universe

MDPI

Review

E_6 GUT and Baryon Asymmetry Generation in the E_6CHM

Roman Nevzorov

I.E.Tamm Theory Department, P.N. Lebedev Physical Institute of the Russian Academy of Sciences, 53 Leninskiy Prospekt, 119991 Moscow, Russia; nevzorovrb@lebedev.ru

Abstract: Grand unified theories (GUTs) may result in the E_6-inspired composite Higgs model (E_6CHM) at low energies, almost stabilizing the electroweak scale. We consider an orbifold GUT in 6 dimensions in which the E_6-gauge group is broken to the gauge symmetry of the standard model (SM) while different multiplets of the SM fermions come from different 27-plets. The strongly coupled sector of the E_6CHM is confined on the brane where E_6 is broken down to its $SU(6)$ subgroup. Near the scale of $f \gtrsim 5$ TeV, this approximate $SU(6)$ symmetry is expected to be further broken down to its $SU(5)$ subgroup, which contains the SM-gauge group. Such a breakdown leads to a set of pseudo-Nambu–Goldstone bosons (pNGBs) that includes an SM-like Higgs doublet. The approximate gauge coupling unification in the E_6CHM takes place at high energies when the right-handed top quark is a composite fermion. To ensure anomaly cancellation, the weakly coupled sector of this model contains extra exotic matter beyond the SM. We discuss the mechanism of the generation of matter–antimatter asymmetry within the variant of the E_6CHM in which the baryon number and CP invariance are violated.

Keywords: unified field theories and models; field theories in dimensions other than four; models beyond the standard model; models of the early universe; composite Higgs; baryon asymmetry generation

check for updates

Citation: Nevzorov, R. E_6 GUT and Baryon Asymmetry Generation in the E_6CHM. *Universe* **2022**, *8*, 33. https://doi.org/10.3390/universe8010033

Academic Editors: Ricardo G. Landim and Marcelo M. Guzzo

Received: 30 November 2021
Accepted: 1 January 2022
Published: 5 January 2022

1. Introduction

It is well known that the standard model (SM) of elementary particles, which involves all known fundamental bosons and fermions, describes rather precisely the major part of all experimental data. At very high energies, the SM can be embedded into grand unified theories (GUTs) [1]. In the simplest GUTs based on the $SU(5)$-gauge group, each SM family of fermions is composed of one antifundamental and one antisymmetric second-rank tensor representation of $SU(5)$, i.e., $\overline{5} + 10$. In the case of the $SO(10)$ GUTs, each family of quarks and leptons fills in a complete single 16-dimensional spinor representation of $SO(10)$. This representation also contains the right-handed neutrino, which may be used for the see-saw mechanism [2,3].

Supersymmetry (SUSY) implies that each supermultiplet includes the same number of bosonic and fermionic degrees of freedom. In $N = 1$ SUSY GUTs with the E_6-gauge group, the fundamental 27 representation of E_6 decomposes under the $SO(10) \times U(1)_\psi$ subgroup as

$$27 \to \left(16, \frac{1}{\sqrt{24}}\right) \oplus \left(10, -\frac{2}{\sqrt{24}}\right) \oplus \left(1, \frac{4}{\sqrt{24}}\right), \tag{1}$$

where the first and second quantities in brackets are the $SO(10)$ representation and its $U(1)_\psi$ charge. As before, the supermultiplet $\left(16, \frac{1}{\sqrt{24}}\right)$ can include one family of quarks and leptons. The doublet of the Higgs bosons may form components of the supermultiplet $\left(10, -\frac{2}{\sqrt{24}}\right)$. The SM-gauge bosons are assigned to the adjoint representation of E_6, i.e., a 78-plet. In $N = 2$ SUSY GUTs based on the E_8-gauge symmetry, all SM bosons and SM fermions may belong to a single 248 representation of E_8 which decomposes under the E_6 subgroup of E_8, as follows:

$$248 \to 78 \oplus 3 \times 27 \oplus 3 \times \overline{27} \oplus 8 \times 1\,. \tag{2}$$

In Equation (2), 3 generations of the SM fermions can be associated with 3 27-plets which may also contain the doublet of the Higgs bosons, while some components of the 78-plet may form the multiplets of the SM-gauge bosons.

The breakdown of gauge symmetry within the SUSY GUTs near some high energy scale $M_X \gtrsim 10^{16}$–10^{17}GeV can result in the gauge group and field content of the SM. In this case, below the scale of M_X, the Higgs scalar potential takes the form

$$V(H) = m_H^2\, H^\dagger H + \lambda\, (H^\dagger H)^2 + \dots\,. \tag{3}$$

In order to ensure that, at low energies, the doublet of the Higgs fields acquires vacuum expectation value (VEV) $\langle H\rangle = v/\sqrt{2} \simeq 174$ GeV breaking the electroweak (EW) symmetry, $|m_H^2|$ is required to be of the order of $(100\ \text{GeV})^2$. On the other hand, most commonly, $|m_H^2|$ is about M_X^2, and an enormous amount of fine tuning is needed to keep $|m_H^2| \sim (100\ \text{GeV})^2$.

Such enormous fine tuning can be avoided if the breakdown of gauge symmetry in SUSY GUTs leads to the extension of the SM with softly broken supersymmetry. The cancellation of quadratic divergences [4–7] within the minimal supersymmetric standard model (MSSM) stabilises the EW scale, solving the hierarchy problem [8,9] (for a review see [10]). $N = 1$ SUSY also facilitates the high-energy convergence of the SM-gauge couplings [11–14] which allows the SM-gauge group, i.e., $SU(3)_C \times SU(2)_W \times U(1)_Y$, to be embedded into SUSY GUTs. Theories with flat [15,16] and warped [17,18] extra spatial dimensions provide new insight into gauge coupling unification [19,20] and also permit the hierarchy between the EW and Planck scales to be explained.

Alternatively, the Higgs boson can be a composite state. Composite Higgs models include two sectors (for a review, see ref. [21]). One of them involves weakly-coupled elementary particles with the quantum numbers of all SM-gauge bosons and SM fermions. The second strongly coupled sector gives rise to a set of bound states that, in particular, contains Higgs doublet. The corresponding idea was proposed in the 1970s [22,23] and 1980s [24–31]. This implies that the EW scale is generated dynamically in a strongly interacting sector, in analogy with the origin of the QCD scale. In general, these models lead to a relatively large quartic coupling λ at the EW scale, and the composite Higgs state tends to be quite heavy. The rather small values of the parameters $\lambda \approx 0.13$ and $m_H^2 \approx -(90\ \text{GeV})^2$ in Equation (3), which are associated with the measured Higgs mass $m_h \simeq 125$–126 GeV, indicate that the Higgs doublet may emerge as a set of pseudo-Nambu–Goldstone bosons (pNGB). The appearance of such pNGB states can be caused by the spontaneous breakdown of an approximate global symmetry of the strongly coupled sector.

In SUSY GUTs with the $E_8 \times G_0$ (or $E_6 \times G_0$)-gauge group, the breakdown of gauge symmetry at high energies down to the $SU(3)_C \times SU(2)_W \times U(1)_Y \times G$ subgroup may lead to the $SU(6)$ global symmetry in the strongly interacting sector [32–37]. Hereafter, it is assumed that fields, which compose the strongly coupled sector, can be charged under both the E_8 (E_6)- and G_0 (G)-gauge symmetries, whereas the elementary states belonging to the weakly coupled sector participate in the E_8 (E_6) interactions only. The scenario mentioned above is realised if, for instance, in the strongly interacting sector, E_8 is broken to E_6, with sequential breakdown of E_6 to its $SU(6)$ subgroup near the GUT scale M_X. The spontaneous $SU(6)$ symmetry, breaking at much lower energies to $SU(5)$, which contains the SM-gauge group, gives rise to the 11 pNGBs in this E_6-inspired composite Higgs model (E_6CHM) [32–37]. The corresponding set of the pNGBs involves the Higgs doublet. It is worth noting that the E_6CHM can not appear as a low energy limit of the heterotic superstring theory with $E_8 \times E_8'$-gauge symmetry. Some phenomenological consequences of the heterotic string model were discussed in [38,39].

This review paper is organised as follows. In the next Section, we briefly review the composite Higgs models and specify the E_6CHM. To suppress the proton decay rate and the Majorana masses of the left-handed neutrino within the E_6CHM, the elementary fermions with different baryon and lepton numbers should stem from different fundamental representations of E_6, whereas all other components of the corresponding 27-plets are expected to gain masses of the order of M_X. In this context, in Section 3, we present a six-dimensional (6D) orbifold GUT model based on the $E_6 \times G_0$-gauge group in which the appropriate splitting of the fundamental representations of E_6 can be achieved. The observed baryon asymmetry in the universe stimulates the exploration of different extensions of the SM. This asymmetry can be created dynamically within the scenarios satisfying Sakharov conditions [40]. A number of such new physics scenarios were proposed, including GUT baryogenesis [41–47], baryogenesis via leptogenesis [48], the Affleck–Dine mechanism [49,50], electroweak baryogenesis [51], etc. In Section 4 we consider the process of the baryon asymmetry generation in the framework of the E_6CHM with explicitly broken $U(1)_B$ baryon symmetry. The sizeable baryon number asymmetry can be induced in this model if CP is violated. Section 5 concludes the paper.

2. Composite Higgs Models and E_6CHM

2.1. Composite Higgs Models—A Brief Review

The strongly interacting sector of the minimal composite Higgs model (MCHM) possesses a global SO(5) $\times U(1)_X$ symmetry [52]. It is expected that, near the scale $f \sim$ 1–10 TeV, this global symmetry is broken down to $SO(4) \times U(1)'_X \cong SU(2)_W \times SU(2)_R \times U(1)'_X$, which includes the SU(2)$_W \times U(1)_Y$-gauge group as a subgroup. Such a breakdown gives rise to a set of pNGBs which form the Higgs doublet. Via the AdS/CFT correspondence, such composite Higgs scenarios are dual to the Randall–Sundrum (RS) extra-dimensional scenarios, with the SM fields in the bulk [52,53]. In the RS scenarios, Kaluza–Klein excitations of the SM fields are associated with the bound states at the compositeness scale, f [52–55].

Thus, the strongly coupled sector in the composite Higgs models should result in a set of massive fields with quantum numbers of all SM particles, which are the so-called composite partners of the SM states. The elementary states from the weakly coupled sector mix with their composite partners. Therefore, at low energies, those states identified with SM fermions (bosons) are superpositions of the elementary fermionic (bosonic) states and their fermionic (bosonic) composite partners. Such partial compositeness [55,56] implies that the SM states couple to the composite Higgs with a strength which is determined by the fraction of the compositeness of this state. As consequence, the effective up- and down-quark Yukawa couplings (y^u_{ij} and y^d_{ij}) are given by

$$y^u_{ij} = s^i_q Y^u_{ij} s^j_u \,, \qquad y^d_{ij} = s^i_q Y^d_{ij} s^j_d \,, \tag{4}$$

where s^i_q, s^j_u and s^j_d are the fractions of compositeness of the left-handed SM quarks as well as the right-handed SM quarks of up- and down-types, whereas $i, j = 1, 2, 3$. In Equation (4), Y^u_{ij} and Y^d_{ij} are the effective Yukawa couplings of the composite Higgs field to the composite partners of the up- and down-quarks. The couplings of the elementary states to the operators of the strongly interacting sector explicitly break its global symmetry. In these models, the Higgs potential arises from loops containing elementary states. This results in the suppression of the effective quartic Higgs coupling. The contributions of the composite partners of the SM states to the EW observables, including the $\hat{S}$ and $\hat{T}$ parameters, were examined in [57–84]. Within the MCHM, the custodial symmetry $SU(2)_{cust} \subset SO(4) \cong SU(2)_W \times SU(2)_R$ [85] protects the Peskin–Takeuchi $\hat{T}$ parameter [86] against the contributions of new composite states.

In the phenomenologically viable composite Higgs models, the fractions of compositeness of the first and second generation fermions should be rather small. If this is the case, the corresponding states have small couplings to the Higgs doublet and therefore tend to be light. In other words, the observed mass hierarchy in the quark and lepton sectors can be reproduced if the couplings of the elementary fermions associated with the first and second generations to the states from the strongly interacting sector are very weak. Such weak couplings result in some suppression of flavour-changing processes and modifications of the W and Z couplings [55,87], playing the role of the generalization of the Glashow–Iliopoulos–Maiani (GIM) mechanism of the SM [88]. Although this generalization of the GIM mechanism reduces the contributions of new composite states to the off-diagonal flavour transitions in the quark and lepton sectors, this suppression is not sufficient. To avoid dangerous flavour-changing processes, the composite Higgs models have to satisfy a set of constraints which were examined in [81–84,89–97]. If the matrices of effective Yukawa couplings in the strongly interacting sector, such as Y^u_{ij} and Y^d_{ij}, are structureless, then the adequate suppression of the non-diagonal flavour transitions can be achieved only if f is larger than 10 TeV [81–83,89–91,94,95]. At the same time, in the composite Higgs models with flavour symmetries [79–81,89,92,93,98–100], under which the third-generation elementary fermions transform as singlets while the first two generations of elementary fermions form different $U(2)$ doublets, the corresponding constraints can be fulfilled even if $f \gtrsim 1$ TeV [92,93].

When $f \ll 10$ TeV, approximate $U(1)_B$ and $U(1)_L$ symmetries, which ensure the conservation of the baryon and lepton numbers, should be imposed in the strongly interacting sector of the composite Higgs models. These symmetries are needed to suppress the operators that give rise to the Majorana masses of the left-handed neutrinos and the baryon number violation. The implications of the composite Higgs models were considered for Higgs physics [73–78,101–124], gauge coupling unification [125,126], dark matter [57,107,127,128] and collider phenomenology [70–73,79,80,89,93,123,124,129–153]. Non-minimal composite Higgs models were explored in [57,101–107,127,128,154–163].

Since the top quark is rather heavy, the left-handed and right-handed top quarks (t and t^c) should have substantial fractions of compositeness. In the case when t^c is in an entirely composite state, the approximate unification of the SM-gauge couplings in the composite Higgs models may take place if all multiplets in the strongly coupled sector form complete representations of $SU(5)$, while the weakly coupled sector includes the following set of matter multiplets [32,164]:

$$(q_i,\, d^c_i,\, \ell_i,\, e^c_i) + u^c_\alpha + \bar{q} + \bar{d^c} + \bar{\ell} + \bar{e^c}\,, \tag{5}$$

where $i = 1, 2, 3$ runs over all 3 generations and $\alpha = 1, 2$ runs over the first 2 generations. Here, we have denoted the right-handed charged leptons and the right-handed down- and up-type quarks by e^c_i, d^c_i and u^c_α, whereas ℓ_i and q_i are associated with the left-handed lepton and quark doublets. In Equation (5), $\bar{e^c}$, $\bar{d^c}$, $\bar{q}$ and $\bar{\ell}$ correspond to the exotic fermions which have opposite $SU(3)_C \times SU(2)_W \times U(1)_Y$ quantum numbers compared to the right-handed charged leptons, right-handed down-type quarks, as well as left-handed quark and left-handed lepton doublets, respectively. The set of elementary states (5) contains all SM fermions except t^c. The particle content of the weakly coupled sector is chosen so that anomaly cancellation takes place.

The phenomenological viability of such composite Higgs models implies that the strongly coupled sector leads to a set of composite fermions that form $\mathbf{10} + \mathbf{\bar{5}}$ multiplets of $SU(5)$. All of them, except the components of the 10-plet associated with the composite t^c, get combined with $\bar{q}$, $\bar{d^c}$, $\bar{\ell}$ and $\bar{e^c}$, composing vector-like states. The composite $SU(3)_C$ triplet identified with t^c survives to the EW scale.

The presence of exotic vector-like fermions facilitates the convergence of the SM-gauge couplings at high energies. In the one-loop approximation, the renormalisation group (RG) flow of the SM-gauge couplings is described by a system of RG equations (RGEs), which can be written in the following form:

$$\frac{d\alpha_i}{dt} = \frac{\beta_i \alpha_i^2}{(2\pi)}\,, \tag{6}$$

where b_i are one-loop beta functions with the index i running from 1 to 3, corresponding to $U(1)_Y$, $SU(2)_W$ and $SU(3)_C$ interactions, $t = \ln(\mu/M_Z)$ and μ is a renormalisation scale. Then, using the solutions of the RGEs (6), one can find $\alpha_3(M_Z)$, for which the exact gauge coupling unification takes place

$$\frac{1}{\alpha_3(M_Z)} = \frac{1}{b_1 - b_2}\left[\frac{b_1 - b_3}{\alpha_2(M_Z)} - \frac{b_2 - b_3}{\alpha_1(M_Z)}\right]. \tag{7}$$

If all bound states in the strongly interacting sector compose complete $SU(5)$ multiplets, they contribute equally to b_i. Because of this, the inclusion of the composite sector fields does not change the differential running. In other words, in the one-loop approximation, $(b_i - b_j)$ are determined by the particle content of the weakly coupled sector (5). Then, from Equation (7), it follows that for $\sin^2\theta_W = 0.231$, $\alpha(M_Z) = 1/127.9$ and the values of $(b_i - b_j)$ corresponding to the elementary particle spectrum (5), the exact gauge coupling unification may be obtained if $\alpha_3(M_Z) \simeq 0.109$. Such unification of the SM-gauge couplings takes place near the scale $M_X \sim 10^{15}$–10^{16} GeV. Despite the fact that $\alpha_3(M_Z) \simeq 0.109$ is considerably lower than the central measured value of this coupling, this estimation indicates that, in the composite Higgs model with a composite t^c, an approximate unification of the SM-gauge couplings can be attained.

