Electroweak Baryogenesis in the Presence of an Isosinglet Quark

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Abstract

We consider the possibility of electroweak baryogenesis in a simple extension of the standard model with an extra singlet complex scalar and a vector-like down quark. We show that in the present model the first-order electroweak phase transition can be strong enough to avoid the baryon asymmetry washout by sphalerons and that the CP-violating effects can be sufficient to explain the observed baryon-to-entropy ratio $n_B/s \sim 10^{-10}$. Other appealing features of the model include the generation of a CKM phase from spontaneous CP breaking at a high energy scale and a possible solution of the strong CP problem through the natural suppression of the parameter $\theta$. 
1 Introduction

Among the various scenarios which attempt to explain the baryon asymmetry of our Universe, the mechanism of electroweak baryogenesis continues being one of the most attractive [1]. One of the motivations to consider the electroweak scenario lies of course in the fact that with the advent of new high-energy colliders, physics at the electroweak scale becomes more accessible and testable. Furthermore, it is very appealing the fact that all the necessary ingredients for a successful baryogenesis [3] (i.e. baryon number violation, $C$ and $CP$ violation and departure from thermal equilibrium) can be easily found in quantum theories of particle interactions.

Unfortunately, the standard model (SM) of electroweak interactions fails in providing the required baryon asymmetry for at least two reasons. Firstly, the electroweak phase transition (EWPT) is not strongly first-order [3, 4, 5] and therefore, any $\Delta B \neq 0$ created during the transition would subsequently be washed out by unsuppressed $B$-violating processes in the broken phase. Secondly, the $CP$-violating effects coming from the Cabibbo-Kobayashi-Maskawa (CKM) mixing matrix are too small [6, 7] to explain the observed baryon-to-entropy ratio $n_B/s \sim 10^{-10}$ [8]. Thus, for electroweak baryogenesis to be feasible, physics beyond the minimal standard model must be invoked.

Several extensions of the SM have been studied in the literature. In particular, in the two Higgs doublet model [9], the singlet majoron model [10] and in the SM with an extra real Higgs singlet [11] or a complex gauge singlet with zero vacuum expectation value (VEV) [4], the first-order EWPT can be strong enough to suppress the sphaleron interactions after the transition. However, it should be stressed that a strong first-order phase transition is not sufficient for a successful baryogenesis; an adequate amount of $CP$ violation is also required. Supersymmetric extensions of the SM also provide a possible framework for electroweak baryogenesis. For instance, in the so-called light stop scenario of the minimal supersymmetric standard model the phase transition can be sufficiently strongly first-order [12]-[15] for values of the lightest Higgs and stop masses consistent with the present experimental bounds. Furthermore, the latter model contains additional sources of $CP$ violation which can account for the observed baryon asymmetry [16]-[19].

In this letter we analyze the mechanism of electroweak baryogenesis in a simple extension of the SM, where the only additional fields are a charge $-1/3$
vector-like quark $D$ and a singlet complex scalar $S$ \cite{20,21}. The addition of extra vector-like quarks to the SM is particularly attractive since they naturally arise in grand unified theories such as $E_6$. After briefly describing the model and explaining how the Higgs sector spontaneously breaks $CP$ through a phase in the VEV of the singlet scalar, we proceed to study the electroweak phase transition. We show that in a wide range of the parameter space and provided that one of the singlet scalar components is light enough, the first-order EWPT can be sufficiently strong so that the sphaleron interactions in the broken phase are too slow to erase the baryon asymmetry created during the transition. We also discuss the possible mechanisms of baryon number violation in the present model. As far as spontaneous baryogenesis \cite{22} (through a space-time dependent complex VEV of the singlet field and/or a radiatively induced $\theta_{\text{strong}}$) is concerned, it turns out that it cannot give the right order of magnitude for the observed $n_B/s$ ratio. Nevertheless, a strong enhancement compared to the SM can be obtained via reflection of quasiparticles on the bubble wall (charge transport baryogenesis \cite{23,24,25}) since the model has new lower dimensional $CP$-violating weak-basis invariants (compared to the SM $CP$-violating invariant).

