

Spectrum of color sextet scalars in realistic $SO(10)$ GUT

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 (Received 19 December 2022; accepted 22 February 2023; published 9 March 2023)

Incorporation of the standard model Yukawa interactions in a grand unified theory (GUT) often predicts varieties of new scalars that couple to the fermions and lead to some novel observational effects. We assess such a possibility for the color sextet diquark scalars within the realistic renormalizable models based on $SO(10)$ GUT. The spectrum consists of five sextets; $\Sigma \sim (6, 1, -\frac{2}{3})$, $S \sim (6, 1, \frac{1}{3})$, $\bar{S} \sim (\bar{6}, 1, -\frac{1}{3})$, $\mathcal{S} \sim (6, 1, \frac{4}{3})$, and $\bar{\mathcal{S}} \sim (\bar{6}, 3, -\frac{1}{3})$. Computing explicitly their couplings with the quarks, we evaluate their contributions to the neutral meson-antimeson mixing and baryon number-violating processes like neutron-antineutron oscillation. The latter arises because of a $B - L$ violating trilinear coupling between the sextets which also contributes to some of the quartic couplings and perturbativity of the same leads to strong limits on the sextet masses. Using the values of the $B - L$ breaking scale and Yukawa couplings permitted in the realistic models, we derive constraints on the masses of these scalars. It is found that Σ along with any of the remaining sextets cannot be lighter than the $B - L$ breaking scale, simultaneously. In the realm of realistic models, this implies no observable $n - \bar{n}$ oscillation in near future experiments. We also point out a possibility in which the sub-GUT scale Σ and a pair of S , allowed by the other constraints, can viably produce the observed baryon asymmetry of the Universe.

DOI: [10.1103/PhysRevD.107.055008](https://doi.org/10.1103/PhysRevD.107.055008)

I. INTRODUCTION

Grand unified theories (GUTs), which provide complete unification of the Standard Model (SM) gauge bosons and full or partial unification of the quarks and leptons, typically predict an enlarged spectrum for spin-0 particles [1–3]. This is in particular the case for the renormalizable versions of $SO(10)$ models constructed on the four-dimensional space-time which provide a unique platform for constructing an explicit, predictive, and realistic model of the grand unification [4–8]. Several scalar fields with varieties of color and electroweak charges are predicted as partners of the electroweak Higgs doublets in these models. Since only the latter are essentially required in the low-energy theory to break the electroweak symmetry, the rest are often assumed as heavy as the GUT scale invoking the so-called minimal survival hypothesis [9]. Nevertheless, if some of these scalars remain lighter than the GUT scale then they can give rise to some phenomenologically interesting effects because of their nontrivial SM charges

and direct couplings with quarks and leptons as predicted by the underlying GUT model. Such effects include flavor anomalies [10–14], distinct signatures for nucleon decays [15], neutron-antineutron oscillation [16–20], baryogenesis [18,19,21–25], precise unification of the SM gauge couplings [18,21,26–28] and some anomalous events in the direct search experiments [27–29].

A complete classification of the scalars that may arise from the most general Yukawa sector of the renormalizable $SO(10)$ models is given in our previous paper [15]. Among the various scalars, the color triplet and sextet fields are of particular phenomenological interest as they all carry non-zero $B - L$, where B (L) denotes baryon (lepton) number. Because of this, they give rise to processes that violate B and/or L which otherwise are good accidental global symmetries of the SM at the perturbative level. Among these, the color triplets are known to induce nucleon decay and they have been comprehensively studied in [15]. Computing explicitly their couplings with the quarks and leptons in the realistic $SO(10)$ GUTs, we derived bounds on their masses arising from various $B - L$ conserving and violating modes of proton and neutron decays. In this paper, we focus on the color sextet scalars with a similar intention to derive the constraints on their spectrum from various phenomenological considerations.

Unlike the color triplet scalars, the sextets do not induce nucleon decay by themselves. However, they can give rise to neutral baryon-antibaryon oscillations if there

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exists $B - L$ violating interaction between the relevant sextets [17,30,31]. The latter is inherent in the renormalizable $SO(10)$ models. One, therefore, expects constraints on the masses of the sextet scalars from neutron-antineutron oscillation experiments [32–35]. Light color sextets can also be constrained from the B conserving but flavor-violating mixings between mesons and antimesons [17,36–38]. Both these constraints primarily depend on (i) the Yukawa couplings of the quarks with the underlying sextet scalars, and (ii) the $B - L$ breaking scale. Unlike in the typical bottom-up approaches, both (i) and (ii) are more or less determined from the low-energy spectrum of the quarks and leptons in the realistic renormalizable $SO(10)$ models. Thus, one obtains more robust and unambiguous bounds on the spectrum of the colored sextet scalars in the top-down approach that we present in this work. Utilizing the spectrum of the sextet scalars allowed within the renormalizable $SO(10)$ models, we also point out a new and self-sufficient possibility of generating the observed baryon asymmetry of the Universe.

The rest of the paper is organized as follows. In the next section, we derive the spectrum and couplings of color sextet scalars in renormalizable $SO(10)$ models.

Various phenomenological implications of these scalars are derived which include flavor violation in Sec. III, neutron-antineutron oscillations in Sec. IV, perturbativity of the effective quartic couplings in Sec. V and baryogenesis in Sec. VI. Constraints from all these observables are analyzed in Sec. VII and the study is concluded in Sec. VIII.

II. COLOR SEXTET SCALARS AND THEIR COUPLINGS

The Yukawa sector of renormalizable $SO(10)$ GUTs comprises scalars in $\mathbf{10}$, $\overline{\mathbf{126}}$, and $\mathbf{120}$ dimensional irreducible representations of the gauge group. Various submultiplets residing in these GUT multiplets, along with their SM and $B - L$ charges and multiplicities, are listed in our previous paper [15]. For convenience, we reproduce the information relevant to $SU(3)_C$ sextet fields in Table I. These scalars arise only from $\overline{\mathbf{126}}_H$ and $\mathbf{120}_H$.