2.2. E_6CHM

Hereafter, we assume that the weakly coupled sector of the E_6CHM involves all elementary states specified in Equation (5). Since the strongly coupled sector of the E_6CHM possesses an approximate global $SU(6)$ symmetry, which is expected to be broken down to its $SU(5)$ subgroup near some scale $f \gg v \simeq 246$ GeV so that the $SU(3)_C \times SU(2)_W \times U(1)_Y$-gauge symmetry remains intact, all composite states must come in complete $SU(5)$ multiplets. Therefore, an approximate gauge coupling unification may be achieved. Because the Lagrangian of the strongly interacting sector of the E_6CHM does not possess any custodial symmetry that may protect the Peskin–Takeuchi $\hat{T}$ parameter against contributions of extra composite states, $|\hat{T}|$ should be of the order $\xi \simeq v^2/f^2$ [57]. In this case, the electroweak precision measurements, which constrain $|\hat{T}| \lesssim 0.002$, result in the lower bound

$$f \gtrsim 5\,\text{TeV}. \tag{8}$$

In the model under consideration, more stringent restrictions on the scale f can be avoided. Indeed, the non-diagonal flavour transitions can be suppressed by imposing approximate flavour symmetry. Due to the mixing between the elementary states and their composite partners, the interactions in the strongly coupled sector may also induce the dimension-5 operators of the form $\ell_i \ell_j HH/f$, which give rise to overly large Majorana neutrino masses, as well as a set of baryon number-violating operators. All these operators are suppressed by the small fractions of compositeness of the SM fermions and by the relatively large scale f. Nevertheless such suppression is not sufficient if $f \ll 10$ TeV. The baryon- and lepton number-violating operators can be forbidden by postulating the conservation of baryon and lepton numbers in the E_6CHM. In principle, the corresponding $U(1)_B$ and $U(1)_L$ symmetries can be part of the symmetries of the composite sector. The G_0-gauge symmetry associated with the strongly coupled sector might be broken down to its subgroup G so that the $U(1)_B$ and $U(1)_L$ symmetries are preserved to very good approximations. As a consequence, at low energies, the Lagrangian of the strongly interacting sector of the E_6CHM respects the approximate $SU(6) \times U(1)_B \times U(1)_L$ global symmetry.

The global $U(1)_L$ symmetry has to be broken down to

$$Z_2^L = (-1)^L, \tag{9}$$

where L is a lepton number. This breakdown allows the left-handed neutrinos to gain non-zero Majorana masses. When Z_2^L remains in almost exact discrete symmetry, it forbids all operators that lead to rapid proton decay.

Near the scale f, the approximate global $SU(6)$ symmetry of the strongly coupled sector is broken down to $SU(5)$ in the E_6CHM. The $SU(6)$ and $SU(5)$ groups have 35 and 24 generators, t^a, respectively, which are normalised so that $\text{Tr}\, t^a t^b = \frac{1}{2}\delta_{ab}$. Here, we denote the 11 broken generators from the coset $SU(6)/SU(5)$ by $T^{\hat{a}}$. The generators of the unbroken $SU(5)$ subgroup of $SU(6)$ are denoted by T^a. The 11 pNGB states can be parameterised in terms of a 6-component unit vector Ω [32]

$$\begin{aligned} \Omega^T &= \Omega_0^T \Sigma^T = e^{i\frac{\phi_0}{\sqrt{15}f}} \left(C\phi_1 \quad C\phi_2 \quad C\phi_3 \quad C\phi_4 \quad C\phi_5 \quad \cos\frac{\tilde{\phi}}{\sqrt{2}f} + \sqrt{\frac{3}{10}} C\phi_0 \right), \\ C &= \frac{i}{\tilde{\phi}} \sin\frac{\tilde{\phi}}{\sqrt{2}f}, \qquad \tilde{\phi} = \sqrt{\frac{3}{10}\phi_0^2 + |\phi_1|^2 + |\phi_2|^2 + |\phi_3|^2 + |\phi_4|^2 + |\phi_5|^2}, \end{aligned} \tag{10}$$

where

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

In Equation (10), ϕ_0 is a real field, whereas $\phi_1\,\phi_2\,\phi_3\,\phi_4$ and ϕ_5 are complex fields. Taking into account that ϕ_0 and $\tilde{\phi}$ are invariant under the $SU(5)$ symmetry transformations, vector Ω can be decomposed into $\mathbf{5} + \mathbf{1}$ under the unbroken $SU(5)$ symmetry. Thus, it is convenient to introduce a 5-component multiplet $\tilde{H} = (\phi_1, \phi_2, \phi_3, \phi_4, \phi_5)$ and $A = \phi_0$, which is a SM singlet field. The first two components of $\tilde{H}$ compose an $SU(2)_W$ doublet. Therefore these components can be associated with the SM-like Higgs doublet H. A total of 3 other components, $T = (\phi_3, \phi_4, \phi_5)$, form an $SU(3)_C$ triplet. Since in the SM the Higgs doublet does not carry any baryon and/or lepton numbers (B and L), all components of the vector Ω have $B = L = 0$. The Lagrangian that describes the interactions of these pNGBs is given by

$$\mathcal{L}_{pNGB} = \frac{f^2}{2} \left| \mathcal{D}_\mu \Omega \right|^2. \tag{11}$$

Integrating out composite partners of the SM states and exotic fermions, one can obtain the pNGB effective potential $V_{eff}(\tilde{H}, T, \phi_0)$. This potential is induced by the interactions of the SM states with their composite partners that break $SU(6)$ symmetry. In the exact $SU(6)$ symmetry limit, it vanishes. The investigation of the pNGB potentials within similar models revealed that there exists a large part of the parameter space where the EW symmetry is broken, while $SU(3)_C$ is preserved [57,127]. Nevertheless, a significant tuning, ∼0.01%, is required in order to get a 125 GeV Higgs state in E_6CHM because $f \gtrsim 5$ TeV. It was shown that the appropriate quadratic term $m_H^2 |H|^2$ in the pNGB effective potential can be induced [127]. The analysis performed in the models, which are similar to the E_6CHM, indicated that, in the corresponding part of the parameter space, the $SU(3)_C$ triplet scalar T is considerably heavier than the SM-like Higgs boson.

As mentioned before, the weakly coupled sector of the E_6CHM includes a set of elementary states (5), whereas the right-handed top quark t^c is a composite state. Such a scenario implies that the dynamics of the strongly interacting sector results in the formation of the composite $\mathbf{10} + \overline{\mathbf{5}}$ multiplets of $SU(5)$. These $SU(5)$ multiplets get combined with $\bar{\ell}$, $\bar{e}^c$, $\bar{q}$ and $\bar{d}^c$, leading to a set of massive vector-like fermions as well as composite t^c. The composite $\mathbf{10} + \overline{\mathbf{5}}$ multiplets of $SU(5)$ may originate from two $\overline{\mathbf{6}}$-plets ($\overline{\mathbf{6}}_1$ and $\overline{\mathbf{6}}_2$) and one $\mathbf{15}$-plet of $SU(6)$. The $\mathbf{15}$-plet and $\overline{\mathbf{6}}$-plet have the following decomposition in terms of $SU(5)$ representations:

$$\mathbf{15} \to \mathbf{10} \oplus \mathbf{5}, \qquad \overline{\mathbf{6}} \to \overline{\mathbf{5}} \oplus \mathbf{1}.$$

The components of $\overline{\mathbf{6}}_1$, $\overline{\mathbf{6}}_2$ and $\mathbf{15}$ decompose under $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_B \times U(1)_L$ as follows:

$$
\begin{aligned}
\mathbf{15} \quad \to \quad & Q = \left(3,\, 2,\, \frac{1}{6},\, -\frac{1}{3},\, 0\right), & \overline{\mathbf{6}}_2 \quad \to \quad & D^c_2 = \left(\bar{3},\, 1,\, \frac{1}{3},\, \frac{1}{3},\, 0\right), \\
& t^c = \left(3^*,\, 1,\, -\frac{2}{3},\, -\frac{1}{3},\, 0\right), & & L_2 = \left(1,\, 2,\, -\frac{1}{2},\, \frac{1}{3},\, 0\right), \\
& E^c = \left(1,\, 1,\, 1,\, -\frac{1}{3},\, 0\right), & & N_2 = \left(1,\, 1,\, 0,\, \frac{1}{3},\, 0\right), \\
& D = \left(3,\, 1,\, -\frac{1}{3},\, -\frac{1}{3},\, 0\right), & \overline{\mathbf{6}}_1 \quad \to \quad & D^c_1 = \left(\bar{3},\, 1,\, \frac{1}{3},\, B_{\overline{6}_1},\, L_{\overline{6}_1}\right), \\
& \overline{L} = \left(1,\, 2,\, \frac{1}{2},\, -\frac{1}{3},\, 0\right); & & L_1 = \left(1, 2, -\frac{1}{2}, B_{\overline{6}_1}, L_{\overline{6}_1}\right), \\
& & & N_1 = \left(1,\, 1,\, 0,\, B_{\overline{6}_1},\, L_{\overline{6}_1}\right).
\end{aligned}
\tag{12}
$$

In Equation (12), 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 the $U(1)_Y$, $U(1)_B$ and $U(1)_L$ charges, respectively. Since the right-handed top quark belongs to the **15**-plet, all components of this multiplet should carry the same baryon and lepton numbers as t^c, i.e., $B_{15} = -1/3$ and $L_{15} = 0$. After the $SU(6)$ symmetry breaking, a $\overline{\mathbf{5}}$-plet from the $\overline{\mathbf{6}}_2$ and **5**-plet from the **15**-plet should compose vector-like states. This can be possible only if $B_{\overline{6}_2} = 1/3$ and $L_{\overline{6}_2} = 0$. Although the baryon and lepton numbers of the components of the $\overline{\mathbf{6}}_1$ multiplet are not fixed in the E_6CHM, the $SU(5)$ singlet components of $\overline{\mathbf{6}}_1$ and $\overline{\mathbf{6}}_2$ may gain mass through the interaction $(\overline{\mathbf{6}}_1\Omega)(\Omega\overline{\mathbf{6}}_2)$ if $B_{\overline{6}_1} = -1/3$ and $L_{\overline{6}_1} = 0$.

As pointed out before, in the composite Higgs models, the elementary fermions acquire masses through mixing with their composite partners. From the conservation of baryon and lepton numbers, it follows that in the E_6CHM, different multiplets of elementary fermions should come from different representations of the GUT-gauge group. All other components of the corresponding GUT multiplets have to be extremely heavy. Therefore, elementary fermions appear at low energies as incomplete GUT multiplets. In the case of the simplest $SU(5)$ GUT, the elementary fermions constitute the following set of incomplete $SU(5)$ multiplets:

$$
\begin{gathered}
u^c_\alpha \in \mathbf{10}^u_\alpha = \left(\mathbf{10},\, -\frac{1}{3},\, 0\right)_\alpha \quad q_i \in \mathbf{10}^q_i = \left(\mathbf{10},\, \frac{1}{3},\, 0\right)_i \quad d^c_i \in \overline{\mathbf{5}}^d_i = \left(\overline{\mathbf{5}},\, -\frac{1}{3},\, 0\right)_i \\
e^c_i \in \mathbf{10}^e_i = \left(\mathbf{10},\, 0,\, -1\right)_i \qquad \ell_i \in \overline{\mathbf{5}}^\ell_i = \left(\overline{\mathbf{5}},\, 0,\, 1\right)_i,
\end{gathered}
\tag{13}
$$

where the first, second and third quantities in brackets are the $SU(5)$ representation and the $U(1)_B$ and $U(1)_L$ charges. In Equation (13), $\alpha = 1, 2$ and $i = 1, 2, 3$. The Higgs doublet h is normally embedded into the fundamental representation of $SU(5)$, i.e., $h \in \mathbf{5}^h$. In this scenario, the Yukawa interactions of the SM, which induce the masses of the up-type quarks at low energies, have the following $SU(5)$ structure:

$$
\mathcal{L}^u_{SU(5)} \simeq h^u_{\alpha i}\, \mathbf{10}^u_\alpha\, \mathbf{10}^q_i\, \mathbf{5}^h. \tag{14}
$$

In the simplest $SU(5)$ models, the masses of the charged leptons and down-type quarks are generated through the Yukawa interactions

$$
\mathcal{L}^d_{SU(5)} \simeq h^e_{ij}\, \mathbf{10}^e_i\, \overline{\mathbf{5}}^\ell_\mathbf{j}\, \overline{\mathbf{5}}^h + h^d_{ij}\, \mathbf{10}^q_i\, \overline{\mathbf{5}}^\mathbf{d}_\mathbf{j}\, \overline{\mathbf{5}}^h\,. \tag{15}
$$

The composite partners of the elementary quarks and leptons must be embedded into the representations of the $SU(6)$ group so that the Yukawa interactions (14) and (15) are allowed. In this case, the Higgs multiplet $\mathbf{5}^h$ has to be replaced by the unit vector Ω. Moreover, instead of $\mathbf{10}^u_\alpha$ and $\mathbf{10}^q_i$ of $SU(5)$, one needs to include two $SU(6)$ multiplets which involve an $SU(5)$ decuplet. The simplest $SU(6)$ multiplet of this type is an antisymmetric second-rank tensor field **15**. The next-simplest $SU(6)$ representation that contains an $SU(5)$ decuplet is a totally antisymmetric third-rank tensor **20** that has the following decomposition in terms of $SU(5)$ representations: $\mathbf{20} = \mathbf{10} \oplus \overline{\mathbf{10}}$. The generalisation of the Yukawa interaction (14) to the case of $SU(6)$ symmetry can be written as

$$\mathcal{L}^u_{SU(6)} \sim \mathbf{20} \times \mathbf{15} \times \mathbf{6}\,. \tag{16}$$

In Equation (16), the **6**-plet has to be identified with the unit vector Ω. Thus, there are two different scenarios. In scenario A, the composite partners of u^c_α and q_i (U_α and Q_i) are components of $\mathbf{15}(U_\alpha)$ and $\mathbf{20}(Q_i)$ representations of $SU(6)$, while in scenario B, the composite partners of u^c_α and q_i belong to $\mathbf{20}(U_\alpha)$ and $\mathbf{15}(Q_i)$, respectively. Below scale f, the mixing between incomplete $SU(5)$ representations ($\mathbf{10}^u_\alpha$ and $\mathbf{10}^q_i$) and their composite partners is induced, and the Yukawa interactions (14) are reproduced.

In scenario A, the $SU(6)$ generalisation of the Yukawa interactions (15) is given by

$$\mathcal{L}^d_{SU(6)} \sim \mathbf{20} \times \overline{\mathbf{15}} \times \overline{\mathbf{6}}'\,, \tag{17}$$

where $\overline{\mathbf{6}}' \equiv \Omega^\dagger$. In Equation (17), the **20**-plet is associated with the $SU(6)$ representations involving the composite partners of q_i, i.e., $\mathbf{20}(Q_i)$, while the $\overline{\mathbf{15}}$-plet involves the $\overline{\mathbf{5}}$ multiplet of the $SU(5)$ group that should include the composite partners of d^c_i, i.e., $\overline{\mathbf{15}} \equiv \overline{\mathbf{15}}(D_i)$.

In scenario B, the $SU(6)$ generalisation of the $SU(5)$ structure of the Yukawa interactions (15) takes the form:

$$\mathcal{L}^d_{SU(6)} \sim \mathbf{15} \times \overline{\mathbf{6}} \times \overline{\mathbf{6}}'\,. \tag{18}$$

In this case, again, the $\overline{\mathbf{6}}'$ in Equation (18) must be identified with $\Omega^\dagger$. The **15**-plet corresponds to $\mathbf{15}(Q_i)$, which contains the composite partners of q_i, whereas the $\overline{\mathbf{6}}$-plet has to involve the composite partners of d^c_i (D_i), i.e., $\overline{\mathbf{6}} \equiv \overline{\mathbf{6}}(D_i)$. After the breakdown of the $SU(6)$ global symmetry near the scale f, Equation (18) should lead to the Yukawa interactions (15).

The interactions (18) can also be used to generate the masses of charged leptons in both scenarios A and B. In this case, $\mathbf{15}(E_i)$ and $\overline{\mathbf{6}}(L_i)$ should contain the composite partners of e^c_i and ℓ_i, respectively. In the simplest $SU(5)$ GUT, the masses of the left-handed neutrinos can be induced through the interactions

$$\mathcal{L}^\nu_{SU(5)} \simeq k_{ij} \left(\overline{\mathbf{5}}^\ell_{\mathbf{i}}\, \mathbf{5}^h \right) \left(\overline{\mathbf{5}}^\ell_{\mathbf{j}}\, \mathbf{5}^h \right). \tag{19}$$

The $SU(6)$ generalisation of the interactions (19) takes the form:

$$\mathcal{L}^\nu_{SU(6)} \simeq \varkappa_{ij} \left(\overline{\mathbf{6}} \times \mathbf{6}' \right) \left(\overline{\mathbf{6}} \times \mathbf{6}' \right), \tag{20}$$

where $\overline{\mathbf{6}}$ should be associated with $\overline{\mathbf{6}}(L_i)$, and $\mathbf{6}' \equiv \Omega$. The inclusion of interactions (20) implies that global $U(1)_L$ symmetry is broken down to $Z^L_2 = (-1)^L$. Since, in the E$_6$CHM, the lepton number is preserved to a very good approximation, $\varkappa_{ij}$ are expected to be very small, giving rise to tiny masses of the left-handed neutrinos.

