Diphoton signature of neutral pseudo–Goldstone boson in the $E_6$CHM at the CERN LHC

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Abstract

The breakdown of the $SU(6)$ global symmetry to its $SU(5)$ subgroup, that contains the standard model (SM) gauge group, in the $E_6$ inspired composite Higgs model ($E_6$CHM) results in a set of pseudo–Nambu–Goldstone bosons (pNGBs). This set, in particular, involves the SM–like Higgs doublet and a SM singlet boson. In the limit when CP is conserved the SM singlet scalar $A$ is a CP–odd state that does not mix with the SM–like Higgs. The interactions of $A$ with exotic matter beyond the SM, which ensures anomaly cancellation and approximate gauge coupling unification, can induce couplings of this pseudoscalar to the SM gauge bosons. As a consequence, we argue that the SM singlet pNGB state can be identified with the 750 GeV diphoton resonance recently reported by the LHC experiments. Such an interpretation requires that either all or some of the exotic states have masses below 1 TeV. We specify the interactions of the SM singlet pNGB state with the exotic vector–like fermions, top quark and SM gauge bosons as well as examining the dependence of the LHC diphoton production cross section, the total width of the pseudoscalar $A$ and its branching ratios on the parameters of the $E_6$CHM.

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1 Introduction

After the discovery of a scalar particle, which is consistent with the Standard Model (SM) Higgs boson, at the LHC Run-1, the focus of experimental and theoretical studies is shifting again towards the investigations of possible new physics phenomena beyond the SM (BSM). So far ATLAS and CMS have collected only a small amount of data at the 13 TeV LHC but this already permitted us to set competitive constraints on the parameters of some BSM models as compared with those obtained at 8 TeV. First results of new physics searches at the LHC Run-2 recently announced by both the ATLAS and CMS collaborations indicate that there is an excess of events in the diphoton channel at an invariant mass, $M$, around 750 GeV. Although the situation is rather uncertain, because of the relatively low significance of the excess that might be due to a statistical fluctuation, the observed signals are compatible both with Run-1 data and between ATLAS and CMS measurements which makes this a compelling hint of a new resonance.

Many papers have already appeared, studying numerous interpretations of the observed excess in terms of new physics scenarios [1]–[4]. The simplest interpretation suggests that the new resonance is a SM singlet scalar (pseudoscalar) state coupled to vector–like fermions. At the moment the data collected at the LHC Run-2 imply that the cross section $\sigma(pp \rightarrow \gamma\gamma)$ associated with the production and sequential diphoton decays of this state should be around $5 - 10$ fb. The analysis of ATLAS measurements favours a large width $\Gamma$ of about 45 GeV, i.e. $\Gamma/M \approx 0.06$, while the best fit to the CMS data has a much narrower width but also a slightly lower significance than the ATLAS case. It is also important that the observed diphoton events are not accompanied by missing energy, leptons or jets. Moreover, no resonance with mass 750 GeV has been detected in any other channels including $pp \rightarrow t\bar{t}, WW, ZZ, b\bar{b}, \tau\bar{\tau}$ and $jj$. This might be an indication that the observed signal is nothing but a statistical fluctuation. On the other hand, if this state is confirmed these observations should set stringent constraints on the new physics scenarios that give rise to such a signature.

In this article we consider the diphoton signature in the $E_6$ inspired composite Higgs model ($E_6$CHM) [5] which leads to possibly the simplest interpretation of the observed excess of diphoton events mentioned above. As in any other composite Higgs model [6] (for a recent review, see [7]), the $E_6$CHM involves weakly–coupled elementary and strongly interacting sectors. The $E_6$SSM implies that at some high energy scale, $M_X$, the $E_6 \times G_0$ gauge group is broken down to the $SU(3)_C \times SU(2)_W \times U(1)_Y \times G$ subgroup where $G_0$ and $G$ are associated with the strongly interacting sector and $SU(3)_C \times SU(2)_W \times U(1)_Y$ is the SM gauge group. Fields belonging to the strongly coupled sector can be charged under both the $E_6$ and $G_0$ ($G$) gauge symmetries. The weakly–coupled sector contains elemen-
tary states that participate in the \( E_6 \) interactions only. In the \( E_6 \) CHM the appropriate suppression of the proton decay rate and the Majorana masses of the left–handed neutrino can be achieved if global \( U(1)_B \) and \( U(1)_L \) symmetries, which ensure the conservation of baryon and lepton numbers, are imposed. Because of the almost exact conservation of the \( U(1)_B \) and \( U(1)_L \) charges, the elementary states with different baryon and/or lepton numbers should stem from different 27–plets, while all other components of these 27–plets have to gain masses of the order of \( M_X \).

The appropriate splitting of the \( E_6 \) fundamental representations may occur within a six–dimensional orbifold GUT model with \( N = 1 \) supersymmetry (SUSY) \([5]\) in which SUSY is broken slightly below the GUT scale \( M_X \). In this model the SM fermions with different baryon and/or lepton numbers are components of different bulk 27–plets. All fields from the strongly interacting sector are localised on the brane, where \( E_6 \) symmetry is broken to the \( SU(6) \times SU(2)_N \) subgroup. It is assumed that the \( E_6 \) gauge symmetry gets broken to the SM gauge group, \( SU(2)_N \) symmetry is entirely broken, while \( SU(6) \), which contains the \( SU(3)_C \times SU(2)_W \times U(1)_Y \) subgroup, remains an approximate global symmetry of the composite sector at low energies.

Below the scale \( f \) (\( f \gtrsim 10 \) TeV) the global \( SU(6) \) symmetry in the \( E_6 \) CHM is expected to be broken down to \( SU(5) \), so that the SM gauge group is preserved, leading to eleven pseudo–Nambu–Goldstone bosons (pNGBs). One of these pNGB states is a SM singlet field, four others form the SM–like Higgs doublet, \( H \), and six pNGB states are associated with an \( SU(3)_C \) triplet, \( T \). None of these pNGB states carry any baryon and/or lepton numbers. The effective potential, which describes the interactions of the pNGB states, is induced by radiative corrections. These corrections are caused by the interactions between the elementary states and their composite partners that violate \( SU(6) \) symmetry. The induced effective potential has a structure that results in the spontaneous breakdown of the electroweak (EW) symmetry, while \( SU(3)_C \) remains intact. In this case the effective quartic Higgs coupling tends to be small and, therefore, may lead to a 125 GeV Higgs.

Although the SM gauge couplings \( \alpha_i(M_X) \) in the orbifold GUT models may not be identical an approximate unification of these couplings is expected to take place near the scale \( M_X \). Such a unification can be attained if the right–handed top quark \( t^c \) is entirely composite and the weakly–coupled sector, together with the SM fields (but without \( t^c \)), involves a set of exotic states. The presence of additional exotic states also guarantees the cancellation of gauge anomalies in the elementary sector. In the \( E_6 \) CHM the exotic states mentioned above get combined with the composite counterparts, which fill complete \( SU(5) \) representations, resulting in a set of vector–like fermion states and composite \( t^c \).

\(^1\)Different phenomenological aspects of the \( E_6 \) inspired models with low-scale supersymmetry breaking were recently studied in \([4], [8]\).
In this article we argue that the interactions of these vector–like fermions with the SM singlet pNGB state can lead to the recently observed diphoton excess if the corresponding pNGB state has a mass around 750 GeV.

The layout of this article is as follows. In the next Section we discuss the matter content of the weakly–coupled elementary sector, the generation of masses of the SM particles, unification of the SM gauge couplings and non–linear realization of the Higgs mechanism in the E₆CHM. The interactions of the exotic elementary states with their composite partners is also considered. In Section 3 we specify the interactions of the SM singlet pNGB state with the exotic vector–like fermions, top quark and SM gauge bosons. In section 4 we examine the the branching ratios and the LHC diphoton production cross section associated with this pNGB state. Section 5 concludes the paper.