The couplings of these colored sextet fields with the SM quarks can be straightforwardly computed using the method discussed and the decompositions given in [15]. For the sextet fields residing in $\overline{\mathbf{126}}_H$, we find

$$\begin{aligned}
-\mathcal{L}_Y^{\overline{\mathbf{126}}} &= F_{AB} \mathbf{16}_A^T C^{-1} \mathbf{16}_B \overline{\mathbf{126}}_H + \text{H.c.} \\
&\supset -\frac{i}{\sqrt{15}} F_{AB} \left(2d_{\alpha A}^{CT} C^{-1} d_{\beta B}^C \Sigma^{\alpha\beta} - \sqrt{2} \epsilon^{\alpha\beta\gamma} u_{\gamma A}^{CT} C^{-1} d_{\sigma B}^C S_{\alpha\beta}^\sigma \right. \\
&\quad \left. + \epsilon^{\alpha\beta\gamma} u_{\sigma A}^{CT} C^{-1} u_{\gamma B}^C S_{\alpha\beta}^\sigma + \sqrt{2} \epsilon_{\alpha\beta\gamma} \epsilon_{ab} q_A^{\alpha a T} C^{-1} q_B^{\sigma c} S_{\sigma c}^{\beta\gamma b} \right) + \text{H.c.}
\end{aligned} \tag{1}$$

In the above, we continue following the notations used by us in the previous work [15] in which the α, β, \dots (a, b, \dots) letters denote $SU(3)_C$ ($SU(2)_L$) indices while A, B, \dots represent three flavors of quarks. Different numerical factors in front of each term in Eq. (1) arise from the Clebsch-Gordan decomposition and canonical normalization of the quark and scalar fields [15].

Analogously for the Yukawa interaction with $\mathbf{120}_H$, we obtain

$$\begin{aligned}
-\mathcal{L}_Y^{\mathbf{120}} &= G_{AB} \mathbf{16}_A^T C^{-1} \mathbf{16}_B \mathbf{120}_H + \text{H.c.} \\
&\supset -\frac{2i}{\sqrt{3}} G_{AB} (\epsilon^{\alpha\beta\gamma} u_{\gamma A}^{CT} C^{-1} d_{\sigma B}^C \tilde{S}_{\alpha\beta}^\sigma \\
&\quad - \epsilon_{\alpha\beta\gamma} u_A^{\gamma T} C^{-1} d_B^\sigma \tilde{S}_\sigma^{\alpha\beta}) + \text{H.c.}
\end{aligned} \tag{2}$$

Here, F and G are symmetric and antisymmetric matrices in the flavor space, respectively. \tilde{S} denotes the color sextet with $Y = 1/3$ residing in $\mathbf{120}_H$ and it is distinguished from S belonging to $\overline{\mathbf{126}}_H$ which has the same quantum numbers. It is noted that \tilde{S} and \mathbb{S} couple to the left-chiral quark fields while the remaining color sextets have interaction vertices

with only the right-chiral quarks. All the interactions in Eqs. (1) and (2) conserve $B - L$.

In the models with both $\overline{\mathbf{126}}_H$ and $\mathbf{120}_H$ present, the fields S and \tilde{S} can mix with each other through gauge-invariant terms like $\mathbf{120}_H \overline{\mathbf{126}}_H^\dagger \mathbf{45}_H$ or $\mathbf{120}_H \overline{\mathbf{126}}_H^\dagger \mathbf{210}_H$. The physical states are then given by linear combinations of S and \tilde{S} . For simplicity, we assume that such linear combinations are parametrized by real parameters and define

TABLE I. Types of colored sextet fields, their charges under the SM gauge group [$SU(3)_C$, $SU(2)_L$, $U(1)_Y$], $B - L$ and multiplicities in $\overline{\mathbf{126}}_H$ and $\mathbf{120}_H$ dimensional scalars of $SO(10)$.

SM charges	Notation	$B - L$	$\overline{\mathbf{126}}_H$	$\mathbf{120}_H$
$(6, 1, \frac{1}{3})$	$S_{\beta\gamma}^\alpha$	$\frac{2}{3}$	1	1
$(\bar{6}, 1, -\frac{1}{3})$	$\tilde{S}_\alpha^{\beta\gamma}$	$-\frac{2}{3}$	0	1
$(6, 1, -\frac{2}{3})$	$\Sigma^{\alpha\beta}$	$\frac{2}{3}$	1	0
$(6, 1, \frac{4}{3})$	$S_{\beta\gamma}^\alpha$	$\frac{4}{3}$	1	0
$(\bar{6}, 3, -\frac{1}{3})$	$\mathbb{S}_{\gamma b}^{\alpha\beta a}$	$-\frac{2}{3}$	1	0

$$S_1 = c_\theta S + s_\theta \tilde{S}, \quad S_2 = -s_\theta S + c_\theta \tilde{S}. \quad (3)$$

Here, $c_\theta = \cos \theta$ and $s_\theta = \sin \theta$. The linear combination S_1 is to be identified with the lighter mass eigenstate, i.e., $M_{S_1} < M_{S_2}$.