3. From E_6 Orbifold GUT to the E$_6$CHM

As previously noted, the nearly exact conservation of the baryon and lepton numbers at low energies requires different multiplets of elementary quarks and leptons to stem from different representations of the GUT group. In this sense, the $U(1)_B$ and $U(1)_L$ charges of the corresponding GUT multiplets are determined by the baryon and lepton numbers of the fermion components of these representations that survive to low energies. All other components of these GUT multiplets must gain huge masses. In this section, we focus on the SUSY GUTs with the E_6-gauge group and assume that elementary fermions originate from different **27**-plets and $\overline{\mathbf{27}}$-plets of E_6. The complete set of the $SU(6)$ and E_6 representations associated with the multiplets of elementary fermions is given by

$$
\begin{aligned}
&\overline{q} \in \overline{\mathbf{15}}^{\overline{q}} = \left(\overline{\mathbf{15}},\, \frac{1}{3},\, 0\right) \in \overline{\mathbf{27}}^{\overline{q}} = \left(\overline{\mathbf{27}},\, \frac{1}{3},\, 0\right),\\
&\overline{e^c} \in \overline{\mathbf{15}}^{\overline{e}} = \left(\overline{\mathbf{15}},\, \frac{1}{3},\, 0\right) \in \overline{\mathbf{27}}^{\overline{e}} = \left(\overline{\mathbf{27}},\, \frac{1}{3},\, 0\right),\\
&\overline{d^c} \in \mathbf{6}^{\overline{d}} = \left(\mathbf{6},\, -B_{\overline{6}_2},\, -L_{\overline{6}_2}\right) \in \overline{\mathbf{27}}^{\overline{d}} = \left(\overline{\mathbf{27}},\, -B_{\overline{6}_2},\, -L_{\overline{6}_2}\right),\\
&\overline{\ell} \in \mathbf{6}^{\overline{\ell}} = \left(\mathbf{6},\, -B_{\overline{6}_2},\, -L_{\overline{6}_2}\right) \in \overline{\mathbf{27}}^{\overline{\ell}} = \left(\overline{\mathbf{27}},\, -B_{\overline{6}_2},\, -L_{\overline{6}_2}\right),\\
&u^c_\alpha \in \mathbf{15}^{u}_\alpha = \left(\mathbf{15},\, -\frac{1}{3},\, 0\right)_\alpha \in \mathbf{27}^{u}_\alpha = \left(\mathbf{27},\, -\frac{1}{3},\, 0\right)_\alpha,\\
&q_i \in \mathbf{15}^{q}_i = \left(\mathbf{15},\, \frac{1}{3},\, 0\right)_i \in \mathbf{27}^{q}_i = \left(\mathbf{27},\, \frac{1}{3},\, 0\right)_i,\\
&e^c_i \in \mathbf{15}^{e}_i = \left(\mathbf{15},\, 0,\, -1\right)_i \in \mathbf{27}^{e}_i = \left(\mathbf{27},\, 0,\, -1\right)_i,\\
&d^c_i \in \overline{\mathbf{6}}^{d}_i = \left(\overline{\mathbf{6}},\, -\frac{1}{3},\, 0\right)_i \in \mathbf{27}^{d}_i = \left(\mathbf{27},\, -\frac{1}{3},\, 0\right)_i,\\
&\ell_i \in \overline{\mathbf{6}}^{\ell}_i = \left(\overline{\mathbf{6}},\, 0,\, 1\right)_i \in \mathbf{27}^{\ell}_i = \left(\mathbf{27},\, 0,\, 1\right)_i.
\end{aligned}
\tag{21}
$$

In Equation (21), the first quantity in brackets is either the $SU(6)$ or E_6 representation, whereas the second and third quantities are the $U(1)_B$ and $U(1)_L$ charges. It is rather problematic to get the desirable splitting of the **27**-plets and $\overline{\mathbf{27}}$-plets within four-dimensional E_6 GUTs. Nevertheless, the appropriate splitting of these E_6 representations can occur in the orbifold GUTs with extra dimensions.

Orbifolding in higher-dimensional theories offers new possibilities for gauge symmetry breaking which have been explored within the SUSY GUT models in five dimensions [165–193] and six dimensions [188–208]. Initially, in the string-motivated work [209–215], it was pointed out that the breakdown of the gauge symmetry can be caused by identifications imposed on the gauge fields under the spacetime symmetries of an orbifold. More recently, it was argued that the orbifold compactifications of the heterotic string may lead to five-dimensional or six-dimensional GUT structures which are similar to orbifold GUT models [216–222]. The unification of gauge couplings in the $5D$ and $6D$ orbifold GUTs was examined in [179–187,206,207]. The models of composite fermions were studied in the context of Sherk–Schwarz compactification in [223].

In this section, an $N = 1$ SUSY GUT in $6D$, which results in a set of elementary fermions given by Equation (5), is considered. This SUSY GUT is based on the $E_6 \times G_0$-gauge group. Near some high energy scale M_X, the E_6 and G_0 groups are broken down to $SU(3)_C \times SU(2)_W \times U(1)_Y$ and G, respectively. The elementary bosons and fermions participate in the E_6 interactions only. Fields from the composite sector are charged under both the G_0 (G) and E_6-gauge symmetries. The phenomenological viability of such a model requires the adequate suppression of operators which give rise to proton decay. In the context of orbifold GUTs, the proton stability was discussed in [176–178,180–185,208]. It was shown that the experimental lower limit on the proton lifetime can be satisfied if $M_X \gtrsim 10^{16}$ GeV [208].

All elementary fermions in the model under consideration are components of the bulk 27-plets. In the four-dimensional $N = 1$ SUSY GUT, the fundamental 27-dimensional representation of E_6 contains components Φ_n ($n = 1, 2, \ldots, 27$) associated with the supermultiplets of one generation of ordinary matter, including the right-handed neutrino (ν^c), i.e., q, ℓ, u^c, d^c, e^c and ν^c. In addition, it also involves the supermultiplets that correspond to the charged $\pm 1/3$ exotic quarks (h and h^c), 2 $SU(2)_W$ doublets (h^d and h^u) as well as a SM singlet s. Each $6D$ fermion state is formed by 2 $4D$ Weyl fermions, ψ and ψ^c. The minimal $N=1$ SUSY in $6D$ implies that each $6D$ superfield involves one $6D$ fermion field and two complex scalars, ϕ and ϕ^c. These fields compose a $4D$ $N=2$ hypermultiplet that contains 2 $4D$ $N=1$ chiral superfields, $\Phi=(\phi,\psi)$ and $\overline{\Phi}=(\phi^c,\psi^c)$, with opposite quantum numbers. Therefore, each $6D$ 27-plet $\widehat{\Phi}_n$ includes two $4D$ $N=1$ supermultiplets, Φ_n (27-plet) and $\overline{\Phi}_n$ ($\overline{27}$-plet). In other words, the $N=1$ SUSY in $6D$ corresponds to $N=2$ supersymmetry in $4D$.

The E_6-gauge supermultiplet which should exist in the bulk must include vector bosons A_M ($M=0,1,2,3,5,6$) as well as $6D$ Weyl fermions (gauginos). Each $6D$ gaugino is formed by 2 $4D$ Weyl fermions, λ and λ'. The components of vector bosons and gauginos can be grouped into chiral and vector supermultiplets of the $N=1$ SUSY in $4D$, i.e.,

$$V=(A_\mu,\lambda)\,,\qquad\qquad \Sigma=\left((A_5+iA_6)/\sqrt{2},\lambda'\right). \tag{22}$$

In Equation (22), $\mu=0,1,2,3$, whereas λ, λ', A_M, V and Σ are matrices in the adjoint E_6 representation. The $N=1$ supermultiplets (22) compose an $N=2$ vector supermultiplet in $4D$.

We assume that 2 extra dimensions $z(=x_6)$ and $y(=x_5)$ are compact with $z\in(-\pi R_6,\pi R_6]$ and $y\in(-\pi R_5,\pi R_5]$ that corresponds to the compactification of extra dimensions on a torus T^2 with fixed radii R_5 and R_6, where R_5 and R_6 are defined by the scale M_X. Using Z_2 symmetry, the orbifold T^2/Z_2 can be obtained. The Z_2 transformation acts on T^2 according to $z\to -z$ and $y\to -y$. The components of the bulk supermultiplets also transform under Z_2 symmetry, while the Lagrangian of the model under consideration has to be invariant under this transformation. The Z_2 symmetry allows the physical region to be reduced to a pillow with the 4 fixed points as corners: $(0,0)$, $(\pi R_5,0)$, $(0,\pi R_6)$ and $(\pi R_5,\pi R_6)$.

3.1. The E_6 Symmetry Breaking to $SU(4)'\times SU(2)_W\times SU(2)_N\times U(1)'$

In this section, we consider a $6D$ SUSY GUT compactified on the orbifold $T^2/(Z_2\times Z_2^I\times Z_2^{II})$, where Z_2, Z_2^I and Z_2^{II} are reflections. In particular, a Z_2 transformation is defined as before, i.e., $z\to -z$ and $y\to -y$. The reflection Z_2^I acts as $z\to -z$ and $y'\to -y'$, where $y'=y-\pi R_5/2$. The reflection Z_2^{II} is defined by $z'\to -z'$ and $y\to -y$, where $z'=z-\pi R_6/2$. The reflection symmetries Z_2^I and Z_2^{II} introduce additional fixed points, resulting in the physical region in which $z\in[0,\pi R_6/2]$ and $y\in[0,\pi R_5/2]$. The irreducible space is a pillow limited by fixed points with 4 $4D$ branes (walls) which are located at its corners.

The Lagrangian of this $6D$ SUSY GUT must be invariant under the transformations of Z_2, Z_2^I and Z_2^{II} symmetries. Each reflection has its own orbifold parity, i.e., P, P_I and P_{II}. The components Φ_n and $\overline{\Phi}_n$ of $6D$ 27-plets transform under Z_2, Z_2^I and Z_2^{II} reflections as follows:

$$\begin{aligned}
&\Phi_n(x,-y,-z)=P_{nn}\Phi_n(x,y,z)\,, && \overline{\Phi}_n(x,-y,-z)=-P_{nn}\overline{\Phi}_n(x,y,z)\,,\\
&\Phi_n(x,-y',-z)=P^I_{nn}\hat{\Phi}_n(x,y',z)\,, && \overline{\Phi}_n(x,-y',-z)=-P^I_{nn}\overline{\Phi}_n(x,y',z)\,,\\
&\Phi_n(x,-y,-z')=P^{II}_{nn}\hat{\Phi}_n(x,y,z')\,, && \overline{\Phi}_n(x,-y,-z')=-P^{II}_{nn}\overline{\Phi}_n(x,y,z')\,.
\end{aligned} \tag{23}$$

In Equation (23), P, P_I and P_{II} are diagonal matrices which have eigenvalues ± 1. The diagonal elements of these matrices can be written as

$$\begin{gathered}
(P)_{ii}=\sigma\exp\{2\pi i\Delta\alpha_i\}\,,\qquad (P^I)_{ii}=\sigma_I\exp\{2\pi i\Delta^I\alpha_i\}\,,\\
(P^{II})_{ii}=\sigma_{II}\exp\{2\pi i\Delta^{II}\alpha_i\}\,,
\end{gathered} \tag{24}$$

where $\sigma, \sigma_I, \sigma_{II} \in \{+,-\}$ are parities of the $6D$ 27-plets and α_i are E_6 weights which are well-known [193]. The gauge shifts Δ, Δ^I and Δ^{II} associated with Z_2, Z_2^I and Z_2^{II} reflections are chosen so that

$$\begin{gathered}\Delta = \left(0,\, 0,\, 0,\, \frac{1}{2},\, 0,\, 0\right), \qquad \Delta^I = \left(\frac{1}{2},\, \frac{1}{2},\, \frac{1}{2},\, \frac{1}{2},\, \frac{1}{2},\, 0\right), \\ \Delta^{II} = \left(\frac{1}{2},\, \frac{1}{2},\, \frac{1}{2},\, 0,\, \frac{1}{2},\, 0\right). \end{gathered} \tag{25}$$

The corresponding orbifold parity assignments are shown in Table 1.

Table 1. Orbifold parity assignments in the bulk 27 supermultiplet with $\sigma = \sigma_I = \sigma_{II} = \sigma_{III} = +1$.

	q	d^c	u^c	ℓ	e^c	ν^c	h^u	h^d	h	h^c	s
Z_2	+	−	+	−	+	−	+	−	+	−	−
Z_2^I	−	+	+	−	+	+	−	−	+	+	+
Z_2^{II}	−	−	+	+	+	−	−	+	+	−	−
Z_2^{III}	+	+	+	+	+	+	+	+	+	+	+

The components of the E_6-gauge supermultiplet (V and Σ) transform under reflections Z_2, Z_2^I and Z_2^{II} as follows:

$$\begin{aligned} &V(x,-y,-z) = PV(x,y,z)P^{-1}, \\ &V(x,-y',-z) = P^I V(x,y',z)(P^I)^{-1}, \\ &V(x,-y,-z') = P^{II} V(x,y,z')(P^{II})^{-1}, \\ &\Sigma(x,-y,-z) = -P\Sigma(x,y,z)P^{-1}, \\ &\Sigma(x,-y',-z) = -P^I\Sigma(x,y',z)(P^I)^{-1}, \\ &\Sigma(x,-y,-z') = -P^{II}\Sigma(x,y,z')(P^{II})^{-1}, \end{aligned} \tag{26}$$

where $\Sigma(x,y,z) = \Sigma^A(x,y,z)T^A$ and $V(x,y,z) = V^A(x,y,z)T^A$, while T^A are the E_6 generators. In the orbifold GUT under consideration, the $4D$ $N=2$ SUSY is broken down to $4D$ $N=1$ supersymmetry because components Φ_n and $\overline{\Phi}_n$ as well as V and Σ transform differently under the reflections. Since P, P^I and P^{II} are not unit matrices, they do not commute with all generators of E_6. As a consequence, the E_6-gauge symmetry is broken as well.

From the P parity assignment, it follows that near the fixed point $y = z = 0$ (brane O) associated with the Z_2 reflection, the E_6-gauge group is broken down to $SU(6) \times SU(2)_N$. Indeed, the fundamental representation of E_6 decomposes under $SU(6) \times SU(2)_N$ as follows:

$$27 \to (15,\, 1) + (\overline{6},\, 2)\,. \tag{27}$$

In Equation (27), the first and second quantities in brackets are the $SU(6)$ and $SU(2)_N$ representations. From Table 1, one can see that $(\overline{6},\, 2)$ is formed by two $SU(2)_W$ doublets ℓ and h^d, two SM singlets ν^c and s as well as two $SU(3)_C$ triplets d^c and h^c. Indeed, these components of the 27-plet transform differently under the Z_2 reflection as compared with the other components of the fundamental representation of E_6 which compose the (15, 1) supermultiplet of $SU(6)$. The unbroken $SU(6)$ group contains a $SU(3)_C \times SU(2)_W \times U(1)_Y$ subgroup. It is assumed that all fields of the strongly interacting sector are localised on the O brane.

At the fixed point $y = \pi R_5/2$, $z = 0$ (brane O_I) of the Z_2^I reflection the E_6 group is broken to its $SU(6)' \times SU(2)_W$ subgroup. According to the P^I parity assignment, all $SU(2)_W$ doublets from the 27-plet transform differently as compared with the other components of the fundamental representation of E_6. These $SU(2)_W$ doublets compose the (6, 2) representation of $SU(6)'$. All other components of the 27-plet form $(\overline{15},\, 1)$ of $SU(6)'$. In this case, the unbroken $SU(6)'$ includes a $SU(3)_C$ subgroup. It is assumed that 2 pairs of

(15, 1) and ($\overline{15}$, 1) of $SU(6)'$ are localised on the brane O_I. They are needed to ensure the appropriate breakdown of the E_6-gauge group to its $SU(3)_C \times SU(2)_W \times U(1)_Y$ subgroup.

Near the fixed point $y = 0$, $z = \pi R_6/2$ (brane O_{II}) associated with the Z_2^{II} reflection the E_6-gauge group is also broken to $SO(10)' \times U(1)'$. Indeed, the P_{II} parity assignment indicates that the 16 components of the fundamental representation of E_6, i.e., q, d^c, ν^c, h^u, h^c and s, are odd, composing a 16-dimensional spinor representation of $SO(10)'$. All other components of the 27-plet are even. Because, in the orbifold GUTs, the mechanism of gauge symmetry breaking preserves the ranks of the group, the unbroken subgroup of the E_6 group has to be $SO(10)' \times U(1)'$. The unbroken $SO(10)'$ contains a $SU(3)_C \times SU(2)_W$ subgroup. The 10 components of the 27-plet, i.e., u^c, ℓ, h^d and h, constitute a 10-dimensional vector representation of $SO(10)'$ while the e^c component of the 27-plet is an $SO(10)'$ singlet. It is worth noting that the spinor representation of the ordinary $SO(10)$ and $SO(10)'$ are composed by different components of the 27-plets. This means that $SO(10)$ and $SO(10)'$ are different subgroups of E_6. We assume that, on the brane O_{II}, a 45-dimensional representation of $SO(10)'$ as well as 3 pairs of e^c_i and $\overline{e^c_i}$ superfields are confined.

At the corner of the physical region, i.e., $y = \pi R_5/2$, $z = \pi R_6/2$, a fourth fixed point (brane O_{III}) is located. It is associated with the Z_2^{III} symmetry, which is obtained by combining reflections Z_2, Z_2^I and Z_2^{II}. The corresponding parity assignment $P_{III} = P\, P_I\, P_{II}$ is just an identity matrix. Therefore, near this fixed point, the E_6-gauge group remains intact, whereas $N = 2$ SUSY is broken to $N = 1$ supersymmetry. We assume that two 27-plets reside on the brane O_{III}.

The intersection of the E_6 subgroups, which remain intact near the branes O, O_I, O_{II} and O_{III}, represents the unbroken gauge group of the effective $4D$ theory. The intersection of the E_6 subgroups $SU(6) \times SU(2)_N$, $SU(6)' \times SU(2)_W$ and $SO(10)' \times U(1)'$ is $SU(4)' \times SU(2)_W \times SU(2)_N \times U(1)'$, where the $SU(4)'$ group includes a $SU(3)_C$ subgroup. The $SU(4)'$ group is also a subgroup of $SU(6)'$ and $SO(10)'$. In Table 2, the charges associated with the $U(1)'$ symmetry are specified.

Table 2. The charges (Q'_i, $\widetilde{Q}_i$, Q^N_i and Q^Y_i) of the components of the 27-plet associated with the $U(1)'$, $\widetilde{U}(1)$, $U(1)_N$ and $U(1)_Y$ symmetries.

	q	d^c	u^c	ℓ	e^c	ν^c	h^u	h^d	h	h^c	s
$\sqrt{24}Q'_i$	1	1	−2	−2	4	1	1	−2	−2	1	1
$\sqrt{24}\widetilde{Q}_i$	1	−1	2	0	0	3	−3	0	−2	−1	3
$\sqrt{40}Q^N_i$	1	2	1	2	1	0	−2	−3	−2	−3	5
$\sqrt{\frac{5}{3}}Q^Y_i$	$\frac{1}{6}$	$\frac{1}{3}$	$-\frac{2}{3}$	$-\frac{1}{2}$	1	0	$\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{3}$	$\frac{1}{3}$	0

3.2. The Breakdown of $SU(4)' \times SU(2)_W \times SU(2)_N \times U(1)'$ to $SU(3)_C \times SU(2)_W \times U(1)_Y$

According to Table 1, the bulk 27-plets contain components which have even and odd parities with respect to the transformations of the Z_2, Z_2^I and Z_2^{II} symmetries. At the same time, only components that have all even parities are allowed to have zero modes which may survive below the scale M_X. In particular, the elementary fermions, u^c_α, e^c_i and $\bar{e^c}$, can stem from the bulk 27-plets $\widehat{\Phi}^u_i$, $\widehat{\Phi}^{\overline{u}}_i$, $\widehat{\Phi}^e_i$ and $\widehat{\Phi}^{\overline{e}}_i$. Hereafter, index $i = 1, 2, 3$ runs over three generations. These bulk 27-plets decompose as follows:

$$\begin{aligned} \widehat{\Phi}^u_i &= (27, +, +, +, +), \qquad & \widehat{\Phi}^{\overline{u}}_i &= (27, -, -, -, -), \\ \widehat{\Phi}^e_i &= (27, +, +, +, +), \qquad & \widehat{\Phi}^{\overline{e}}_i &= (27, -, -, -, -), \end{aligned} \tag{28}$$

In Equation (28), the quantities in brackets are the parities of the bulk 27-plets σ, σ_I, σ_{II} and σ_{III} as well as the E_6 representations of these $6D$ supermultiplets. The parities of $\widehat{\Phi}^u_i$ are such that only the components u^c_i, e^c_i and h_i of Φ^u_i have zero modes. Since the parities of the components of Φ^u_i and $\overline{\Phi}^u_i$ are opposite, the $N = 1$ supermultiplet $\overline{\Phi}^u_i$ does not lead to zero modes. On the other hand, the Kaluza–Klein (KK) expansion of $\widehat{\Phi}^{\overline{u}}_i$ involves only the zero modes of components $\overline{u^c}_i$, $\overline{e^c}_i$ and $\overline{h}_i$. The bulk 27-plets $\widehat{\Phi}^e_i$ and $\widehat{\Phi}^{\overline{e}}_i$ result in a similar set of zero modes.