2 The E₆CHM

2.1 Gauge coupling unification and SU(6) symmetry breaking

As already mentioned, approximate gauge coupling unification can be achieved in the E₆CHM. This scenario is realised if the right–handed top quark, \( t^c \), is composite and the sector of weakly–coupled elementary states contains the following set of matter multiplets (see also [9]):

\[
(q_i, d_i^c, \ell_i, e_i^c) + u_\alpha^c + \bar{q} + \bar{d}^c + \bar{\ell} + \bar{e}^c + \eta ,
\]

where \( \alpha = 1, 2 \) runs over the first two generations and \( i = 1, 2, 3 \) runs over all three. In Eq. (1) \( q_i \) and \( \ell_i \) represent left-handed quark and lepton doublets, \( u_\alpha^c, d_i^c \) and \( e_i^c \) correspond to the right-handed up- and down-type quarks and charged leptons, while \( \bar{q}, \bar{d}^c, \bar{\ell} \) and \( \bar{e}^c \) are associated with the exotic states that have exactly opposite \( SU(3)_C \times SU(2)_W \times U(1)_Y \) quantum numbers to the left-handed quark doublets, right-handed down-type quarks, left-handed lepton doublets and right-handed charged leptons, respectively. An extra exotic state, \( \eta \), with spin 1/2 does not participate in the \( SU(3)_C \times SU(2)_W \times U(1)_Y \) gauge interactions. It is included to ensure the phenomenological viability of this model. The set of fermion states (1) is chosen so that the weakly–coupled elementary sector involves all SM fermions except \( t^c \) and anomaly cancellation takes place.

At low energies (\( E \lesssim 4\pi f \)) the strongly interacting sector in the composite Higgs models [6] leads to a set of bound states that includes the pNGB states as well as massive fields with quantum numbers of all SM particles. These are the so–called composite partners of the SM fermions and bosons. The contributions of these new states to electroweak precision observables were analysed in Refs. [10]–[18]. In the E₆CHM the composite bound states fill complete \( SU(6) \) representations, that involve the composite partners of
quarks, leptons and gauge bosons. The elementary states couple to the appropriate operators of the strongly interacting sector that give rise to the mixing between these states and their composite partners. The partial compositeness of the SM bosons and fermions makes possible the generation of their masses caused by non-zero vacuum expectation value (VEV) of the pNGB state associated with the Higgs boson. The couplings of the SM states to the composite Higgs are determined by the fractions of the compositeness of these states. In most cases, especially for the first and second generations of fermions, the corresponding fractions are sufficiently small. Therefore the non-diagonal flavor transitions and the modifications of the $W$ and $Z$ couplings associated with the light SM fermions are suppressed. Within the composite Higgs models the constraints that come from flavour-changing processes in the quark and lepton sectors were examined in Refs. [16]–[21] and [21]–[24], respectively. In particular, it was shown that in the case when the matrices of effective Yukawa couplings in the strong sector are structureless, i.e anarhich matrices, adequate suppression of flavor changing neutral currents (FCNCs) can be obtained only if $f$ is larger than 10 TeV [16]–[17], [19]–[20], [23]². The implications of the composite Higgs models were also studied for Higgs physics [13]–[14], [26]–[29], gauge coupling unification [30]–[31], dark matter [11], [27], [31]–[32] and collider phenomenology [12]–[13], [15], [19], [22], [29], [33]. Non–minimal composite Higgs models were considered in Refs. [11], [26]–[27], [31]–[32], [34].

The presence of additional exotic states in the E6 CHM facilitates the convergence of the SM gauge couplings at very high energies. Indeed, all states in the strongly coupled sector come in complete SU(6) and SU(5) representations which contribute equally to the one–loop beta functions of the SU(3)$_C$, SU(2)$_W$ and U(1)$_Y$ interactions. Thus composite sector fields should not spoil the convergence of the SM gauge couplings in the leading approximation, which is determined by the matter content of the elementary sector. Using the one–loop renormalisation group equations (RGEs) one can find that for $\alpha(M_Z) = 1/127.9$, $\sin^2 \theta_W = 0.231$ and the elementary particle spectrum given by Eq. (1), the exact unification of the SM gauge couplings takes place if $\alpha_3(M_Z) \approx 0.109$. The corresponding gauge coupling unification scale is somewhat close to $M_X \sim 10^{15} - 10^{16}$ GeV.

Although $\alpha_3(M_Z) \approx 0.109$ is considerably smaller than the central measured low energy value of this coupling, this estimation demonstrates that an approximate gauge coupling unification may be achieved in the E6 CHM. Moreover, it was also argued that the inclusion of higher order effects may improve the unification of the SM gauge couplings [9], [31].

In the E6 CHM the global SU(6) symmetry of the strongly interacting sector is broken

²This bound on the scale $f$ can be significantly alleviated in the composite Higgs models with flavour symmetries [15]–[16], [19], [21]–[22], [23].
down to $SU(5)$ below the scale $f$. We denote the generators of the $SU(5)$ subgroup of $SU(6)$ by $T^a$, whereas the eleven generators from the coset $SU(6)/SU(5)$ associated with the pNGB states are denoted by $T^a_{\tilde{\alpha}}$. Here the $SU(6)$ generators are normalised so that $\text{Tr} T^a T^b = \frac{1}{2} \delta_{ab}$. It is convenient to use the non–linear representation of the pNGB states in terms of a 6–component unit vector $\Omega$, which is a fundamental representation of $SU(6)$, that is

$$\Omega^T = \frac{1}{\sqrt{f}} \begin{pmatrix} C \phi_1 & C \phi_2 & C \phi_3 & C \phi_4 & C \phi_5 & \cos \frac{\phi}{\sqrt{2f}} + \sqrt{\frac{3}{10}} C \phi_0 \end{pmatrix},$$

(2)

where

$$\Omega_0^T = (0 \ 0 \ 0 \ 0 \ 0 \ 1), \quad \Sigma = e^{i \Pi/f}, \quad \Pi = \Pi^a T^a.$$

The fields $\phi_1, \phi_2, \phi_3, \phi_4$ and $\phi_5$ are complex while $\phi_0$ is a real field. Vector $\Omega$ transforms as $5 + 1$ under the transformation of the unbroken $SU(5)$ subgroup where $5 = \tilde{H} \sim (\phi_1 \phi_2 \phi_3 \phi_4 \phi_5)$ and $1 = \phi_0$ is a SM singlet field. The first two components of $\tilde{H}$ transform as an $SU(2)_W$ doublet, $H \sim (\phi_1 \phi_2)$, and correspond to the SM–like Higgs doublet. Three other components of $\tilde{H}$, $T \sim (\phi_3 \phi_4 \phi_5)$, are associated with the $SU(3)_C$ triplet. In the $E_6$CHM the components of vector $\Omega$ do not carry any baryon and/or lepton numbers. In the leading approximation the Lagrangian, that describes their interactions, is given by

$$L_{pNGB} = \frac{f^2}{2} \left| D_\mu \Omega \right|^2.$$

(3)