2 The model

We consider a simple extension of the SM where, in addition to the usual field content, we introduce in the fermion sector a vector-like down quark $D$, singlet under $SU(2)$ and in the Higgs sector, an extra complex scalar $S$, singlet under $SU(2) \times U(1)$ \cite{20}. As shown in ref. \cite{20}, this minimal Higgs structure can spontaneously break $CP$ through a phase in the VEV of the singlet $S$. Furthermore, this phase induces a non-vanishing phase in the CKM matrix, which is not suppressed by a small ratio $v/\sigma = |\langle \phi \rangle|/|\langle S \rangle| \ll 1$. We emphasize that small values of the latter ratio are desirable to naturally suppress not only the existing flavour-changing neutral currents, but also the one-loop finite contributions to the parameter $\bar{\theta}$ associated with strong $CP$ violation \cite{21}. It turns out that the most stringent bounds on the VEV of the scalar singlet, $\sigma$, come from the new contributions to the neutron electric dipole moment. Nevertheless, the present experimental limit $|d_n| \lesssim 10^{-25}\text{e} \text{cm}$ \cite{23} implies $\sigma \gtrsim \text{few TeV}$ for acceptable values of the couplings and parameters of the model \cite{24,21}.
The field content of the model is given by
\[
(u \ d)_L^i, \ u_R^i, \ d_R^i, \ D_L, \ \phi, \ S, \quad i = 1, 2, 3, \quad \alpha = 1, \ldots, 4,
\]
where \(i, \alpha\) are family indices and \(\phi, S\) denote the standard Higgs doublet and the extra complex scalar singlet, respectively. We impose \(CP\) invariance at the Lagrangian level and also introduce a \(Z_2\) discrete symmetry under which \(D_L\) and \(S\) are odd, while all the other fields are even. The role of the \(Z_2\) symmetry is to forbid quark bare mass terms of the form \(\bar{D}_L d^i_R\) and therefore, to guarantee the vanishing of \(\theta\)-strong at the tree level. However, we will allow a soft breaking of such symmetry in the scalar sector. As we shall see below, soft-breaking terms, linear and cubic in the singlet scalar field, will play a crucial role in avoiding the baryon asymmetry washout by sphalerons during the EWPT.

The most general \(SU(2) \times U(1) \times Z_2\) invariant potential can be written in the form
\[
V_0 = -m^2 \phi^\dagger \phi + \lambda (\phi^\dagger \phi)^2 - m_S^2 S^* S + \lambda_S (S^* S)^2 + \beta (\phi^\dagger \phi) (S^* S)
+ (\mu^2 + \beta' (\phi^\dagger \phi) + \lambda' (S^* S)) (S^2 + S'^2) + \lambda'' (S^4 + S'^4).
\] (1)
All the parameters are assumed to be real, so that the Lagrangian is \(CP\) invariant. In addition to \(V_0\), we shall introduce the following \(Z_2\)-breaking terms in the direction of the real component of the singlet scalar field
\[
V' = \sqrt{2} \xi (\phi^\dagger \phi) (S + S^*) - \frac{\alpha}{6\sqrt{2}} (S + S^*)^3,
\] (2)
with \(\xi\) and \(\alpha\) real. The \(SU(2) \times U(1) \times Z_2\) invariant Yukawa interactions are
\[
\mathcal{L}_Y = -\sqrt{2}(\bar{u} \ \bar{d})_L^i \left( g_{ij} \phi d_R^j + h_{ij} \phi u_R^j \right) - \mu_D \bar{D}_L D_R
- \sqrt{2} (f_i S + f_i' S^*) \tilde{D}_L d_R^i + \text{h.c.,}
\] (3)
where \(\tilde{\phi} = i \sigma_2 \phi^*, D_R \equiv d_R^4\) and all interaction constants are real because of \(CP\) invariance. The down quark mass matrix \(\mathcal{M}_d\) can be then written as follows:
\[
\mathcal{M}_d = \begin{pmatrix} M_d & 0 \\ M_D & \mu_D \end{pmatrix}
\] (4)
with

\begin{align*}
(M_d)_{ij} &= \sqrt{2} g_{ij} \langle \phi^0 \rangle, \\
(M_D)_{i} &= \sqrt{2} (f_i \langle S \rangle + f'_i \langle S^* \rangle);
\end{align*}

(5)

\langle \phi^0 \rangle \text{ and } \langle S \rangle \text{ are the VEV for the neutral Higgs doublet and the singlet scalar, respectively. The up quark mass matrix is the same as in the SM,}

\[ M_u = (M_u)_{ij} = \sqrt{2} h_{ij} \langle \phi^0 \rangle. \]  

(6)

For our purposes it will be more convenient to express the complex scalar \( S \) in terms of its real and imaginary parts. Let us write

\[ S = \frac{1}{\sqrt{2}} (S_1 + iS_2) \]

and denote by \( h = \sqrt{2} \phi^0 \) the neutral component of the Higgs field. In terms of these fields, the potential \( V = V_0 + V' \) reads

\[ V = -\frac{1}{2} m^2 h^2 + \frac{\lambda}{4} h^4 + \frac{\beta_1}{2} h^2 S^2 + \xi h^2 S_1 - \frac{1}{2} \mu_1^2 S_1^2 - \frac{\alpha}{3} S_1^3 + \frac{\lambda_1}{4} S_1^4 \\
+ \frac{\beta_2}{2} h^2 S_2^2 + \frac{1}{2} \gamma S_1^2 S_2^2 - \frac{1}{2} \mu_2^2 S_2^2 + \frac{\lambda_2}{4} S_2^4, \]

(7)

with the obvious identifications:

\[ \mu_1^2 = m_S^2 - 2\mu^2, \quad \mu_2^2 = m_S^2 + 2\mu^2, \]
\[ \lambda_1 = \lambda_S + 2\lambda' + 2\lambda'', \quad \lambda_2 = \lambda_S - 2\lambda' + 2\lambda'', \]
\[ \beta_1 = \frac{1}{2} (\beta + 2\beta'), \quad \beta_2 = \frac{1}{2} (\beta - 2\beta'), \quad \gamma = \lambda_S - 6\lambda''. \]

In what follows we suppose \( m^2, \lambda, \mu_k^2, \lambda_k > 0; \ \beta_k, \gamma \geq 0 \ (k = 1, 2) \).

