



Baryon asymmetry generation in the E_6 CHM



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ABSTRACT

In the E_6 inspired composite Higgs model (E_6 CHM) the strongly interacting sector possesses an $SU(6)$ global symmetry which is expected to be broken down to its $SU(5)$ subgroup at the scale $f \gtrsim 10$ TeV. This breakdown results in a set of pseudo-Nambu–Goldstone bosons (pNGBs) that includes one Standard Model (SM) singlet scalar, a SM-like Higgs doublet and an $SU(3)_C$ triplet of scalar fields, T . In the E_6 CHM the Z_2^L symmetry, which is a discrete subgroup of the $U(1)_L$ associated with lepton number conservation, can be used to forbid operators which lead to rapid proton decay. The remaining baryon number violating operators are sufficiently strongly suppressed because of the large value of the scale f . We argue that in this variant of the E_6 CHM a sizeable baryon number asymmetry can be induced if CP is violated. At the same time, the presence of the $SU(3)_C$ scalar triplet with mass in the few TeV range may give rise to spectacular new physics signals that may be detected at the LHC in the near future.

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1. Introduction

Although the new scalar particle discovered at the LHC in 2012 is consistent with the Standard Model (SM) Higgs boson, it could in principle be composed of more fundamental degrees of freedom. The idea of a composite Higgs boson was proposed in the 70's [1] and 80's [2]. It implies the presence of a strongly interacting sector in which electroweak (EW) symmetry breaking (EWSB) is generated dynamically. Generically, in models of this type the composite Higgs tends to have a large quartic coupling $\lambda \gtrsim 1$. At the same time, the observed SM-like Higgs boson is relatively light and corresponds to $\lambda \simeq 0.13$. This indicates that the discovered Higgs state could possibly be a pseudo-Nambu–Goldstone boson (pNGB) originating from the spontaneous breakdown of an approximate global symmetry of the strongly interacting sector.

The minimal composite Higgs model (MCHM) [3] contains two sectors (for a review, see Ref. [4]). One of them involves weakly-coupled elementary particles, including all SM gauge bosons and SM fermions. The second, strongly interacting sector gives rise to a set of bound states that, in particular, include composite partners of the elementary particles; that is, massive fields with quantum numbers of all SM particles.

The composite sector of the MCHM possesses a global $SO(5) \times U(1)_X$ symmetry which is broken down at the scale f to $SO(4) \times U(1)'_X \cong SU(2)_W \times SU(2)_R \times U(1)'_X$, which in turn contains the $SU(2)_W \times U(1)_Y$ gauge group as a subgroup. This breakdown results in a set of pNGB states that form the Higgs doublet. The global custodial symmetry $SU(2)_{cust} \subset SU(2)_W \times SU(2)_R$ [5] protects the Peskin–Takeuchi \hat{T} parameter [6], which is strongly constrained by experimental data [7], against the contributions induced by the composite states. The contributions of these new states to electroweak observables were examined in Refs. [8–16]. The implications of the composite Higgs models were also considered for Higgs physics [11,12], [17–20], gauge coupling unification [21,22], dark matter [9], [18], [22,23] and collider phenomenology [10,11], [13], [20], [24–26]. Non-minimal composite Higgs models were studied in Refs. [9], [17,18], [22,23], [27].

In these models the elementary and composite states with the same quantum numbers mix, so that at low energies those states associated with the SM fermions (bosons) are superpositions of the corresponding elementary fermion (boson) states and their composite fermion (boson) partners. In this partial compositeness framework [28,29] the SM fields couple to the composite states, including the Higgs boson, with a strength which is proportional to the compositeness fraction of each SM field. In this case the mass hierarchy in the quark and lepton sectors can be reproduced if the compositeness fractions of the first and second generation fermions are rather small. This also leads to the suppression of off-diagonal flavor transitions, as well as modifications of the W and Z

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couplings associated with these light quark and lepton fields [28], [30].