The pNGB effective potential $V_{eff}(\tilde{H}, T, \phi_0)$ can be obtained by integrating out the exotic fermions and heavy resonances of the composite sector. It is induced by the interactions of the elementary fermions and gauge bosons with their composite partners, that break $SU(6)$ global symmetry and vanish in the exact $SU(6)$ symmetry limit. The analysis of the pNGB effective potentials in the composite Higgs models, which are similar to the $E_6$CHM, revealed that there is a considerable part of the parameter space where the $SU(2)_W \times U(1)_Y$ gauge symmetry is broken to $U(1)_{em}$, associated with electromagnetism, while $SU(3)_C$ remains intact. Since in the $E_6$CHM the scale $f \gtrsim 10$ TeV, a significant tuning, $\sim 0.01\%$, is needed to get the quadratic term $m_{\tilde{H}}^2 |H|^2$ in $V_{eff}(\tilde{H}, T, \phi_0)$ with the appropriate value of the parameter $m_{\tilde{H}}^2$ that leads to a 125 GeV Higgs state. It was shown that such tuning can be accomplished by cancelling two different contributions to $m_{\tilde{H}}^2$ associated with the gauge fields and exotic fermions which appear with different signs in such composite Higgs models. In these models the $SU(3)_C$ triplet scalar $T$ tends to be substantially heavier than the SM–like Higgs boson.
2.2 Exotic fermion states

The mechanism responsible for the generation of the masses of the SM particles discussed above requires $t^c$ to either have a sizeable fraction of compositeness or to be almost completely composite, because the top quark is so heavy and precision measurements, like $Z \rightarrow b\bar{b}$, indicate that the left–handed $b$–quark, and hence left–handed $t$, should have a reasonably small admixture of composite states. The phenomenological viability of the scenario under consideration implies that the dynamics of the strongly coupled sector below the $SU(6)$ breaking scale $f$ results in the composite $10 + \overline{5} + 1$ multiplets of $SU(5)$. The components of these $SU(5)$ multiplets decompose under $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_B \times U(1)_L$ as follows:

\[
\begin{align*}
10 & \rightarrow Q = (U, D) = \left(3, 2, \frac{1}{6}, -\frac{1}{3}, 0\right), \\
t^c & = \left(3^*, 1, -\frac{2}{3}, \frac{1}{3}, 0\right), \\
E^c & = \left(1, 1, 1, -\frac{1}{3}, 0\right); \\
\overline{5} & \rightarrow D^c = \left(3, 1, \frac{1}{3}, \pm\frac{1}{3}, 0\right), \\
L & = (N, E) = \left(1, 2, -\frac{1}{2}, \pm\frac{1}{3}, 0\right); \\
1 & \rightarrow \eta = \left(1, 1, 0, \mp\frac{1}{3}, 0\right). 
\end{align*}
\]

The first and second quantities in brackets are the $SU(3)_C$ and $SU(2)_W$ representations, while the third, fourth and fifth quantities are $U(1)_Y$, $U(1)_B$ and $U(1)_L$ charges respectively. The conservation of the $U(1)_B$ and $U(1)_L$ charges implies that all components of the 10–plet specified above carry the same baryon and lepton numbers as $t^c$. The multiplets $\overline{5}$ and 1 are allowed to have baryon charges $-1/3$ and $+1/3$ [5].

The composite $SU(5)$ multiplets [11] are expected to get combined with the elementary exotic states $\bar{q}$, $\bar{e}^c$, $\bar{d}^c$, $\bar{\ell}$ and $\eta$ resulting in a set of vector–like fermion states. The only exceptions are the components of the 10–plet associated with the composite $t^c$, which survive down to the EW scale. In the E$_6$CHM the exotic states $\bar{q}$, $\bar{e}^c$, $\bar{d}^c$, $\bar{\ell}$ and $\eta$ constitute the following incomplete multiplets of $SU(6)$ at low energies

\[
\begin{align*}
\bar{q} \in \mathbf{15'}, & \quad \bar{e}^c \in \mathbf{15'}, & \quad \bar{d}^c \in \mathbf{6'}, & \quad \bar{\ell} \in \mathbf{6'}, & \quad \eta \in \mathbf{1''}.
\end{align*}
\]

As discussed in Ref. [5], the composite partners of the up type quarks can stem from either the totally antisymmetric third–rank tensor $20$ or the antisymmetric second–rank tensor $15$. These $SU(6)$ representations decompose under unbroken $SU(5)$ as follows: $15 = 10 \oplus 5$ and $20 = 10 \oplus \overline{10}$. Thus the 10–plet associated with the composite $t^c$ may
belong to either a 15–plet or 20–plet of $SU(6)$. Further, we assume that this 10–plet is a linear superposition of the corresponding components of the 15–plet ($15'$) and 20–plet ($20'$). On the other hand, the composite $\mathbf{5}$ can originate from either the $\mathbf{15}$–plet ($\mathbf{15}$) or $\mathbf{6}$–plet ($\mathbf{6}$). Therefore it seems to be natural to expect that the composite $\mathbf{5}$, which survives below scale $f$, is a superposition of the appropriate components of $\mathbf{15}$ and $\mathbf{6}$.

Here we also assume that the composite SM singlet state, $\bar{\eta}$, that arises below the scale $f$, is a linear combination of 1 and the corresponding component of $\mathbf{6}$ of $SU(6)$. In the most general case, the interactions of the $SU(6)$ composite multiplets mentioned above with the incomplete $SU(6)$ representations that involve elementary exotic states (5) can be written as

$$L_{\text{exotic}} = \mu_Q \mathbf{15}' \mathbf{15}' + \sigma_Q f \mathbf{15}' \mathbf{15}' + \mu_E \mathbf{15} \mathbf{15}' + \sigma_E f \mathbf{15} \mathbf{15}' + \mu_D \mathbf{6} \mathbf{6}'$$

$$+ \sigma_D f \mathbf{15} \mathbf{6} \mathbf{6}' + \mu_L \mathbf{6} \mathbf{6}' + \sigma_L f \mathbf{15} \mathbf{6} \mathbf{6}' + \mu_1 \mathbf{1} \mathbf{1} + \sigma_1 f \mathbf{1} \mathbf{1} + \text{h.c.} \, .$$

### 3 Couplings of the SM singlet pNGB state

Using Eqs. (2) and (6) one can obtain the explicit analytical expressions for the masses of exotic fermion states and their couplings to the SM singlet field $\phi_0$. Throughout the rest of this article we assume that $\phi_0$ can be identified with the state which results in the excess of diphoton events at an invariant mass around 750 GeV recently reported by the LHC experiments. This SM singlet state could potentially mix with the Higgs boson which would give rise to large partial decay widths of the 750 GeV resonance to pairs of $Z$-bosons, $W$-bosons and $t\bar{t}$, leading to the strong suppression of the branching ratios of its decays to photons. At the same time ATLAS and CMS measurements indicate that these partial decay widths should be at least of the same order of magnitude as or even smaller than, the diphoton decay rate of this resonance.

Here, motivated by the nonobservation of CP violation beyond the SM, invariance under CP transformation is imposed and that permits us to forbid the mixing between $\phi_0$ and the SM–like Higgs state. Indeed, if all couplings in Eq. (6) are real, then $\phi_0$ manifests itself in the Yukawa interactions with fermions as a pseudoscalar field. As a consequence $\phi_0$ can not mix with the Higgs boson because of the almost exact CP–conservation.

In the leading approximation, the Lagrangian that describes the interactions between $\phi_0 = A$ and other states is given by

$$L_A = \frac{y_t}{\Lambda_4} A(i\bar{t}_L H t_R + \text{h.c.}) + A \left( i\kappa_D \bar{d}_L D^c + i\kappa_Q \bar{q} Q + i\lambda_L \bar{L} L + i\lambda_E \bar{E} E^c + i\lambda_\eta \bar{\eta} \eta + \text{h.c.} \right)$$

$$+ \frac{\alpha_Y}{16\pi\Lambda_1} A B_{\mu\nu} \tilde{B}^{\mu\nu} + \frac{\alpha_2}{16\pi\Lambda_2} A W^\alpha_{\mu\nu} \tilde{W}^{\mu\nu} + \frac{\alpha_3}{16\pi\Lambda_3} A G^\sigma_{\mu\nu} \tilde{G}^{\sigma\mu\nu} \, .$$

(7)
where $B_{\mu\nu}, W^a_{\mu\nu}, G^\sigma_{\mu\nu}$ are field strengths for the $U(1)_Y$, $SU(2)_W$ and $SU(3)_C$ gauge interactions, respectively, whereas $\tilde{G}^{\sigma\mu\nu} = \frac{1}{2} \epsilon^{\mu\nu\lambda\rho} G^\sigma_{\lambda\rho}$, $\tilde{W}^{a\mu\nu} = \frac{1}{2} \epsilon^{\mu\nu\lambda\rho} W^a_{\lambda\rho}$ and $\tilde{B}^{\mu\nu} = \frac{1}{2} \epsilon^{\mu\nu\lambda\rho} B_{\lambda\rho}$.