After the spontaneous symmetry breaking we have

\[ \langle h \rangle = v, \quad \langle S \rangle = \frac{\sigma e^{i\eta}}{\sqrt{2}} = \frac{1}{\sqrt{2}} (s_1 + is_2). \]

(8)

If the absolute minimum of the potential [7] is at \( v \neq 0, \ \sigma \neq 0 \) and \( \eta \neq k\pi/2 \ (k = 0, 1, \ldots) \), then the vacuum breaks in general both \( T \) and \( CP \) invariance. Such spontaneous \( CP \) breaking is expected to occur at a high energy scale (\( \sigma \gtrsim \text{few TeV} \)) and at least one of the VEV, \( s_1 \) or \( s_2 \), should be of the order of this scale. On the other hand, in order to avoid the baryon asymmetry washout by the sphalerons during the EWPT, one of
the singlet components must be light enough to produce a strong first-order phase transition.

We shall assume that the field component \( S_2 \) is heavy, \( m_{S_2} \sim \text{few TeV} \), while \( S_1 \) is light, \( m_{S_1} \gtrsim m_H \). In this case, the \( S_2 \) field will decouple from the theory and the tree-level potential is reduced to the SM one, plus a real singlet field. We stress however that, although not relevant in our analysis of the phase transition, the \( S_2 \) field will play an essential role (through the phase \( \eta = \arctan(s_2/s_1) \)) in what concerns the magnitude of \( CP \) violation and production of a net baryon asymmetry, \( \Delta B \neq 0 \), during the EWPT. Notice also that to have \( \nu \ll \sigma \) without any further fine tuning, the parameters \( \beta_2 \) and \( \gamma \) in Eq.(7) should be taken small enough. Here we will assume, for simplicity, \( \beta_2 = \gamma = 0 \).

Under the above assumptions, the minimization of \( \rho \) implies \( s_2 = \mu_2/\sqrt{\lambda_2} \gg v, s_1 \), while for \( v, s_1 \) one obtains the system of equations:

\[
\begin{align*}
\nu \{ -m^2 + \lambda v^2 + \beta_1 s_1^2 + 2 \xi s_1 \} &= 0, \\
\lambda_1 s_1^3 - \alpha s_1^2 - (\mu_1^2 - \beta_1 v^2) s_1 + \xi v^2 &= 0.
\end{align*}
\]

(9) (10)

The system (9)-(10) can have up to six solutions (with \( v \geq 0 \)), two of which can be local minima. The field-dependent scalar masses are given by

\[
\begin{align*}
m_{\chi}^2 &= -m^2 + \lambda v^2 + \beta_1 s_1^2 + 2 \xi s_1, \\
m_{h,s_1}^2 &= \frac{1}{2} \left[ -m^2 - \mu_1^2 + (3 \lambda + \beta_1) v^2 + (\beta_1 + 3 \lambda_1) s_1^2 + 2 (\xi - \alpha) s_1 \right] \\
&\quad \pm \frac{1}{2} \left\{ \left[ -m^2 + \mu_1^2 + (3 \lambda - \beta_1) v^2 + (\beta_1 - 3 \lambda_1) s_1^2 + 2 (\xi + \alpha) s_1 \right]^2 \\
&\quad + 16 v^2 (\beta_1 s_1 + \xi s_1)^2 \right\}^{1/2},
\end{align*}
\]

(11)

where \( m_{\chi} \) corresponds to the Goldstone bosons. The gauge boson masses remain the same as in the SM. At the global minimum, which satisfies Eqs.(9)-(10), we have the following relations for the physical masses:

\[
\begin{align*}
m_h^2 + m_{S_1}^2 &= 2 \lambda v_0^2 + 2 \lambda_1 s_1^2 - \xi v_0^2 / s_1 - \alpha s_1, \\
m_h^2 m_{S_1}^2 &= 2 \lambda v_0^2 \left( 2 \lambda_1 s_1^2 - \xi v_0^2 / s_1 - \alpha s_1 \right) - 4 v_0^2 (\beta_1 s_1 + \xi)^2,
\end{align*}
\]

(13) (14)

with \( v_0 \approx 246 \text{ GeV} \) in order to reproduce the experimental values of the gauge boson masses, \( m_W \approx 80 \text{ GeV} \), \( m_Z \approx 91 \text{ GeV} \).
3 The electroweak phase transition