Even though this partial compositeness considerably reduces the contributions of composite states to dangerous flavour-changing processes, this suppression is not sufficient to satisfy all constraints. Within the composite Higgs models the constraints that stem from the off-diagonal flavour transitions in the quark and lepton sectors were explored in Refs. [14–16], [24], [31,32] and [25], [32–34], respectively. In particular, it was argued that in the case when the matrices of effective Yukawa couplings in the strongly interacting sector are structureless, i.e. anarchic matrices, the scale f has to be larger than 10 TeV [14,15], [24], [31], [33]. This bound can be considerably alleviated in composite Higgs models with flavour symmetries [13,14], [24,25], [32], [35], in which the constraints originating from the Kaon and B systems can be fulfilled if $f \gtrsim 1$ TeV. For such low values of the scale f adequate suppression of the baryon number violating operators and the Majorana masses of the left-handed neutrinos can be attained provided global $U(1)_B$ and $U(1)_L$ symmetries, which guarantee the conservation of the baryon and lepton numbers respectively, are imposed.

In this note we focus on an E_6 inspired composite Higgs model (E_6 CHM) in which the strongly interacting sector is invariant under the transformations of an $SU(6) \times U(1)_L$ global symmetry [36]. In the weakly-coupled sector $U(1)_L$ is broken down to a Z_2^L discrete symmetry which stabilises the proton. Since the composite sector in the E_6 CHM does not possess any flavour or custodial symmetry, $SU(6)$ is expected to be broken down to $SU(5)$, which in turn contains the SM gauge group, at a sufficiently high scale, $f \gtrsim 10$ TeV. This breakdown of the $SU(6)$ symmetry gives rise to a set of pNGBs that involves the SM-like Higgs doublet, scalar coloured triplet and a SM singlet boson. Because the scale f is so high, all baryon number violating operators are sufficiently strongly suppressed so that the existing experimental constraints are satisfied. Nevertheless, we argue that in the E_6 CHM, with explicitly broken $U(1)_B$ baryon symmetry, the observed matter–antimatter asymmetry can be induced if CP is broken. The pNGB scalar coloured triplet plays a key role in this process and leads to a distinct signature that can be observed at the LHC.

The layout of this article is as follows. In the next Section we discuss the E_6 CHM with broken baryon symmetry. In Section 3 we consider the process through which the baryon asymmetry is generated and present our estimate of its value. Section 4 concludes the paper.

2. E_6 CHM with baryon number violation

The gauge and global symmetries of the E_6 CHM, as well as its particle content, can originate from a Grand Unified Theory (GUT) based on the $E_6 \times G_0$ gauge group. At some high energy scale, M_X , the $E_6 \times G_0$ gauge symmetry is broken down to the $SU(3)_C \times SU(2)_W \times U(1)_Y \times G$ subgroup. The gauge groups G_0 and G are associated with the strongly interacting sector, whereas $SU(3)_C \times SU(2)_W \times U(1)_Y$ is the SM gauge group. Multiplets from the strongly coupled sector are charged under both the E_6 and G_0 (G) gauge symmetries. The weakly-coupled sector comprises fields that participate in the E_6 interactions only. It is expected that in this sector different multiplets of the elementary quarks and leptons come from different fundamental 27-dimensional representations of E_6 . All other components of these 27-plets acquire masses of the order of M_X . The corresponding splitting of the 27-plets can occur within a six-dimensional orbifold GUT model

with $N = 1$ supersymmetry (SUSY) [36] in which SUSY is broken somewhat below the GUT scale M_X .²

In this orbifold GUT model the elementary quarks and leptons are components of different bulk 27-plets, while all fields from the strongly interacting sector are localised on the brane, where the E_6 symmetry is broken down to the $SU(6) \times SU(2)_N$ subgroup. In the model under consideration E_6 is broken down to the SM gauge group and $SU(2)_N$ symmetry is entirely broken. Furthermore, $SU(6)$ can remain an approximate global symmetry of the strongly coupled sector. We assume that around the scale $f \gtrsim 10$ TeV the $SU(6)$ global symmetry is broken down to $SU(5)$. That, in turn, contains the SM gauge group as a subgroup, leading to a set of pNGB states which includes the SM-like Higgs doublet.