Replacing S and \tilde{S} with the physical states $S_{1,2}$ using Eq. (3) and converting the quarks fields into their physical basis using $f \rightarrow U_f f$, the Yukawa couplings between the various sextet fields, $\Phi = \Sigma, S_i, \mathcal{S}, \mathbb{S}, \tilde{\mathbb{S}}$, and quarks can be rewritten as

$$\begin{aligned} -\mathcal{L}_\Phi = & Y_{AB}^\Sigma d_{\alpha A}^{CT} C^{-1} d_{\beta B}^C \Sigma^{\alpha\beta} + Y_{AB}^{S_i} \epsilon^{\alpha\beta\gamma} u_{\gamma A}^{CT} C^{-1} d_{\sigma B}^C S_{i\alpha\beta}^\sigma \\ & + Y_{AB}^{\mathcal{S}} \epsilon^{\alpha\beta\gamma} u_{\sigma A}^{CT} C^{-1} u_{\gamma B}^C \mathcal{S}_{\alpha\beta}^\sigma \\ & + Y_{AB}^{\mathbb{S}} \epsilon_{\alpha\beta\gamma} \epsilon_{ab} q_A^{a\alpha T} C^{-1} q_B^{\sigma c} \mathbb{S}_{\sigma c}^{\beta\gamma b} \\ & + Y_{AB}^{\tilde{\mathbb{S}}} \epsilon_{\alpha\beta\gamma} u_A^{\gamma T} C^{-1} d_B^\sigma \tilde{\mathbb{S}}_{\sigma}^{\alpha\beta} + \text{H.c.}, \end{aligned} \quad (4)$$

where $i = 1, 2$. The 3×3 matrices Y^Φ obtained from Eqs. (1) and (2) as

$$\begin{aligned} Y^\Sigma &= -\frac{2i}{\sqrt{15}} U_{d^c}^T F U_{d^c}, \\ Y^{S_1} &= \frac{\sqrt{2}i}{\sqrt{3}} \left(\frac{1}{\sqrt{5}} c_\theta U_{u^c}^T F U_{d^c} - \sqrt{2} s_\theta U_{u^c}^T G U_{d^c} \right), \\ Y^{S_2} &= -\frac{\sqrt{2}i}{\sqrt{3}} \left(\frac{1}{\sqrt{5}} s_\theta U_{u^c}^T F U_{d^c} + \sqrt{2} c_\theta U_{u^c}^T G U_{d^c} \right), \\ Y^{\mathcal{S}} &= -\frac{i}{\sqrt{15}} U_{u^c}^T F U_{u^c}, \\ Y^{\mathbb{S}} &= -\frac{\sqrt{2}i}{\sqrt{15}} U_q^T F U_{q'}, \quad Y^{\tilde{\mathbb{S}}} = \frac{2i}{\sqrt{3}} U_u^T G U_d. \end{aligned} \quad (5)$$

Here $q, q' = u, d$. The matrices $Y^\Sigma, Y^{\mathcal{S}}$ are symmetric in the flavor space while $Y^{\mathbb{S}}$ is symmetric when $q = q'$. Note that the matrices U_u and U_d determine the quark mixing matrix, $U_u^\dagger U_d \equiv V_{\text{CKM}}$. Therefore, $U_u \sim U_d$ serves as a good approximation and $Y^{\mathbb{S}}$ can be considered symmetric at the leading order.

The advantage of deriving the expressions in Eq. (5) is that all the Y^Φ can be explicitly computed in the realistic $SO(10)$ models in which the fundamental couplings F, G ,

and the diagonalizing matrices U_f are determined from the fermion mass fits. We now use the above couplings to determine various phenomenologically relevant processes involving sextet scalars in the subsequent sections.

III. QUARK FLAVOR VIOLATION

The Yukawa sector of viable $SO(10)$ models must consist of more than one GUT scalar. This implies that Y^Φ are not diagonal matrices and, therefore, the sextet scalars can lead to a new source of flavor violation in the quark sector. The strongest constraints on this type of new physics come from the $|\Delta F| = 2$ processes involving neutral meson-antimeson oscillations. Following the effective theory approach, we estimate the sextets-induced contributions to $K^0-\bar{K}^0, B_{d/s}^0-\bar{B}_{d/s}^0$, and $D^0-\bar{D}^0$ mixing at tree and one-loop levels.

Integrating out various sextets from Eq. (4) and parametrizing the effective Lagrangian as

$$\mathcal{L}_{\text{eff}}^{\Delta F=2} = \sum_{q=d,u} (c_q \mathcal{O}_q + \tilde{c}_q \tilde{\mathcal{O}}_q) + \text{H.c.}, \quad (6)$$

we find the following independent operators

$$\begin{aligned} \tilde{\mathcal{O}}_d &= \left(\bar{d}_{RA}^\alpha \gamma^\mu d_{RC}^\alpha \right) \left(\bar{d}_{RB}^\beta \gamma_\mu d_{RD}^\beta \right), \\ \tilde{\mathcal{O}}_u &= \left(\bar{u}_{RA}^\alpha \gamma^\mu u_{RC}^\alpha \right) \left(\bar{u}_{RB}^\beta \gamma_\mu u_{RD}^\beta \right), \\ \mathcal{O}_d &= \left(\bar{d}_{LA}^\alpha \gamma^\mu d_{LC}^\alpha \right) \left(\bar{d}_{LB}^\beta \gamma_\mu d_{LD}^\beta \right), \\ \mathcal{O}_u &= \left(\bar{u}_{LA}^\alpha \gamma^\mu u_{LC}^\alpha \right) \left(\bar{u}_{LB}^\beta \gamma_\mu u_{LD}^\beta \right). \end{aligned} \quad (7)$$

Here, we use $q = q_L$ and $q^C = C q_R^*$ for $q = u, d$ to obtain the above operators in the usual left- and right-chiral notations. $\mathcal{O}_{u,d}$ and $\tilde{\mathcal{O}}_{u,d}$ are related by $L \leftrightarrow R$. Operators \mathcal{O}_u and \mathcal{O}_d induce flavor-changing neutral meson-antimeson oscillations in the up-type and down-type quark sectors, respectively.

