel, the dashed lines are spaced by 0.5, while the dotted ones by 0.1. On the right, we zoomed near the SM point.

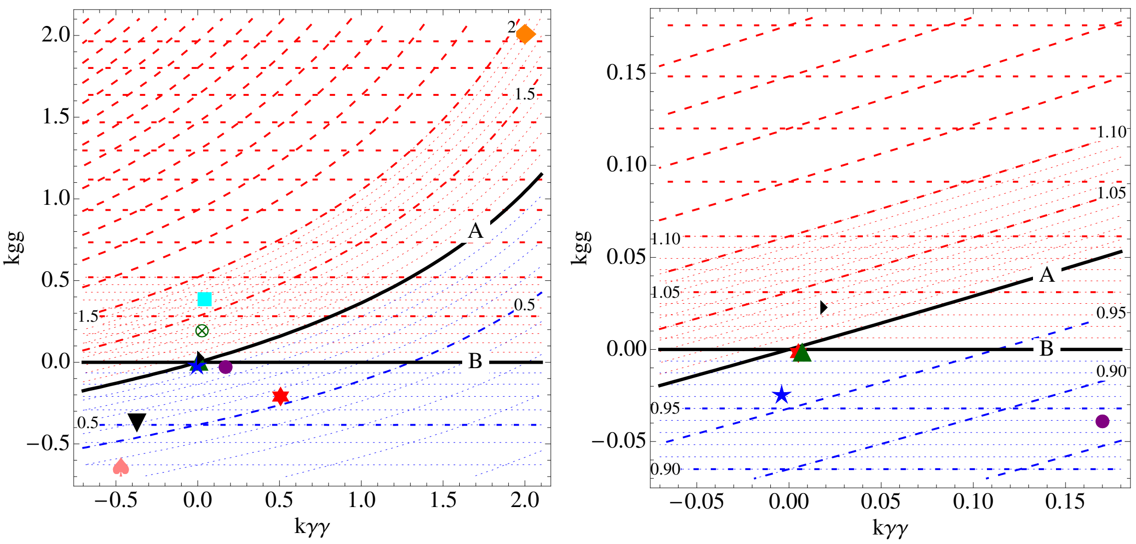

Figure 4: $\kappa_{\gamma\gamma}$ and κ_{gg} at the LHC for a Higgs near the WW threshold ($m_H = 150$ GeV). The two solid lines correspond to the SM values of the inclusive $\gamma\gamma$ channel (**A**), and the inclusive V^*V channel ($V = W, Z$) (**B**). On the left panel, the dashed lines are spaced by 0.5, while the dotted ones by 0.1. On the right, we zoomed near the SM point.

exception is supersymmetry, for which in general large tree level modifications of the $H \to b\bar{b}$ are possible. However we showed that in most of the parameters space in which the $H \to \gamma\gamma$ channel relevant to our formalism can still be applied. On more general grounds the inclusion of tree-level effects which differ from the SM can be taken care of by introducing new κs. We avoided doing this in order to keep the parametrisation as simple as possible and in practice many models are

PoS(Kruger 2010)056

actually SM-like in the sense discussed above.

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PoS(Kruger 2010)056