In the E_6 CHM the $U(1)_L$ global symmetry, which ensures the conservation of lepton number, can be used to suppress the operators in the strongly interacting sector that may induce too large Majorana masses of the left-handed neutrino. In the weakly-coupled elementary sector this symmetry is supposed to be broken down to

$$Z_2^L = (-1)^L, \quad (1)$$

where L is a lepton number, to guarantee that the left-handed neutrinos gain non-zero Majorana masses. If Z_2^L is an almost exact discrete symmetry it also forbids potentially dangerous operators that give rise to rapid proton decay. All other baryon number violating operators in the model under consideration are sufficiently strongly suppressed by the relatively large value of the scale f , as well as the rather small mixing between elementary states and their composite partners. Indeed, in the SM the effective operators responsible for $\Delta B = 2$ and $\Delta L = 0$ are given by

$$\mathcal{L}_{\Delta B=2} = \frac{1}{\Lambda^5} \left[q_i q_j q_k q_m (d_n^c d_l^c)^* + u_i^c d_j^c d_k^c u_m^c d_n^c d_l^c \right], \quad (2)$$

where q_i are doublets of the left-handed quarks, u_i^c and d_j^c are the right-handed up- and down-type quarks and the generation indices are $i, j, k, m, n, l = 1, 2, 3$.

The $n - \bar{n}$ mixing mass can be deduced from this operator by simple dimensional analysis to be $\delta m \simeq \kappa \Lambda_{QCD}^6 / \Lambda^5$, where κ is of order one and $\Lambda_{QCD} \simeq 200$ MeV. For $\Lambda \sim \text{few} \times 100$ TeV one finds the free $n - \bar{n}$ oscillation time to be $\tau_{n-\bar{n}} \simeq 1/\delta m \simeq 10^8$ s, which is rather close to the present experimental limit [39,40]. A similar bound on the scale Λ comes from the rare nuclear decay searches looking for the annihilation of the two nucleons $NN \rightarrow KK$, which may be also induced by the operators (2). In this case one obtains a lower limit on Λ of around 200–300 TeV. On the other hand, in the composite Higgs models the small mixing between elementary states and their composite partners leads to $\Lambda \gtrsim \text{few} \times 100$ TeV when $f \gtrsim 10$ TeV.

Thus, to ensure the phenomenological viability of the E_6 CHM, the Lagrangian of the strongly coupled sector of this model should respect the $SU(6) \times U(1)_L$ global symmetry. Here we also assume that the low energy effective Lagrangian of the E_6 CHM is invariant with respect to an approximate Z_2^B symmetry, which is a discrete subgroup of $U(1)_B$, i.e.

$$Z_2^B = (-1)^{3B}, \quad (3)$$

where B is the baryon number. The Z_2^B discrete symmetry does not forbid baryon number violating operators (2) but it does provide an additional mechanism for the suppression of the proton decay.

² Different phenomenological aspects of the E_6 inspired models with low-scale SUSY breaking were recently explored in [37,38].

In order to embed the E_6 CHM into a Grand Unified Theory (GUT) based on the $E_6 \times G_0$ gauge group, the SM gauge couplings extrapolated to high energies using the renormalisation group equations (RGEs) should converge to some common value near the scale M_X . Such an approximate unification of the SM gauge couplings can be achieved if the right-handed top quark t^c is entirely composite and the sector of weakly-coupled elementary states involves [36], [41]

$$(q_i, d_i^c, \ell_i, e_i^c) + u_\alpha^c + \bar{q} + \bar{d}^c + \bar{\ell} + \bar{e}^c, \quad (4)$$

where $\alpha = 1, 2$ and $i = 1, 2, 3$. Here we have denoted the left-handed lepton doublet by ℓ_i , the right-handed charged leptons by e_i^c , while the extra exotic states in Eq. (4), \bar{q} , \bar{d}^c , $\bar{\ell}$ and \bar{e}^c , 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. This scenario also implies that the strongly coupled sector gives rise to the composite $\mathbf{10} + \bar{\mathbf{5}}$ multiplets of $SU(5)$. These multiplets get combined with \bar{q} , \bar{d}^c , $\bar{\ell}$ and \bar{e}^c , resulting in a set of vector-like states. The only exceptions are the components of the 10-plet that correspond to t^c , which survive down to the EW scale.³

