SU(5) Octet Scalar at the LHC

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Color scalars are salient features of non-minimal SU(5) model, where Higgs sector is extended by 45-dimensional multiplet. We show that gauge coupling unification can be realized in this model with TeV octet scalar and intermediate (∼ 10^6 GeV) color-triplet scalar at scale larger than 10^{15} GeV. We analyze the possible LHC signatures of these TeV octet scalars. We emphasize that multi-(b)-jet final states provide significant signal for direct probe of octet scalars at the LHC.

I. INTRODUCTION

The major goal of the Large Hadron Collider (LHC) experiments at CERN is to probe new physics beyond the Standard Model (SM) at TeV energy scales. The grand unified theory based on SU(5) gauge symmetry is one of the most appealing scenarios for possible extension of the SM. The minimal SU(5) accommodates the matter fields in 5* and 10 dimensional representations, while the scalar sector consists of 24 and 5 dimensional Higgs multiplets [1].

However, the minimal SU(5) suffers from several sever problems. For instance, it does not seem to unify the SM gauge couplings. In addition it predicts a wrong fermions mass relation: m_{µ(e)} = m_{e(µ)} that contradicts the experiment measurements. A possible approach to overcome some of these problems, is to introduce an extra Higgs multiplet with 45_H dimensional representation [2-4]. The 45_H transforms under the SM gauge as

\[ 45_H = (8, 2)_{1/2} \oplus (1, 1)_{1/2} \oplus (3, 1)_{-1/3} \oplus (3, 3)_{-1/3} \oplus (6^*, 1)_{-1/3} \oplus (3^*, 2)_{-7/6} \oplus (3^*, 1)_{4/3}. \]  

It also satisfies the following constraints: 45_H^{αβ} = -45_H^{βα} and \( \sum_α (45_H)^{αβ} = 0 \). Thus, the electroweak symmetry SU(2)_L × U(1)_Y can be spontaneously broken into U(1)_em through the non-vanishing Vacuum Expectation Value (VEV) of 5_H and 45_H, namely

\[ \langle 5_H \rangle = v_5, \quad \langle 45_H \rangle^{15} = \langle 45_H \rangle^{25} = \langle 45_H \rangle^{35} = v_{45}, \quad \langle 45_H \rangle^{45} = -3v_{45}. \]  

The 45_H-doublet is defined as [4]

\[ D \equiv (1, 2)_{1/2} = \left( \frac{(45_H)_5^{54}}{(45_H)_4^{54}}, \frac{(45_H)_5^{45}}{(45_H)_4^{45}} \right) \equiv \left( -\frac{(45_H)_5^{45}}{(45_H)_4^{45}}, \frac{(45_H)_4^{45}}{(45_H)_5^{45}} \right) = \left( -D^+, D^0 \right). \]

While, the 45_H-color octet scalars are given by [3]

\[ S_j^{iat} \equiv (8, 2)_{1/2} = (45_H)_j^{iat} - \frac{1}{3} \delta_j^{iat} \langle 45_H \rangle_m = \left( \begin{array}{c} S_i^+ \\ S_R^0 + i S_I^0 \end{array} \right) = S^A T^A, \]  

where \( i, j = 1, 2, 3, \ A = 1, \ldots, 8 \), and \( T^A \) are the SU(3) generators. It is clear that the octet scalars are defined such that they have vanishing VEVs. Moreover, one can define the other components of the 45_H as follows:

\[ \langle 6^*, 1 \rangle_k^{ij} \equiv (45_H)_k^{ij} = \Phi_k^{ij}, \quad \langle 3^*, 2 \rangle_c^{ij} \equiv (45_H)_c^{ij} = \Phi_c^{ij}, \quad \langle 3^*, 1 \rangle_k^{ab} \equiv (45_H)_k^{ab} = \Phi_k^{ab}, \quad \langle (3, 1) \rangle_c^{ij} \equiv (45_H)_c^{ij} = \Phi_c^{ij}. \]

These scalars acquire masses from the potential \( V(5_H, 24_H, 45_H) \) after breaking the SU(5) into SU(3)_C × SU(2)_L × U(1)_Y through the VEV of 24_H and the electroweak symmetry through the VEV of the doublets of 5_H and 45_H.
II. GAUGE COUPLING UNIFICATION