The coefficient of the operator \mathcal{O}_d is obtained as

$$\begin{aligned} \tilde{c}_d[A, B, C, D] = & \frac{1}{M_\Sigma^2} \left[\frac{1}{2} Y_{AB}^\Sigma (Y_{CD}^\Sigma)^* - \frac{6}{64\pi^2} \left[(Y_2^\Sigma)_{AC} (Y_2^\Sigma)_{BD} + (Y_2^\Sigma)_{AD} (Y_2^\Sigma)_{BC} \right] \right] \\ & - \frac{1}{M_{S_1}^2} \frac{6}{16\pi^2} \left[(\bar{Y}_2^{S_1})_{AC} (\bar{Y}_2^{S_1})_{BD} + (\bar{Y}_2^{S_1})_{AD} (\bar{Y}_2^{S_1})_{BC} \right], \end{aligned} \quad (8)$$

where $Y_2^\Phi = Y^\Phi Y^{\Phi\dagger}$ and $\bar{Y}_2^\Phi = Y^{\Phi\dagger} Y^\Phi$ are hermitian matrices. The first term in the above expression denotes the tree-level contribution mediated by Σ while the second and third terms are contributions generated at one loop by the scalars Σ and S_1 , respectively. Note that we have not included the contribution from S_2 , which is of a similar kind as S_1 , as it is generically expected to be suppressed since $M_{S_2} > M_{S_1}$.

Analogously, we find the following coefficients of the remaining operators.

$$\begin{aligned} \tilde{c}_u[A, B, C, D] = & \frac{1}{M_S^2} \left[Y_{AB}^S (Y_{CD}^S)^* - \frac{6}{16\pi^2} \left[(Y_2^S)_{AC} (Y_2^S)_{BD} + (Y_2^S)_{AD} (Y_2^S)_{BC} \right] \right] \\ & - \frac{1}{M_{S_1}^2} \frac{6}{16\pi^2} \left[(Y_2^{S_1})_{AC} (Y_2^{S_1})_{BD} + (Y_2^{S_1})_{AD} (Y_2^{S_1})_{BC} \right], \end{aligned} \quad (9)$$

$$\begin{aligned} c_d[A, B, C, D] = & \frac{1}{M_S^2} \left[Y_{AB}^S (Y_{CD}^S)^* - \frac{30}{16\pi^2} \left[(Y_2^S)_{AC} (Y_2^S)_{BD} + (Y_2^S)_{AD} (Y_2^S)_{BC} \right] \right] \\ & - \frac{1}{M_{\bar{S}}^2} \frac{6}{16\pi^2} \left[(\bar{Y}_2^{\bar{S}})_{AC} (\bar{Y}_2^{\bar{S}})_{BD} + (\bar{Y}_2^{\bar{S}})_{AD} (\bar{Y}_2^{\bar{S}})_{BC} \right], \end{aligned} \quad (10)$$

and

$$\begin{aligned} c_u[A, B, C, D] = & \frac{1}{M_S^2} \left[Y_{AB}^S (Y_{CD}^S)^* - \frac{30}{16\pi^2} \left[(Y_2^S)_{AC} (Y_2^S)_{BD} + (Y_2^S)_{AD} (Y_2^S)_{BC} \right] \right] \\ & - \frac{1}{M_{\bar{S}}^2} \frac{6}{16\pi^2} \left[(Y_2^{\bar{S}})_{AC} (Y_2^{\bar{S}})_{BD} + (Y_2^{\bar{S}})_{AD} (Y_2^{\bar{S}})_{BC} \right]. \end{aligned} \quad (11)$$

It can be seen from the above expressions of $\tilde{c}_{u,d}$ and $c_{u,d}$ that only the sextet Σ (\mathcal{S}) gives rise to flavor violations in the down-type (up-type) quark sectors while \mathbb{S} contributes in both the sectors at tree level. Contributions of S_1 and \bar{S} to the flavor violation arise only at the loop level. Nevertheless, the latter can be of a similar order as that of the tree level depending on the hierarchical structure of the Yukawa couplings.

For a quantitative estimation of the constraint on the masses and couplings of sextet scalars, relevant $c_{u,d}$ and $\tilde{c}_{u,d}$ need to be evolved from the mass scale of the integrated-out scalar, i.e., $\mu = M_\Phi$ where Eqs. (8)–(11) hold, down to the scale at which meson-antimeson mixing is experimentally determined. For K^0 - \bar{K}^0 oscillation, the relevant scale is $\mu = 2$ GeV and the renormalisation group evolution (RGE) evolved effective coefficient for $M_\Phi > m_t$ is given by [39]

$$C_K^1 = \left(0.82 - 0.016 \frac{\alpha_s(M_\Phi)}{\alpha_s(m_t)} \right) \left(\frac{\alpha_s(M_\Phi)}{\alpha_s(m_t)} \right)^{0.29} c_d[1, 1, 2, 2](M_\Phi). \quad (12)$$

The same equation is obtained for \tilde{C}_K^1 by replacing c_d with \tilde{c}_d in the above expression. Note that C_K^1 and \tilde{C}_K^1 are coefficients of effective operators, \mathcal{Q}_1 and $\tilde{\mathcal{Q}}_1$, given in [39] which are identical to our \mathcal{O}_d and $\tilde{\mathcal{O}}_d$, respectively, with $A = B = 1$ and $C = D = 2$. These operators do not mix with the other operators through RGE evolution and, therefore, one obtains a relatively simple expression, Eq. (12).