In the simplest case the composite $\mathbf{10} + \bar{\mathbf{5}}$ multiplets of $SU(5)$ stem from one $\mathbf{15}$ -plet and two $\bar{\mathbf{6}}$ -plets ($\bar{\mathbf{6}}_1$ and $\bar{\mathbf{6}}_2$) of $SU(6)$. These $SU(6)$ representations have the following decomposition in terms of $SU(5)$ representations: $\mathbf{15} = \mathbf{10} \oplus \mathbf{5}$ and $\bar{\mathbf{6}} = \bar{\mathbf{5}} \oplus \mathbf{1}$. The components of these $\mathbf{15}$, $\bar{\mathbf{6}}_1$ and $\bar{\mathbf{6}}_2$ multiplets decompose under $SU(3)_C \times SU(2)_W \times U(1)_Y$ as follows:

$$\begin{aligned} \mathbf{15} &\rightarrow Q = \left(3, 2, \frac{1}{6}\right), \\ t^c &= \left(3^*, 1, -\frac{2}{3}\right), & \bar{\mathbf{6}}_\alpha &\rightarrow D_\alpha^c = \left(\bar{3}, 1, \frac{1}{3}\right), \\ E^c &= \left(1, 1, 1\right), & L_\alpha &= \left(1, 2, -\frac{1}{2}\right), \\ D &= \left(3, 1, -\frac{1}{3}\right), & N_\alpha &= \left(1, 1, 0\right), \\ \bar{L} &= \left(1, 2, \frac{1}{2}\right); \end{aligned} \quad (5)$$

where $\alpha = 1, 2$. The first and second quantities in brackets are the $SU(3)_C$ and $SU(2)_W$ representations, while the third are the $U(1)_Y$ charges. The large mass of the top quark can be generated only if t^c is Z_2^B -odd. As a consequence all components of the $\mathbf{15}$ -plet have to be odd under the Z_2^B symmetry. After the $SU(6)$ symmetry breaking a $\mathbf{5}$ -plet from the $\mathbf{15}$ -plet and $\bar{\mathbf{5}}$ -plet from the $\bar{\mathbf{6}}_2$ form vector-like states. The corresponding mass terms are allowed if all components of $\bar{\mathbf{6}}_2$ are Z_2^B -odd. In principle, the components of $\bar{\mathbf{6}}_1$ multiplet could be either even or odd under the Z_2^B symmetry. Hereafter we assume that D_1^c , L_1 and N_1 are Z_2^B -even.

The breakdown of the $SU(6)$ symmetry also induces the Majorana masses for the SM singlet states N_1 and N_2 . The mixing of these states is suppressed because of the approximate Z_2^B symmetry. As discussed above, the remaining components of the $SU(6)$

multiplets $\mathbf{15}$ and $\bar{\mathbf{6}}_1$ get combined with \bar{q} , \bar{d}^c , $\bar{\ell}$ and \bar{e}^c leading to the composite right-handed top quark t^c and a set of vector-like states. In general all extra exotic fermions tend to gain masses which are a few times larger than f . Therefore it is unlikely that these states will be detected at the LHC in the near future. In the next section we consider the phenomenological implications of this variant of the E_6 CHM, assuming that N_1 is considerably lighter than other exotic fermion states and has a mass which is somewhat smaller than f .

3. Generation of baryon asymmetry

As mentioned earlier, the breakdown of the $SU(6)$ to its $SU(5)$ subgroup gives rise to a set of pNGB states. The masses of all pNGB states are expected to be considerably lower than $f \gtrsim 10$ TeV, so that these resonances are the lightest composite states. The corresponding set involves eleven pNGB states. One of them, ϕ_0 , is a real SM singlet scalar. The collider signatures associated with the presence of such a scalar, in the limit where CP is conserved, were examined in Ref. [42]. Ten other pNGB states form a fundamental representation of unbroken $SU(5)$, i.e. $\tilde{H} = \mathbf{5}$. The first two components of \tilde{H} transform as an $SU(2)_W$ doublet and correspond to the SM-like Higgs doublet, H , whereas three other components of \tilde{H} are associated with the $SU(3)_C$ triplet of scalar fields T . The pNGB effective potential $V_{eff}(\tilde{H}, T, \phi_0)$ is induced by the interactions of the elementary states with their composite partners that explicitly violate the $SU(6)$ global symmetry. In the model under consideration substantial tuning, $\sim 0.01\%$, is required to get $v \ll f$ and a 125 GeV Higgs boson, because the scale f is so large. Nevertheless, it has been shown that in models similar to the E_6 CHM there exists a part of the parameter space where the $SU(2)_W \times U(1)_Y$ gauge symmetry is broken to $U(1)_{em}$, whilst $SU(3)_C$ remains intact [9], [22]. In these composite Higgs models the $SU(3)_C$ triplet scalar, T , tends to be substantially heavier than the Higgs scalar.