In quantum field theory, the gauge couplings are functions of the energy at which they are measured and their Renormalization Group Equations (RGE), at one loop, are given by

$$\frac{d\alpha_i(t)}{dt} = \frac{b_i}{2\pi} \alpha_i^2(t), \quad i = 1, 2, 3$$

(5)

where \( t = \ln \mu \) with \( \mu \) is the running scale from the Electroweak scale, \( \mu_Z \), up to the scale of grand unification theory, \( \mu_{GUT} \). The couplings \( \alpha_i \) are defined as \( \alpha_i(t) = \frac{\alpha_i^0}{\alpha_i(t)} \) and the \( U(1)_Y \) coupling \( \alpha_1 \) is normalized by the factor 5/3. One can solve these RGEs to obtain \( \alpha_i(\mu) \) at a scale \( \mu \) for a given \( \alpha_i(\mu_0) \),

$$\alpha_i(\mu)^{-1} = \alpha_i(\mu_0)^{-1} - \frac{b_i}{2\pi} \ln \left( \frac{\mu}{\mu_0} \right).$$

(6)

The one loop coefficients \( b_i \) are defined as

$$b = -\frac{11}{3} T_G(R) + \frac{2}{3} T_F(R) + \frac{1}{3} T_B(R),$$

(7)

where \( G, F, B \) stand for gauge bosons, fermions and bosons respectively. The Dynkin index \( T(R) \) is defined for a representation \( R \) as \( T(R) = T(R) T_a, \), where \( T_a \) are the generators of this representation. In \( SU(N) \), \( T(\text{vector}) = \frac{1}{2} \) and \( T(\text{adjoint}) = N \). Therefore, the \( b_i \) in the SM are given by

$$b_1^{SM} = \frac{41}{10}, \quad b_2^{SM} = -\frac{19}{6}, \quad b_3^{SM} = -7.$$  

(8)

In this case, one can easily show that using the experimental values of the coupling constants at electroweak scale: \( \alpha_{\text{ew}}^{-1} = 127.9, \sin^2 \theta_W = 0.2329 \) and \( \alpha_3 = 0.110 \), the three coupling constants will never meet at single point.

If one assumes that the mass of \( SU(5) \)-Higgs doublet is of order \( \mu_Z \sim \mathcal{O}(100) \) GeV, while the octet scalar masses are of order \( \mu_S \sim \mathcal{O}(1) \) TeV and the triplet scalar masses are of order \( \mu_T \gg \mu_S \), then the following three regions for gauge coupling evolution are obtained: From \( \mu_Z \) to \( \mu_S \), from \( \mu_S \) to \( \mu_T \) and from \( \mu_T \) to \( \mu_{GUT} \). In these regions the beta function coefficients are given by

$$b_i^{EW} = \left( \frac{43}{10}, -\frac{17}{6}, -7 \right), \quad b_i^S = \left( \frac{51}{10}, -\frac{3}{2}, -5 \right), \quad b_i^T = \left( \frac{193}{30}, \frac{1}{2}, -\frac{25}{6} \right).$$

(9)

Therefore the RGE solution in Eq. (6) takes the form

$$\alpha_i^{-1}(\mu_Z) = \alpha_i^{-1}(\mu_{GUT}) - \frac{b_i^{EW}}{2\pi} \ln \left( \frac{\mu_S}{\mu_Z} \right) + \frac{b_i^S}{2\pi} \ln \left( \frac{\mu_T}{\mu_S} \right) + \frac{b_i^T}{2\pi} \ln \left( \frac{\mu_{GUT}}{\mu_T} \right).$$

(10)

One can eliminate \( \alpha_i^{-1} \) from the above equations and consider the following two equations in terms of the unknown scales: \( \mu_S, \mu_T \) and \( \mu_{GUT} \):

$$\alpha_1^{-1}(\mu_Z) - \alpha_2^{-1}(\mu_Z) = \frac{1}{2\pi} \left( b_1^{EW} - b_1^S \right) \ln \left( \frac{\mu_S}{\mu_Z} \right) + \frac{1}{2\pi} \left( b_2^S - b_2^T \right) \ln \left( \frac{\mu_T}{\mu_S} \right) + \frac{1}{2\pi} \left( b_1^T - b_1^{EW} \right) \ln \left( \frac{\mu_{GUT}}{\mu_T} \right).$$

