Electroweak Symmetry Breaking and the Scalar Boson: Confronting Theories at Colliders

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We will discuss implications of the recently discovered scalar boson at 125 GeV on the physics Beyond Standard Model(BSM). The proceedings will be focused on the constraints coming from the measurements of the Higgs couplings at LHC on the composite and supersymmetric theories

1 Introduction

Higgs[1] field is the only missing element of the Electroweak Symmetry breaking mechanism. Recently both collaborations at LHC reported discovery [2,3] of a new resonance, which might be the Higgs boson of the Standard Model(SM). However the mass of the Higgs boson in the SM suffers from the hierarchy problem, which predicts its natural value to be of the order of the Planck scale. This issue can be addressed within various BSM scenarios, where the new resonances at the scale of a few TeV stabilize the radiative corrections to the Higgs mass. Direct searches for the new states at LHC put strong constraints on these theories. Discovery of the new boson at LHC provides a new way of testing these ideas: generically the BSM Higgs boson couplings are modified compared to the SM values and observation of the significant deviation will be a new physics signal. In this proceedings we will discuss bounds on supersymmetry and composite Higgs models coming from the measurements of the Higgs couplings at LHC[a]. The paper is organized as follows : first we will discuss the single Higgs effective lagrangian, then we will review the constraints on the Composite Higgs models and supersymmetric theories and after that we will conclude.

2 Single Higgs effective theory

As we mentioned in the previous section, the generic BSM theories predict modifications of the Higgs couplings. We can parametrize these couplings using electroweak chiral lagrangian with all the possible additional interactions involving a singlet scalar h[5]. LEP constraints[6,7] on $\Delta\rho$ parameter force our lagrangian to be symmetric under custodial $SU(2)_L \times SU(2)_R$ symmetry, and we will assume that the electroweak symmetry breaking pattern is $SU(2)_L \times SU(2)_R \to SU(2)_V$. Then the Nambu-Goldstone (NG) bosons of $SU(2)_L \times SU(2)_R/SU(2)_V$ symmetry breaking can be described by the 2×2 matrix field

$$\Sigma(x) = \exp\left(i\sigma^a \chi^a(x)/v\right) , \tag{1}$$

[a] See for an example [4] for exhaustive list of the literature on the subject

where σ^a, χ^a are the Pauli matrices and the corresponding Goldstones bosons and $v = 246\,\text{GeV}$. The chiral lagrangian can then be systematically expanded in numbers of derivatives acting on the Σ field. In our case, when we have an additional custodial singlet scalar h we should add all the possible interactions between this field and the chiral lagrangian, which can be written as

$$\mathcal{L} = -V(h) + \mathcal{L}^{(2)} + \mathcal{L}^{(4)} + \ldots, \tag{2}$$

where $\mathcal{L}^{(n)}$ includes the terms with n derivatives and $V(h)$ is the potential for h. At the level of two derivatives we have

$$\begin{aligned} \mathcal{L}^{(2)} = \quad & \tfrac{1}{2}(\partial_\mu h)^2 + \tfrac{v^2}{4} Tr\left(D_\mu \Sigma^\dagger D^\mu \Sigma\right)\left(1 + 2a\,\tfrac{h}{v} + b\,\tfrac{h^2}{v^2} + \cdots\right) \\ - \quad & \tfrac{v}{\sqrt{2}}\lambda^u_{ij}\,(\bar{u}^{(i)}_L, \bar{d}^{(i)}_L)\,\Sigma\,(u^{(i)}_R, 0)^T\left(1 + c_u\,\tfrac{h}{v} + c_{2u}\,\tfrac{h^2}{v^2} + \cdots\right) + h.c. \\ + \quad & (u_R \Leftrightarrow d_R, c_u \Leftrightarrow c_d) + ((u, d \Leftrightarrow \nu, e), c_u \Leftrightarrow c_l) \end{aligned} \tag{3}$$