In the interactions with other SM particles the Higgs boson manifests itself as a Z_2^B -even state. Therefore all other pNGB states should be also even under the Z_2^B symmetry. The gauge and Z_2^B symmetries allow the decays of the scalar triplet T into up and down antiquarks. At the same time, the decays of the $SU(3)_C$ scalar triplet into a charged lepton and an up quark as well as into a neutrino and a down quark are forbidden by the almost exact Z_2^L symmetry. Since the fractions of compositeness of the first and second generation quarks are rather small, the decay mode $T \rightarrow \bar{t}\bar{b}$ tends to be the dominant one. At the energies $E \lesssim f$ almost all resonances of the composite sector, except the pNGB states, can be integrated out and all baryon number violating operators are strongly suppressed, so that baryon number is conserved to a very good approximation. In this limit T manifests itself in the interactions with other SM bosons and fermions as a diquark, i.e. an $SU(3)_C$ scalar triplet with $B = -2/3$.⁴

The presence of this exotic $SU(3)_C$ scalar triplet with mass m_T in the few TeV range makes possible the generation of the baryon asymmetry via the out-of-equilibrium decays of N_1 , provided N_1 is the lightest exotic fermion in the spectrum.⁵ Indeed, the Majorana mass of N_1 is set by f , while $m_T \ll f$. As a result the decays

³ The composite nature of the right-handed top quark manifests itself through operators which are suppressed by inverse powers of the scale f . Since in the E_6 CHM $f \gtrsim 10$ TeV all couplings of t^c are expected to be extremely close to those predicted by the SM. Therefore it seems rather problematic to test such weak interactions at the LHC.

⁴ Because the couplings of the diquark T to the first and second generation quarks are very small, the LHC production cross section of single T is strongly suppressed. Nonetheless this $SU(3)_C$ scalar triplet can be pair produced at the LHC. In this case the process $pp \rightarrow T\bar{T}$ leads to a final state involving four heavy quarks $t\bar{t}b\bar{b}$. This signature is somewhat similar to that associated with the lightest squark in R-parity violating SUSY models. Recent LHC searches for such squarks set stringent lower bounds on their masses [43]. As a result, scenarios with the mass of the diquark T below 700 GeV are already strongly disfavoured.

⁵ In different contexts baryogenesis caused by the decays of the neutral fermion into quark and scalar diquark was discussed in [44].

$N_1 \rightarrow T + \bar{d}_i$ and $N_1 \rightarrow T^* + d_i$ are kinematically allowed. Since at low energies $E \lesssim f$ the $SU(3)_C$ scalar triplet, T , manifests itself in the interactions with other SM states as a diquark, the Majorana fermion N_1 can decay into final states with baryon numbers ± 1 . The interactions of N_1 and N_2 with the pNGB state T and down-type quarks are described by the Lagrangian

$$\mathcal{L}_N = \sum_{i=1}^3 \left(g_{i1}^* T d_i^c N_1 + g_{i2}^* T d_i^c N_2 + h.c. \right). \quad (6)$$

In the exact Z_2^B symmetry limit, the couplings g_{i1} have to vanish. Therefore the approximate Z_2^B symmetry ensures that the g_{i1} couplings are somewhat suppressed in comparison with g_{i2} , i.e. $|g_{i1}| \ll |g_{i2}|$.

The process of the baryon asymmetry generation is controlled by the flavour CP (decay) asymmetries $\varepsilon_{1,k}$ that appear on the right-hand side of Boltzmann equations. There are three decay asymmetries associated with three quark flavours d , s and b . These are given by

$$\varepsilon_{1,k} = \frac{\Gamma_{N_1 d_k} - \Gamma_{N_1 \bar{d}_k}}{\sum_m (\Gamma_{N_1 d_m} + \Gamma_{N_1 \bar{d}_m})}, \quad (7)$$

where $\Gamma_{N_1 d_k}$ and $\Gamma_{N_1 \bar{d}_k}$ are partial decay widths of $N_1 \rightarrow d_k + T^*$ and $N_1 \rightarrow \bar{d}_k + T$ with $k, m = 1, 2, 3$. At the tree level the CP asymmetries (7) vanish because (see [38])

$$\Gamma_{N_1 d_k} = \Gamma_{N_1 \bar{d}_k} = \frac{3|g_{k1}|^2}{32\pi} M_1, \quad (8)$$