Here $\alpha_Y = 3 \alpha_1 / 5$ while $\alpha_1$, $\alpha_2$ and $\alpha_3$ are (GUT normalised) gauge couplings of $U(1)_Y$, $SU(2)_W$ and $SU(3)_C$ interactions.

In Eq. (7) all interactions between the pseudoscalar $A$ and SM fermions except the coupling of $A$ to the top quarks were ignored, because the masses of the SM quarks and leptons are negligibly small as compared with the scale $f$ and the top quark mass $m_t$.

The first term in Eq. (7) stems from the interactions

$$L_t = g_t 15(Q) \Omega 20^t + \tilde{g}_t 20(Q) \Omega 15^t + h.c. \quad (8)$$

in the strongly interacting sector, where the linear combination of the appropriate components of the 10–plets of $SU(5)$ from $15(Q)$ and $20(Q)$ corresponds to the composite partners of the third generation left–handed quark doublet. Then the mixing between the third generation left–handed quark doublet from the weakly–coupled elementary sector and its composite partners leads to the first term in Eq. (7). In the case when $t^c$ is predominantly the appropriate component of $20^t$ of $SU(6)$, the scale $\Lambda_t = \sqrt{15} f$. If $t^c$ is mainly a component of $15^t$ of $SU(6)$ then $\Lambda_t = \sqrt{\frac{60}{49}} f \simeq 1.1 f$. Since our analysis is motivated by the recently observed diphoton excess associated with the 750 GeV resonance and this resonance has not been detected in other channels, including $pp \to t\bar{t}$, we are going to focus on the scenario in which the decay rate $A \to t\bar{t}$ is strongly suppressed. Such a suppression can be attained if $\Lambda_t \gg 10$ TeV. This scenario is realised when $t^c$ is predominantly a component of $20^t$, so that $\Lambda_t \gg f$.

The second term in Eq. (7) originates from the Lagrangian (6). The interactions specified in Eq. (6) also give rise to the mass terms of the exotic fermions

$$L_{mass} = \mu_D \tilde{D}^c D^c + \mu_Q \tilde{Q} Q + \mu_L \tilde{L} L + \mu_E \tilde{E} E + \mu_\eta \tilde{\eta} \eta + h.c. \quad (9)$$

The masses $\mu_i$ in Eq. (9) are linear combinations of two contributions. One of these contributions is proportional to $\tilde{\mu}_i$ whereas another is induced as a result of the breakdown of $SU(6)$ global symmetry to its $SU(5)$ subgroup, that contains the $SU(3)_C \times SU(2)_W \times U(1)_Y$ gauge symmetry and is therefore proportional to $\tilde{\sigma}_i f$. The approximate gauge coupling unification requires exotic fermions to be substantially lighter than 10 TeV. This can be always achieved by adjusting the mass parameters $\tilde{\mu}_i$ in Eq. (9).

On the other hand, the Yukawa couplings $\kappa_i$ and $\lambda_i$ in Eq. (7) are proportional to $\tilde{\sigma}_i$ only. Thus the masses $\mu_i$ and couplings $\kappa_i$ and $\lambda_i$ are completely independent parameters, which are not constrained by either $SU(6)$ or $SU(5)$ global symmetries.

The lightest exotic fermion state in the $E_6$CHM must be stable. Indeed, because of baryon number conservation the low energy effective Lagrangian of the $E_6$CHM is also
invariant under the transformations of the baryon triality \([11]\) which is defined as

\[ \Psi \rightarrow e^{2\pi i B_3/3} \Psi, \quad B_3 = (3B - n_C)_{\text{mod } 3}, \]  

(10)

where \(B\) is the baryon number of the given multiplet \(\Psi\) and \(n_C\) is the number of colour indices (\(n_C = 1\) for the colour triplet and \(n_C = -1\) for \(\overline{3}\)). All SM particles have \(B_3 = 0\), while exotic states carry either \(B_3 = 1\) or \(B_3 = 2\). Because of this, the lightest exotic state can not decay into SM particles. In this context it is worth noting that models with stable charged exotic particles are ruled out by various experiments \([35]\). Moreover, the coupling of the neutral Dirac fermion, which is absolutely stable, to the \(Z\)–boson should be extremely suppressed. Otherwise such particles would scatter on nuclei leading to unacceptably large spin–independent cross sections (for a recent analysis see \([36]\)). As a consequence, only a Dirac fermion that involves mostly \(\eta\) and \(\bar{\eta}\), can be the lightest exotic state in the \(E_6\) CHM.

The last three terms in Eq. (7) are induced by the composite partners of the SM fermions. Therefore \(\Lambda_1, \Lambda_2\) and \(\Lambda_3\) are expected to be of the order of scale \(f\). In order to obtain a sufficiently large diphoton production cross section, we also assume that \(\mu_D, \mu_Q, \mu_L\) and \(\mu_E\) are larger than 375 GeV, so that the on-shell decays of \(A\) into the corresponding exotic fermions are not kinematically allowed. Integrating out the heavy exotic states, which appear in the usual triangle loop diagrams, one gets the effective Lagrangian that describes the interactions of pseudoscalar \(A\) with the top quark, SM gauge bosons and dark matter particles (see also \([2]–[3]\))

\[ L_{\text{eff}}^A = c_1 AB_{\mu \nu} \tilde{B}^{\mu \nu} + c_2 AW_{\mu \nu}^{a} \tilde{W}^{a \mu \nu} + c_3 AG_{\mu \nu}^{\sigma} \tilde{G}^{\sigma \mu \nu} + \frac{y_t}{\Lambda_t} A(i\bar{t}_L H t_R + h.c.) + i\lambda_\eta A(\bar{\eta}\eta + h.c.), \]  

(11)

where

\[ c_1 = \frac{\alpha_Y}{16\pi} \left[ \frac{2 \kappa_D}{3 \mu_D} B(x_D) + \frac{\kappa_Q}{3 \mu_Q} B(x_Q) + \frac{\lambda_L}{\mu_L} B(x_L) + 2 \frac{\lambda_E}{\mu_E} B(x_E) \right] + \frac{\alpha_Y}{16\pi \Lambda_1}, \]

(12a)

\[ c_2 = \frac{\alpha_2}{16\pi} \left[ \frac{3 \kappa_Q}{\mu_Q} B(x_Q) + \frac{\lambda_L}{\mu_L} B(x_L) \right] + \frac{\alpha_2}{16\pi \Lambda_2}, \]

(12b)

\[ c_3 = \frac{\alpha_3}{16\pi} \left[ \frac{\kappa_D}{\mu_D} B(x_D) + 2 \frac{\kappa_Q}{\mu_Q} B(x_Q) \right] + \frac{\alpha_3}{16\pi \Lambda_3}, \]

\[ B(x) = 2x \arcsin^2[1/\sqrt{x}], \quad \text{for} \quad x \geq 1. \]

In Eq. (12) \(x_D = 4 \mu_D^2/m_A^2, x_Q = 4 \mu_Q^2/m_A^2, x_L = 4 \mu_L^2/m_A^2, x_E = 4 \mu_E^2/m_A^2\) and \(m_A \approx 750\) GeV is the mass of the SM singlet pNGB state \(A\). The last terms in the analytical expressions for \(c_i\) (see Eq. (12)) tend to be rather small, because the scenario under consideration implies that \(\Lambda_i \sim f\) and \(\mu_D \sim \mu_Q \sim \mu_L \sim \mu_E \ll f\), and, thereby, can be neglected in our analysis.
Since we focus here on the diphoton decays of the pseudoscalar $A$ that may lead to the 750 GeV diphoton excess, it is convenient to derive an explicit analytical expression for the coupling of the SM singlet pNGB state $A$ to the electromagnetic field. Using Eqs. (11)–(12) one finds

$$L_{\text{eff}}^{A\gamma\gamma} = c_{\gamma} A F_{\mu\nu} \tilde{F}^{\mu\nu}, \quad (13)$$

where $F_{\mu\nu}$ is a field strength associated with the electromagnetic interaction, 

$$\tilde{F}^{\mu\nu} = \frac{1}{2} \epsilon^{\mu\nu\lambda\rho} F_{\lambda\rho}$$

and $\theta_W$ is the weak mixing (Weinberg) angle.