SM corresponds to the point where all $a = b = c_i = 1$ and $c_{2i} = 0$. At the level of four derivatives the bosonic operators that lead to cubic and quartic vertices of NG bosons and gauge fields with one or two Higgs bosons are:

$$\begin{gathered} O_1 = \mathrm{Tr}\Big[(D_\mu\Sigma)^\dagger(D^\mu\Sigma)\Big]\,(\partial_\nu F_1(h))^2, \quad O_2 = \mathrm{Tr}\Big[(D_\mu\Sigma)^\dagger(D_\nu\Sigma)\Big]\,\partial^\mu F_2(h)\partial^\nu F_{2'}(h), \\ O_{GG} = G_{\mu\nu}G^{\mu\nu}\,F_{GG}(h), \quad O_{BB} = B_{\mu\nu}B^{\mu\nu}\,F_{BB}(h), \\ O_W = D_\mu W^a_{\mu\nu}\,\mathrm{Tr}\Big[\Sigma^\dagger\sigma^a i \overleftrightarrow{D}_\nu\,\Sigma\Big]\,F_W(h), \quad O_B = -\partial_\mu B_{\mu\nu}\,\mathrm{Tr}\Big[\Sigma^\dagger i \overleftrightarrow{D}_\nu\,\Sigma\,\sigma^3\Big]\,F_B(h), \\ O_{WH} = i\,W^a_{\mu\nu}\,\mathrm{Tr}\Big[(D^\mu\Sigma)^\dagger\sigma^a D^\nu\Sigma\Big]\,F_{WH}(h), \quad O_{BH} = -i\,B_{\mu\nu}\,\mathrm{Tr}\Big[(D^\mu\Sigma)^\dagger(D^\nu\Sigma)\sigma^3\Big]\,F_{BH}(h), \\ O_{W\partial H} = \tfrac{1}{2}\,W^a_{\mu\nu}\,\mathrm{Tr}\left[\Sigma^\dagger\sigma^a i \overleftrightarrow{D}^\mu\,\Sigma\right]\partial^\nu F_{W\partial H}(h), \quad O_{B\partial H} = -\tfrac{1}{2}\,B_{\mu\nu}\,\mathrm{Tr}\left[\Sigma^\dagger i \overleftrightarrow{D}^\mu\,\Sigma\sigma^3\right]\partial^\nu F_{W\partial B}(h), \\ \text{where} \qquad F_i(h) = \alpha_i^{(0)} + \alpha_i^{(1)}\,h + \alpha_i^{(2)}\,h^2 + \ldots \end{gathered} \tag{4}$$

We can see that at the level of four derivatives we have ten operators, and all of them have a priori arbitrary coefficients, however some of them are irrelevant for discussing Higgs physics at LHC. For example O_1 and O_2 operators always involve at least two Higgses bosons, thus we can safely ignore them now. $O_{GG}(O_{BB})$ are very important because they contribute to the Higgs production from gluon fusion and decays to gluons/photons. O_W and O_B are constrained from the electroweak precision S parameter[6,7], so that we expect them to be small. O_{WH}/O_{BH} involve at least three W bosons and we can ignore them. $O_{W\partial H}/O_{B\partial H}$ contribute to the Higgs couplings to $Z\gamma$,ZZ and WW and are phenomenologically important, however these operators are generated at least at one loop level so we expect their effects in WW/ZZ couplings to be suppressed compared to the a parameter of Eq.3, and experimental $Z\gamma$ constraints are still weak[8].