where M_1 is the Majorana mass of N_1 . However, if CP invariance is broken the non-zero contributions to the CP asymmetries arise from the interference between the tree-level amplitudes of the N_1 decays and the one-loop corrections to them. The tree-level and one-loop diagrams that give contributions to the CP asymmetries associated with the decays $N_1 \rightarrow \bar{d}_k + T$ can be found in [38]. Assuming that the $SU(3)_C$ scalar triplet T is much lighter than N_1 and N_2 , the direct calculation of the appropriate one-loop diagrams gives⁶

$$\varepsilon_{1,i} = \frac{1}{(8\pi)} \frac{1}{(\sum_{m=1}^3 |g_{m1}|^2)} \times \left[\sum_{n=1}^3 \text{Im}(g_{i1}^* g_{i2} g_{n1}^* g_{n2}) \sqrt{x} \left(\frac{3}{2(1-x)} + 1 - (1+x) \ln \frac{1+x}{x} \right) + \sum_{n=1}^3 \text{Im}(g_{i1}^* g_{i2} g_{n1} g_{n2}^*) \frac{3}{2(1-x)} \right], \quad (9)$$

where $x = (M_2/M_1)^2$ and M_2 is the Majorana mass of N_2 .

In order to calculate the total baryon asymmetries induced by the decays of N_1 , the system of Boltzmann equations that describe the evolution of baryon number densities have to be solved. The corresponding solution should be somewhat similar to the solutions of the Boltzmann equations for thermal leptogenesis; so that in the first approximation the generated baryon asymmetry can be estimated using an approximate formula given in Ref. [46]⁷

$$Y_{\Delta B} \sim 10^{-3} \left(\sum_{k=1}^3 \varepsilon_{1,k} \eta_k \right), \quad (10)$$

where $Y_{\Delta B}$ is the baryon asymmetry relative to the entropy density, i.e.

$$Y_{\Delta B} = \frac{n_B - n_{\bar{B}}}{s} \Big|_0 = (8.75 \pm 0.23) \times 10^{-11}.$$

In Eq. (10) η_k are efficiency factors. A thermal population of N_1 decaying completely out of equilibrium without washout effects would lead to $\eta_k = 1$. However washout processes reduce the induced asymmetries by the factors η_k , where η_k varies from 0 to 1.

To simplify our analysis we assume that the pNGB state T couples primarily to the b -quarks, i.e. $|g_{31}| \gg |g_{21}|, |g_{11}|$ and $|g_{32}| \gg |g_{22}|, |g_{12}|$, and N_1 is substantially lighter than all other exotic fermions. In particular, we set $M_2 = 10 \cdot M_1$. The imposed hierarchical structure of the Yukawa couplings implies that the decay asymmetries $\varepsilon_{1,2}$ and $\varepsilon_{1,1}$ are much smaller than $\varepsilon_{1,3}$ and can be neglected. If $g_{31} = |g_{31}| e^{i\varphi_{31}}$ and $g_{32} = |g_{32}| e^{i\varphi_{32}}$ then in the limit $x \gg 1$ one finds

$$\varepsilon_{1,3} \simeq -\frac{1}{(4\pi)} \frac{|g_{32}|^2}{\sqrt{x}} \sin 2\Delta\varphi, \quad \Delta\varphi = \varphi_{32} - \varphi_{31}. \quad (11)$$

The CP asymmetry (11) vanishes when all Yukawa couplings are real, i.e. CP invariance is preserved. The decay asymmetry $\varepsilon_{1,3}$ attains its maximum absolute value when $\Delta\varphi = \pm\pi/4$, i.e. $\sin 2\Delta\varphi$ is equal to ± 1 .

In order to estimate the efficiency factor η_3 , we concentrate on the so-called strong washout scenario (see, for example [46]) for which

$$\eta_3 \simeq H(T = M_1) / \Gamma_3, \quad \Gamma_3 = \Gamma_{N_1 d_3} + \Gamma_{N_1 \bar{d}_3} = \frac{3|g_{31}|^2}{16\pi} M_1, \quad H = 1.66 g_*^{1/2} \frac{T^2}{M_{Pl}}, \quad (12)$$

where H is the Hubble expansion rate and $g_* = n_b + \frac{7}{8} n_f$ is the number of relativistic degrees of freedom in the thermal bath. Within the SM $g_* = 106.75$, whereas in the E_6 CHM $g_* = 113.75$ for $T \lesssim f$. Eqs. (12) indicate that η_3 increases with diminishing of $|g_{31}|$. Thus this coupling of N_1 to the pNGB state T can be adjusted so that η_3 becomes relatively close to unity. In particular, from Eqs. (12) it follows that for $|g_{31}| \simeq 10^{-6}$ and $M_1 \simeq 10$ TeV the efficiency factor η_3 is around 0.25.