The 45-dimensional representation of $SO(10)'$ that resides on the brane O_{II} involves one component, φ, corresponding to the generator of the $\widetilde{U}(1)$ subgroup of $SO(10)'$, the charges of which are specified in Table 2. It is assumed that φ acquires a non-zero vacuum expectation value (VEV), φ_0, and couples to $\widehat{\Phi}^e_i$ and $\widehat{\Phi}^{\overline{e}}_i$. This VEV, which is smaller than the scale M_X, breaks the $SU(4)' \times SU(2)_W \times SU(2)_N \times U(1)'$-gauge symmetry to $SU(3)_C \times SU(2)_W \times SU(2)_N \times \widetilde{U}(1) \times U(1)'$, generating masses of the the zero modes u^c_i, $\overline{u^c}_i$, h_i and $\overline{h}_i$ of supermultiplets $\widehat{\Phi}^e_i$ and $\widehat{\Phi}^{\overline{e}}_i$. In this case, only zero modes associated with the components e^c_i and $\overline{e^c}_i$ remain massless. The couplings of other $6D$ supermultiplets to φ are expected to be forbidden. Because supermultiplet φ is confined on the brane O_{II}, the VEV φ_0 does not directly break the $SU(6) \times SU(2)_N$ global symmetry of the strongly interacting sector which is localised on the brane O.

At the same time, the superfield e^c_i, which resides on brane O_{II}, can be combined with the corresponding zero modes of $\widehat{\Phi}^{\overline{u}}_i$, giving rise to vector-like states with masses of order of M_X. In the same way, $\overline{e^c_i}$ and the appropriate zero modes of $\widehat{\Phi}^u_i$ can form vector-like states that gain masses set by the scale M_X. As a result, only zero modes of the components of $\widehat{\Phi}^{\overline{u}}_i$ and $\widehat{\Phi}^u_i$ with the quantum numbers of $\overline{u^c}_i$, $\overline{h}_i$, u^c_i and h_i remain massless.

The zero modes with the quantum numbers of q_i, d^c_i, $\bar{q}$ and $\bar{d}^c$ can originate from the bulk supermultiplets

$$\begin{aligned} \widehat{\Phi}^q_i &= (27, +, -, -, +), \qquad & \widehat{\Phi}^{\overline{q}}_i &= (27, -, +, +, -), \\ \widehat{\Phi}^d_i &= (27, -, +, -, +), \qquad & \widehat{\Phi}^{\overline{d}}_i &= (27, +, -, +, -). \end{aligned} \tag{29}$$

Using the parity assignments given in Table 1, it is easy to check that all parities of d^c_i, h^c_i, s_i and ν^c_i components of $\widehat{\Phi}^d_i$, $\overline{d^c_i}$, $\overline{h^c_i}$, $\overline{s_i}$ and $\overline{\nu^c_i}$ components of $\widehat{\Phi}^{\overline{d}}_i$, h^u_i and q_i components of $\widehat{\Phi}^q_i$ as well as $\overline{h^u_i}$ and $\overline{q_i}$ and components of $\widehat{\Phi}^{\overline{q}}_i$ are positive. Therefore, the corresponding KK expansions include zero modes.

In order to obtain the zero modes associated with the ℓ_i and $\bar{\ell}$ components of the 27-plets and $\overline{27}$-plet, the set of the $6D$ supermultiplets has to be supplemented by

$$\widehat{\Phi}^\ell_i = (27, -, -, +, +), \qquad \widehat{\Phi}^{\overline{\ell}}_i = (27, +, +, -, -). \tag{30}$$

Again, one can check that all parities of h^d_i and ℓ_i components of $\widehat{\Phi}^\ell_i$ as well as $\overline{h^d_i}$ and $\overline{\ell_i}$ components of $\widehat{\Phi}^{\overline{\ell}}_i$ are positive, resulting in the corresponding set of zero modes. The full set of the bulk supermultiplets, as well as their zero modes, which remain massless below $\langle\varphi\rangle = \varphi_0$, are given in Table 3. We assume here that the mass terms associated with the zero modes with opposite quantum numbers are forbidden.

Table 3. The bulk supermultiplets and their zero modes that remain massless below the scales M_X, φ_0 and ϕ_0.

	$\widehat{\Phi}^u_i$	$\widehat{\Phi}^e_i$	$\widehat{\Phi}^q_i$	$\widehat{\Phi}^d_i$	$\widehat{\Phi}^\ell_i$	$\widehat{\Phi}^{\overline{u}}_i$	$\widehat{\Phi}^{\overline{e}}_i$	$\widehat{\Phi}^{\overline{q}}_i$	$\widehat{\Phi}^{\overline{d}}_i$	$\widehat{\Phi}^{\overline{\ell}}_i$
$E \lesssim M_X$	u^c_i, e^c_i, h_i	u^c_i, e^c_i, h_i	q_i, h^u_i	d^c_i, ν^c_i, h^c_i, s_i	ℓ_i, h^d_i	$\overline{u^c}_i$, $\overline{e^c}_i$, $\overline{h}_i$	$\overline{u^c}_i$, $\overline{e^c}_i$, $\overline{h}_i$	$\overline{q_i}$, $\overline{h^u_i}$	$\overline{d^c_i}$, $\overline{\nu^c_i}$, $\overline{h^c_i}$, $\overline{s_i}$	$\overline{\ell_i}$, $\overline{h^d_i}$
$E \lesssim \varphi_0$	u^c_i, h_i	e^c_i	q_i, h^u_i	d^c_i, ν^c_i, h^c_i, s_i	ℓ_i, h^d_i	$\overline{u^c}_i$, $\overline{h}_i$	$\overline{e^c}_i$	$\overline{q_i}$, $\overline{h^u_i}$	$\overline{d^c_i}$, $\overline{\nu^c_i}$, $\overline{h^c_i}$, $\overline{s_i}$	$\overline{\ell_i}$, $\overline{h^d_i}$
$E \lesssim \phi_0$	u^c_i	e^c_i	q_i	d^c_i	ℓ_i	$\overline{u^c}$	$\overline{e^c}$	$\overline{q}$	$\overline{d^c}$	$\overline{\ell}$

Below the scale φ_0, the $6D$ supermultiplets result in the set of zero modes which involves 3 pairs of $N = 1$ chiral 27 and $\overline{27}$-plets. The model under consideration implies that 2 $\overline{27}$-plets associated with $i = 1, 2$ and 2 27 supermultiplets which are confined on the O_{III} brane compose vector-like states with masses $M_0 \lesssim \varphi_0$. It is also expected that

2 pairs of $\overline{15}$ and 15 of $SU(6)'$ that reside on the brane O_I acquire VEVs of the order of $\phi_0 \lesssim M_0$. The VEVs of the $\overline{\nu^c}$ and ν^c components of one pair of 15 and $\overline{15}$ break the $SU(3)_C \times SU(2)_W \times SU(2)_N \times \widetilde{U}(1) \times U(1)'$-gauge group to $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_N$. Different phenomenological aspects of the SUSY extensions of the SM based on the $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_N$-gauge symmetry were explored in [224–250]. The VEVs of the $\bar{s}$ and s components of another pair of 15 and $\overline{15}$ break the $SU(3)_C \times SU(2)_W \times U(1)_Y \times U(1)_N$-gauge group to $SU(3)_C \times SU(2)_W \times U(1)_Y$. The VEVs of $\overline{\nu^c}$, ν^c, $\bar{s}$ and s also induce the set of the mass terms in the superpotential as follows:

$$\begin{aligned} \delta W_{mass} &= M^{\eta}_{ij} h_i h^c_j + M^{\zeta}_{ij} h^u_i h^d_j + M^{\xi}_{ij} s_i s_j + M^{\nu}_{ij} \nu^c_i \nu^c_j + \overline{M}^{\eta}\, \overline{h}\, \overline{h^c} \\ &+ \overline{M}^{\zeta}\, \overline{h^u}\, \overline{h^d} + \overline{M}^{\xi}\, \bar{s}^2 + \overline{M}^{\nu}\, \overline{\nu^c}^2 , \end{aligned} \tag{31}$$

where $i, j = 1, 2, 3$ and $M^{\eta}_{ij} \sim M^{\zeta}_{ij} \sim M^{\xi}_{ij} \sim M^{\nu}_{ij} \sim \overline{M}^{\eta} \sim \overline{M}^{\zeta} \sim \overline{M}^{\xi} \sim \overline{M}^{\nu} \sim \phi_0$. The pairs of 15 and $\overline{15}$ of $SU(6)'$ compose vector-like states with masses which are close to ϕ_0. These states reside on the brane O_I and do not interact directly with the supermultiplets which are localised on the brane O. Therefore, the $SU(6)$ global symmetry of the strongly interacting sector may remain unbroken.

In the six-dimensional orbifold GUTs, there are two types of anomalies: bulk anomalies [251–255] and $4D$ anomalies at orbifold fixed points [256–258]. Bulk anomalies are induced by box diagrams, the contributions of which are proportional to the trace of four generators. This trace includes a part which may be reduced to the product involving traces of two generators. It corresponds to the reducible anomaly that can be canceled by the Green–Schwarz mechanism [259]. Another nonfactorizable part is associated with the irreducible gauge anomaly. The E_6 orbifold GUT model in six dimensions does not have irreducible anomaly [254,255]. The $4D$ anomalies at the fixed points reduces to the anomalies associated with the unbroken subgroup of the E_6-gauge group in the vicinity of such points. These anomalies are determined by the sum of the contributions that come from the zero modes confined on the brane [251–253,260–262]. In the orbifold GUT under consideration, the corresponding brane anomalies are cancelled automatically.

Finally, near the scale M_S, which is somewhat lower than ϕ_0, SUSY gets broken. As a consequence, the scalar components of all superfields gain masses of the order of M_S. At the scale of M_S, the SM singlet superfield S develops a non-zero VEV. This superfield interacts only with the components of $\widehat{\Phi}^{u}_3$ and $\widehat{\Phi}^{\overline{u}}_3$. The interactions of the superfield S with the components of other bulk supermultiplets can be forbidden by the discrete Z^S_2 symmetry, under which only $\widehat{\Phi}^{\overline{u}}_3$ and S are odd, whereas all other bulk 27-plets are even. The VEV of S gives rise to the masses of zero modes u_3 and $\overline{u_3}$. Thus, below scale M_S, the weakly coupled elementary sector includes a set of the fermion states given by Equation (5). Since different fermion multiplets are the zero modes of different bulk supermultiplets, they are allowed to have different baryon and lepton numbers in this case.

4. Generation of Matter–Antimatter Asymmetry in the E$_6$CHM

In this section, we restrict our consideration to the scenarios with $f \gtrsim 10\,\text{TeV}$. The $U(1)_B$ symmetry can be explicitly broken in the strongly coupled sector in this case, since Z^L_2 can be nearly exact, forbidding all operators which lead to the proton decay. When scale f is so high, all other operators that violate baryon number are sufficiently strongly suppressed. This suppression is caused by the small mixing between elementary fermions and their composite partners as well as by the large value of f. For instance, the effective operators in the SM, which give rise to the processes with $\Delta B = 2$ and $\Delta L = 0$, are given by

$$\mathcal{L}_{\Delta B=2} = \frac{1}{\Lambda^5} \left[q_i q_j q_k q_m (d^c_n d^c_l)^* + u^c_i d^c_j d^c_k u^c_m d^c_n d^c_l \right] , \tag{32}$$

where $i, j, k, m, n, l = 1, 2, 3$ are the generation indices. The $n - \bar{n}$ mixing mass δm and $n - \bar{n}$ oscillation time $\tau_{n-\bar{n}}$ can be estimated as

$$\delta m \simeq \varkappa \frac{\Lambda_{QCD}^6}{\Lambda^5}\,, \qquad\qquad \tau_{n-\bar{n}} \simeq \frac{1}{\delta m}\,. \tag{33}$$

In Equation (33), $\varkappa \sim 1$ and $\Lambda_{QCD} \simeq 200\,\text{MeV}$. The $n - \bar{n}$ oscillation time becomes close to the experimental limit 10^8 s. Refs. [263,264] for $\Lambda \sim$ few $\times$ 100 TeV. The operators (32) also induce the process of the annihilation of the two nucleons $NN \to KK$, resulting in rare nuclear decays. The searches for such decays set a lower bound on Λ of around 200–300 TeV. At the same time, in the E_6CHM with $f \gtrsim 10\,\text{TeV}$, the value of $\Lambda \gtrsim$ few $\times$ 100 TeV.

Here we assume that the effective Lagrangian of the E_6CHM possesses an approximate Z_2^B symmetry. This discrete symmetry is a subgroup of $U(1)_B$, i.e.,

$$Z_2^B = (-1)^{3B}\,. \tag{34}$$

In Equation (34), B is the baryon number of the multiplet. The Z_2^B discrete symmetry forbids proton decay, but it does not suppress the baryon-number-violating operators (32) Thus, in this case, the Lagrangian of the strongly interacting sector of the E_6CHM respects the approximate $SU(6) \times U(1)_L \times Z_2^B$ symmetry.

In the scenario under consideration, after the breakdown of the $SU(6)$ global symmetry to its $SU(5)$ subgroup, all composite states and exotic fermions, including the components of the $SU(6)$ multiplets $\overline{\mathbf{6}}_1$, $\overline{\mathbf{6}}_2$, $\mathbf{15}$ as well as $\bar{q}$, $\bar{d}^c$, $\bar{\ell}$ and $\bar{e}^c$, gain masses which are several times larger than f. The only exceptions are the components of the 15-plet which are identified with t^c. These components survive to the EW scale. As follows from Equation (12), all components of the $\mathbf{15}$-plet and $\overline{\mathbf{6}}_2$ multiplet as well as $\bar{q}$ and $\bar{e}^c$ are odd under the Z_2^B symmetry. The components of the $\overline{\mathbf{6}}_1$ multiplet, $\bar{q}$ and $\bar{\ell}$ can be either odd or even under the Z_2^B symmetry. Hereafter, it is assumed that these fermions are Z_2^B-even.

The N_1 and N_2 components of $\overline{\mathbf{6}}_1$ and $\overline{\mathbf{6}}_2$ multiplets acquire Majorana masses through the interactions $(\overline{\mathbf{6}}_1\Omega)(\Omega\overline{\mathbf{6}}_1)$ and $(\overline{\mathbf{6}}_2\Omega)(\Omega\overline{\mathbf{6}}_2)$, respectively. These operators are allowed by the approximate Z_2^B symmetry. Nevertheless, this symmetry suppresses the mixing between N_1 and N_2. We further assume that N_1 is substantially lighter than other composite and exotic fermions and has a mass of order of f.

The pNGB states have masses which are considerably lower than $f \gtrsim 10\,\text{TeV}$. Therefore these resonances are the lightest composite particles in the E_6CHM spectrum. All pNGB states are even under the Z_2^B symmetry because the Higgs boson manifests itself in interactions with SM particles as a Z_2^B-even state. The Z_2^B and gauge symmetries permit the decays of the $SU(3)_C$ triplet of scalar fields T into up and down antiquarks. On the other hand, almost exact Z_2^L symmetry forbids the decays of T into either a neutrino and a down quark or a charged lepton and an up quark. The decay mode $T \to \bar{t}\bar{b}$ should be the dominant one because the first and second generation quarks have quite small fractions of compositeness. For $E \lesssim f$, all operators that violate baryon number are suppressed, and T manifests itself in the interactions with other particles as a diquark with $B = -2/3$. At the LHC, the $SU(3)_C$ triplet can be pair produced, resulting in four heavy quarks in the final state, i.e., $pp \to T\bar{T} \to t\bar{t}b\bar{b}$. A somewhat similar signature arises in the R-parity-violating SUSY models. It is associated with the lightest squark in these models. Nowadays, scenarios with the mass of the $SU(3)_C$ triplet T below 700 GeV are disfavored by the LHC constraints on the masses of such squarks [265].

Although at low energies $E \lesssim f$ the baryon number violating processes are suppressed within the E_6CHM, a sizeable baryon number asymmetry may still be generated via the out-of-equilibrium decays of N_1 if N_1 has a mass which is substantially lower than the masses of all other composite and exotic fermions. This can happen if CP is violated and the mass of the $SU(3)_C$ triplet T (m_T) is in the multi TeV range, provided $m_T \ll m_{N_1} \sim f$

and the decays $N_1 \to T^* + d_i$ and $N_1 \to T + \bar{d}_i$ are allowed. The Lagrangian that describes the decays of N_1 and N_2 into down-type quarks and the pNGB state T is given by

$$\mathcal{L}_N = \sum_{i=1}^{3} \left(g_{i1}^* T d_i^c N_1 + g_{i2}^* T d_i^c N_2 + h.c. \right). \tag{35}$$

When Z_2^B symmetry is exact, the coupling g_{i1} vanishes. Therefore, one can expect that, in the case of the approximate Z_2^B symmetry, $|g_{i1}| \ll |g_{i2}|$. Because the pNGB state T decays mostly into $\bar{t}\bar{b}$, the decays of the Majorana fermion N_1 lead to the final states with $B = \pm 1$. The baryon asymmetry generation via the neutral fermion decays into scalar diquark and quark was considered in [266–281].