Let us now study how the electroweak phase transition proceeds in the present model. In the presence of a thermal bath, the effective potential (i.e. the free energy density) must include the interactions between the fields and the hot plasma. At one loop in perturbation theory, this amounts to add to the ground state energy, the free energy of a gas of noninteracting particles at finite temperature. The result is well known \[26\]

\[
\Delta V_T = \frac{T^4}{2\pi^2} \left\{ \sum_B g_B I_- \left( \frac{m_B}{T} \right) + \sum_F g_F I_+ \left( \frac{m_F}{T} \right) \right\},
\]

\[
I_\pm(y) = \pm \int_0^\infty x^2 \ln \left( 1 \mp e^{-\sqrt{x^2+y^2}} \right) dx,
\]

(15)

where \(g_B\) (\(g_F\)) are the boson (fermion) degrees of freedom; \(m_B\) (\(m_F\)) is the mass of a boson (fermion) in the presence of background fields. In the high temperature limit, when \(m \lesssim T\), Eq.(15) can be expanded in powers of \(m/T\) and one obtains \[26\]

\[
\Delta V_T = \sum_{i=B,F} g_i \left\{ A_i \frac{m_i^2 T^2}{48} - B_i \frac{m_i^3 T}{12\pi} - D_i \frac{m_i^4}{64\pi^2} \ln \left( \frac{m_i^2}{T^2} \right) - C_i \right\},
\]

(16)

where \(A_i = 2\ (1)\), \(B_i = 1\ (0)\), \(D_i = 1\ (-1)\), for bosons (fermions) and \(C_B \approx 5.41\), \(C_F \approx 2.46\).

At this point it is worthwhile making a few comments. First we notice that in the high temperature limit the \(T\)-independent vacuum fluctuations \(\sim m^4 \ln m^2\) in Eq.(14) are exactly cancelled by similar contributions coming from the one-loop effective potential at \(T = 0\). As far as the \(m^4 \ln T^2\) fluctuations are concerned, they tend to modify the quartic terms, but are subdominant provided the couplings are small. In what follows we shall assume that the latter conditions apply and therefore we neglect these terms. Similarly, we shall also neglect all one-loop scalar self-interactions. Finally, we note that since for the scalar component \(S_2\), \(m_{S_2} \gg T\), its contribution to the finite temperature effective potential is Boltzmann suppressed. The same argument applies to the vector-like down quark \(D\) which is assumed to have a mass \(m_D > m_t\).

Thus, adding the zero-temperature potential (7) and the finite-temperature terms (13) for all the relevant particles involved (Higgs, Goldstone and gauge...
bosons, singlet scalar and top quark), finally we obtain the approximate one-loop effective potential at high temperatures

\[
V_T = \frac{1}{2} m^2(T) h^2 - \frac{1}{3} \delta(T) h^3 + \frac{\lambda}{4} h^4 + \frac{1}{2} \beta_1 h^2 S_1^2 + \xi h^2 S_1 + (4 \xi - \alpha) \frac{T^2}{12} S_1 + \frac{1}{2} \mu_1^2(T) S_1^2 - \frac{\alpha}{3} S_3^3 + \frac{\lambda_1}{4} S_1^4,
\]

(17)

where

\[
m^2(T) = -m^2 + \left( \frac{\beta_1}{3} + 2\lambda + \epsilon \right) \frac{T^2}{4},
\]

(18)

\[
\mu_1^2(T) = -\mu_1^2 + \left( \frac{\beta_1}{3} + \frac{\lambda_1}{4} \right) T^2,
\]

(19)

\[
\epsilon = \frac{2m_W^2 + m_Z^2 + 2m_t^2}{v_0^2} \approx 1.37,
\]

(20)

\[
\delta(T) = \frac{(2m_W^2 + m_Z^2) T}{2\pi v_0^3} \approx 0.02 T,
\]

(21)

and we have used the experimental central value for the top quark mass \( m_t \approx 176 \text{ GeV} \).

The minimization of the free energy density (17) with respect to the fields \( h \) and \( S_1 \) yields the system of equations

\[
v \left\{ m^2(T) - \delta(T) v + \lambda v^2 + \beta_1 s_1^2 + 2\xi s_1 \right\} = 0,
\]

(22)

\[
\lambda_1 s_1^3 - \alpha s_1^2 + \left( \mu_1^2(T) + \beta_1 v^2 \right) s_1 + \xi v^2 + (4 \xi - \alpha) \frac{T^2}{12} = 0.
\]

(23)

Solving this system we can find \( v = v(T) \) and \( s_1 = s_1(T) \). Note that as \( T \to \infty \), the solution of (22)-(23) is

\[
v = 0, \ s_1 \to \frac{\alpha - 4\xi}{4\beta_1 + 3\lambda_1}.
\]

(24)

Therefore, the electroweak symmetry restoration does take place at high temperature, while in general \( s_1 \neq 0 \). As the Universe cools down, a new minimum will appear with \( v \neq 0 \) and the electroweak symmetry will be spontaneously broken.