4 750 GeV diphoton excess

At the LHC the SM singlet pNGB state $A$ is predominantly produced through gluon fusion. The corresponding production cross section is determined by the effective coupling $|c_3|^2$. Since the partial width of the decays $A \to gg$ is also proportional to $|c_3|^2$ it is convenient to present the LHC production cross section of the pseudoscalar $A$ in the following form [2]

$$\sigma_{gg} \simeq \frac{C_{gg}}{m_A} \frac{\Gamma(A \to gg)}{\Gamma_A} \simeq 7.3 \text{ fb} \times \left( \frac{\Gamma(A \to gg)}{m_A} \times 10^6 \right), \quad (14)$$

where $C_{gg} \simeq 3163$, $\sqrt{s} \simeq 13$ TeV and the partial decay width $\Gamma(A \to gg)$ is given by

$$\Gamma(A \to gg) = \frac{2m_A^3}{\pi} |c_3|^2. \quad (15)$$

Then the cross section $\sigma_{\gamma\gamma} = \sigma(pp \to A \to \gamma\gamma)$ for given total width of the SM singlet pNGB state $\Gamma_A$ and partial decay width $\Gamma(A \to \gamma\gamma)$ can be approximately estimated as

$$\sigma_{\gamma\gamma} = \sigma_{gg} \frac{\Gamma(A \to \gamma\gamma)}{\Gamma_A} \simeq 7.3 \text{ fb} \times \left( \frac{\Gamma(A \to gg)}{\Gamma_A m_A} \frac{\Gamma(A \to \gamma\gamma)}{m_A} \times 10^6 \right). \quad (16)$$

In Eq. (16) the partial decay width associated with the decays of the SM singlet pNGB state into a pair of photons is given by

$$\Gamma(A \to \gamma\gamma) = \frac{m_A^3}{4\pi} |c_\gamma|^2. \quad (17)$$

If all exotic fermions were lighter than 375 GeV then the tree-level decays of the pseudoscalar $A$ into pairs of such fermions would be kinematically allowed and as a consequence the total decay width $\Gamma_A$ would be rather large, because the Yukawa couplings $\kappa_D$, $\kappa_Q$, $\lambda_L$, $\lambda_E$ and $\lambda_\eta$ in Eq. (7) are expected to be of the order of unity. This would result in a strong suppression of the cross section (16). Therefore in our analysis we restrict
our consideration to the case when all charged exotic fermions are heavier than half the mass of the 750 GeV resonance. At the same time we allow the lightest exotic fermion state to be either lighter or heavier than 375 GeV. If $\mu_\eta > 375$ GeV then all decays of the pseudoscalar $A$ proceed through usual triangle loop diagrams and the corresponding total decay width $\Gamma_A$ tends to be sufficiently small. When $\mu_\eta < 375$ GeV the tree-level decays of $A$ into $\eta\bar{\eta}$ are kinematically allowed so that $\Gamma_A$ can be rather large. Here we examine these two possibilities separately. In both cases we assume that the mass of the lightest exotic fermion is quite close to half the mass of the 750 GeV resonance and the Yukawa coupling $\lambda_\eta \gtrsim 1$ giving rise to sufficiently large annihilation cross section for $\eta\bar{\eta} \to$ SM particles. This should enable the lightest exotic fermion state to account for all or some of the observed cold dark matter density.

4.1 Narrow width case

Let us first consider the scenario with $\mu_\eta > 375$ GeV. Since no indication of the presence of the 750 GeV resonance has been detected in the channels $pp \to t\bar{t}, WW, ZZ, b\bar{b}, \tau\bar{\tau}$ and $jj$, we initially identify the scenario that results in the appropriate suppression of the decay rates $A \to t\bar{t}, WW, ZZ, b\bar{b}, \tau\bar{\tau}$. The partial decay width, that corresponds to the decay mode $A \to t\bar{t}$, is given by

$$\Gamma(A \to t\bar{t}) = \frac{3m_A m_t^2}{8\pi \Lambda_t^2} \sqrt{1 - \frac{4m_t^2}{m_A^2}}.$$ (18)

The requirement $\Gamma(A \to t\bar{t}) \lesssim \Gamma(A \to \gamma\gamma)$ can be fulfilled only if $\Lambda_t \gg 10$ TeV. Thus the appropriate suppression of the decay width (18) can be obtained when the scale $f$ is rather large, i.e. $f \gtrsim 10$ TeV. Although $\Gamma(A \to t\bar{t})$ can become negligibly small, when $\Lambda_t$ and $f$ increase, this region of the parameter space is strongly disfavoured by naturalness arguments. Indeed, the masses of the exotic fermions tend to be of the order of $\kappa_i f$ and $\lambda_i f$. Because it is expected that $\kappa_i \sim \lambda_i \sim 1$ some fine-tuning is required to keep exotic fermions substantially lighter than the scale $f$ which is needed to ensure the approximate unification of the SM gauge couplings and to get sufficiently large LHC production cross section of the pseudoscalar $A$. In this case an extremely large fine–tuning is also needed to obtain the EW scale $v \ll f$. In our numerical analysis we set $\Lambda_t \simeq 80$ TeV. When $t^c$ is predominantly formed by the appropriate components of the composite $20'$ representation of $SU(6)$ so large value of $\Lambda_t$ corresponds to $f \simeq 20$ TeV. For so large values of the scale $f$ the decay rates $A \to b\bar{b}$ and $A \to \tau\bar{\tau}$ become negligibly small and can be ignored in the leading approximation.

The analytical expressions for the partial widths associated with the decays
$A \to WW, ZZ, \gamma Z$ can be presented in the following form

$$\Gamma(A \to WW) = \frac{m_A^3}{2\pi} |c_2|^2 \left(1 - \frac{4M_W^2}{m_A^2}\right)^{3/2}, \quad (19)$$

$$\Gamma(A \to ZZ) = \frac{m_A^3}{4\pi} \left|c_1 \sin^2 \theta_W + c_2 \cos^2 \theta_W\right|^2 \left(1 - \frac{4M_Z^2}{m_A^2}\right)^{3/2}, \quad (20)$$

$$\Gamma(A \to \gamma Z) = \frac{m_A^3}{8\pi} \sin^2 2\theta_W |c_1 - c_2|^2 \left(1 - \frac{M_Z^2}{m_A^2}\right)^3. \quad (21)$$

In order to obtain $\Gamma(A \to WW) \lesssim \Gamma(A \to \gamma \gamma)$ the absolute value of $c_2$ should be quite small. When $|c_2| \ll |c_1|$ the partial decay widths $\Gamma(A \to ZZ)$ and $\Gamma(A \to \gamma Z)$ tend to be considerably smaller than $\Gamma(A \to \gamma \gamma)$ because the value of $\sin^2 2\theta_W$ is quite small. On the other hand if $\mu_D \sim \mu_Q \sim \mu_L \sim \mu_E$ then from Eqs. (2) it follows that $|c_2|$ is substantially larger than $|c_1|$. This happens because $c_2$ and $c_1$ are proportional to the $SU(2)_W$ and $U(1)_Y$ gauge couplings respectively and at low energies $c_2$ is considerably bigger than $\alpha_Y$. The absolute value of $c_1$ becomes substantially bigger than $|c_2|$ only when the exotic fermions that form $SU(2)_W$ doublets are much heavier than the ones that do not participate in the $SU(2)_W$ gauge interactions.