Similar arguments also follow for $B_{d/s}^0$ - $\bar{B}_{d/s}^0$ mixing. In this case, one obtains the relevant Wilson coefficients as

$$\begin{aligned} C_{B_d}^1 &= \left(0.865 - 0.017 \frac{\alpha_s(M_\Phi)}{\alpha_s(m_t)} \right) \left(\frac{\alpha_s(M_\Phi)}{\alpha_s(m_t)} \right)^{0.29} c_d[1, 1, 3, 3](M_\Phi), \\ C_{B_s}^1 &= \left(0.865 - 0.017 \frac{\alpha_s(M_\Phi)}{\alpha_s(m_t)} \right) \left(\frac{\alpha_s(M_\Phi)}{\alpha_s(m_t)} \right)^{0.29} c_d[2, 2, 3, 3](M_\Phi), \end{aligned} \quad (13)$$

at $\mu = m_b$ [40]. In the case of the charm mixing governed by D^0 - \bar{D}^0 oscillations, one finds [41]

$$C_D^1 = \left(0.837 - 0.016 \frac{\alpha_s(M_\Phi)}{\alpha_s(m_t)} \right) \left(\frac{\alpha_s(M_\Phi)}{\alpha_s(m_t)} \right)^{0.29} c_u[1, 1, 2, 2](M_\Phi), \quad (14)$$

at $\mu = 2.8$ GeV. Various coefficients then can be compared with the present limits obtained from a fit to the experimental data by UTFit Collaboration [41]. The present limits are

$$\begin{aligned} |C_K^1| &< 9.6 \times 10^{-13}, & |C_{B_d}^1| &< 2.3 \times 10^{-11}, \\ |C_{B_s}^1| &< 1.1 \times 10^{-9}, & |C_D^1| &< 7.2 \times 10^{-13}. \end{aligned} \quad (15)$$

The same upper bounds are also applicable on the corresponding \tilde{C}^1 .

IV. NEUTRON-ANTINEUTRON OSCILLATIONS

The color sextet fields can also induce a transition between the neutral baryons and their antiparticles. Unlike the flavor transitions discussed in the previous section, this requires a source of baryon number violation. In renormalizable $SO(10)$, the latter naturally arises from gauge invariant quartic couplings between three sextets and an SM singlet residing in $\overline{\mathbf{126}}_H$, namely σ , which carries $B - L = -2$ (see Table I in [15]). The vacuum expectation value (VEV) of σ breaks $B - L$ and generates masses for the right-handed neutrinos. It also gives rise to $B - L$ violating trilinear couplings between various sextet scalars which can induce neutral $n - \bar{n}$ oscillations through dimension-9 six-fermion operators [42–44].

To derive the effective operators relevant for neutral baryon-antibaryon oscillations, we first write the most general quartic interaction terms between three sextet scalars and σ allowed by the SM gauge symmetry and $B - L$. They are found as

$$\eta_{ij} \sigma \Sigma^{\alpha\beta} S_{i\alpha\gamma}^\gamma S_{j\beta\theta}^\theta, \quad (16)$$

$$\eta_2 \epsilon_{\alpha\gamma\theta} \sigma \Sigma^{\alpha\beta} \Sigma^{\gamma\sigma} S_{\beta\sigma}^\theta, \quad (17)$$

$$\eta_3 \sigma^* \Sigma_{\alpha\beta}^* S_{\sigma a}^{\alpha\sigma a} S_{\rho a}^{\beta\rho b}, \quad (18)$$

$$\eta_4 \sigma^* \Sigma_{\alpha\beta}^* \bar{S}_\sigma^{\alpha\sigma} \bar{S}_\rho^{\beta\rho}. \quad (19)$$

The terms in the first two lines above can arise from a quartic term $(\overline{\mathbf{126}}_H)^4$ and mixing between S and \tilde{S} . The third and fourth line terms can result from the gauge-invariant terms $(\overline{\mathbf{126}}_H^\dagger \overline{\mathbf{126}}_H)^2$ and $(\overline{\mathbf{126}}_H^\dagger \mathbf{120}_H)^2$, respectively. In Eqs. (16)–(19), all the sextet fields are written in a mass basis. Equation (16) denotes three distinct operators corresponding to $i, j = 1, 2$ with $\eta_{12} = \eta_{21}$. Since $M_{S_1} < M_{S_2}$, the dominant contribution to the neutral baryon-antibaryon transition generically arises from the term corresponding to η_{11} in Eq. (16). Therefore, we consider only the S_1 induced operators in the following.

Integrating out various color sextet fields from the operators listed in Eqs. (16)–(19) and their Hermitian conjugate terms and using the diquark couplings evaluated in Eqs. (1) and (2), we find the following effective Lagrangian that gives rise to the baryon-antibaryon oscillation at the leading order,

$$\mathcal{L}_{\text{eff}}^{|\Delta B|=2} = \sum_{i=1}^3 c_i \mathcal{O}_i + \text{H.c.}, \quad (20)$$

with

$$\begin{aligned} \mathcal{O}_1 &= \frac{1}{2} \epsilon^{\alpha\beta\gamma} \epsilon^{\sigma\rho\eta} (u_{\alpha A}^{C\dagger} C^{-1} d_{\beta B}^{C*}) (u_{\sigma C}^{C\dagger} C^{-1} d_{\rho D}^{C*}) (d_{\gamma E}^{C\dagger} C^{-1} d_{\eta F}^{C*}), \\ \mathcal{O}_2 &= \epsilon^{\alpha\beta\gamma} \epsilon^{\sigma\rho\eta} (d_{\alpha A}^{C\dagger} C^{-1} d_{\sigma B}^{C*}) (d_{\beta C}^{C\dagger} C^{-1} d_{\rho D}^{C*}) (u_{\gamma E}^{C\dagger} C^{-1} u_{\eta F}^{C*}), \\ \mathcal{O}_3 &= \frac{1}{2} \epsilon^{\alpha\beta\gamma} \epsilon^{\sigma\rho\eta} (u_{\alpha A}^\dagger C^{-1} d_{\beta B}^*) (u_{\sigma C}^\dagger C^{-1} d_{\rho D}^*) (d_{\gamma E}^{CT} C^{-1} d_{\eta F}^C). \end{aligned} \quad (21)$$

Here, the quark fields are written in a physical basis. The operators \mathcal{O}_1 and \mathcal{O}_2 arise from the quartic terms Eqs. (16) and (17), respectively. The remaining terms, Eqs. (18) and (19), both lead to a single operator denoted by \mathcal{O}_3 .