3 Bounds on the Composite Higgs

One of the most attractive solutions to the SM hierarchy problem is provided in the Composite Higgs models, where the Higgs boson appears as a composite state[9] of some new strong dynamics. The Higgs can be made much lighter than the rest of the composite resonances by implementing Pseudo-Goldstone Mechanism[9]. This symmetry is explicitly broken by the SM gauge and Yukawa interactions, which lead to the generation of the Higgs mass. Note that the same symmetry arguments constrains heavily the operators, which can appear in the effective Higgs theory once all the composite resonances are integrated out. For the Higgs interactions with vector bosons we get

$$\begin{aligned} \mathcal{L} &= \frac{f^2}{4}\sin^2\left(\frac{h}{f}\right) W^+_\mu W^{\mu-} \\ a &= \cos\left(\frac{< h >}{f}\right) = \sqrt{1-\xi}, \end{aligned} \tag{5}$$

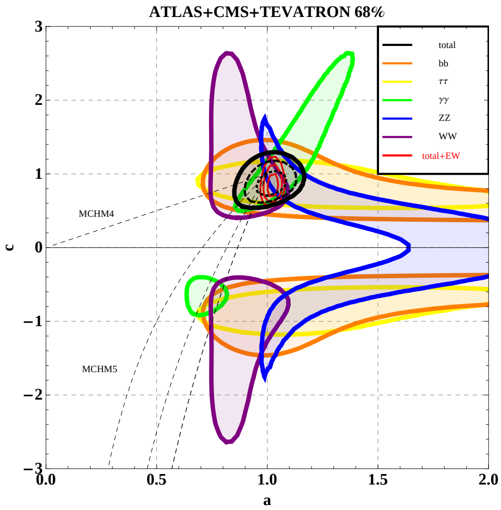

Figure 1: Fit of the Higgs couplings in the (a, c) plane, coloured contours indicate the bounds coming from individual contributions, black- $68, 95, 99\%$ probability contours from the LHC and Tevatron, red-the same with the data from LEP

where f is an analogue of the pion decay constant. Embedding of the SM fermions in the composite framework is model dependent and within this paper we will consider only the theories, where the SM fermion masses are generated using the partial compositeness mechanism [10], which protects the flavor violating observables (see for example [11] for the most recent discussion of the constraints). The Higgs interactions with fermions within this framework are rescaled by some trigonometric function

$$\mathcal{L} \propto \bar{\psi}_L \psi_R \sin\left(\frac{h}{f}\right)^{1+2m} \cos\left(\frac{h}{f}\right)^{n}$$
$$c_\psi = \frac{1 + 2m - (1 + 2m + n)\xi}{\sqrt{1-\xi}}, \tag{6}$$

where the parameters (m, n) are fixed by the representations of the global group and are model dependent[b]. In the case when top, bottom and tau mix with the same representations of the composite fermions, all the modifications of the Higgs couplings are characterized by only two parameters a, c and the model space can be represented as contours in the (a, c) plane. The Fig.1 shows fits of the couplings in the (a, c) plane as well contours for various composite Higgs models. The point $a = c = 1$ corresponds to the SM and is achieved in the decoupling limit when $f \to \infty$. We see from the Fig.1 that electroweak precision observables push towards the $a = 1$ point, which require

$$\frac{v^2}{f^2} < 0.04 \to \quad f > 1.2\text{TeV} \tag{7}$$

note that the strongest constraints are still from the electroweak precisions and are indirect, thus they can be overcome with the contributions of the new resonances.

4 Constraints on Supersymmetry

The other solution to the Planck weak hierarchy problem is realized within the models with TeV scale supersymmetry. These type of models also predict generic modifications of the Higgs

[b]Note that there are additional terms in the expression for the fermion mass, however they are suppressed by the additional powers of the mass of the composite resonance M_*, so the Eq. 6 becomes exact in the limit $f/M_* \ll 1$