(11)

$$\alpha_2^{-1}(\mu_Z) - \alpha_3^{-1}(\mu_Z) = \frac{1}{2\pi} \left( b_2^{EW} - b_2^S \right) \ln \left( \frac{\mu_S}{\mu_Z} \right) + \frac{1}{2\pi} \left( b_3^S - b_3^T \right) \ln \left( \frac{\mu_T}{\mu_S} \right) + \frac{1}{2\pi} \left( b_2^T - b_2^{EW} \right) \ln \left( \frac{\mu_{GUT}}{\mu_T} \right).$$

(12)
FIG. 1: The running of the gauge couplings in low energy effective $SU(5)$ model with TeV scale octet scalars and intermediate scale ($\sim 10^6$ GeV) triplet scalars.

From these equations one finds that three gauge couplings are unified at $\mu_{GUT} = 2 \times 10^{16}$ GeV if $\mu_s \sim 4 \times 10^6$ GeV and $\mu_T \sim 4 \times 10^{11}$ GeV. Nevertheless, as shown in Fig. [11] gauge coupling unification with a lower scale of octet scalars (of order TeV) is possible at $\mu_{GUT} \simeq 4 \times 10^{15}$ Gev. In this case, the scale of triplet scalars becomes of order $10^9$ GeV. It is also worth noting that the recent CMS and ATLAS experimental results, based on searches for dijet pair signatures at $\sqrt{s} = 7$ TeV, imposed stringent constraint on the octet scalar masses: $m_S \gtrsim 2$ TeV [6].

III. OCTET SCALAR AT THE LHC

TeV scale octet scalars can be produced copiously at the LHC and lead to very interesting signatures. In the rest of the paper, we discuss the phenomenology of these scalars, in particular their production and decay in the LHC. The interaction lagrangian of colored octet scalars with SM fermions is given by [4]

$$\mathcal{L}_{int} = 2(Y_2)_{ij} \bar{d}_{Ri} Q_{Lj} S^i + 4\epsilon_{\alpha\beta}(Y_4^T - Y_4)_{ij} \bar{u}_{Ri} Q_{Lj} S^j + h.c.,$$

(13)

where the Yukawa couplings $Y_2$ and $Y_4$, along with the Yukawa couplings $Y_1$ and $Y_3$ of 5-plet Higgs, define the fermion masses as follows [4, 8]

$$M_E = Y_1^T v_5 - 6 Y_2^T v_{45}^*,$$

(14)

$$M_D = Y_1 v_5^* + 2 Y_2 v_{45}^*,$$

(15)

$$M_U = 4(Y_3 + Y_4^T) v_5 - 8(Y_4^T - Y_4) v_{45}. $$

(16)

Therefore, the Yukawa coupling $Y_2$ can be written in terms of charged lepton and down quark masses

$$Y_2 = \frac{M_D - M_E}{8v_{45}}.$$

However, the situation for $Y_4$ is more involved. It is clear that one cannot relate $Y_4$ directly to the up-quark masses. Moreover, in the basis where $M_U$ is diagonal, i.e., the quark mixing is emerged from down quark sector only and the rotational matrices are given by $V_L^u = V_R^u = I$ and $V_L^d = V_{CKM}$, then $Y_3$ and $Y_4$ matrices should be also diagonal, unless there is a significant fine tuning between them. In this case one finds that the couplings of up-quarks with octet scalars vanish identically. This conclusion can be also obtained if $Y_4$ is a symmetric matrix. Thus, the interaction lagrangian in the physical mass basis is given by

$$\mathcal{L}^S = \bar{d} \left[ P_L \left( \frac{m_D V_{CKM}^\dagger - m_E}{4v_{45}} \right) \right] S^- u + \bar{u} \left[ P_R \left( \frac{V_{CKM} m_D - m_E}{4v_{45}} \right) \right] S^+ d - \frac{S_0}{4v_{45}} \left[ \bar{d} \left( P_L m_E V_{CKM} + P_R V_{CKM}^\dagger m_E \right) \right] d - m_D \bar{d} d - \frac{i S_0}{4v_{45}} \left[ \bar{d} \left( P_L m_E V_{CKM} - P_R V_{CKM}^\dagger m_E \right) \right] d - m_D \bar{d} \gamma_5 d \right].$$