The interesting issue is, of course, the nature of such a transition. In order to make electroweak baryogenesis possible, a strong enough first-order EWPT
is required, so that any baryon asymmetry produced during the transition can survive without being diluted by the sphalerons. In the SM (with only one Higgs doublet), the requirement for the survival of the baryon asymmetry produced at the electroweak scale is given by [28]

$$\frac{E_{\text{sph}}(T_c)}{T_c} = \frac{4\pi B(\lambda/g_w^2)v(T_c)}{g_w T_c} \gtrsim 45,$$

(25)

where $E_{\text{sph}}(T_c)$ is the sphaleron energy, $T_c$ is the critical temperature and $1.56 \leq B(\lambda/g_w^2) \leq 2.72$ for $0 \leq \lambda/g_w^2 < \infty$. Thus, a large jump in the Higgs VEV is required during the transition,

$$\frac{v(T_c)}{T_c} \gtrsim 1,$$

(26)

which in turn translates into an upper bound on the Higgs mass.

There are convincing arguments [3, 4, 5] that the first-order phase transition in the minimal SM is too weak to yield an acceptable baryon asymmetry for Higgs masses consistent with the present experimental bounds [25]. In the model we are considering, one expects however new effects to come from the linear and cubic terms in the singlet field $S_1$, which can lead to an enhancement of the first-order EWPT. To show that this is indeed the case, we looked for the coexistence of two degenerate minima in the potential of Eq.(17) at some (critical) temperature by solving the extrema Eqs.(22)-(23). Then we verified that the constraint given by Eq.(26) was satisfied. The parameter space was chosen in the region of validity of Eq.(17) and of compatibility with the present experimental limits on the Higgs mass. In the $(v, s_1)$-plane we found a wide parameter range where there exist two degenerate minima, namely $(0, s_+(T_c))$ and $(v(T_c), s_-(T_c))$, and for which the constraint of Eq.(26) is satisfied. The first-order phase transition occurs as the global minimum tunnels from $(0, s_+(T_c))$ (symmetric phase) to $(v(T_c), s_-(T_c))$ (broken phase). For $T > T_c$, we also find in general another first-order phase transition in the $s_1$-direction [11], but this transition is not necessary for electroweak baryogenesis to succeed.

In Fig.1 we plot the ratio $v(T_c)/T_c$ as a function of the $Z_2$-breaking parameter $\xi$ for the particular values of $\alpha = 1.5$, $\beta = 0.01$, $\lambda = \lambda_1 = 0.07$ and for the

$1$ Notice that the requirement of Eq.(26) also applies to models with additional singlet scalar fields since the sphaleron energy in such models is smaller than the corresponding one of the standard model [10].
Higgs mass values $m_H = 80, 85, 90$ GeV. Fig. 2 shows the allowed region in the plane $(\xi, \alpha)$ where the constraint (26) is satisfied for $\beta = 0.01, \lambda = \lambda_1 = 0.07$ and $m_H = 80$ GeV.

Figure 1: Dependence of the ratio $v(T_c)/T_c$ on the soft-breaking parameter $\xi$ for the particular values of $\alpha = 1.5$, $\beta = 0.01$, $\lambda = \lambda_1 = 0.07$ and for the Higgs mass values $m_H = 80, 85, 90$ GeV.

We see that there exists a wide range in the parameter space where the EWPT is strongly first-order and thus baryogenesis is possible. We also notice that the soft-breaking terms (cf. Eq.(2)) play a crucial role in this respect.

4 CP-violation and baryogenesis

A strong enough first-order EWPT is a necessary but not a sufficient condition for baryogenesis to succeed. To explain the observed baryon asymmetry of the Universe a right amount of $CP$ violation is also required.

Two of the appealing features of the model under consideration are the possibility of spontaneous $CP$ violation and natural suppression of strong
Figure 2: The allowed region in the \((\xi, \alpha)\)-plane for which the requirement 
\(v(T_c)/T_c > 1\) is satisfied (shaded-in area). The plot is given for the particular 
values of \(\beta = 0.01, \lambda = \lambda_1 = 0.07\) and the Higgs mass \(m_H = 80\) GeV.