If the efficiency factor η_3 is sufficiently large, i.e. $\eta_3 \sim 0.1-1$, the baryon asymmetry is determined by the induced decay asymmetry $\varepsilon_{1,3}$. Indeed, from Eqs. (9) and (11) one can see that in the limit $g_{21} = g_{11} \rightarrow 0$ the CP asymmetries $\varepsilon_{1,2}$ and $\varepsilon_{1,1}$ vanish while $\varepsilon_{1,3}$ does not depend on the absolute value of the Yukawa coupling g_{31} . Therefore, for a given ratio M_2/M_1 the CP asymmetry $\varepsilon_{1,3}$ is set by $|g_{32}|$ and the combination of the CP-violating phases $\Delta\varphi$. The dependence of the absolute value of $\varepsilon_{1,3}$ on these parameters is examined in Fig. 1, where we fix $(M_2/M_1) = 10$. Since the Yukawa coupling of N_2 to $SU(3)_C$ scalar triplet and b -quark is not suppressed by the Z_2^B symmetry, $|g_{32}|$ is expected to be relatively large, i.e. $|g_{32}| \gtrsim 0.1$. In Fig. 1a we plot the absolute value of $\varepsilon_{1,3}$ as a function of $\Delta\varphi$ for $|g_{32}| = 0.1$ and $|g_{32}| = 1$. Fig. 1a illustrates that the CP asymmetry $\varepsilon_{1,3}$ attains its maximum absolute value $\sim 10^{-4}-10^{-2}$ for $\Delta\varphi \simeq \pi/4$. Thus a larger value of $|\varepsilon_{1,3}|$ can lead to a phenomenologically acceptable baryon density only for sufficiently small values of efficiency factor, $\eta_3 = 10^{-5}-10^{-3}$. When this factor is reasonably large, i.e. $\eta_3 \sim 0.1-1$, and $|g_{32}| \simeq 0.1$ a phenomenologically acceptable value

⁶ These calculations are very similar to the ones performed in the case of thermal leptogenesis [45] (for the review see [46]).

⁷ The induced baryon asymmetry is partially converted into lepton asymmetry due to $(B+L)$ -violating sphaleron interactions [47]. Here we ignore sphaleron processes.

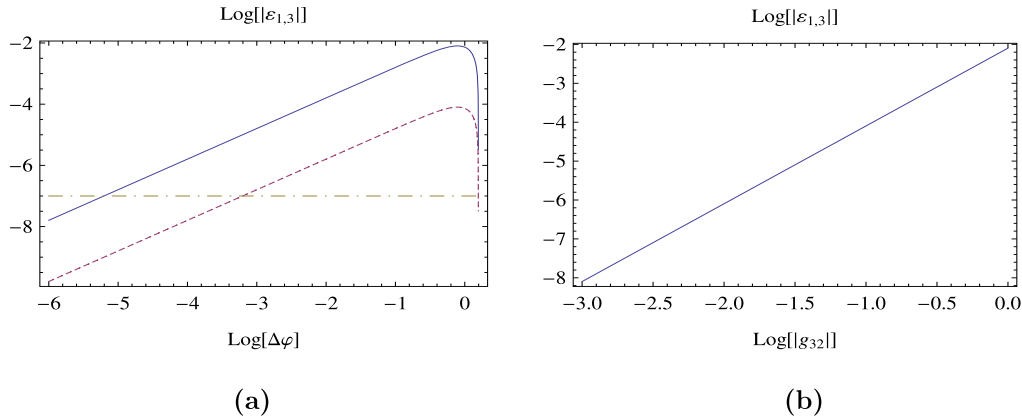


Fig. 1. Logarithm (base 10) of the absolute value of the CP asymmetry $\varepsilon_{1,3}$ for $g_{11} = g_{21} = g_{12} = g_{22} = 0$ and $M_2 = 10 \cdot M_1$. In (a) the absolute value of $\varepsilon_{1,3}$ is given as a function of logarithm (base 10) of $\Delta\varphi$ for $|g_{32}| = 1$ (solid line) and $|g_{32}| = 0.1$ (dashed line). In (b) the maximal absolute value of $\varepsilon_{1,3}$ is presented as a function of logarithm (base 10) of $|g_{32}|$ for $\Delta\varphi = \pi/4$.