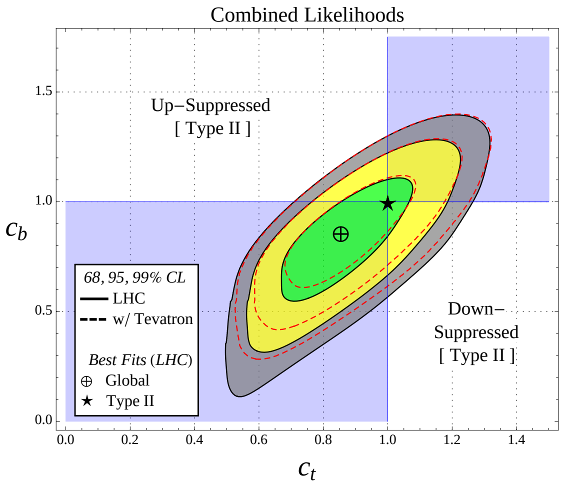

Figure 2: Fit of the Higgs couplings in the (a, c_t, c_b) plane. The plot is result of the marginalization over the a parameter in the $[0, 1]$ range

couplings and the constraints from LHC already put interesting bounds on the susy parameter space. Note that supersymmetric theories require the presence of the second doublet by the holomorphy of the superpotential as well as by the anomaly freedom of the theory. The holomorphy of the superpotential requires that one of the Higgs doublets couples only to the up quarks and the other doublet couples only to the down quarks, as a result couplings to down and up quarks are modified in a different way, so we need to analyse the couplings in the three dimensional parameter space (a, c_b, c_τ)

$$a = \sin(\beta - \alpha),\ c_t = \frac{\cos\alpha}{\sin\beta},\ c_b = c_\tau = -\frac{\sin\alpha}{\cos\beta},$$

where $\tan\beta = v_u/v_d$ is ratio of the vevs of two doublets and α is the mixing angle between two CP even scalars. Generically we have also independent contributions to the O_{GG} and O_{BB} operators from the loops with stops and staus, which we will ignore for the beginning. This approximation is justified in the limit, when both stops and staus are much heavier than the second Higgs doublet, the contribution to the $c_{\gamma\gamma}$ of the charged Higgs is generically small and we will ignore it here. The results of the three dimensional scan are shown on the Fig.2, where we have marginalized over the a -Higgs coupling to the vector bosons. Note that supersymmetric models can populate only the white regions of the plot and the MSSM is located almost always in the region with up suppressed and down enhanced Higgs couplings. Another way to look at the Higgs couplings in supersymmetry is to assume that "supersymmetric" parametrization of the couplings in Eq.(8) is correct and look for the preferred direction in the $(\sin\alpha, \tan\beta)$ space(see Fig.3). In this parametrization the SM point becomes a line $(\beta - \alpha = \pi/2)$, which is the MSSM in the decoupling limit, i.e. when the masses of the all new particles tend to infinity. The plot on on Fig.3, shows generic preference for the decoupling limit, without any specific preference for any $\tan\beta$. Note that the shape of the contours on the Fig.3 are dominated by the constraints on c_b parameter, indeed expressing α in terms of the decoupling parameter m_Z^2/m_A^2 we get

$$\begin{aligned} a &\sim 1 - \frac{2m_Z^4}{m_A^4}\frac{\cos^2 2\beta \sin^4\beta}{\tan^2\beta} \\ c_b &\sim 1 - \frac{2m_Z^2}{m_A^2}\cos 2\beta \sin^2\beta \\ c_t &\sim 1 + \frac{2m_Z^2}{m_A^2}\frac{\cos 2\beta \sin^2\beta}{\tan^2\beta}. \end{aligned} \tag{8}$$

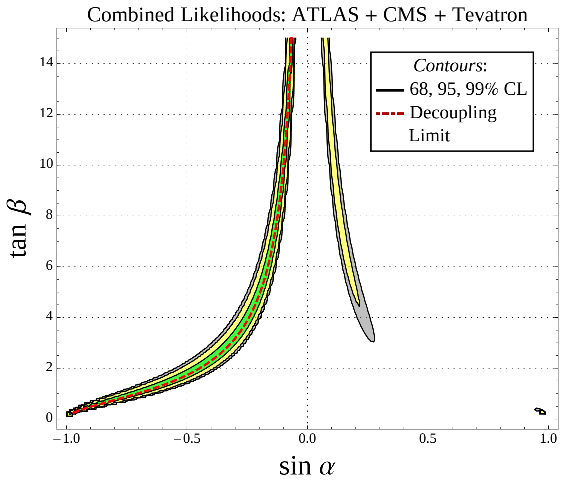

Figure 3: Fit of the Higgs couplings in the $\sin\alpha, \tan\beta$ plane. Red line indicates decoupling limit, which corresponds to the SM