\[\begin{align*}
\text{Figure 2:} & \quad \text{The allowed region in the \((\xi, \alpha)\)-plane for which the requirement } v(T_c)/T_c > 1 \text{ is satisfied (shaded-in area). The plot is given for the particular values of } \\
& \quad \beta = 0.01, \lambda = \lambda_1 = 0.07 \text{ and the Higgs mass } m_H = 80 \text{ GeV.}
\end{align*}\]

\[\begin{align*}
CP \text{ violation } [20, 21]. & \quad \text{The model satisfies the Nelson-Barr criteria [29] which guarantee the tree-level vanishing of the } \theta_{\text{strong}} \text{ parameter. Moreover, due to the presence of vector-like quarks, there are new sources of } CP \text{ violation. In } \\
& \quad \text{general in an extension of the SM with } n \text{ standard generations and } n_d \quad Q = \\
& \quad -1/3 \quad \text{ isosinglet quarks, the generalized CKM matrix consists of the first } n \text{ lines of a } (n + n_d)-\text{dimensional unitary matrix and it has been shown [30] that there are } \\
& \quad (n - 1)(\frac{n - 2}{2} + n_d) \text{ physical } CP\text{-violating phases in the CKM matrix. Hence, with } n = 3 \text{ and } n_d = 1, \text{ there are three } CP\text{-violating phases. More } \\
& \quad \text{importantly, it has been shown [31] that in this minimal extension of the SM, there is } CP \text{ violation even in the chiral limit where } m_u = m_d = m_s = m_c = 0. \\
& \quad \text{This implies that } CP\text{-violating effects relevant for baryogenesis at the EWPT will not be suppressed by the smallness of the light quark masses. At this } \\
& \quad \text{point, it is worth recalling that in the three generation SM, the strength of }
\end{align*}\]
CP violation is controlled by the invariant \[^{32, 33}\]

\[ I_{\text{SM}} = \det \left[ (M_u M_u^\dagger), (M_d M_d^\dagger) \right] = \frac{1}{3} \text{Tr} \left[ (M_u M_u^\dagger), (M_d M_d^\dagger) \right]^3 \]

\[ = -2i \left( \Delta_{ij}^2 \Delta_{tu}^2 \Delta_{cu}^2 \right) \left( \Delta_{ij}^2 \Delta_{bd}^2 \Delta_{sd}^2 \right) \text{Im} \left( V_{ud} V_{dc} V_{us}^* V_{cd}^* \right), \quad (27) \]

where \( M_u, M_d \) are the up and down quark mass matrices and \( \Delta_{ij}^2 = m_i^2 - m_j^2 \). A natural question is whether in the present model there are new CP violating weak-basis invariants which may play a role in the generation of the observed baryon asymmetry\[^{31}\]. In ref.\[^{31}\] a complete set of necessary and sufficient conditions for CP invariance were given, expressed in terms of weak-basis invariants. In particular, it was shown that in the presence of one vector-like down quark the lowest weak-basis invariant is:

\[ I_{\text{VL}} = \text{ImTr}[M_u M_u^\dagger M_d M_d^\dagger M_d M_d^\dagger], \quad (28) \]

where \( M_u, M_d \) and \( M_D \) matrices are defined as in Eqs.(4)-(6). It can also be shown that in the chiral limit \( m_{u,d,s,c} = 0 \), all weak-basis invariants which measure CP violation in the present model, are proportional to \( I_{\text{VL}} \). Therefore, \( I_{\text{VL}} \) plays the same role that \( I_{\text{SM}} \) plays in the SM, giving the strength of CP violation relevant for baryogenesis at the EWPT.

In the SM, based on dimensional analysis, a naive estimate for the baryon asymmetry can be given as \[^{28}\]

\[ \frac{n_B}{s} \sim \frac{I_{\text{SM}}}{T_c^{12}} \sim 10^{-20}, \quad (29) \]

where \( T_c \) is the critical temperature (natural mass scale) at the EWPT.

In fact, a more refined analysis shows that the critical temperature is not the only scale at the phase transition. Taking into account the effects of strong interactions on the lifetime of quasiparticles, a better estimate of the baryon asymmetry in the SM can be given \[^{11}\]:

\[ \frac{n_B}{s} \sim \frac{I_{\text{SM}}}{T_c^6 \ell^{-6}} \sim 10^{-25}, \quad (30) \]

where \( \ell \) is the coherence length of the quasiparticles (i.e. the distance over which the quasiparticles propagate during their lifetime). Thus, there are

\[^{2}\]For previous discussions on the sources of CP violation in models with vector-like quarks and their role in electroweak baryogenesis, see e.g. refs.\[^{34}\].
two relevant scales, namely, the critical temperature $T_c$ and the inverse of the coherence length $\ell^{-1}$.

In our model two different mechanisms of baryogenesis can occur:

- *Spontaneous baryogenesis* through a space-time dependent complex VEV of the singlet $S$ and/or through a radiatively induced $\theta_{\text{strong}}$,

- *Charge transport baryogenesis* through reflection of quasiparticles off the oncoming bubble wall during the phase transition.

As we shall see below, the first mechanism cannot give the right order of magnitude for the baryon asymmetry in this type of model. Nevertheless, a strong enhancement compared to the SM results can be obtained via reflection of quasiparticles on the bubble wall.