The coefficients c_i defined in Eq. (20) are determined as

$$\begin{aligned} c_1 &= \frac{2\eta_{11} v_\sigma}{M_\Sigma^2 M_{S_1}^4} (U_{d^c}^\dagger F^* U_{d^c})_{EF} \left[\frac{4ic_\theta^2}{15\sqrt{15}} (U_{u^c}^\dagger F^* U_{d^c})_{AB} (U_{u^c}^\dagger F^* U_{d^c})_{CD} + \frac{8is_\theta^2}{3\sqrt{15}} (U_{u^c}^\dagger G^* U_{d^c})_{AB} (U_{u^c}^\dagger G^* U_{d^c})_{CD} \right. \\ &\quad \left. - \frac{8\sqrt{2}ic_\theta s_\theta}{15\sqrt{3}} (U_{u^c}^\dagger F^* U_{d^c})_{AB} (U_{u^c}^\dagger G^* U_{d^c})_{CD} \right], \\ c_2 &= \frac{4i}{15\sqrt{15}} \frac{\eta_2 v_\sigma}{M_\Sigma^2 M_S^2} (U_{d^c}^\dagger F^* U_{d^c})_{AB} (U_{d^c}^\dagger F^* U_{d^c})_{CD} (U_{u^c}^\dagger F^* U_{u^c})_{EF}, \\ c_3 &= \frac{2v_\sigma}{M_\Sigma^2} (U_{d^c}^T F U_{d^c})_{EF} \left[\frac{24i\eta_3}{15\sqrt{15}M_S^4} (U_u^\dagger F^* U_d)_{AB} (U_u^\dagger F^* U_d)_{CD} - \frac{8i\eta_4}{3\sqrt{15}M_S^4} (U_u^\dagger G^* U_d)_{AB} (U_u^\dagger G^* U_d)_{CD} \right]. \end{aligned} \quad (22)$$

Here, $v_\sigma = \langle \sigma \rangle$ which breaks $B - L$ by two units. The unitary matrices U_f and U_{f^c} with $f = u, d$ denote rotations in the flavor space and they can be explicitly computed from the corresponding quark mass matrices as described in detail in [15]. It can be noted that c_1 and c_3 receive contributions from both $\overline{\mathbf{126}}_H$ and $\mathbf{120}_H$.

To identify the above operators with the ones listed in [45,46], we rewrite various quark fields in the left- and right-chiral notations using the following relations:

$$\psi^T C^{-1} \chi = \overline{(\psi_L)^c} \chi_L, \quad \psi^{C\dagger} C^{-1} \chi^{C*} = \overline{(\psi_R)^c} \chi_R. \quad (23)$$

Substituting the above in Eq. (21), we obtain

$$\begin{aligned} \mathcal{O}_1 &= \frac{1}{2} \epsilon^{\alpha\beta\gamma} \epsilon^{\sigma\rho\eta} (\overline{(u_{R\alpha A})^c} d_{R\beta B}) (\overline{(u_{R\sigma C})^c} d_{R\rho D}) (\overline{(d_{R\gamma E})^c} d_{R\eta F}), \\ \mathcal{O}_2 &= \epsilon^{\alpha\beta\gamma} \epsilon^{\sigma\rho\eta} (\overline{(d_{R\alpha A})^c} d_{R\sigma B}) (\overline{(d_{R\beta C})^c} d_{R\rho D}) (\overline{(u_{R\gamma E})^c} u_{R\eta F}), \\ \mathcal{O}_3^* &= \frac{1}{2} \epsilon^{\alpha\beta\gamma} \epsilon^{\sigma\rho\eta} (\overline{(u_{L\alpha A})^c} d_{L\beta B}) (\overline{(u_{L\sigma C})^c} d_{L\rho D}) (\overline{(d_{R\gamma E})^c} d_{R\eta F}). \end{aligned} \quad (24)$$

For $n-\bar{n}$ oscillations, one substitutes $A = B = C = D = E = F = 1$ in the above expressions of \mathcal{O}_i and c_i . In this case, the operators \mathcal{O}_1 and \mathcal{O}_3^* can straightforwardly be identified with \mathcal{O}_1 and \mathcal{O}_3 as listed in [45,46], respectively. Our \mathcal{O}_2 is proportional to \mathcal{O}_{RRR}^1 defined in [46] which in turn is a linear combination of the operators \mathcal{O}_1 and \mathcal{O}_4 also defined in [46]. Explicitly,

$$\mathcal{O}_2 = \frac{1}{4} \mathcal{O}_{RRR}^1 = \frac{1}{5} (\mathcal{O}_4 - 12\tilde{\mathcal{O}}_1). \quad (25)$$

Notably, \mathcal{O}_4 has vanishing nuclear matrix element [47] while $\tilde{\mathcal{O}}_1$ has nuclear matrix element identical to that of \mathcal{O}_1 . Therefore, the operator \mathcal{O}_2 is directly related to the operator \mathcal{O}_1 in our case. This leaves only two linearly-independent operators, \mathcal{O}_1 and \mathcal{O}_3 , as listed in Eq. (24).