The generation of the baryon asymmetry is determined by the flavour CP asymmetries $\varepsilon_{1,k}$

$$\varepsilon_{1,k} = \frac{\Gamma_{N_1 d_k} - \Gamma_{N_1 \bar{d}_k}}{\sum_m \left(\Gamma_{N_1 d_m} + \Gamma_{N_1 \bar{d}_m} \right)}, \tag{36}$$

where $k, m = 1, 2, 3$, whereas $\Gamma_{N_1 \bar{d}_k}$ and $\Gamma_{N_1 d_k}$ are partial decay widths of $N_1 \to \bar{d}_k + T$ and $N_1 \to d_k + T^*$. There are three CP (decay) asymmetries that correspond to three quark flavours, i.e., d, s and b. At the tree level,

$$\Gamma_{N_1 d_k} = \Gamma_{N_1 \bar{d}_k} = \frac{3|g_{k1}|^2}{32\pi} M_1, \tag{37}$$

and the decay asymmetries (36) vanish in this approximation. In Equation (37), M_1 is the mass of N_1. The interference of the tree-level decay amplitudes of N_1 with the one-loop corrections to them yields the non-zero values of the CP asymmetries (36) if CP invariance is violated. The corresponding tree-level and one-loop diagrams can be found in [250]. Assuming that $m_T \ll m_{N_1}$, and using the results obtained in the case of thermal leptogenesis [48,282–286], the direct calculation of all these diagrams gives

$$\begin{aligned} \varepsilon_{1,i} \;=\; & \frac{1}{(8\pi)} \frac{1}{(\sum_{m=1}^{3} |g_{m1}|^2)} \left[\sum_{n=1}^{3} \mathrm{Im}(g_{i1}^* g_{i2} g_{n1}^* g_{n2}) \sqrt{x} \left(\frac{3}{2(1-x)} + 1 \right.\right. \\ & \left.\left. -(1+x) \ln \frac{1+x}{x} \right) + \sum_{n=1}^{3} \mathrm{Im}(g_{i1}^* g_{i2} g_{n1} g_{n2}^*) \frac{3}{2(1-x)} \right], \end{aligned} \tag{38}$$

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

To compute the baryon asymmetries generated by the decays of N_1, one needs to solve the system of Boltzmann equations that determine the evolution of baryon number densities. Since the corresponding solution has to be similar to the solutions in the case of thermal leptogenesis, the induced baryon asymmetry relative to the entropy density $Y_{\Delta B}$ can be approximately estimated as (see [286])

$$Y_{\Delta B} \sim 10^{-3} \left(\sum_{k=1}^{3} \varepsilon_{1,k} \eta_k \right), \qquad Y_{\Delta B} = \frac{n_B - n_{\bar{B}}}{s} \bigg|_0 = (8.75 \pm 0.23) \times 10^{-11}. \tag{39}$$

where η_k are efficiency factors that vary from 0 to 1, and s is the entropy density. In the limit, when washout processes can be neglected, $\eta_k = 1$. Here we ignore sphaleron processes that partially convert baryon asymmetry into lepton asymmetry.

It is expected that $|g_{32}| \gg |g_{22}|, |g_{12}|$ and $|g_{31}| \gg |g_{21}|, |g_{11}|$ because s- and d-quarks are almost elementary fermions which couple very weakly to the operators of the strongly interacting sector. Such hierarchical structure of the Yukawa couplings ensures that $\varepsilon_{1,2}$

and $\varepsilon_{1,1}$ are negligibly small. To simplify our analysis, we set $M_2 = 10 \cdot M_1$. In the limit $x \gg 1$ for $\varepsilon_{1,3}$, one finds

$$\varepsilon_{1,3} \simeq -\frac{1}{(4\pi)} \frac{|g_{32}|^2}{\sqrt{x}} \sin 2\Delta\varphi\,, \qquad \Delta\varphi = \varphi_{32} - \varphi_{31}\,, \tag{40}$$

where the phases φ_{32} and φ_{31} are defined as $g_{32} = |g_{32}|e^{i\varphi_{32}}$ and $g_{31} = |g_{31}|e^{i\varphi_{31}}$. When CP invariance is preserved, i.e., all Yukawa couplings are real, the decay asymmetry (40) goes to zero. The maximum absolute value of $\varepsilon_{1,3}$ is attained for $\Delta\varphi = \pm\pi/4$.

The efficiency factor η_3 can be of the order of unity in the E_6CHM. Indeed, in the strong washout scenario η_3 may be estimated as follows (see, for example [286])

$$\begin{gathered}\eta_3 \simeq H(T = M_1)/\Gamma_3\,,\\ \Gamma_3 = \Gamma_{N_1 d_3} + \Gamma_{N_1 \bar{d}_3} = \frac{3|g_{31}|^2}{16\pi} M_1\,, \qquad H = 1.66 g_*^{1/2} \frac{T^2}{M_{Pl}}\,,\end{gathered} \tag{41}$$

where $g_* = n_b + \frac{7}{8} n_f$ is the number of relativistic degrees of freedom in the thermal bath and H is the Hubble expansion rate. In the SM, $g_* = 106.75$, while in the E_6CHM, $g_* = 113.75$ for $T \lesssim f$. From Equation (41), one can see that the efficiency factor η_3 increases with diminishing of $|g_{31}|$, and for sufficiently small values of $|g_{31}|$, it may become close to unity. For instance, in the case when $M_1 \simeq 10\,\text{TeV}$ and $|g_{31}| \simeq 10^{-6}$, this factor is around 0.25.

When $\eta_3 \sim 1$, the generated baryon asymmetry is defined by the CP asymmetry $\varepsilon_{1,3}$ which does not depend on $|g_{31}|$. From Equation (40), it follows that for a given ratio M_2/M_1 and negligibly small absolute values of the Yukawa couplings $|g_{11}|$, $|g_{21}|$, $|g_{12}|$ and $|g_{22}|$, the decay asymmetry $\varepsilon_{1,3}$ is set by the combination of phases $\Delta\varphi$ and $|g_{32}|$. Because g_{32} is not suppressed by the Z_2^B discrete symmetry, it is expected that $|g_{32}| \gtrsim 0.1$. In Figure 1, the dependence of $|\varepsilon_{1,3}|$ on $\Delta\varphi$ is shown for two different values of $|g_{32}|$, i.e., $|g_{32}| = 1$ and $|g_{32}| = 0.1$. Figure 1 illustrates that $|\varepsilon_{1,3}|$ increases when $|g_{32}|$ grows and attains its maximal possible value for $\Delta\varphi \simeq \pi/4$. Near its maximum, the value of $|\varepsilon_{1,3}|$ is so large that a phenomenologically acceptable baryon density is induced only for $\eta_3 \lesssim 10^{-3}$. If $|g_{32}| \gtrsim 0.1$ and $\eta_3 \sim 1$, then the appropriate matter–antimatter asymmetry, corresponding to $\varepsilon_{1,3} \lesssim 10^{-7} - 10^{-6}$, can be obtained only in the limit when $\Delta\varphi \lesssim 0.01$. This demonstrates that, in the E_6CHM, the observed baryon density can be induced even if CP is approximately preserved. For $\Delta\varphi \sim 1$ and $(M_2/M_1) = 10$, phenomenologically acceptable matter–antimatter asymmetry can be also generated when $|g_{32}|$ varies from 0.01 to 0.1.

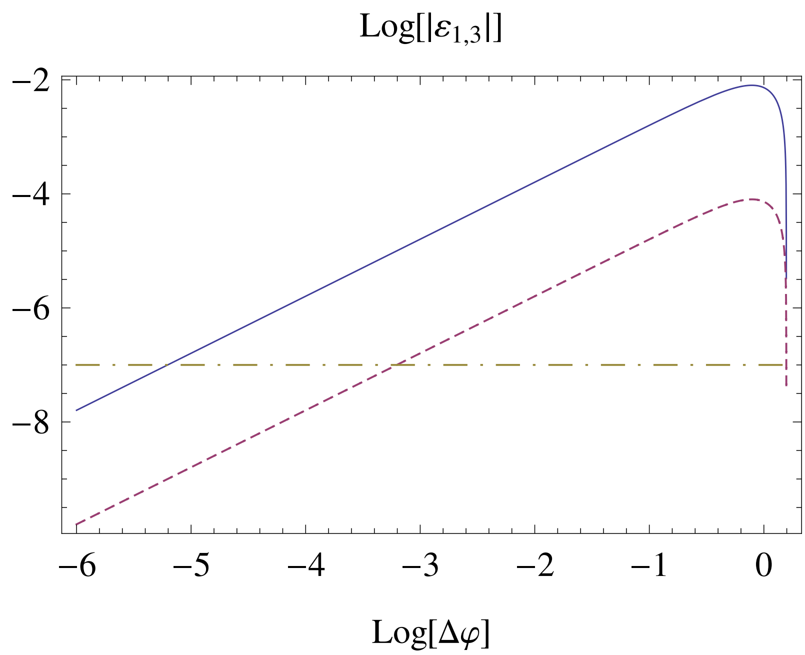

Figure 1. Logarithm (base 10) of the absolute value of the decay asymmetry $\varepsilon_{1,3}$ as a function of logarithm (base 10) of $\Delta\varphi$ for $|g_{32}| = 1$ (solid line) and $|g_{32}| = 0.1$ (dashed line) in the case when $g_{11} = g_{21} = g_{12} = g_{22} = 0$ and $M_2 = 10 \cdot M_1$.

5. Conclusions

The breakdown of gauge symmetry within GUTs can lead, at low energies, to the E_6-inspired composite Higgs model (E_6CHM), which almost allows the mass hierarchy to be stabilized. In particular, the E_6CHM can originate from the SUSY GUT based on the $E_6 \times G_0$-gauge symmetry. In the vicinity of some high energy scale M_X, the $E_6 \times G_0$ group can be broken down to the $SU(3)_C \times SU(2)_W \times U(1)_Y \times G$ subgroup, where G and G_0 are associated with the strongly interacting sector. We consider a six-dimensional orbifold SUSY GUT in which all fields of the strongly coupled sector reside on the brane where E_6 is broken down to $SU(6)$. It is expected that, in the E_6CHM $SU(6)$, there still remains an approximate symmetry of the strongly interacting sector even at low energies, and this gets spontaneously broken around the scale $f \gtrsim 5\,\text{TeV}$ to $SU(5)$, which contains the $SU(3)_C \times SU(2)_W \times U(1)_Y$-gauge group as a subgroup. The E_6 orbifold GUTs in six dimensions do not have irreducible bulk anomaly. In the orbifold SUSY GUT under consideration, brane anomalies are cancelled. Within this E_6 model, different multiplets of the elementary quarks and leptons stem from different bulk 27-plets. All other components of these 27-plets gain huge masses which are somewhat close to M_X. As a consequence, the low energy Lagrangian of the E_6CHM can be invariant with respect to the global $U(1)_L$ and $U(1)_B$ symmetries, which guarantee the conservation of the lepton and baryon charges to a very good approximation. To ensure that the left-handed neutrinos acquire non-zero Majorana masses, the $U(1)_L$ symmetry should be broken down to its Z_2^L discrete subgroup, which forbids all operators giving rise to rapid proton decay.

The $SU(6)$ symmetry breaking to $SU(5)$ in the E_6CHM results in eleven pNGB states. Four of these states form the SM-like Higgs doublet H. One of these pNGBs is a SM singlet boson A. Six others are associated with the $SU(3)_C$ triplet of scalar field T. The pNGB states mentioned above do not carry any lepton and/or baryon numbers. A significant fine-tuning, $\sim$0.01%, is required to obtain a Higgs boson with mass around 125 GeV in this model because $v \ll f$.

The masses of the SM fermions in the E_6CHM are induced through the mixing between elementary states and their composite partners. There are two different scenarios of quark mass generation. In scenario A, the composite partners of the right-handed down-type quarks, left-handed quarks and right-handed up-type quarks are components of $\overline{\mathbf{15}}$, **20** and **15** representations of $SU(6)$. Scenario B implies that the composite partners of the right-handed down-type quarks, left-handed quarks and right-handed up-type quarks belong to $\overline{\mathbf{6}}$, **15** and **20** representations of the $SU(6)$ group. In the case of the lepton sector, the corresponding masses can be generated if the composite partners of the right-handed charged leptons and left-handed leptons are components of the **15** and $\overline{\mathbf{6}}$ representations of $SU(6)$.

The embedding of the E_6CHM into an orbifold GUT with the E_6-gauge group implies that, at some high energy scale, the SM-gauge couplings are approximately equal. This can be achieved when the dynamics of the strongly coupled sector lead to the composite right-handed top quark t^c. In addition to the SM fields (without t^c), the weakly coupled sector in this case must involve a set of exotic fermions that also permits anomalies to be canceled. In particular, this set of exotic particles contains two SM singlet Majorana fermions N_2 and N_1. In general, all exotic fermions and all composite resonances except the pNGB states gain masses which are a few times larger than f. The pNGB states have masses which tend to be considerably lower than f. Therefore, they are the lightest composite resonances in the E_6CHM spectrum. In our analysis, N_1 is assumed to be the lightest exotic fermion; it has a mass around 10 TeV. The discrete Z_2^B symmetry, which is a subgroup of $U(1)_B$, forbids all couplings that allow N_1 to decay. When Z_2^B is an approximate symmetry, N_1 can be a long-lived composite state.

When the $SU(6)$ symmetry-breaking scale $f \gtrsim 10$ TeV and Z_2^L symmetry are almost exact, all operators that violate baryon number are sufficiently strongly suppressed even if $U(1)_B$ is explicitly broken. In this variant of the E_6CHM, the out-of-equilibrium decays $N_1 \to T^* + b$ and $N_1 \to T + \bar{b}$ can induce the observed baryon asymmetry if CP is violated. This scenario implies that the lifetime of N_1 is less than 10^{-15} s. Phenomenologically acceptable matter–antimatter asymmetry can be obtained, even in the limit when all CP-violating phases are small ($\lesssim$0.01). The electric dipole moments (EDMs) of atoms, neutrons and elementary states, which have not been observed in different experiments, are suppressed if CP invariance is approximately preserved. These EDMs, as well as baryon-number-violating processes, such as neutron-antineutron oscillations, are going to be searched for in the near future [263,264].

The lightest exotic fermion N_1 becomes absolutely stable if Z_2^L and Z_2^B are exact symmetries. In this limit, N_1 can account for all or some of the observed cold dark matter density if it has a mass M_1 which is much smaller than the scale f. In particular, when M_1 is close to half the mass of the SM singlet boson A, the annihilation cross section for $N_1N_1 \to$ SM particles can be relatively large, resulting in the cold dark matter density which is smaller than its measured value.

For a large $SU(6)$ symmetry-breaking scale, i.e., $f \gtrsim$ 5–10 TeV, all exotic fermions and almost all composite resonances are too heavy to be observed at the LHC. Because the deviations of the couplings of the SM-like Higgs boson to the SM particles within the E_6CHM are determined by v^2/f^2, the modifications of the appropriate Higgs branching fractions are negligibly small in this case. Therefore, it is going to be rather problematic to probe such small deviations at the LHC and future e^+e^- collider. Too large a value of f also implies that the interactions of the top quark with the SM particles are very similar to the ones which are predicted by the SM. This makes the top quark with its significant admixture of composite components basically indistinguishable from the corresponding SM state. Nonetheless, the spectrum of the E_6CHM must contain the $SU(3)_C$ triplet of scalar fields T ($T^\dagger$) with an electric charge $-1/3$ ($+1/3$) and a mass which is significantly lower than f. This state predominantly decays into $T \to \bar{t} + \bar{b}$. If this $SU(3)_C$ triplet T has a mass in the few TeV range, then it can be pair produced at the LHC, resulting in some enhancement of the cross section of $pp \to t\bar{t}b\bar{b}$. The discovery of such a colour state will provide a smoking gun signal of the composite Higgs model under consideration.

Funding: This research received no external funding.

Institutional Review Board Statement: Not applicable.

Informed Consent Statement: Not applicable.

Data Availability Statement: Data sharing is not applicable to this article.

Acknowledgments: R.N. thanks A.W. Thomas for sharing his valuable ideas in connection with composite Higgs models. R.N. acknowledges fruitful discussions with E. Akhmedov, E. Boos, M. Dubinin, S. Duplij, D. Gorbunov, D. Kazakov, V. Novikov, V. Rubakov, M. Vasiliev, M. Vysotsky and E. Zhemchugov.

Conflicts of Interest: The author declares no conflict of interest.