of the baryon density, corresponding to $\varepsilon_{1,3} \lesssim 10^{-7} - 10^{-6}$, is generated only if the combination of the CP-violating phases $\Delta\varphi$ is rather small, i.e. $\Delta\varphi \lesssim 0.01$. This demonstrates that the appropriate baryon asymmetry can be obtained within the E_6 CHM even if CP is approximately preserved.

In Fig. 1b the dependence of the maximum value of $|\varepsilon_{1,3}|$ on $|g_{32}|$ is studied. From Eq. (11) and Fig. 1b it follows that the maximum absolute value of this CP asymmetry grows monotonically with increasing of $|g_{32}|$. Fig. 1b also indicates that the appropriate baryon density associated with $\varepsilon_{1,3} \gtrsim 10^{-7} - 10^{-6}$ can be obtained even if the absolute value of the corresponding Yukawa coupling varies from 0.01 to 0.1.

4. Conclusions

In the E_6 inspired composite Higgs model (E_6 CHM) the approximate $SU(6)$ global symmetry of the strongly coupled sector is supposed to be broken down at the scale $f \gtrsim 10$ TeV to its $SU(5)$ subgroup, which incorporates the $SU(3)_C \times SU(2)_W \times U(1)_Y$ gauge symmetry. Within this model the operators that may result in rapid proton decay can be suppressed by a Z_2^L discrete symmetry. Since the scale f is so large all baryon number violating operators, which are not forbidden by the Z_2^L symmetry, are sufficiently strongly suppressed. Nonetheless, this variant of the E_6 CHM leads to baryon number violating processes, like neutron-antineutron oscillations, that are going to be searched for in future experiments [39,40]. To ensure the approximate unification of the SM gauge couplings, that makes possible the embedding of the E_6 CHM into a suitable GUT, this model involves extra matter. Additional matter multiplets give rise to a composite right-handed top quark and a set of exotic fermions that, in particular, includes two SM singlet Majorana states N_1 and N_2 . In general all exotic fermions acquire masses which are somewhat larger than f . In our analysis we assumed that N_1 is the lightest exotic fermion, with a mass around 10 TeV.

The pNGB states, which originate from the breakdown of $SU(6)$ to its $SU(5)$ subgroup, are the lightest composite resonances in the E_6 CHM. The corresponding set of states contains one SM singlet scalar, a SM-like Higgs doublet and an $SU(3)_C$ triplet of scalar fields, T . The masses of all these resonances tend to be substantially lower than f . At energies $E \lesssim f$ baryon number is preserved to a very good approximation and the $SU(3)_C$ scalar triplet T manifests itself in the interactions with the SM particles as a diquark. We argued that in this variant of the E_6 CHM the baryon asymmetry can be generated via the out-of equilibrium decays of N_1 into final states with baryon numbers ± 1 , i.e., $N_1 \rightarrow T + \bar{d}_i$ and

$N_1 \rightarrow T^* + d_i$, provided CP is violated. Moreover, if the absolute value of the Yukawa coupling of N_2 to T and b -quark varies in the range 0.1 to 1 a phenomenologically acceptable baryon density may be obtained, even when all CP-violating phases are quite small ($\lesssim 0.01$). In this case the approximate CP conservation leads to suppression of the electric dipole moments (EDMs) of the neutron, elementary states and atoms that have not been observed in numerous experiments but can be measured in the near future (see [40]). Since the couplings of N_1 , N_2 and T to the first and second generation quarks are tiny, their contributions to the baryon number violating processes, like $n - \bar{n}$ oscillations, are sufficiently strongly suppressed. On the other hand, the $SU(3)_C$ scalar triplet T , with mass in the few TeV range, can be pair produced at the LHC and predominantly decays into $T \rightarrow \bar{t} + \bar{b}_i$, leading to some enhancement of the cross section of $pp \rightarrow t\bar{t}b\bar{b}$. Thus the scenario under consideration emphasises the importance of the complementarity of different experiments.

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