The absence of new particles carrying colour with masses below 1 TeV, which should have sufficiently large LHC production cross section, implies that all exotic, coloured fermions and scalar coloured triplet in the E6CHM have to be quite heavy. To simplify our analysis here we focus on the scenario with $\mu_D = \mu_Q = \mu_L = \mu_0 \gtrsim \mu_E$, $\kappa_D = \kappa_Q = \lambda_L = \lambda_E = \sigma$ and $A_1 \simeq 80$ TeV. In the case when $\mu_0 \gg \mu_E$ this scenario should lead to rather suppressed decay rates for $A \to t \bar{t}, WW, ZZ, b \bar{b}$ and $\tau \bar{\tau}$. At the same time $\mu_0$ should be sufficiently small so that reasonable large LHC production cross section of the SM singlet pNGB state $A$ can be obtained. The results of our numerical analysis indicates that in this part of the parameter space the pseudoscalar $A$ still mainly decays into a pair of gluons. These in turn are very difficult to detect at the LHC if the corresponding cross section is sufficiently small. Since in this scenario $\Gamma_A \approx \Gamma(A \to gg)$ the cross section (16) is basically determined by the partial decay width $\Gamma(A \to \gamma \gamma)$. Then the mass of the $SU(2)_W$ singlet charge $\pm 1$, vector-like exotic fermion with baryon number $\mp 1/3$ can be adjusted so that the ratio $\left(\Gamma(A \to \gamma \gamma)/m_A\right) \sim 10^{-6}$ giving rise to $\sigma(pp \to A \to \gamma \gamma) \simeq 5 - 10$ fb.

Our results are summarised in Fig. 1 and 2. In Figs. 1a and 1b we consider the dependence of the cross section $\sigma(pp \to A \to \gamma \gamma)$ on $\sigma$ and $\mu_0$ respectively. One can see that this cross section decreases very rapidly with increasing $\mu_E$. From Fig. 1a it follows $\sigma(pp \to A \to \gamma \gamma)$ can be larger than 5 fb for $\mu_E = 400$ GeV ($\mu_E = 500$ GeV) when $\sigma \gtrsim 1.5$ ($\sigma \gtrsim 2$), even if all other exotic fermions are extremely heavy $\mu_0 \simeq 5$ TeV, i.e. beyond the
Figure 1: The cross section of $\sigma(pp \rightarrow A \rightarrow \gamma\gamma)$ and the ratio $\Gamma_A/m_A$ are shown for $\mu_E = 400$ GeV (solid lines), $\mu_E = 500$ GeV (dashed lines) and $\mu_E = 800$ GeV (dashed–dotted lines), for the case $\Lambda_t = 80$ TeV and $\kappa_D = \kappa_Q = \lambda_L = \lambda_E = \sigma$. In (a) we show the cross section of $\sigma(pp \rightarrow A \rightarrow \gamma\gamma)$ as a function of $\sigma$ for $\mu_Q = \mu_D = \mu_L = 5$ TeV. In (b) we give the cross section of $\sigma(pp \rightarrow A \rightarrow \gamma\gamma)$ as a function of $\mu_Q = \mu_D = \mu_L = \mu_0$ for $\sigma = 1.5$. In (c) the ratio $\Gamma_A/m_A$ is shown as a function of $\sigma$ for $\mu_Q = \mu_D = \mu_L = 5$ TeV. In (d) the ratio $\Gamma_A/m_A$ is presented as a function of $\mu_Q = \mu_D = \mu_L = \mu_0$ for $\sigma = 1.5$. 
LHC reach. For $\mu_E \simeq 700 - 800$ GeV and $\mu_0 = 5$ TeV the cross section $\sigma(pp \rightarrow A \rightarrow \gamma\gamma)$ can reach 5 fb only if the coupling $\sigma$ is very large, i.e. close to the upper bound ($\sigma \sim \sqrt{4\pi}$) beyond which perturbation theory is not valid. Fig. 1c indicates that in all these cases the total width $\Gamma_A$ of the SM singlet pNGB state $A$ is quite narrow, i.e. $\Gamma_A/m_A \sim 10^{-5}$.

This corresponds to the LHC production cross section (14) around $100 - 200$ fb. When $\mu_E \simeq 500 - 800$ GeV the cross section $\sigma(pp \rightarrow A \rightarrow \gamma\gamma)$ $\simeq 5 - 10$ fb can also be obtained for $\sigma \simeq 1.5$, provided other exotic fermions have masses about 1 TeV, i.e. $\mu_0 = 1$ TeV (see Fig. 1b). For $\mu_E \lesssim 500$ GeV and $\sigma \simeq 1.5$ the same values of the cross section (16) can be reproduced for larger $\mu_0$. Fig. 1d demonstrates that the total width $\Gamma_A$ of the pseudoscalar $A$ increases considerably when $\mu_0$ decreases. If $\mu_0 \simeq 1$ TeV the ratio $\Gamma_A/m_A$ becomes of the order of $10^{-4}$, which leads to an LHC production cross section for the pNGB state $A$ of order 1 pb.

In Fig. 2 the dependence of the branching ratios of the SM singlet pNGB state $A$ on $\sigma$ and $\mu_0$ is examined. The branching ratios of the pseudoscalar $A$ are given by

$$BR(A \rightarrow gg) = \frac{\Gamma(A \rightarrow gg)}{\Gamma_A}, \quad BR(A \rightarrow \gamma\gamma) = \frac{\Gamma(A \rightarrow \gamma\gamma)}{\Gamma_A},$$

$$BR(A \rightarrow t\bar{t}) = \frac{\Gamma(A \rightarrow t\bar{t})}{\Gamma_A}, \quad BR(A \rightarrow Z\gamma) = \frac{\Gamma(A \rightarrow Z\gamma)}{\Gamma_A},$$

$$BR(A \rightarrow ZZ) = \frac{\Gamma(A \rightarrow ZZ)}{\Gamma_A}, \quad BR(A \rightarrow WW) = \frac{\Gamma(A \rightarrow WW)}{\Gamma_A}.$$ (22)

Figs. 2a and 2b show the dependence of the branching ratios (22) on the dimensionless coupling $\sigma$ for $\mu_E \simeq 400$ GeV and $\mu_E \simeq 500$ GeV, respectively. One can see that the branching ratio associated with the decays $A \rightarrow t\bar{t}$ decreases with increasing $\sigma$, whereas all other branching ratios become bigger. In particular, for $\mu_E \simeq 400$ GeV and $\sigma > 1.5$ as well as for $\mu_E \simeq 500$ GeV and $\sigma > 2$ the branching fraction corresponding to $A \rightarrow \gamma\gamma$ is larger than $BR(A \rightarrow t\bar{t})$ because $\Lambda_t$ is sufficiently large. However the branching ratio of $A \rightarrow \gamma\gamma$ is not the dominant one. In the part of the parameter space under consideration the biggest branching fraction is associated with the decays of the pseudoscalar $A$ into a pair of gluons. $BR(A \rightarrow \gamma\gamma)$ is the second largest one. Since $\mu_0 \simeq 5$ TeV the branching fractions $BR(A \rightarrow WW)$, $BR(A \rightarrow ZZ)$ and $BR(A \rightarrow Z\gamma)$ are considerably smaller than $BR(A \rightarrow \gamma\gamma)$ and that may explain why the decays $A \rightarrow WW, ZZ, Z\gamma$ have not been observed yet. On the other hand, the decays of the pseudoscalar $A$ into a pair of gluons are quite difficult to detect since the total LHC production cross section of this state remains rather small.