References

1. Georgi, H.; Glashow, S.L. Unity Of All Elementary Particle Forces. *Phys. Rev. Lett.* **1974**, *32*, 438. [CrossRef]
2. Minkowski, P. $\mu \to e\gamma$ at a Rate of One Out of 10^9 Muon Decays? *Phys. Lett. B* **1977**, *67*, 421. [CrossRef]
3. Mohapatra, R.N.; Senjanovic, G. Neutrino Mass and Spontaneous Parity Nonconservation. *Phys. Rev. Lett.* **1980**, *44*, 912. [CrossRef]
4. Witten, E. Dynamical Breaking of Supersymmetry. *Nucl. Phys. B* **1981**, *188*, 513. [CrossRef]
5. Sakai, N. Naturalness in Supersymmetric GUTs. *Z. Phys. C* **1981**, *11*, 153. [CrossRef]
6. Dimopoulos, S.; Georgi, H. Softly Broken Supersymmetry and SU(5). *Nucl. Phys. B* **1981**, *193*, 150. [CrossRef]
7. Kaul, R.K.; Majumdar, P. Cancellation of Quadratically Divergent Mass Corrections in Globally Supersymmetric Spontaneously Broken Gauge Theories. *Nucl. Phys. B* **1982**, *199*, 36. [CrossRef]
8. Gildener, E.; Weinberg, S. Symmetry Breaking and Scalar Bosons. *Phys. Rev. D* **1976**, *13*, 3333. [CrossRef]
9. Gildener, E. Gauge Symmetry Hierarchies. *Phys. Rev. D* **1976**, *14*, 1667. [CrossRef]
10. Chung, D.J.H.; Everett, L.L.; Kane, G.L.; King, S.F.; Lykken, J.; Wang, L.T. The soft supersymmetry-breaking Lagrangian: Theory and applications. *Phys. Rept.* **2005**, *407*, 1. [CrossRef]
11. Ellis, J.R.; Kelley, S.; Nanopoulos, D.V. Probing the desert using gauge coupling unification. *Phys. Lett. B* **1991**, *260*, 131. [CrossRef]
12. Langacker, P.; Luo, M.X. Implications of precision electroweak experiments for M_t, ρ_0, $\sin^2\theta_W$ and grand unification. *Phys. Rev. D* **1991**, *44*, 817. [CrossRef] [PubMed]
13. Amaldi, U.; de Boer, W.; Furstenau, H. Comparison of grand unified theories with electroweak and strong coupling constants measured at LEP. *Phys. Lett. B* **1991**, *260*, 447. [CrossRef]
14. Anselmo, F.; Cifarelli, L.; Peterman, A.; Zichichi, A. The Effective experimental constraints on $M_(susy)$ and $M_(gut)$. *Nuovo Cim. A* **1991**, *104*, 1817. [CrossRef]
15. Arkani-Hamed, N.; Dimopoulos, S.; Dvali, G.R. The Hierarchy problem and new dimensions at a millimeter. *Phys. Lett. B* **1998**, *429*, 263. [CrossRef]
16. Antoniadis, I.; Arkani-Hamed, N.; Dimopoulos, S.; Dvali, G.R. New dimensions at a millimeter to a Fermi and superstrings at a TeV. *Phys. Lett. B* **1998**, *436*, 257. [CrossRef]
17. Randall, L.; Sundrum, R. A Large mass hierarchy from a small extra dimension. *Phys. Rev. Lett.* **1999**, *83*, 3370. [CrossRef]
18. Randall, L.; Sundrum, R. An Alternative to compactification. *Phys. Rev. Lett.* **1999**, *83*, 4690. [CrossRef]
19. Dienes, K.R.; Dudas, E.; Gherghetta, T. Extra space-time dimensions and unification. *Phys. Lett. B* **1998**, *436*, 55. [CrossRef]
20. Dienes, K.R.; Dudas, E.; Gherghetta, T. Grand unification at intermediate mass scales through extra dimensions. *Nucl. Phys. B* **1999**, *537*, 47. [CrossRef]
21. Bellazzini, B.; Csáki, C.; Serra, J. Composite Higgses. *Eur. Phys. J. C* **2014**, *74*, 2766. [CrossRef]
22. Terazawa, H.; Akama, K.; Chikashige, Y. Unified Model of the Nambu-Jona-Lasinio Type for All Elementary Particle Forces. *Phys. Rev. D* **1977**, *15*, 480. [CrossRef]
23. Terazawa, H. Subquark Model of Leptons and Quarks. *Phys. Rev. D* **1980**, *22*, 184. [CrossRef]
24. Dimopoulos, S.; Preskill, J. Massless Composites With Massive Constituents. *Nucl. Phys. B* **1982**, *199*, 206. [CrossRef]
25. Kaplan, D.B.; Georgi, H. $SU(2) \times U(1)$ Breaking by Vacuum Misalignment. *Phys. Lett. B* **1984**, *136*, 183. [CrossRef]
26. Kaplan, D.B.; Georgi, H.; Dimopoulos, S. Composite Higgs Scalars. *Phys. Lett. B* **1984**, *136*, 187. [CrossRef]
27. Georgi, H.; Kaplan, D.B.; Galison, P. Calculation of the Composite Higgs Mass. *Phys. Lett. B* **1984**, *143*, 152. [CrossRef]
28. Banks, T. Constraints on $SU(2) \times U(1)$ Breaking by Vacuum Misalignment. *Nucl. Phys. B* **1984**, *243*, 125. [CrossRef]
29. Georgi, H.; Kaplan, D.B. Composite Higgs and Custodial $SU(2)$. *Phys. Lett. B* **1984**, *145*, 216. [CrossRef]
30. Dugan, M.J.; Georgi, H.; Kaplan, D.B. Anatomy of a Composite Higgs Model. *Nucl. Phys. B* **1985**, *254*, 299. [CrossRef]
31. Georgi, H. A Tool Kit for Builders of Composite Models. *Nucl. Phys. B* **1986**, *266*, 274. [CrossRef]
32. Nevzorov, R.; Thomas, A.W. E_6 inspired composite Higgs model. *Phys. Rev. D* **2015**, *92*, 075007. [CrossRef]
33. Nevzorov, R.; Thomas, A.W. LHC signatures of neutral pseudo-Goldstone boson in the E_6CHM. *J. Phys. G* **2017**, *44*, 075003. [CrossRef]
34. Nevzorov, R.; Thomas, A.W. E_6 inspired composite Higgs model and 750 GeV diphoton excess. *EPJ Web Conf.* **2016**, *125*, 02021. [CrossRef]
35. Nevzorov, R.; Thomas, A.W. Baryon asymmetry generation in the E_6CHM. *Phys. Lett. B* **2017**, *774*, 123. [CrossRef]
36. Nevzorov, R.; Thomas, A.W. Generation of baryon asymmetry in the E_6CHM. *EPJ Web Conf.* **2018**, *191*, 02004. [CrossRef]

37. Nevzorov, R.; Thomas, A.W. E_6 Inspired Composite Higgs Model and Baryon Asymmetry Generation. *Phys. Part. Nucl.* **2020**, *51*, 709. [CrossRef]
38. Khlopov, M.Y.; Shibaev, K.I. New physics from superstring phenomenology. *Grav. Cosmol. Suppl.* **2002**, *8*, 45.
39. Khlopov, M.Y. What comes after the Standard model? *Prog. Part. Nucl. Phys.* **2021**, *116*, 103824. [CrossRef]
40. Sakharov, A.D. Violation of CP invariance, C asymmetry, and baryon asymmetry of the universe. *JETP Lett.* **1967**, *5*, 24.
41. Ignatiev, A.Y.; Krasnikov, N.V.; Kuzmin, V.A.; Tavkhelidze, A.N. Universal CP Noninvariant Superweak Interaction and Baryon Asymmetry of the Universe. *Phys. Lett. B* **1978**, *76*, 436. [CrossRef]
42. Yoshimura, M. Unified Gauge Theories and the Baryon Number of the Universe. *Phys. Rev. Lett.* **1978**, *41*, 281. [CrossRef]
43. Toussaint, D.; Treiman, S.B.; Wilczek, F.; Zee, A. Matter - Antimatter Accounting, Thermodynamics, and Black Hole Radiation. *Phys. Rev. D* **1979**, *19*, 1036. [CrossRef]
44. Weinberg, S. Cosmological production of baryons. *Phys. Rev. Lett.* **1979**, *42*, 850. [CrossRef]
45. Yoshimura, M. Origin of Cosmological Baryon Asymmetry. *Phys. Lett. B* **1979**, *88*, 294. [CrossRef]
46. Barr, S.M.; Segre, G.; Weldon, H.A. The Magnitude of the Cosmological Baryon Asymmetry. *Phys. Rev. D* **1979**, *20*, 2494. [CrossRef]
47. Nanopoulos, D.V.; Weinberg, S. Mechanisms for Cosmological Baryon Production. *Phys. Rev. D* **1979**, *20*, 2484. [CrossRef]
48. Fukugita, M.; Yanagida, T. Baryogenesis Without Grand Unification. *Phys. Lett. B* **1986**, *174*, 45. [CrossRef]
49. Affleck, I.; Dine, M. A New Mechanism for Baryogenesis. *Nucl. Phys. B* **1985**, *249*, 361. [CrossRef]
50. Dine, M.; Randall, L.; Thomas, S.D. Baryogenesis from flat directions of the supersymmetric standard model. *Nucl. Phys. B* **1996**, *458*, 291. [CrossRef]
51. Riotto, A.; Trodden, M. Recent progress in baryogenesis. *Ann. Rev. Nucl. Part. Sci.* **1999**, *49*, 35. [CrossRef]
52. Agashe, K.; Contino, R.; Pomarol, A. The Minimal composite Higgs model. *Nucl. Phys. B* **2005**, *719*, 165. [CrossRef]
53. Contino, R.; Nomura, Y.; Pomarol, A. Higgs as a holographic pseudoGoldstone boson. *Nucl. Phys. B* **2003**, *671*, 148. [CrossRef]
54. Agashe, K.; Delgado, A.; May, M.J.; Sundrum, R. RS1, custodial isospin and precision tests. *J. High Energy Phys.* **2003**, *308*, 50. [CrossRef]
55. Contino, R.; Kramer, T.; Son, M.; Sundrum, R. Warped/composite phenomenology simplified. *J. High Energy Phys.* **2007**, *705*, 74. [CrossRef]
56. Kaplan, D.B. Flavor at SSC energies: A New mechanism for dynamically generated fermion masses. *Nucl. Phys. B* **1991**, *365*, 259. [CrossRef]
57. Frigerio, M.; Serra, J.; Varagnolo, A. Composite GUTs: Models and expectations at the LHC. *J. High Energy Phys.* **2011**, *1106*, 29. [CrossRef]
58. Agashe, K.; Contino, R. The Minimal composite Higgs model and electroweak precision tests. *Nucl. Phys. B* **2006**, *742*, 59. [CrossRef]
59. Agashe, K.; Contino, R.; Da Rold, L.; Pomarol, A. A Custodial symmetry for $Zb\bar{b}$. *Phys. Lett. B* **2006**, *641*, 62. [CrossRef]
60. Giudice, G.F.; Grojean, C.; Pomarol, A.; Rattazzi, R. The Strongly-Interacting Light Higgs. *J. High Energy Phys.* **2007**, *706*, 45. [CrossRef]
61. Barbieri, R.; Bellazzini, B.; Rychkov, V.S.; Varagnolo, A. The Higgs boson from an extended symmetry. *Phys. Rev. D* **2007**, *76*, 115008. [CrossRef]
62. Lodone, P. Vector-like quarks in a 'composite' Higgs model. *J. High Energy Phys.* **2008**, *0812*, 029. [CrossRef]
63. Gillioz, M. A Light composite Higgs boson facing electroweak precision tests. *Phys. Rev. D* **2009**, *80*, 055003. [CrossRef]
64. Anastasiou, C.; Furlan, E.; Santiago, J. Realistic Composite Higgs Models. *Phys. Rev. D* **2009**, *79*, 075003. [CrossRef]
65. Panico, G.; Wulzer, A. The Discrete Composite Higgs Model. *J. High Energy Phys.* **2011**, *1109*, 135. [CrossRef]
66. De Curtis, S.; Redi, M.; Tesi, A. The 4D Composite Higgs. *J. High Energy Phys.* **2012**, *1204*, 42. [CrossRef]
67. Marzocca, D.; Serone, M.; Shu, J. General Composite Higgs Models. *J. High Energy Phys.* **2012**, *1208*, 013. [CrossRef]
68. Orgogozo, A.; Rychkov, S. The S parameter for a Light Composite Higgs: A Dispersion Relation Approach. *J. High Energy Phys.* **2013**, *1306*, 14. [CrossRef]
69. Grojean, C.; Matsedonskyi, O.; Giuliano, P. Light top partners and precision physics. *J. High Energy Phys.* **2013**, *1310*, 160. [CrossRef]
70. Carena, M.; Ponton, E.; Santiago, J.; Wagner, C.E.M. Light Kaluza Klein States in Randall-Sundrum Models with Custodial SU(2). *Nucl. Phys. B* **2006**, *759*, 202. [CrossRef]
71. Pomarol, A.; Serra, J. Top Quark Compositeness: Feasibility and Implications. *Phys. Rev. D* **2008**, *78*, 074026. [CrossRef]
72. Pappadopulo, D.; Thamm, A.; Torre, R. A minimally tuned composite Higgs model from an extra dimension. *J. High Energy Phys.* **2013**, *1307*, 58. [CrossRef]
73. Bellazzini, B.; Csaki, C.; Hubisz, J.; Serra, J.; Terning, J. Composite Higgs Sketch. *J. High Energy Phys.* **2012**, *1211*, 3. [CrossRef]
74. Gillioz, M.; Grober, R.; Grojean, C.; Muhlleitner, M.; Salvioni, E. Higgs Low-Energy Theorem (and its corrections) in Composite Models. *J. High Energy Phys.* **2012**, *1210*, 4. [CrossRef]
75. Azatov, A.; Galloway, J. Electroweak Symmetry Breaking and the Higgs Boson: Confronting Theories at Colliders. *Int. J. Mod. Phys. A* **2013**, *28*, 1330004. [CrossRef]
76. Falkowski, A.; Riva, F.; Urbano, A. Higgs at last. *J. High Energy Phys.* **2013**, *1311*, 111. [CrossRef]
77. Azatov, A.; Contino, R.; Di Iura, A.; Galloway, J. New Prospects for Higgs Compositeness in $h \to Z\gamma$. *Phys. Rev. D* **2013**, *88*, 075019. [CrossRef]

78. Gillioz, M.; Gröber, R.; Kapuvari, A.; Mühlleitner, M. Vector-like Bottom Quarks in Composite Higgs Models. *J. High Energy Phys.* **2014**, *1403*, 37. [CrossRef]
79. Barbieri, R.; Isidori, G.; Pappadopulo, D. Composite fermions in Electroweak Symmetry Breaking. *J. High Energy Phys.* **2009**, *902*, 29. [CrossRef]
80. Matsedonskyi, O. On Flavour and Naturalness of Composite Higgs Models. *J. High Energy Phys.* **2015**, *1502*, 154. [CrossRef]
81. Barbieri, R.; Buttazzo, D.; Sala, F.; Straub, D.M.; Tesi, A. A 125 GeV composite Higgs boson versus flavour and electroweak precision tests. *J. High Energy Phys.* **2013**, *1305*, 69. [CrossRef]
82. Csaki, C.; Falkowski, A.; Weiler, A. The Flavor of the Composite Pseudo-Goldstone Higgs. *J. High Energy Phys.* **2008**, *809*, 8. [CrossRef]
83. Agashe, K.; Azatov, A.; Zhu, L. Flavor Violation Tests of Warped/Composite SM in the Two-Site Approach. *Phys. Rev. D* **2009**, *79*, 056006. [CrossRef]
84. Vignaroli, N. Δ F=1 constraints on composite Higgs models with LR parity. *Phys. Rev. D* **2012**, *86*, 115011. [CrossRef]
85. Sikivie, P.; Susskind, L.; Voloshin M.B.; Zakharov, V.I. Isospin Breaking in Technicolor Models. *Nucl. Phys. B* **1980**, *173*, 189. [CrossRef]
86. Peskin, M.E.; Takeuchi, T. Estimation of oblique electroweak corrections. *Phys. Rev. D* **1992**, *46*, 381. [CrossRef]
87. Agashe, K.; Perez, G.; Soni, A. Flavor structure of warped extra dimension models. *Phys. Rev. D* **2005**, *71*, 016002. [CrossRef]
88. Glashow, S.L.; Iliopoulos, J.; Maiani, L. Weak Interactions with Lepton-Hadron Symmetry. *Phys. Rev. D* **1970**, *2*, 1285. [CrossRef]
89. Redi, M.; Weiler, A. Flavor and CP Invariant Composite Higgs Models. *J. High Energy Phys.* **2011**, *1111*, 108. [CrossRef]
90. Blanke, M.; Buras, A.J.; Duling, B.; Gori, S.; Weiler, A. Δ F = 2 Observables and Fine-Tuning in a Warped Extra Dimension with Custodial Protection. *J. High Energy Phys.* **2009**, *903*, 1. [CrossRef]
91. Gedalia, O.; Isidori, G.; Perez, G. Combining Direct & Indirect Kaon CP Violation to Constrain the Warped KK Scale. *Phys. Lett. B* **2009**, *682*, 200.
92. Barbieri, R.; Buttazzo, D.; Sala, F.; Straub, D.M. Flavour physics from an approximate $U(2)^3$ symmetry. *J. High Energy Phys.* **2012**, *1207*, 181. [CrossRef]
93. Redi, M. Leptons in Composite MFV. *J. High Energy Phys.* **2013**, *1309*, 60. [CrossRef]
94. Agashe, K.; Blechman, A.E.; Petriello, F. Probing the Randall-Sundrum geometric origin of flavor with lepton flavor violation. *Phys. Rev. D* **2006**, *74*, 053011. [CrossRef]
95. Csaki, C.; Grossman, Y.; Tanedo, P.; Tsai, Y. Warped penguin diagrams. *Phys. Rev. D* **2011**, *83*, 073002. [CrossRef]
96. Csaki, C.; Delaunay, C.; Grojean, C.; Grossman, Y. A Model of Lepton Masses from a Warped Extra Dimension. *J. High Energy Phys.* **2008**, *0810*, 055. [CrossRef]
97. del Aguila, F.; Carmona, A.; Santiago, J. Neutrino Masses from an A4 Symmetry in Holographic Composite Higgs Models. *J. High Energy Phys.* **2010**, *1008*, 127. [CrossRef]
98. Cacciapaglia, G.; Csaki, C.; Galloway, J.; Marandella, G.; Terning, J.; Weiler, A. A GIM Mechanism from Extra Dimensions. *J. High Energy Phys.* **2008**, *804*, 6. [CrossRef]
99. Redi, M. Composite MFV and Beyond. *Eur. Phys. J. C* **2012**, *72*, 2030. [CrossRef]
100. König, M.; Neubert, M.; Straub, D.M. Dipole operator constraints on composite Higgs models. *Eur. Phys. J. C* **2014**, *74*, 2945. [CrossRef]
101. Gripaios, B.; Pomarol, A.; Riva, F.; Serra, J. Beyond the Minimal Composite Higgs Model. *J. High Energy Phys.* **2009**, *904*, 70. [CrossRef]
102. Mrazek, J.; Pomarol, A.; Rattazzi, R.; Redi, M.; Serra, J.; Wulzer, A. The Other Natural Two Higgs Doublet Model. *Nucl. Phys. B* **2011**, *853*, 1. [CrossRef]
103. Redi, M.; Tesi, A. Implications of a Light Higgs in Composite Models. *J. High Energy Phys.* **2012**, *1210*, 166. [CrossRef]
104. Bertuzzo, E.; Ray, T.S.; de Sandes, H.; Savoy, C.A. On Composite Two Higgs Doublet Models. *J. High Energy Phys.* **2013**, *1305*, 153. [CrossRef]
105. Montull, M.; Riva, F. Higgs discovery: The beginning or the end of natural EWSB? *J. High Energy Phys.* **2012**, *1211*, 18. [CrossRef]
106. Chala, M. $h \to \gamma\gamma$ excess and Dark Matter from Composite Higgs Models. *J. High Energy Phys.* **2013**, *1301*, 122. [CrossRef]
107. Frigerio, M.; Pomarol, A.; Riva, F.; Urbano, A. Composite Scalar Dark Matter. *J. High Energy Phys.* **2012**, *1207*, 15. [CrossRef]
108. Contino, R.; Grojean, C.; Moretti, M.; Piccinini, F.; Rattazzi, R. Strong Double Higgs Production at the LHC. *J. High Energy Phys.* **2010**, *1005*, 89. [CrossRef]
109. Low, I.; Vichi, A. On the production of a composite Higgs boson. *Phys. Rev. D* **2011**, *84*, 045019. [CrossRef]
110. Contino, R.; Marzocca, D.; Pappadopulo, D.; Rattazzi, R. On the effect of resonances in composite Higgs phenomenology. *J. High Energy Phys.* **2011**, *1110*, 81. [CrossRef]
111. Azatov, A.; Galloway, J. Light Custodians and Higgs Physics in Composite Models. *Phys. Rev. D* **2012**, *85*, 055013. [CrossRef]
112. Contino, R.; Ghezzi, M.; Moretti, M.; Panico, G.; Piccinini, F.; Wulzer, A. Anomalous Couplings in Double Higgs Production. *J. High Energy Phys.* **2012**, *1208*, 154. [CrossRef]
113. Contino, R.; Ghezzi, M.; Grojean, C.; Muhlleitner, M.; Spira, M. Effective Lagrangian for a light Higgs-like scalar. *J. High Energy Phys.* **2013**, *1307*, 35. [CrossRef]
114. Delaunay, C.; Grojean, C.; Perez, G. Modified Higgs Physics from Composite Light Flavors. *J. High Energy Phys.* **2013**, *1309*, 90. [CrossRef]
115. Banfi, A.; Martin, A.; Sanz, V. Probing top-partners in Higgs+jets. *J. High Energy Phys.* **2014**, *1408*, 53. [CrossRef]