(17)
We now consider the searches for the octet scalars at the LHC. The single neutral octet scalars can be produced at tree level from quark-anti-quark annihilation: $qq \rightarrow S^0_{R,I}$, with $q = u, d$ or $S$ and at one loop level through the gluon fusion process: $gg \rightarrow S^0_{R,I}$, with $b$-quark or $S^0/\bar{S}^0_{R,I}$ exchanges. While the charged octet scalars are singly produced at tree level only from $qq^\prime \rightarrow S^\pm$. The one-loop triangle diagram of gluon-gluon fusion usually gives the dominant contribution for neutral single production, however since our octet scalar is quite heavy, this contribution is rather suppressed and becomes even smaller than the tree level production for $m_{S^0} \gtrsim 2$ TeV. The pair productions of neutral and charged octet scalars occur at tree level as shown in Fig. 2. The gluon interaction with the octet scalars is one of their most relevant interactions with the SM particles. It is given by

$$L_{\text{gluon}}^S = ig_s \text{Tr} \left[ S^A - G^B \partial_\mu S^{D \mu +} + S^A_{R} G^B \partial_\mu S^{D \mu 0} + S^A_{I} G^B \partial_\mu S^{D I 0} \right] F^{ABD} + g_s^2 \text{Tr} \left[ S^A - G^B G^C S^{D \mu +} + S^A_{R} G^B G^C S^{D \mu 0} + S^A_{I} G^B G^C S^{D I 0} \right] F^{ABCD} + h.c.,$$

where

$$F^{ABD} = \text{tr}[T^A T^B T^D] = 1/4 (d^{ABD} + i f^{ABD}),$$

$$F^{ABDE} = \text{tr}[T^A T^B T^D t^E] = 2/9 d^{AB} g^{DE} + 1/8 \{ i d^{ABC} d^{DEC} + i d^{ABC} f^{DEC} + i f^{ABC} d^{DEC} - f^{ABC} f^{DEC} \},$$

with $d^{ABC}$ and $f^{ABC}$ are the $SU(3)$ symmetric and antisymmetric structure constants, respectively. Therefore, the cross section of the partonic pair production of octet scalars is given by

$$\frac{d\sigma}{dt}(gg \rightarrow SS) = \frac{\pi \alpha_s^2}{s^2} \left( \frac{27}{32} + \frac{9(u-t)}{32s^2} \right) \left( 1 + \frac{2m_s^2}{u-m_s^2} + \frac{2m_s^2}{t-m_s^2} \right) + \frac{2m_s^4}{(u-m_s^2)^2} + \frac{2m_s^4}{(t-m_s^2)^2} + \frac{4m_s^4}{(u-m_s^2)(t-m_s^2)},$$

which implies that

$$\sigma(gg \rightarrow SS) = \frac{\pi \alpha_s^2}{s} \left( \frac{15k}{16} + \frac{51km_s^2}{8s} + \frac{9m_s^2(s-m_s^2)}{2s^2} \ln \left( \frac{1-k}{1+k} \right) \right),$$

where $k = (1 - \frac{4m_s^2}{s})^{1/2}$. Note that the initial kinematics threshold of this process is given by $s = 4m_s^2$. We have used Feynrules [10] to generate the model files and MadEvent5 [11] to calculate the numerical values of the cross sections of neutral and charged octet scalar. We assume $\sqrt{s} = 14$ TeV. In Fig. 3 we show different cross sections of single and pair productions of neutral and charged octet scalars in terms of universal octet scalar mass. As can be seen from this figure that for octet scalar mass of order $O(2)$ TeV, the single production cross section is about one order of magnitude larger than the pair production cross sections. This result is consistent with the findings in other extensions of the SM with heavy octet scalars [12,13]. Therefore, one may expect that the single production of octet scalar at LHC would have a higher possibility.