The operators $\mathcal{O}_{1,3}$ need to be evolved from the scale of sextet masses, namely $\mu = M_\Phi$, down to $\mu_0 = 2$ GeV, where nuclear matrix elements are computed using the lattice calculations. The noteworthy feature about the basis in which $\mathcal{O}_{1,3}$ are written is that they do not mix through renormalization group evolution. The mean lifetime of the $n-\bar{n}$ transition can be computed in terms of the effective operators as

$$\begin{aligned} \tau_{n\bar{n}}^{-1} &= \left| \sum_{i=1,3} \langle \bar{n} | \mathcal{O}_i(\mu_0) | n \rangle c_i(\mu_0) \right|, \\ &= \left| \sum_{i=1,3} \langle \bar{n} | \mathcal{O}_i(\mu_0) | n \rangle \left(\frac{\alpha_s^{(4)}(m_b)}{\alpha_s^{(4)}(\mu_0)} \right)^{\frac{3\gamma_i^{(0)}}{50}} \left(\frac{\alpha_s^{(5)}(m_t)}{\alpha_s^{(5)}(m_b)} \right)^{\frac{3\gamma_i^{(0)}}{46}} \right. \\ &\quad \left. \times \left(\frac{\alpha_s^{(6)}(M_\Phi)}{\alpha_s^{(6)}(m_t)} \right)^{\frac{\gamma_i^{(0)}}{14}} c_i(M_\Phi) \right|, \end{aligned} \quad (26)$$

where $\gamma_i^{(0)}$ is the leading-order anomalous dimension of operator \mathcal{O}_i and $\alpha_s^{(n_f)}$ is the strong coupling constant with n_f flavors of the light quarks. We have $\gamma_1^{(0)} = 4$ and $\gamma_3^{(0)} = 0$ from [46]. Using the lattice calculation results from [47] and parametrizing the running effects as in [45], one finally finds

$$\begin{aligned} \tau_{n\bar{n}}^{-1} &= \left| 0.760 \left(\frac{\alpha_s(M_\Phi)}{\alpha_s(10^5 \text{ GeV})} \right)^{2/7} \left(c_1 - \frac{12}{5} c_2 \right) \right. \\ &\quad \left. + 1.08 c_3^* \right| \Lambda_{\text{QCD}}^6. \end{aligned} \quad (27)$$

The $n-\bar{n}$ oscillation time in the given model can be explicitly computed by substituting $c_{1,3}$ from Eq. (22) in the above expression.

Note that $(U_u^\dagger G^* U_d^*)_{11} \neq 0$ despite of G being antisymmetric in the flavor space. This follows from the fact that $U_u \neq U_d$ in general, which in fact is necessarily required by the realistic quark mixing. Therefore, the color sextet scalars from $\mathbf{120}_H$ can induce nonvanishing contribution to $n-\bar{n}$ oscillation.

V. PERTURBATIVITY OF THE EFFECTIVE QUARTIC COUPLINGS

It can be seen from the expressions of c_i in Eq. (22) that the maximization of $n-\bar{n}$ transition rate would require large v_σ and at least two color sextet fields at the low scale. However, this possibility is known to lead to large negative effective quartic couplings for the light scalars [48]. It arises from the correction induced by the trilinear terms which get generated when σ acquires VEV in Eqs. (16)–(19). To quantify the constraint on the masses of the underlying scalars, we consider the first term, Eq. (16), and compute the correction to the quartic coupling of Σ arising from the diagram shown in Fig. 1.

The effective quartic coupling, arising from a tree-level coupling λ and the higher-order corrections, can be written as

$$-i\lambda_{\text{eff}} = -i\lambda + i\delta\lambda. \quad (28)$$

$\delta\lambda$ computed from the diagram shown in Fig. 1, for the vanishing external momentum, is given by

$$i\delta\lambda = 4 \sum_{i,j,k,l} \eta_{ij} \eta_{jk}^* \eta_{kl} \eta_{li}^* v_\sigma^4 \underbrace{\int \frac{d^4 p}{(2\pi)^4} \frac{1}{(p^2 - M_{S_i}^2)(p^2 - M_{S_j}^2)(p^2 - M_{S_k}^2)(p^2 - M_{S_l}^2)}}_{\equiv I_{ijkl}}. \quad (29)$$

A straightforward computation for $i, j, k, l = 1, 2$ leads to

$$\begin{aligned} i \frac{\delta\lambda}{4v_\sigma^4} &= |\eta_{11}|^4 I_{1111} + |\eta_{22}|^4 I_{2222} \\ &+ 4|\eta_{12}|^2 (|\eta_{11}|^2 I_{1112} + |\eta_{22}|^2 I_{2221}) \\ &+ (2|\eta_{12}|^4 + 4\text{Re}[\eta_{11}\eta_{22}\eta_{12}^*\eta_{12}^*]) I_{1122}, \end{aligned} \quad (30)$$

where the different integrals are determined as [49]

$$\begin{aligned} I_{iiii} &= \frac{i}{4\pi^2} \frac{1}{24M_{S_i}^4}, \\ I_{iiij} &= \frac{i}{4\pi^2} \frac{M_{S_i}^4 - M_{S_j}^4 + 2M_{S_i}^2 M_{S_j}^2 \log(M_{S_i}^2/M_{S_j}^2)}{8M_{S_i}^2 (M_{S_i}^2 - M_{S_j}^2)^3}, \\ I_{iiij} &= \frac{i}{4\pi^2} \frac{1}{4(M_{S_i}^2 - M_{S_j}^2)^2} \left(\frac{M_{S_i}^2 + M_{S_j}^2}{M_{S_i}^2 - M_{S_j}^2} \log\left(\frac{M_{S_i}^2}{M_{S_j}^2}\right) - 2 \right). \end{aligned} \quad (31)$$

It can be seen that all integrals divided by i are positive. The only negative contribution to $\delta\lambda$ can come from the last term in Eq. (30). It is apparent that this term cannot cancel completely the positive contributions coming from the remaining terms. Therefore, $\delta\lambda$ is always positive. For a hierarchical $M_{S_1} \ll M_{S_2}$, one finds