116. Montull, M.; Riva, F.; Salvioni, E.; Torre, R. Higgs Couplings in Composite Models. *Phys. Rev. D* **2013**, *88*, 095006. [CrossRef]
117. Contino, R.; Grojean, C.; Pappadopulo, D.; Rattazzi, R.; Thamm, A. Strong Higgs Interactions at a Linear Collider. *J. High Energy Phys.* **2014**, *1402*, 6. [CrossRef]
118. Flacke, T.; Kim, J.H.; Lee, S.J.; Lim, S.H. Constraints on composite quark partners from Higgs searches. *J. High Energy Phys.* **2014**, *1405*, 123. [CrossRef]
119. Grojean, C.; Salvioni, E.; Schlaffer, M.; Weiler, A. Very boosted Higgs in gluon fusion. *J. High Energy Phys.* **2014**, *1405*, 22. [CrossRef]
120. Carena, M.; Da Rold, L.; Pontón, E. Minimal Composite Higgs Models at the LHC. *J. High Energy Phys.* **2014**, *1406*, 159. [CrossRef]
121. Carmona, A.; Goertz, F. A naturally light Higgs without light Top Partners. *J. High Energy Phys.* **2015**, *1505*, 2. [CrossRef]
122. Buchalla, G.; Cata, O.; Krause, C. A Systematic Approach to the SILH Lagrangian. *Nucl. Phys. B* **2015**, *894*, 602. [CrossRef]
123. Pomarol, A.; Riva, F. The Composite Higgs and Light Resonance Connection. *J. High Energy Phys.* **2012**, *1208*, 135. [CrossRef]
124. Matsedonskyi, O.; Panico, G.; Wulzer, A. Light Top Partners for a Light Composite Higgs. *J. High Energy Phys.* **2013**, *1301*, 164. [CrossRef]
125. Agashe, K.; Delgado, A.; Sundrum, R. Grand unification in RS1. *Ann. Phys.* **2003**, *304*, 145. [CrossRef]
126. Gherghetta, T. Partly supersymmetric grand unification. *Phys. Rev. D* **2005**, *71*, 065001. [CrossRef]
127. Barnard, J.; Gherghetta, T.; Ray, T.S.; Spray, A. The Unnatural Composite Higgs. *J. High Energy Phys.* **2015**, *1501*, 67. [CrossRef]
128. Asano, M.; Kitano, R. Partially Composite Dark Matter. *J. High Energy Phys.* **2014**, *1409*, 171. [CrossRef]
129. Agashe, K.; Belyaev, A.; Krupovnickas, T.; Perez, G.; Virzi, J. LHC Signals from Warped Extra Dimensions. *Phys. Rev. D* **2008**, *77*, 015003. [CrossRef]
130. Lillie, B.; Randall, L.; Wang, L.T. The Bulk RS KK-gluon at the LHC. *J. High Energy Phys.* **2007**, *0709*, 074. [CrossRef]
131. Agashe, K.; Davoudiasl, H.; Gopalakrishna, S.; Han, T.; Huang, G.Y.; Perez, G.; Si, Z.G.; Soni, A. LHC Signals for Warped Electroweak Neutral Gauge Bosons. *Phys. Rev. D* **2007**, *76*, 115015. [CrossRef]
132. Carena, M.; Medina, A.D.; Panes, B.; Shah, N.R.; Wagner, C.E.M. Collider phenomenology of gauge-Higgs unification scenarios in warped extra dimensions. *Phys. Rev. D* **2008**, *77*, 076003. [CrossRef]
133. Contino, R.; Servant, G. Discovering the top partners at the LHC using same-sign dilepton final states. *J. High Energy Phys.* **2008**, *0806*, 026. [CrossRef]
134. Agashe, K.; Gopalakrishna, S.; Han, T.; Huang, G.Y.; Soni, A. LHC Signals for Warped Electroweak Charged Gauge Bosons. *Phys. Rev. D* **2009**, *80*, 075007. [CrossRef]
135. Aguilar-Saavedra, J.A. Identifying top partners at LHC. *J. High Energy Phys.* **2009**, *0911*, 030. [CrossRef]
136. Mrazek, J.; Wulzer, A. A Strong Sector at the LHC: Top Partners in Same-Sign Dileptons. *Phys. Rev. D* **2010**, *81*, 075006. [CrossRef]
137. Agashe, K.; Azatov, A.; Han, T.; Li, Y.; Si, Z.G.; Zhu, L. LHC Signals for Coset Electroweak Gauge Bosons in Warped/Composite PGB Higgs Models. *Phys. Rev. D* **2010**, *81*, 096002. [CrossRef]
138. Dissertori, G.; Furlan, E.; Moortgat, F.; Nef, P. Discovery potential of top-partners in a realistic composite Higgs model with early LHC data. *J. High Energy Phys.* **2010**, *1009*, 19. [CrossRef]
139. Vignaroli, N. Early discovery of top partners and test of the Higgs nature. *Phys. Rev. D* **2012**, *86*, 075017. [CrossRef]
140. Cacciapaglia, G.; Deandrea, A.; Panizzi, L.; Perries, S.; Sordini, V. Heavy Vector-like quark with charge 5/3 at the LHC. *J. High Energy Phys.* **2013**, *1303*, 4. [CrossRef]
141. De Simone, A.; Matsedonskyi, O.; Rattazzi, R.; Wulzer, A. A First Top Partner Hunter's Guide. *J. High Energy Phys.* **2013**, *1304*, 4. [CrossRef]
142. Li, J.; Liu, D.; Shu, J. Towards the fate of natural composite Higgs model through single t' search at the 8 TeV LHC. *J. High Energy Phys.* **2013**, *1311*, 47. [CrossRef]
143. Redi, M.; Sanz, V.; de Vries, M.; Weiler, A. Strong Signatures of Right-Handed Compositeness. *J. High Energy Phys.* **2013**, *1308*, 8. [CrossRef]
144. Delaunay, C.; Flacke, T.; Gonzalez-Fraile, J.; Lee, S.J.; Panico, G.; Perez, G. Light Non-degenerate Composite Partners at the LHC. *J. High Energy Phys.* **2014**, *1402*, 55. [CrossRef]
145. Matsedonskyi, O.; Riva, F.; Vantalon, T. Composite Charge 8/3 Resonances at the LHC. *J. High Energy Phys.* **2014**, *1404*, 59. [CrossRef]
146. Cheng, H.C.; Gu, J. Top seesaw with a custodial symmetry, and the 126 GeV Higgs. *J. High Energy Phys.* **2014**, *1410*, 2. [CrossRef]
147. Gripaios, B.; Müller, T.; Parker, M.A.; Sutherland, D. Search Strategies for Top Partners in Composite Higgs models. *J. High Energy Phys.* **2014**, *1408*, 171. [CrossRef]
148. Azatov, A.; Panico, G.; Perez, G.; Soreq, Y. On the Flavor Structure of Natural Composite Higgs Models & Top Flavor Violation. *J. High Energy Phys.* **2014**, *1412*, 82.
149. Backović, M.; Flacke, T.; Kim, J.H.; Lee, S.J. Boosted Event Topologies from TeV Scale Light Quark Composite Partners. *J. High Energy Phys.* **2015**, *1504*, 82. [CrossRef]
150. Kanemura, S.; Kaneta, K.; Machida, N.; Shindou, T. New resonance scale and fingerprint identification in minimal composite Higgs models. *Phys. Rev. D* **2015**, *91*, 115016. [CrossRef]
151. Thamm, A.; Torre, R.; Wulzer, A. Future tests of Higgs compositeness: Direct vs indirect. *J. High Energy Phys.* **2015**, *1507*, 100. [CrossRef]
152. Azatov, A.; Chowdhury, D.; Ghosh, D.; Ray, T.S. Same sign di-lepton candles of the composite gluons. *J. High Energy Phys.* **2015**, *1508*, 140. [CrossRef]

153. Serra, J. Beyond the Minimal Top Partner Decay. *J. High Energy Phys.* **2015**, *1509*, 176. [CrossRef]
154. Barnard, J.; Gherghetta, T.; Ray, T.S. UV descriptions of composite Higgs models without elementary scalars. *J. High Energy Phys.* **2014**, *1402*, 2, [arXiv:1311.6562 [hep-ph]]. [CrossRef]
155. Ferretti, G.; Karateev, D. Fermionic UV completions of Composite Higgs models. *J. High Energy Phys.* **2014**, *1403*, 077. [CrossRef]
156. Cacciapaglia, G.; Sannino, F. Fundamental Composite (Goldstone) Higgs Dynamics. *J. High Energy Phys.* **2014**, *1404*, 111. [CrossRef]
157. Hietanen, A.; Lewis, R.; Pica, C.; Sannino, F. Fundamental Composite Higgs Dynamics on the Lattice: SU(2) with Two Flavors. *J. High Energy Phys.* **2014**, *1407*, 116. [CrossRef]
158. Ferretti, G. UV Completions of Partial Compositeness: The Case for a SU(4) Gauge Group. *J. High Energy Phys.* **2014**, *1406*, 142. [CrossRef]
159. Parolini, A. Phenomenological aspects of supersymmetric composite Higgs models. *Phys. Rev. D* **2014**, *90*, 115026. [CrossRef]
160. Geller, M.; Telem, O. Holographic Twin Higgs Model. *Phys. Rev. Lett.* **2015**, *114*, 191801. [CrossRef]
161. Gripaios, B.; Nardecchia, M.; Renner, S.A. Composite leptoquarks and anomalies in *B*-meson decays. *J. High Energy Phys.* **2015**, *1505*, 6. [CrossRef]
162. Low, M.; Tesi, A.; Wang, L.T. Twin Higgs mechanism and a composite Higgs boson. *Phys. Rev. D* **2015**, *91*, 095012. [CrossRef]
163. Golterman, M.; Shamir, Y. Top quark induced effective potential in a composite Higgs model. *Phys. Rev. D* **2015**, *91*, 094506. [CrossRef]
164. Agashe, K.; Contino, R.; Sundrum, R. Top compositeness and precision unification. *Phys. Rev. Lett.* **2005**, *95*, 171804. [CrossRef]
165. Kawamura, Y. Triplet doublet splitting, proton stability and extra dimension. *Prog. Theor. Phys.* **2001**, *105*, 999. [CrossRef]
166. Altarelli, G.; Feruglio, F. SU(5) grand unification in extra dimensions and proton decay. *Phys. Lett. B* **2001**, *511*, 257. [CrossRef]
167. Hebecker, A.; March-Russell, J. A Minimal $S^1/(Z_2 \times Z_2')$ orbifold GUT. *Nucl. Phys. B* **2001**, *613*, 3. [CrossRef]
168. Barbieri, R.; Hall, L.J.; Nomura, Y. Softly broken supersymmetric desert from orbifold compactification. *Phys. Rev. D* **2002**, *66*, 045025. [CrossRef]
169. Haba, N.; Shimizu, Y.; Suzuki, T.; Ukai, K. Fermion mass hierarchy in the grand unified theory on $S_1/(Z_2 \times Z_2')$ orbifold. *Prog. Theor. Phys.* **2002**, *107*, 151. [CrossRef]
170. Barr, S.M.; Dorsner, I. Unifying flipped SU(5) in five-dimensions. *Phys. Rev. D* **2002**, *66*, 065013. [CrossRef]
171. Hebecker, A.; March-Russell, J. The Flavor hierarchy and seesaw neutrinos from bulk masses in 5-d orbifold GUTs. *Phys. Lett. B* **2002**, *541*, 338. [CrossRef]
172. Paccetti Correia, F.; Schmidt, M.G.; Tavartkiladze, Z. 5-D SUSY orbifold SU(6) GUT and pseudoGoldstone Higgs doublets. *Phys. Lett. B* **2002**, *545*, 153. [CrossRef]
173. Hebecker, A.; March-Russell, J.; Yanagida, T. Higher dimensional origin of heavy sneutrino domination and low scale leptogenesis. *Phys. Lett. B* **2003**, *552*, 229. [CrossRef]
174. Kim, H.D.; Raby, S. Neutrinos in 5-D SO(10) unification. *J. High Energy Phys.* **2003**, *0307*, 14. [CrossRef]
175. Bhattacharyya, G.; Branco, G.C.; Silva-Marcos, J.I. CP Violation and Flavour Mixings in Orbifold GUTs. *Phys. Rev. D* **2008**, *77*, 011901. [CrossRef]
176. Kobakhidze, A.B. Proton stability in TeV scale GUTs. *Phys. Lett. B* **2001**, *514*, 131. [CrossRef]
177. Hebecker, A.; March-Russell, J. Proton decay signatures of orbifold GUTs. *Phys. Lett. B* **2002**, *539*, 119. [CrossRef]
178. Shafi, Q.; Tavartkiladze, Z. Neutrino democracy and other phenomenology from 5-D SO(10). *Nucl. Phys. B* **2003**, *665*, 469. [CrossRef]
179. Hall, L.J.; Nomura, Y. Gauge unification in higher dimensions. *Phys. Rev. D* **2001**, *64*, 055003. [CrossRef]
180. Nomura, Y. Strongly coupled grand unification in higher dimensions. *Phys. Rev. D* **2002**, *65*, 085036. [CrossRef]
181. Hall, L.J.; Nomura, Y. Gauge coupling unification from unified theories in higher dimensions. *Phys. Rev. D* **2002**, *65*, 125012. [CrossRef]
182. Dermisek, R.; Mafi, A. SO(10) grand unification in five-dimensions: Proton decay and the mu problem. *Phys. Rev. D* **2002**, *65*, 055002. [CrossRef]
183. Hall, L.J.; Nomura, Y. A Complete theory of grand unification in five-dimensions. *Phys. Rev. D* **2002**, *66*, 075004. [CrossRef]
184. Kim, H.D.; Raby, S. Unification in 5-D SO(10). *J. High Energy Phys.* **2003**, *0301*, 56. [CrossRef]
185. Dorsner, I. Flipping SU(5) towards five-dimensional unification. *Phys. Rev. D* **2004**, *69*, 056003. [CrossRef]
186. Kim, H.-D.; Kim, J.E.; Lee, H.M. Top - bottom mass hierarchy, s - mu puzzle and gauge coupling unification with split multiplets. *Eur. Phys. J. C* **2002**, *24*, 159. [CrossRef]
187. Correia, F.P.; Schmidt, M.G.; Tavartkiladze, Z. Gauge coupling unification and phenomenology of selected orbifold 5-D N=1 SUSY models. *Nucl. Phys. B* **2003**, *649*, 39. [CrossRef]
188. Hall, L.J.; Nomura, Y.; Tucker-Smith, D. Gauge Higgs unification in higher dimensions. *Nucl. Phys. B* **2002**, *639*, 307. [CrossRef]
189. Hall, L.J.; March-Russell, J.; Okui, T.; Tucker-Smith, D. Towards a theory of flavor from orbifold GUTs. *J. High Energy Phys.* **2004**, *0409*, 26. [CrossRef]
190. Babu, K.S.; Barr, S.M.; Kyae, B.-S. Family unification in five-dimensions and six-dimensions. *Phys. Rev. D* **2002**, *65*, 115008. [CrossRef]
191. Kim, H.D.; Raby, S.; Schradin, L. Quark and lepton masses in 5-D SO(10). *J. High Energy Phys.* **2005**, *0505*, 36. [CrossRef]
192. Forste, S.; Nilles, H.P.; Wingerter, A. Geometry of rank reduction. *Phys. Rev. D* **2005**, *72*, 026001. [CrossRef]