In our model, the neutral scalars decay dominantly into $b\bar{b}$, while the charged octet scalars decay into $b\bar{t}$ or $t\bar{b}$. If $m_{S^0} > m_{S^\pm}$, then one may consider the decay channel: $S^0 \rightarrow S^\pm W^\mp$. Recall that the octet scalar interactions with
FIG. 3: The hadronic production cross sections of $q\bar{q} \to S_{R,I}^0$, $q\bar{q} \to S^\pm$, gluonic fusion ($gg \to S_{R,I}^0$), and the pair production $pp \to S^0/\pm S^0/\mp$ as a function of $m_S$ at CME $\sqrt{s} = 14$ TeV.

electroweak gauge bosons are obtained from the kinetic term of $45_H$ \cite{5}: $\text{Tr}[(D^\mu 45_H)^\dagger (D_\mu 45_H)]$

$$\begin{align*}
\mathcal{L} &= -\frac{im_H}{v} S^0 W^\mu \partial_\mu S^0 + \frac{ig}{\sqrt{2}} S^0 W^\mu \partial_\mu S^0 - \frac{ig}{\sqrt{2}} S^0 W^\mu \partial_\mu S^0 + \frac{im_H}{v} S^0 Z^\mu \partial_\mu S^0 \\
&\quad - \frac{1}{4} \left[ \frac{m_{H_1}^2}{v^2} S^0 A^\mu A_\mu S^0 - 2 \sqrt{2} g m_{H_1} S^0 A^\mu W^\mu_\mu S^0 - 2 g^2 S^0 W^\mu W^- W^+ S^0 + 2 \sqrt{2} g m_{H_1} S^0 W^\mu Z^\mu S^0 \\
&\quad - 2 \sqrt{2} g m_{H_1} S^0 W^\mu Z^- S^0 - 2 g^2 S^0 W^\mu W^- S^0 + 2 \sqrt{2} g m_{H_1} S^0 Z^\mu W^- S^0 - \frac{m_{H_1}^2}{v^2} S^0 Z^\mu Z^\mu S^0 \right].
\end{align*}$$

(23)

It is important to mention that phenomenological analysis of octet scalars has been recently studied in literature \cite{13}. However, most of these analysis was based on the octet scalars given in Manohar-Wise model \cite{14}, where the octet scalars have free coupling with SM up and down quarks as well as the SM gauge bosons. Here our octet scalars interactions are limited with the effective non-minimal $SU(5)$ model as explained in detail in Ref. [4]. Moreover, we also impose the stringent experimental constraint that neutral/charged octet scalars mass $\sim 2$ TeV, which is consistent with our conclusions of the previous section for having a gauge coupling unification at scale larger than $10^{15}$ GeV.

As advocated above, the process $pp \to S^0 \to b\bar{b}$ has a large cross section, however it turns out that the SM background of this process exceeds any possible signature even if one applies a large $P_T$ cuts on the outgoing jets. In Fig. 4 (left panel) we plot the number of reconstructed events per bin of the invariant mass of $b\bar{b}$ of this process for $S$ signal and SM background at $P_T$ cut $> 800$ GeV with $m_{S^0} = 2$ TeV and $\sqrt{s} = 14$ TeV. This figure shows that it is not possible to extract a good significance for the octet signal in this channel. In addition, we also consider the process: $pp \to S^0 \to S^+ W^- \to t\bar{b} W^- \to 2l + 2b +$ missing energy. This process has a small cross section in SM, hence applying small cuts on the final states suppresses the SM background. Nevertheless, these cuts also suppress our signal. Therefore, as can be seen from Fig. 4 (right panel), one cannot get a good significance for $S$ signal for this channel as well.