In Fig. 2c and 2d the variations of the branching fractions of the SM singlet pNGB state $A$ as a function of $\mu_0$ is explored for $\mu_E \simeq 400$ GeV and $\sigma = 1.5$, as well as for $\mu_E \simeq 500$ GeV and $\sigma = 2$, respectively. The decay mode $A \rightarrow gg$ remains the dominant decay channel for any value of $\mu_0$ from 1 TeV to 5 TeV. The corresponding branching
Figure 2: The branching ratios of the decays of the pNGB state $A$ into $t\bar{t}$ (dashed–dotted lines), $gg$ (highest solid lines), $\gamma\gamma$ (highest dashed lines), $WW$ (thick solid lines), $ZZ$ (lowest solid lines) and $\gamma Z$ (lowest dashed lines) for $\Lambda_t = 80$ TeV and $\kappa_D = \kappa_Q = \lambda_L = \lambda_E = \sigma$. In (a) the branching fractions of the decays of $A$ are given as a function of $\sigma$ for $\mu_E = 400$ GeV and $\mu_Q = \mu_D = \mu_L = 5$ TeV. In (b) the branching ratios of the decays of $A$ are shown as a function of $\sigma$ for $\mu_E = 500$ GeV and $\mu_Q = \mu_D = \mu_L = 5$ TeV. In (c) the branching fractions of the decays of $A$ are presented as a function of $\mu_Q = \mu_D = \mu_L = \mu_0$ for $\mu_E = 400$ GeV and $\sigma = 1.5$. In (d) the branching ratios of the decays of $A$ as a function of $\mu_Q = \mu_D = \mu_L = \mu_0$ are shown for $\mu_E = 500$ GeV and $\sigma = 2.0$. 
fraction is always close to 100%. $\text{BR}(A \to \gamma\gamma)$ and $\text{BR}(A \to t\bar{t})$ decreases whereas $\text{BR}(A \to ZZ)$ and $\text{BR}(A \to WW)$ increases with decreasing $\mu_0$. In this part of the parameter space $\text{BR}(A \to t\bar{t})$ and $\text{BR}(A \to \gamma Z)$ are considerably smaller than $\text{BR}(A \to \gamma\gamma)$. The branching ratio associated with the decay $A \to \gamma Z$ tends to be the lowest branching fraction and vanishes at some values of $\mu_0$. When $\mu_0$ is around 1 TeV the values of $\text{BR}(A \to ZZ)$ and $\text{BR}(A \to WW)$ are a few times bigger than $\text{BR}(A \to \gamma\gamma)$. Although the branching ratios of $A \to WW$ and $A \to ZZ$ can be substantially larger than the branching ratio of $A \to \gamma\gamma$, their experimental detection is more problematic because the $W$ and $Z$ decay mainly into quarks.

4.2 Large width scenario

Let us now explore the scenario with $\mu_\eta \lesssim 375$ GeV, so that the decays of $A$ into $\eta\bar{\eta}$ are kinematically allowed. Fig. 3a indicates that the partial width associated with the decays $A \to \eta\bar{\eta}$ tends to be rather large in this case if the Yukawa coupling $\lambda_\eta$ is considerably larger than unity. Therefore the parameters of the $E_6$CHM can be always adjusted so that the total width of the SM singlet pNGB state $A$ is of order 45 GeV, or even larger. Further, we assume that the mass of the lightest exotic fermion is rather close to half the mass of the 750 GeV resonance, so that the cross section for $\eta\bar{\eta} \to \text{SM}$ particles is sufficiently large. The large total width $\Gamma_A$ of the pseudoscalar $A$ leads to the significant suppression of the diphoton production cross section (16). Because of this, relatively large partial decay widths $\Gamma(A \to gg)$ and $\Gamma(A \to \gamma\gamma)$ are needed in order to get $\sigma(pp \to A \to \gamma\gamma) \simeq 5\text{--}10$ fb. When $\kappa_D = \kappa_Q = \lambda_L = \lambda_E = \sigma$, this can be achieved only if $\sigma \gtrsim 2$ (see Fig. 3b). In the scenario with $\sigma \simeq 2$ or slightly bigger a reasonably large diphoton production cross section can only be obtained if the masses of all charged exotic fermions are very close to 375 GeV. Since all exotic fermions, including the lightest one, are almost degenerate in this case, the energy of the jets that appear in the final states associated with the decays of these fermions can be very low. If the energy of these jets are so low that their detection at the LHC is not possible, the decays of the charged exotic fermions basically result in invisible final states. Although there is a chance that such ”invisible” exotic states could remain so far undetected, the presence of exotic coloured fermions with masses around 400 GeV should, in general, give rise to a rather large cross section for the process $pp \to j + E_T^{\text{miss}}$, which has not been observed by experiments at the LHC [37].

A more realistic scenario would correspond to the case where the total width $\Gamma_A$ is somewhat smaller than 45 GeV and the diphoton production cross section (16) is a bit lower than 5 fb. In this case the exotic coloured fermions can be considerably heavier than 400 GeV. For example, as follows from Fig. 3b the cross section $\sigma(pp \to A \to \gamma\gamma) \simeq 4$ fb can be obtained for $\mu_Q = \mu_D \simeq 1$ TeV if $\mu_L = \mu_E = 400$ GeV, $\sigma = \sqrt{4\pi}$ and $\Gamma_A = 20$ GeV.
Figure 3: The ratio $\Gamma(A \rightarrow \eta\bar{\eta})/m_A$, the cross sections $\sigma(pp \rightarrow A \rightarrow \gamma\gamma)$ and $\sigma(pp \rightarrow A)$ as well as the branching ratios of the decays of the pNGB state $A$ are presented for $\mu_\eta \lesssim 375$ GeV, $\mu_L = \mu_E = 400$ GeV and $\kappa_D = \kappa_Q = \lambda_L = \lambda_E = \sigma$. In (a) the ratio $\Gamma(A \rightarrow \eta\bar{\eta})/m_A$ is shown as a function of the mass of the lightest exotic fermion $\mu_\eta$ for $\lambda_\eta = 2.5$ (solid line), $\lambda_\eta = 2$ (dashed line) and $\lambda_\eta = 1.4$ (dashed–dotted line). In (b) the cross section of $\sigma(pp \rightarrow A \rightarrow \gamma\gamma)$ is shown as a function of $\mu_Q = \mu_D$ for $\Gamma_A = 45$ GeV (solid lines) and $\Gamma_A = 20$ GeV (dashed lines). The upper solid and dashed lines correspond to $\sigma = \sqrt{4\pi}$, whereas lower solid and dashed lines are associated with $\sigma = 2$. In (c) the production cross section of $A$ at 13 TeV LHC is given as a function of $\mu_Q = \mu_D$ for $\sigma = \sqrt{4\pi}$ (upper solid line) and $\sigma = 2$ (lower solid line). In (d) we show the branching ratios of the decays of $A$ into $t\bar{t}$ (dashed–dotted lines), $gg$ (highest solid lines), $\gamma\gamma$ (highest dashed lines), $WW$ (thick solid lines), $ZZ$ (lowest solid lines) and $\gamma Z$ (lowest dashed lines) as a function of $\mu_Q = \mu_D$ for $\Lambda_t = 20$ TeV, $\sigma = \sqrt{4\pi}$ and $\Gamma_A = 45$ GeV.
Fig. 3c shows that in order to get $\sigma(pp \to A \to \gamma\gamma) \gg 1 \text{ fb}$ the LHC production cross section of the SM singlet pNGB state should be of order of 10 pb. This is because the branching ratio $A \to \gamma\gamma$ is extremely small, i.e. $\sim 10^{-3}$ (see Fig. 3d). Since $A$ decays predominantly into a pair $\eta\bar{\eta}$ that leads to the $E_T^{\text{miss}}$ in the final state, the large LHC production cross section of the pseudoscalar $A$ should also result in a non-negligible cross section for $pp \to j + E_T^{\text{miss}}$. In this context the scenarios with TeV scale exotic coloured fermions are more favorable. Indeed, as one can see from Fig. 3c the LHC production cross section of the pseudoscalar $A$ decreases with increasing of masses of the exotic coloured fermions. As a consequence, the cross section of $pp \to j + E_T^{\text{miss}}$ also becomes considerably smaller if $\mu_Q$ and $\mu_D$ increase. When $\mu_Q = \mu_D \simeq 1 \text{ TeV}$ and $\sigma \simeq \sqrt{4\pi}$ the LHC production cross section of the SM singlet pNGB state reduces to about 4 pb, which is an order of magnitude smaller than that calculated for $\mu_Q = \mu_D \simeq 400 \text{ GeV}$.