193. Braam, F.; Knochel, A.; Reuter, J. An Exceptional SSM from E_6 Orbifold GUTs with intermediate LR symmetry. *J. High Energy Phys.* **2010**, *1006*, 13. [CrossRef]
194. Li, T.-J. GUT breaking on $M^4 \times T^2$ / $(Z_2 \times Z_2')$. *Phys. Lett. B* **2001**, *520*, 377. [CrossRef]
195. Asaka, T.; Buchmuller, W.; Covi, L. Gauge unification in six-dimensions. *Phys. Lett. B* **2001**, *523*, 199. [CrossRef]
196. Li, T.-J. N=2 supersymmetric GUT breaking on T^2 orbifolds. *Nucl. Phys. B* **2001**, *619*, 75. [CrossRef]
197. Haba, N.; Kondo, T.; Shimizu, Y. Fermion mass hierarchy in six-dimensional SO(10) grand unified theory on a T^2/Z_2 orbifold. *Phys. Lett. B* **2002**, *531*, 245. [CrossRef]
198. Watari, T.; Yanagida, T. Higher dimensional supersymmetry as an origin of the three families for quarks and leptons. *Phys. Lett. B* **2002**, *532*, 252. [CrossRef]
199. Haba, N.; Kondo, T.; Shimizu, Y. Fermion mass hierarchy in six- dimensional SO(10) SUSY GUT. *Phys. Lett. B* **2002**, *535*, 271. [CrossRef]
200. Watari, T.; Yanagida, T. Geometric origin of large lepton mixing in a higher dimensional space-time. *Phys. Lett. B* **2002**, *544*, 167. [CrossRef]
201. Asaka, T.; Buchmuller, W.; Covi, L. Exceptional coset spaces and unification in six-dimensions. *Phys. Lett. B* **2002**, *540*, 295. [CrossRef]
202. Hebecker, A.; Ratz, M. Group theoretical aspects of orbifold and conifold GUTs. *Nucl. Phys. B* **2003**, *670*, 3. [CrossRef]
203. Asaka, T.; Buchmuller, W.; Covi, L. Quarks and leptons between branes and bulk. *Phys. Lett. B* **2003**, *563*, 209. [CrossRef]
204. Buchmuller, W.; Kersten, J.; Schmidt-Hoberg, K. Squarks and sleptons between branes and bulk. *J. High Energy Phys.* **2006**, *0602*, 69. [CrossRef]
205. Buchmuller, W.; Covi, L.; Emmanuel-Costa, D.; Wiesenfeldt, S. CP Violation and Neutrino Masses and Mixings from Quark Mass Hierarchies. *J. High Energy Phys.* **2007**, *0712*, 30. [CrossRef]
206. Hall, L.J.; Nomura, Y.; Okui, T.; Tucker-Smith, D. SO(10) unified theories in six-dimensions. *Phys. Rev. D* **2002**, *65*, 035008. [CrossRef]
207. Lee, H.M. Gauge coupling unification in six dimensions. *Phys. Rev. D* **2007**, *75*, 065009. [CrossRef]
208. Buchmuller, W.; Covi, L.; Emmanuel-Costa, D.; Wiesenfeldt, S. Flavour structure and proton decay in 6D orbifold GUTs. *J. High Energy Phys.* **2004**, *0409*, 4. [CrossRef]
209. Candelas, P.; Horowitz, G.T.; Strominger, A.; Witten, E. Vacuum Configurations for Superstrings. *Nucl. Phys. B* **1985**, *258*, 46. [CrossRef]
210. Witten, E. Symmetry Breaking Patterns in Superstring Models. *Nucl. Phys. B* **1985**, *258*, 75. [CrossRef]
211. Dixon, L.J.; Harvey, J.A.; Vafa, C.; Witten, E. Strings on Orbifolds. *Nucl. Phys. B* **1985**, *261*, 678. [CrossRef]
212. Breit, J.D.; Ovrut, B.A.; Segre, G.C. E_6 Symmetry Breaking in the Superstring Theory. *Phys. Lett. B* **1985**, *158*, 33. [CrossRef]
213. Dixon, L.J.; Harvey, J.A.; Vafa, C.; Witten, E. Strings on Orbifolds 2. *Nucl. Phys. B* **1986**, *274*, 285. [CrossRef]
214. Sen, A. Naturally Light Higgs Doublet in Supersymmetric E_6 Grand Unified Theory. *Phys. Rev. Lett.* **1985**, *55*, 33. [CrossRef] [PubMed]
215. Ibanez, L.E.; Kim, J.E.; Nilles, H.P.; Quevedo, F. Orbifold Compactifications with Three Families of SU(3) $\times$ SU(2) $\times$ U(1)^n. *Phys. Lett. B* **1987**, *191*, 282. [CrossRef]
216. Kobayashi, T.; Raby, S.; Zhang, R.-J. Constructing 5-D orbifold grand unified theories from heterotic strings. *Phys. Lett. B* **2004** *593*, 262. [CrossRef]
217. Kobayashi, T.; Raby, S.; Zhang, R.-J. Searching for realistic 4d string models with a Pati-Salam symmetry: Orbifold grand unified theories from heterotic string compactification on a Z(6) orbifold. *Nucl. Phys. B* **2005**, *704*, 3. [CrossRef]
218. Buchmuller, W.; Hamaguchi, K.; Lebedev, O.; Ratz, M. Supersymmetric standard model from the heterotic string. *Phys. Rev. Lett.* **2006**, *96*, 121602. [CrossRef] [PubMed]
219. Lebedev, O.; Nilles, H.P.; Raby, S.; Ramos-Sanchez, S.; Ratz, M.; Vaudrevange, P.K.S.; Wingerter, A. A Mini-landscape of exact MSSM spectra in heterotic orbifolds. *Phys. Lett. B* **2007**, *645*, 88. [CrossRef]
220. Buchmuller, W.; Hamaguchi, K.; Lebedev, O.; Ratz, M. Supersymmetric Standard Model from the Heterotic String (II). *Nucl. Phys. B* **2007**, *785*, 149. [CrossRef]
221. Buchmuller, W.; Ludeling, C.; Schmidt, J. Local SU(5) Unification from the Heterotic String. *J. High Energy Phys.* **2007**, *0709*, 113. [CrossRef]
222. Lebedev, O.; Nilles, H.P.; Raby, S.; Ramos-Sanchez, S.; Ratz, M.; Vaudrevange, P.K.S.; Wingerter, A. The Heterotic Road to the MSSM with R parity. *Phys. Rev. D* **2008**, *77*, 046013. [CrossRef]
223. Chaichian, M.; Chkareuli, J.L.; Kobakhidze, A. Composite quarks and leptons in higher space-time dimensions. *Phys. Rev. D* **2002**, *66*, 095013. [CrossRef]
224. Ma, E. Neutrino masses in an extended gauge model with E_6 particle content. *Phys. Lett. B* **1996**, *380*, 286. [CrossRef]
225. Keith, E.; Ma, E. Generic consequences of a supersymmetric U(1) gauge factor at the TeV scale. *Phys. Rev. D* **1997**, *56*, 7155. [CrossRef]
226. Suematsu, D. Neutralino decay in the μ problem solvable extra U(1) models. *Phys. Rev. D* **1998**, *57*, 1738. [CrossRef]
227. Daikoku, Y.; Suematsu, D. Mass bound of the lightest neutral Higgs scalar in the extra U(1) models. *Phys. Rev. D* **2000**, *62*, 095006. [CrossRef]
228. King, S.F.; Moretti, S.; Nevzorov, R. Theory and phenomenology of an exceptional supersymmetric standard model. *Phys. Rev. D* **2006**, *73*, 035009. [CrossRef]

229. King, S.F.; Moretti, S.; Nevzorov, R. Exceptional supersymmetric standard model. *Phys. Lett. B* **2006**, *634*, 278. [CrossRef]
230. King, S.F.; Moretti, S.; Nevzorov, R. Gauge coupling unification in the exceptional supersymmetric standard model. *Phys. Lett. B* **2007**, *650*, 57. [CrossRef]
231. King, S.F.; Moretti, S.; Nevzorov, R. E_6SSM. *AIP Conf. Proc.* **2007**, *881*, 138.
232. Athron, P.; King, S.F.; Miller, D.J.; Moretti, S.; Nevzorov, R. Predictions of the Constrained Exceptional Supersymmetric Standard Model. *Phys. Lett. B* **2009**, *681*, 448. [CrossRef]
233. Athron, P.; King, S.F.; Miller, D.J.; Moretti, S.; Nevzorov, R. The Constrained Exceptional Supersymmetric Standard Model. *Phys. Rev. D* **2009**, *80*, 035009. [CrossRef]
234. Athron, P.; Hall, J.P.; Howl, R.; King, S.F.; Miller, D.J.; Moretti, S.; Nevzorov, R. Aspects of the exceptional supersymmetric standard model. *Nucl. Phys. Proc. Suppl.* **2010**, *200–202*, 120. [CrossRef]
235. Hall, J.P.; King, S.F.; Nevzorov, R.; Pakvasa, S.; Sher, M. Novel Higgs Decays and Dark Matter in the E_6SSM. *Phys. Rev. D* **2011**, *83*, 075013. [CrossRef]
236. Athron, P.; King, S.F.; Miller, D.J.; Moretti, S.; Nevzorov, R. LHC Signatures of the Constrained Exceptional Supersymmetric Standard Model. *Phys. Rev. D* **2011**, *84*, 055006. [CrossRef]
237. Athron, P.; King, S.F.; Miller, D.J.; Moretti, S.; Nevzorov, R. Constrained Exceptional Supersymmetric Standard Model with a Higgs Near 125 GeV. *Phys. Rev. D* **2012**, *86*, 095003. [CrossRef]
238. Nevzorov, R. E_6 inspired supersymmetric models with exact custodial symmetry. *Phys. Rev. D* **2013**, *87*, 015029. [CrossRef]
239. Nevzorov, R. Quasifixed point scenarios and the Higgs mass in the E_6 inspired supersymmetric models. *Phys. Rev. D* **2014**, *89*, 055010. [CrossRef]
240. Nevzorov, R.; Pakvasa, S. Exotic Higgs decays in the E_6 inspired SUSY models. *Phys. Lett. B* **2014**, *728*, 210. [CrossRef]
241. Athron, P.; Mühlleitner, M.; Nevzorov, R.; Williams, A.G. Non-Standard Higgs Decays in U(1) Extensions of the MSSM. *J. High Energy Phys.* **2015**, *1501*, 153. [CrossRef]
242. Athron, P.; Harries, D.; Nevzorov, R.; Williams, A.G. E_6 Inspired SUSY benchmarks, dark matter relic density and a 125 GeV Higgs. *Phys. Lett. B* **2016**, *760*, 19. [CrossRef]
243. Nevzorov, R.; Pakvasa, S. Nonstandard Higgs decays in the E_6 inspired SUSY models. *Nucl. Part. Phys. Proc.* **2016**, *273-275*, 690. [CrossRef]
244. King, S.F.; Nevzorov, R. 750 GeV Diphoton Resonance from Singlets in an Exceptional Supersymmetric Standard Model. *J. High Energy Phys.* **2016**, *1603*, 139. [CrossRef]
245. Athron, P.; Harries, D.; Nevzorov, R.; Williams, A.G. Dark matter in a constrained E_6 inspired SUSY model. *J. High Energy Phys.* **2016**, *1612*, 128. [CrossRef]
246. Nevzorov, R. Leptogenesis as an origin of hot dark matter and baryon asymmetry in the E_6 inspired SUSY models. *Phys. Lett. B* **2018**, *779*, 223. [CrossRef]
247. Nevzorov, R. E_6 inspired SUSY models with custodial symmetry. *Int. J. Mod. Phys. A* **2018**, *33*, 1844007. [CrossRef]
248. King, S.F.; Moretti, S.; Nevzorov, R. A Review of the Exceptional Supersymmetric Standard Model. *Symmetry* **2020**, *12*, 557. [CrossRef]
249. Nevzorov, R. Higgs Boson with Mass around 125 GeV in SUSY Extensions of the SM. *Phys. Atom. Nucl.* **2020**, *83*, 338. [CrossRef]
250. King, S.F.; Luo, R.; Miller, D.J.; Nevzorov, R. Leptogenesis in the Exceptional Supersymmetric Standard Model: Flavour dependent lepton asymmetries. *J. High Energy Phys.* **2008**, *0812*, 042. [CrossRef]
251. Asaka, T.; Buchmuller, W.; Covi, L. Bulk and brane anomalies in six-dimensions. *Nucl. Phys. B* **2003**, *648*, 231. [CrossRef]
252. von Gersdorff, G.; Quiros, M. Localized anomalies in orbifold gauge theories. *Phys. Rev. D* **2003**, *68*, 105002. [CrossRef]
253. Scrucca, C.A.; Serone, M. Anomalies in field theories with extra dimensions. *Int. J. Mod. Phys. A* **2004**, *19*, 2579. [CrossRef]
254. Borghini, N.; Gouverneur, Y.; Tytgat, M.H.G. Anomalies and fermion content of grand unified theories in extra dimensions. *Phys. Rev. D* **2002**, *65*, 025017. [CrossRef]
255. von Gersdorff, G. Anomalies on Six Dimensional Orbifolds. *J. High Energy Phys.* **2007**, *703*, 83. [CrossRef]
256. Adler, S.L. Axial vector vertex in spinor electrodynamics. *Phys. Rev.* **1969**, *177*, 2426. [CrossRef]
257. Adler, S.L.; Bardeen, W.A. Absence of higher order corrections in the anomalous axial vector divergence equation. *Phys. Rev.* **1969**, *182*, 1517. [CrossRef]
258. Bell, J.S.; Jackiw, R. A PCAC puzzle: Pi0 –> gamma gamma in the sigma model. *Nuovo Cim. A* **1969**, *60*, 47. [CrossRef]
259. Green, M.B.; Schwarz, J.H. Anomaly Cancellation in Supersymmetric D=10 Gauge Theory and Superstring Theory. *Phys. Lett. B* **1984**, *149*, 117. [CrossRef]
260. Arkani-Hamed, N.; Cohen, A.G.; Georgi, H. Anomalies on orbifolds. *Phys. Lett. B* **2001**, *516*, 395. [CrossRef]
261. Scrucca, C.A.; Serone, M.; Silvestrini, L.; Zwirner, F. Anomalies in orbifold field theories. *Phys. Lett. B* **2002**, *525*, 169. [CrossRef]
262. Barbieri, R.; Contino, R.; Creminelli, P.; Rattazzi, R.; Scrucca, C.A. Anomalies, Fayet-Iliopoulos terms and the consistency of orbifold field theories. *Phys. Rev. D* **2002**, *66*, 024025. [CrossRef]
263. Phillips, D.G., II; Snow, W.M.; Babu, K.; Banerjee, S.; Baxter, D.V.; Berezhiani, Z.; Bergevin, M.; Bhattacharya, S.; Brooijmans, G.; Castellanos, L. Neutron-Antineutron Oscillations: Theoretical Status and Experimental Prospects. *Phys. Rept.* **2016**, *612*, 1. [CrossRef]
264. Kronfeld, A.S.; Tschirhart, R.S.; Al-Binni, U.; Altmannshofer, W.; Ankenbrandt, C.; Babu, K.; Banerjee, S.; Bass, M.; Batell, B.; Baxter, D.V. Project X: Physics Opportunities. *arXiv* **2013**, arXiv:1306.5009.

265. The ATLAS collaboration. A search for pair-produced resonances in four-jet final states at $\sqrt{s}$ = 13 TeV with the ATLAS detector. *Eur. Phys. J. C* **2018**, *78*, 250.
266. Cline, J.M.; Raby, S. Gravitino induced baryogenesis: A Problem made a virtue. *Phys. Rev. D* **1991**, *43*, 1781. [CrossRef]
267. Scherrer, R.J.; Cline, J.M.; Raby, S.; Seckel, D. Gravitino induced baryogenesis, primordial nucleosynthesis, and the Tremaine-Gunn limit. *Phys. Rev. D* **1991**, *44*, 3760. [CrossRef] [PubMed]
268. Ma, E.; Raidal, M. Three active and two sterile neutrinos in an E_6 model of diquark baryogenesis. *J. Phys. G* **2002**, *28*, 95. [CrossRef]
269. Ma, E. Multiplicative conservation of baryon number and baryogenesis. *Phys. Lett. B* **2008**, *661*, 273. [CrossRef]
270. Ma, E. Common Origin of $(-)^L$, $(-)^{3B}$, and Strong CP Conservation. *Phys. Rev. D* **2008**, *78*, 047701. [CrossRef]
271. Ma, E. Axionic Extensions of the Supersymmetric Standard Model. *Mod. Phys. Lett. A* **2009**, *24*, 1335. [CrossRef]
272. Kohri, K.; Mazumdar, A.; Sahu, N. Inflation, baryogenesis and gravitino dark matter at ultra low reheat temperatures. *Phys. Rev. D* **2009**, *80*, 103504. [CrossRef]
273. Cui, Y.; Sundrum, R. Baryogenesis for weakly interacting massive particles. *Phys. Rev. D* **2013**, *87*, 116013. [CrossRef]
274. Krauss, L.M.; Long, A.J.; Sabharwal, S. Gravitino Leptogenesis. *Phys. Rev. D* **2014**, *89*, 043503. [CrossRef]
275. Rompineve, F. Weak Scale Baryogenesis in a Supersymmetric Scenario with R-parity violation. *J. High Energy Phys.* **2014**, *1408*, 014. [CrossRef]
276. Boucenna, S.M.; Morisi, S. Theories relating baryon asymmetry and dark matter: A mini review. *Front. Phys.* **2014**, *1*, 33. [CrossRef]
277. Ishiwata, K.; Jeong, K.S.; Takahashi, F. Moduli-induced Baryogenesis. *J. High Energy Phys.* **2014**, *1402*, 62. [CrossRef]
278. Dhuria, M.; Hati, C.; Sarkar, U. Explaining the CMS excesses, baryogenesis and neutrino masses in E_6 motivated $U(1)_N$ model. *Phys. Rev. D* **2016**, *93*, 015001. [CrossRef]
279. Cui, Y. A Review of WIMP Baryogenesis Mechanisms. *Mod. Phys. Lett. A* **2015**, *30*, 1530028. [CrossRef]
280. Farina, M.; Monteux, A.; Shin, C.S. Twin mechanism for baryon and dark matter asymmetries. *Phys. Rev. D* **2016**, *94*, 035017. [CrossRef]
281. Cui, Y.; Okui, T.; Yunesi, A. LHC Signatures of WIMP-triggered Baryogenesis. *Phys. Rev. D* **2016**, *94*, 115022. [CrossRef]
282. Luty, M.A. Baryogenesis Via Leptogenesis. *Phys. Rev. D* **1992**, *45*, 455. [CrossRef] [PubMed]
283. Flanz, M.; Paschos, E.A.; Sarkar, U. Baryogenesis from a lepton asymmetric universe. *Phys. Lett. B* **1995**, *345*, 248. [CrossRef]
284. Plumacher, M. Baryogenesis and lepton number violation. *Z. Phys. C* **1997**, *74*, 549. [CrossRef]
285. Buchmuller, W.; Plumacher, M. CP asymmetry in Majorana neutrino decays. *Phys. Lett. B* **1998**, *431*, 354. [CrossRef]
286. Davidson, S.; Nardi, E.; Nir, Y. Leptogenesis. *Phys. Rept.* **2008**, *466*, 105. [CrossRef]