Let us first discuss the spontaneous baryogenesis in the framework of our model. From the form of the down quark matrix given in Eq. (4) it follows that $\text{Im} \det M_d = 0$. During the electroweak phase transition, $M_d$ is a space-time dependent function of the VEV of the scalar Higgs doublet and the singlet field. The matrix $M_d$ can be diagonalized by unitary matrices belonging to the $SU(N)$ group. The kinetic term for fermions is not invariant under the field redefinition which diagonalizes $M_d$. However, all the currents induced by this redefinition are orthogonal to the baryonic charge. Thus, imposing $\theta_{\text{strong}}$ to be zero at the tree level implies in a natural way that the spontaneous baryogenesis mechanism is inefficient to produce the baryon asymmetry of the Universe.

Similar conclusions can be drawn for the radiatively induced $\theta_{\text{strong}}$. Indeed, the detailed balance principle allows us to relate the baryon density $n_B$ to the chemical potential $\mu_B = N \dot{\theta}$ associated to the baryonic charge ($N \sim O(1)$ is a constant),

$$ n_B = -N_f^2 \int \frac{\Gamma_{\text{sph}}(T)}{T} \mu_B dt, $$

where $N_f$ is the number of flavours, $\Gamma_{\text{sph}}(T)$ is the sphaleron rate and the integration goes until the time when the anomalous processes are out of thermal equilibrium. Given that the entropy density is $s = 2\pi^2 g^* T^3 / 45$ and that on dimensional grounds $\Gamma_{\text{sph}} = \kappa (\alpha_w T)^4$ in the symmetric phase, one
finds for the baryon-to-entropy ratio
\[
\frac{n_B}{s} \approx \frac{45N_f^2\kappa\alpha_w^4N\theta_{\text{strong}}}{2\pi^2g^*} \lesssim 10^{-18},
\] (32)
with \(g^* \sim 100\) is the effective number of degrees of freedom at the electroweak scale, \(0.1 \lesssim \kappa \lesssim 1\) and \(\alpha_w = g_w^2/4\pi\). We remark that the most severe constraint on \(\theta_{\text{strong}}\) comes from the experimental limits on the neutron electric dipole moment, \(|d_n| \lesssim 10^{-25}\text{ecm}\) [25], which implies \(|\theta_{\text{strong}}| < 10^{-10}\).

To evaluate the effects of \(CP\)-violation due to reflection of quasiparticles on the bubble wall we shall follow the approach of ref.\[7\]. The first-order EWPT proceeds via nucleation of bubbles, inside which the VEV of the Higgs field \(v \neq 0\) and which expand until they fill up the symmetric (\(v = 0\)) Universe. Due to \(CP\)-violating effects, quarks and antiquarks reflect asymmetrically on the bubble wall and therefore their distributions are different inside and outside the bubble. The fast \(B\)-violating processes will cause \(\langle \Delta B \rangle\) to relax to zero in the symmetric phase. The net baryon number produced is equal to (minus) the thermal average of the baryon number in the unbroken phase (i.e. the sum of the excess of baryons from the symmetric phase reflected off the wall plus the excess of baryons transmitted from the broken phase into the symmetric one).

Using unitarity and \(CPT\) invariance, and assuming that the sphaleron rate is equal to infinity in the unbroken phase and equal to zero in the broken phase, the ratio \(n_B/s\) can be given as [7]
\[
\frac{n_B}{s} = -\frac{1}{3} \frac{45}{2\pi^2g^*} \int \frac{d\omega}{2\pi} \frac{1}{\pi} \left( n_{L,0}^L(\omega) - n_{R,0}^R(\omega) \right) \text{Tr}[R_{LR}^\dagger R_{LR} - R_{RL}^\dagger R_{RL}],
\] (33)
where the matrix (in flavour space) \(R_{LR}\) is the probability of a left-handed quark to be reflected on the bubble wall into a right-handed quark; \(n_{L,0}^L(\omega)\) are Fermi-Dirac distributions boosted to the bubble wall frame, \(n[\gamma(\omega + \vec{v}_w,\vec{p})] = 1/(e^{\gamma(\omega + \vec{v}_w,\vec{p})/T} + 1)\), \(\gamma\) is the relativistic factor, \(\vec{v}_w\) is the bubble wall velocity and \((\omega,\vec{p})\) is the quasiparticle four-momentum. For small wall velocities, we can expand Eq.(33) in powers of \(\vec{v}_w\) to obtain:
\[
\frac{n_B}{s} = \frac{15}{2\pi^2g^*} \int \frac{d\omega}{2\pi} n(\omega) \left( 1 - n(\omega) \right) \frac{\vec{v}_w \cdot (\vec{p}_L - \vec{p}_R)}{T} \times \text{Tr}[R_{LR}^\dagger R_{LR} - R_{RL}^\dagger R_{RL}] + \mathcal{O}(v_w^2). \] (34)
The reflection coefficients can be evaluated using an expansion in the quark masses \[7\]. For the dominant contribution to the baryon asymmetry, finally we find

\[
\frac{n_B}{s} \approx -10^{-3} v_w \frac{I_{\text{VL}}}{T_c^8},
\]  

(35)

where \(I_{\text{VL}}\) is defined in Eq.(28) and we have assumed \(T_c \approx 1/\ell\).