Fig. 3d demonstrates that in the case when all exotic fermions have masses below 1 TeV the branching ratio of $A \to gg$ is the second largest, while the branching ratios of $A \to WW$ and $A \to ZZ$ are the third and fourth largest, respectively. These branching fractions are considerably larger than BR($A \to \gamma\gamma$). Despite this, the decay modes $A \to gg$, $WW$ and $ZZ$ result in final states that predominantly contain jets and therefore their experimental identification tends to be rather problematic. This may why the 750 GeV resonance has been observed in the $\gamma\gamma$ channel only. Fig. 3d also indicates that other decay rates, including $A \to Z\gamma$ and $A \to t\bar{t}$, can be significantly smaller in this case as compared with the one associated with $A \to \gamma\gamma$. All branching ratios except BR($A \to \eta\bar{\eta}$) and BR($A \to t\bar{t}$) decrease when the masses of the exotic coloured fermions grow (see Fig. 3d). In particular, for $\sigma = \sqrt{4\pi}$ and $\Gamma_A = 45 \text{ GeV}$ the branching fractions BR($A \to gg$), BR($A \to WW$) and BR($A \to ZZ$) reduce from 16%, 0.82% and 0.28% to 1%, 0.16% and 0.06%, respectively, if $\mu_Q = \mu_D$ vary from 400 GeV to 1 TeV.

The presence of exotic fermions with masses below 1 TeV should lead to remarkable signatures which have already been mentioned in [5]. The production processes of exotic coloured fermions involve gluon–induced QCD interactions, so that these states are doubly produced. Assuming that such exotic coloured states couple most strongly to the third generation fermions, they decay into a pair of third generation quarks and the lightest neutral exotic state resulting in the enhancement of the cross sections for

$$pp \to t\bar{t}b\bar{b} + E_T^{\text{miss}} + X$$
and
$$pp \to b\bar{b}b\bar{b} + E_T^{\text{miss}} + X . \quad (23)$$

Although the production cross sections of other exotic fermions tend to be somewhat smaller, their pair production can also result in an enhancement of the cross sections for the processes with the final states (23). Thus the exotic fermion states should be observable in dedicated searches at Run 2 of the LHC.
5 Conclusions

In the $E_6$ inspired composite Higgs model ($E_6$CHM) the strongly interacting sector possesses an $SU(6) \times U(1)_B \times U(1)_L$ global symmetry. Global $U(1)_B$ and $U(1)_L$ symmetries ensure the conservation of baryon and lepton numbers. Near the scale $f \gtrsim 10 \text{ TeV}$ the $SU(6)$ global symmetry is broken down to its $SU(5)$ subgroup, that contains the SM gauge group. Such breakdown of the $SU(6)$ symmetry leads to eleven pNGB states. This set of pNGB states involves, in particular, the SM–like Higgs doublet and SM singlet boson. If CP is conserved then the SM singlet scalar is CP–odd and its mixing with the SM–like Higgs boson is forbidden. Below scale $f$ the particle spectrum of the $E_6$CHM also includes extra matter beyond the SM that leads to the approximate unification of the SM gauge couplings and gives rise to a set of exotic vector–like fermions as well as a composite, right-handed top quark. Baryon number conservation guarantees that the lightest exotic fermion state is stable. Because of this the phenomenological viability of the model under consideration requires the lightest exotic fermion to be a SM singlet, so that it can play the role of dark matter.

In this paper we focused on the scenario in which the SM singlet pNGB boson $A$ can be identified with the 750 GeV diphoton resonance recently reported by ATLAS and CMS. Since $f \gtrsim 10 \text{ TeV}$, the decay rate $A \to t\bar{t}$ tends to be sufficiently strongly suppressed so that it may explain why the 750 GeV resonance has not been observed in the $pp \to t\bar{t}$ channel. This scenario also implies that all charged exotic fermions are heavier than 375 GeV, so that the decays of the pseudoscalar $A$ into pairs of these exotic states are not kinematically allowed. Nevertheless, the interactions of charged exotic fermions with $A$ induce couplings of this SM singlet pNGB state to the SM gauge bosons. We considered the interactions of $A$ with exotic fermions and specified the couplings of this pseudoscalar boson to $gg$, $\gamma\gamma$, $WW$, $ZZ$ and $Z\gamma$. Also we explored the dependence of the LHC diphoton production cross section, total width of the SM singlet pNGB state and its branching ratios on the $E_6$CHM parameters. In particular, we identified the part of the parameter space where the decay rates of $A \to WW, ZZ$ and $\gamma Z$ are quite suppressed while the partial widths associated with the decays $A \to gg$ and $A \to \gamma\gamma$ are comparable. In this case one charged exotic fermion, that does not carry colour charge, and the lightest SM singlet exotic fermion have masses just above 375 GeV, whereas all other exotic states can be considerably heavier, lying even beyond the projected LHC sensitivity. The total width of the pseudoscalar $A$ is rather narrow in this case, i.e. $\Gamma_A/m_A \sim 10^{-5}$.

In the $E_6$CHM there is also a part of the parameter space in which the SM singlet pNGB state has a relatively large width, i.e. $\Gamma_A/m_A \gtrsim 0.02$. In this case the pseudoscalar $A$ decays mainly into the lightest exotic fermions resulting in invisible final states. In order
to get a sufficiently large diphoton production cross section for $\Gamma_A/m_A \simeq 0.06$, the masses of all other charged exotic fermions should be around 400 GeV. If the total width of the 750 GeV resonance is 20 GeV or narrower, $\sigma(pp \rightarrow A \rightarrow \gamma\gamma) \simeq 4 - 5 \text{ fb}$ can be obtained provided the masses of exotic coloured fermions are around 1 TeV or larger. In general this scenario leads to quite a large cross section of the process $pp \rightarrow j + E_T^{\text{miss}}$. Moreover, the branching ratios of the pseudoscalar decays $A \rightarrow gg, WW, ZZ$ are considerably larger than the one associated with $A \rightarrow \gamma\gamma$. Although the final states of the decays $A \rightarrow gg, WW, ZZ$ predominantly contain jets, that makes their experimental detection more problematic, the observation of these decay modes should be possible in Run 2 at the LHC. In addition, the presence of exotic coloured fermions with masses below 1 TeV should give rise to an enhancement of the cross sections $pp \rightarrow t\bar{t}b\bar{b} + E_T^{\text{miss}} + X$ and $pp \rightarrow b\bar{b}b\bar{b} + E_T^{\text{miss}} + X$, which should be also observable at the LHC.

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