We now turn to possible signals from the pair production processes. Although, these processes have smaller cross sections, one finds that by imposing suitable cuts a significant signal can be obtained. In Fig. 5 (left panel) we plot the number of events of $pp \to S^0 S^0 \to b\bar{b} b\bar{b}$ per bin at the parton level at 14 TeV center of mass energy versus the invariant mass of the $b\bar{b}$ pair. While in Fig. 5 (right panel), we present the partonic center of mass energy of the process $pp \to S^+ S^- \to b\bar{t} t\bar{b}, (t \to W^- b), (t \to W^+ b)$, with the possibility that $W$ boson decays leptonically into $l\nu$ (30 %) or hadronically into $q \bar{q}$ (70 %). So the following three modes are available: (a) $(W^- \to l^- \bar{\nu}), (W^+ \to q\bar{q})$ with final states, $l + 2(b\bar{b}) + 2$ jets + missing energy, (b) $(W^- \to l^- \bar{\nu}), (W^+ \to l^+ \nu)$ with final states, $2l + 2(b\bar{b}) +$ missing energy and (c) $(W^- \to q\bar{q}), (W^+ \to q\bar{q})$. 
FIG. 4: The number of reconstructed events per bin of the invariant mass of jet-pair for $S$ signal and the SM background for $P_T > 800$ GeV, $m_{S^0} = 2$ TeV and $\sqrt{s} = 14$ TeV for (Left panel) $pp \rightarrow S^0 \rightarrow b\bar{b}$ and $pp \rightarrow S^0 \rightarrow W^- S^+ \rightarrow l^- l^+ b\bar{b} \nu \bar{\nu}$ (Right panel). Here the bin size is 30 GeV.

FIG. 5: (Left panel) The number of events per bin of the invariant mass of $b\bar{b}$ pair in $pp \rightarrow S^0 S^0 \rightarrow b\bar{b} b\bar{b}$ process at $\sqrt{s} = 14$ TeV. (Right panel) The partonic center of mass energy for $pp \rightarrow S^+ S^- \rightarrow b\bar{t} t \bar{b}$ (Left panel) and $pp \rightarrow S^0 \rightarrow S^+ W^-$ (Right panel) according to the decay chains (a),(b) and (c) for each channel as declared in the text. The size of the bin is 30 GeV.

From these figures, it is clear that the $b\bar{b}b\bar{b}$ final state channel gives the best signal for probing the neutral octet scalar at the LHC. To investigate the 4b-tagged jets, a strong cut should be imposed to suppress the SM background. We apply cut on the transverse momentum $P_T$ of the produced 4-jets to be $P_T > 800$ GeV. This is an acceptable cut, since we expect high energetic jets produced by the octet scalars. This cut enhances the significance $S = N_{signal}/\sqrt{N_{background}}$, where $\sigma_{SM}$ becomes $3.967 \times 10^{-4}$ $fb$ and $\sigma_S = 1.827 \times 10^{-3}$ $fb$. Decreasing the $P_T$ cut to $> 700$ GeV implies a reduction in the significance of our signal, since $\sigma_{SM}$ becomes $1.66 \times 10^{-3}$ $fb$ and $\sigma_S = 3.954 \times 10^{-3}$ $fb$.

In Fig. 6 we display the number of events per bin of the highest $P_T$ GeV jets. Here we use pythia and PGS detector simulator for the octet scalars signal and SM background at integrated luminosity $200 fb^{-1}$ and $\sqrt{s} = 14$ TeV [15]. We also used MadAnalysis5 [16] to plot our results.

IV. CONCLUSIONS

In this paper we have studied some phenomenological aspects of the non-minimal $SU(5)$ model. We emphasized that the low energy Higgs sector of this class of models consists of two SM-like Higgs doublets, two neutral and one charged octet scalars, in addition to two triplet colored scalars. We have shown that the gauge coupling unification could be realized at scale larger than $10^{15}$ GeV, if the octet scalar masses are of order TeV and the triplet scalar...
masses are of order an intermediate scale $\sim 10^6$ GeV. We have also analyzed the possible LHC signatures of these TeV octet scalars. We showed that they can be singly produced at tree level from $q\bar{q}$ annihilation (neutral octet scalar) or $q\bar{q'}$ annihilation (charged octet scalar) and in pairs through the process $gg \to S^0(\pm)S^0(\mp)$. We found that the process $pp \to S^0 \to b\bar{b}$ has the largest cross section, however the SM background is quite large and exceeds any possible signature even if one applies a large $p_T$ cut on the outgoing jets. We argued that applying a high $p_T$ cut of order $\gtrsim 800$ GeV on the outgoing jets, the best channel which can provide a good signal for our heavy octet scalars is the $4b$-tagged jets final states. Although the other channels of multi-jets with associated leptons and missing energy have significant cross sections, they can not lead to good significance due to the high $p_T$ cut that one should implement in order to suppress the SM background.

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