The above equation can be expressed in terms of the mixing angles and the physical quark masses. Then, using the constraints coming from FCNC and the measured elements of the CKM matrix, an estimate on the \(D\) quark mass can be given in order to reproduce the observed baryon asymmetry.

In the chiral limit \(m_{u,d,s,c} = 0\), all \(CP\)-violating effects are due to physics beyond the SM. In this limit, the lowest invariant \(28\) can be easily written in terms of mixing angles and quark masses \[31\]. We have

\[
\sigma_{\text{VL}} = -m_t^2 m_D^2 m_b^2 (m_D^2 - m_b^2) \text{Im} \left( V_{tb} V^*_{4b} V_{4D} V^*_{tD} \right),
\]  

(36)

where \(V_{\alpha\beta}\) are matrix elements of the \(4 \times 4\) unitary matrix \(V\) which diagonalizes \(M_d M_d^\dagger\) in the weak-basis where \(M_u\) is diagonal. Therefore,

\[
\frac{n_B}{s} \approx 10^{-3} v_w \frac{m_t^2 m_D^2 m_b^2}{T_c^8} \text{Im} \left( V_{tb} V^*_{4b} V_{4D} V^*_{tD} \right).
\]  

(37)

The observed value for the cosmological baryon asymmetry is \(n_B/s \sim 10^{-10}\). The main constraint on the mixing angles comes from the FCNC and is given by \(|\text{Im} V_{tb} V^*_{4b} V_{4D} V^*_{tD}| \lesssim 10^{-4}\) \[31\]. Using the values \(T_c \approx 1/\ell \approx 120\) GeV, \(v_w = 0.1\), \(m_t \approx 95\) GeV, \(m_b \approx 3\) GeV, and assuming that the mixing angles are of the order of the above bound, we derive that the required mass of the \(D\) quark should be about \(m_D \approx 200\) GeV. Of course, this value of \(m_D\) should be taken only as a rough estimate, in view of uncertainties in the evaluation of the baryon-to-entropy ratio (in particular concerning the properties of the matter during the phase transition and the propagation of the bubble wall). In spite of these uncertainties, one may safely conclude that in the present model a \(D\) quark mass of the order of a few hundred GeV is needed in order to produce the baryon asymmetry of the Universe at the electroweak scale.

A final comment is in order. It is clear that the sign of the quantity \(\text{Im} Q_1 \equiv \text{Im}(V_{tb} V^*_{4b} V_{4D} V^*_{tD})\) is dictated by the sign of the baryon asymmetry
\[ n_B. \text{ Now, due to unitarity constraints, } \text{Im}Q_1 \text{ is related to the imaginary parts of other rephasing invariant quartets} \ [31]. \text{ Indeed one obtains:} \]
\[ \text{Im}Q_1 = \text{Im}Q_2 - \text{Im}Q_3 + \text{Im}Q_4 - \text{Im}Q_5, \quad (38) \]

where
\[
Q_2 = V_{ud}V_{tb}V_{td}^*V_{ub}^*, \quad Q_3 = V_{us}^*V_{tb}V_{ub}V_{ts}, \\
Q_4 = V_{cd}V_{tb}V_{td}^*V_{cb}^*, \quad Q_5 = V_{cs}^*V_{tb}V_{cb}V_{ts}. \quad (39)
\]

Note that \( \text{Im}Q_j \ (j = 2, 3, 4, 5) \) in Eq.\((39)\) involve CKM matrix elements connecting only standard quarks. In the SM, as a result of \(3 \times 3\) unitarity of \(V_{\text{CKM}}\), all \( \text{Im}Q_j \ (j = 2, 3, 4, 5) \) are equal and therefore the r.h.s. of Eq.\((38)\) vanishes. It is clear that \( \text{Im}Q_1 \) can in principle be obtained from low energy data signaling deviations from the SM predictions.

5 Conclusions

In this letter we have considered the possibility of generating at the electroweak scale the observed baryon asymmetry of the Universe in a simple extension of the SM, in which a complex singlet scalar field and an isosinglet down quark are added to the theory. We have shown that in the presence of soft \(Z_2\)-breaking terms, linear and cubic in the singlet field, the first-order electroweak phase transition can be strong enough to suppress the \(B\)-violating sphaleron processes after the transition, thus avoiding the washout of any \(B\) asymmetry generated during the transition. Moreover, we have seen that in the presence of isosinglet quarks, the theory contains new lower dimensional \(CP\)-violating weak-basis invariants, which can significantly contribute to the produced baryon asymmetry at the electroweak scale via the charge transport mechanism.

Of course, there are still uncertainties in the evaluation of the baryon-to-entropy ratio. Nevertheless, relations \((37)-(39)\) show that the sign of \(CP\)-violation in the \(K\)- and \(B\)-meson systems could in principle be related to the sign of the baryon asymmetry of the Universe. In our opinion this is a nice feature of the present model.
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