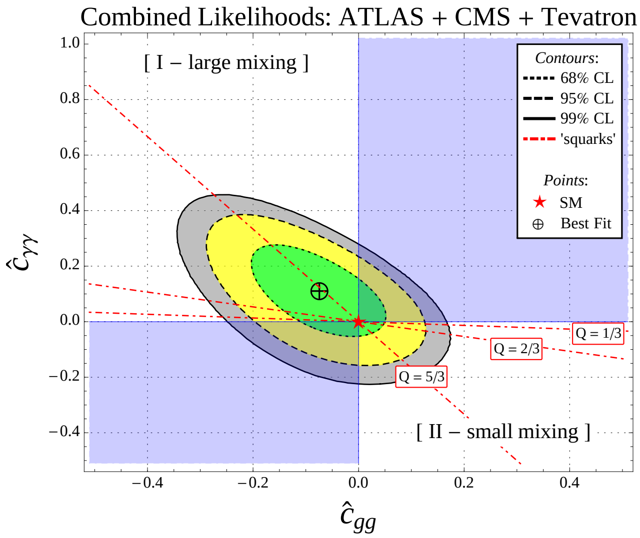

Figure 4: Fit of the Higgs couplings in the $c_{gg}, c_{\gamma\gamma}$ plane. New contributions are normalized to the SM values. Red lines indicate indicate the contours of the contributions of new fields, which are in the fundamentals of the QCD and have various electrc charges Q.

The modification to the a is suppressed by the m_Z^4/m_A^4 and to c_t by $m_Z^2/m_A^2/\tan^2\beta$, so that c_b is the most sensitive parameter to the scale of new physics.

4.1 Loops of new fields

So far we have been ignoring completely direct contributions of the new fields to the O_{GG} and O_{BB} operators. However in the case of the light stops (or any other charged/colored resonances) these might be the dominant modifications to the Higgs couplings. The contribution to the $c_{gg}/c_{\gamma\gamma}$ coefficients can be easily calculated using the Higgs Low Energy Theorems [12] and for example in the case of stops we get

$$\tilde{m}_t^2 \sim \begin{pmatrix} m_Q^2 + y_t^2 v^2 s_\beta^2 & y_t s_\beta v X_t \\ y_t s_\beta v X_t & m_u^2 + y_t^2 v^2 s_\beta^2 \end{pmatrix}$$

$$c_{gg} \propto \frac{\partial\, \mathrm{Det} \log \tilde{m}_t^2}{\partial \log v} \sim \frac{1}{4}\left(\frac{m_t^2}{m_{\tilde{t}_1}^2} + \frac{m_t^2}{m_{\tilde{t}_2}^2} - \frac{m_t^2 X_t^2}{m_{\tilde{t}_1}^2 m_{\tilde{t}_2}^2}\right), \tag{9}$$

where $c_{gg} = 1$ corresponds to the contribution by the infinitely heavy top quark. The results of the fit are shown on the Fig.4, where both c_{gg} and $c_{\gamma\gamma}$ are normalized in terms of their SM contributions. In the case, when we have only one new particle in the loop only two regions (unshaded on Fig.4) are possible, where region I(II) corresoponds to the large(small) mixing ((X_t) term). We can see from the Fig.4 that the SM prediction is within 68% probability contour.

5 Summary

We have reviewed the current status of the fits of the Higgs couplings, and the corresponding constraints on the new physics models. We have discussed the composite Higgs and supersymmetric models, both of these models provide a solution to the hierarchy problem and can be also tested by the measurements of the Higgs couplings. So far the experimental data prefers SM limit of these models.

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