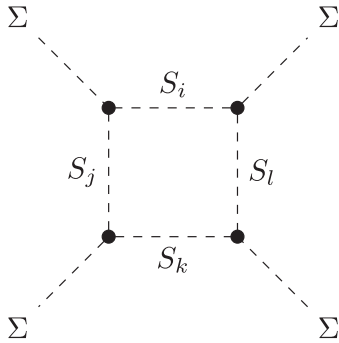


FIG. 1. Diagram representing a typical correction to the quartic coupling induced by trilinear coupling between three color sextet scalars. The vertex denoted by the bullet point is a trilinear coupling which arises from the VEV of $B - L$ charged scalar σ .

$$\lambda_{\text{eff}} \simeq \lambda - \frac{|\eta_{11}|^4}{24\pi^2} \frac{v_\sigma^4}{M_{S_1}^4}. \quad (32)$$

Since both λ_{eff} and λ are required to be positive and perturbative, the above leads to a constraint

$$M_{S_1} \geq \frac{|\eta_{11}| v_\sigma}{(24\pi^2)^{1/4}}, \quad (33)$$

if $M_\Sigma < v_\sigma$. For $M_\Sigma > v_\sigma$, the above constraint does not apply as Σ has to be integrated out first and the effective theory below v_σ does not contain a quartic term for Σ .

It is straightforward to generalize the above discussion for the remaining sextet scalars and their interactions given in Eqs. (16)–(19). Generically, two or more sextet scalars coupled through $B - L$ violating vertex leads to unstable potential if they are all well below the $B - L$ breaking scale. To evade this situation, at least one of these sextets is required to be heavier than the $B - L$ breaking scale and this in turn leads to relatively suppressed $n - \bar{n}$ transition rate in the models with high $B - L$ breaking scale.

VI. BARYOGENESIS

We now point out a possibility of generating baryon asymmetry using the color sextet scalars and their interactions predicted within this framework. The mechanism relies on the fact that the $B - L$ violating interactions present in the model can generate baryon asymmetry that avoids washout by the electroweak sphalerons [50,51]. The sextet scalars relevant for this are Σ , S , \tilde{S} (or $S_{1,2}$ in the physical basis) and their interactions extracted from Eqs. (4) and (16) are summarized as

$$\begin{aligned} -\mathcal{L} \supset & Y_{AB}^\Sigma d_{\alpha A}^{CT} C^{-1} d_{\beta B}^C \Sigma^{\alpha\beta} + Y_{AB}^{S_i} \epsilon^{\alpha\beta\gamma} u_{\gamma A}^{CT} C^{-1} d_{\sigma B}^C S_{i\alpha\beta}^\sigma \\ & + \eta_{ij} \sigma \Sigma^{\alpha\beta} S_{i\alpha\gamma}^\gamma S_{j\beta\theta}^\theta + \text{H.c.} \end{aligned} \quad (34)$$

The above Lagrangian contains all the necessary conditions for baryogenesis [52]. It inherently violates P while CP violation can arise from the phases in η_{ij} as described in detail below. The VEV of σ gives rise to $B - L$ violation as mentioned earlier. Departure from the thermal equilibrium

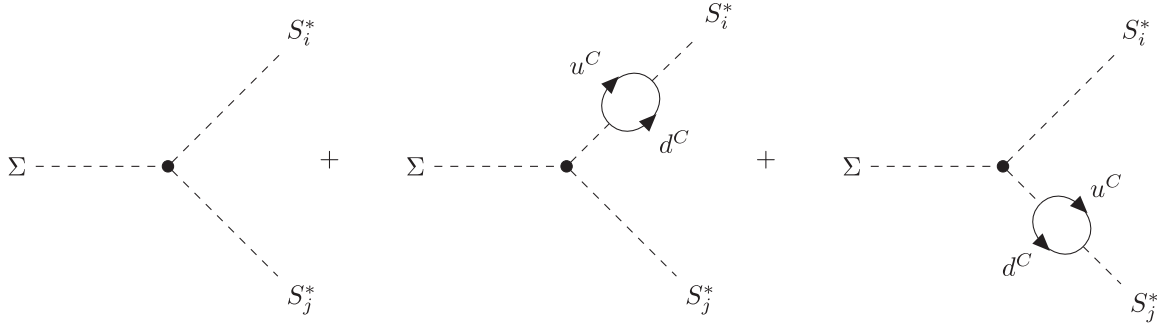


FIG. 2. The tree and one-loop diagrams involving decays of the color sextet scalar give rise to CP asymmetry. The vertex denoted by the bullet point is a trilinear coupling induced by the VEV of $B - L$ charged scalar σ .

is arranged through out-of-equilibrium decays of Σ in the expanding universe as discussed below.

We assume the mass hierarchy $M_\Sigma \gg M_{S_2} > M_{S_1} \gg m_t$. The relevant processes for baryogenesis, as can be read from Eq. (34), are categorized as the following:

- (i) B conserving: decay $\Sigma \rightarrow d_A^C d_B^C$ and scatterings $\Sigma S_i^* \rightarrow d_A^C u_B^C$, $\Sigma u^C \rightarrow S_i d^C$ and $\Sigma d_A^C \rightarrow S_i u_B^C$,
- (ii) B violating: decay $\Sigma \rightarrow S_i^* S_j^*$ and scatterings $S_i^* S_j^* \rightarrow d_A^C d_B^C$, $\Sigma S_i \rightarrow d_A^C u_B^C$, $S_i^* d_A^C \rightarrow S_j d_B^C$, $\Sigma u_A^C \rightarrow S_i^* d_B^C$ and $\Sigma d_A^C \rightarrow S_i^* u_B^C$,

along with their CP conjugate and inverse processes.

Concentrating on the B violating decay modes of Σ , the CP asymmetry in the decay $\Sigma \rightarrow S_i^* S_j^*$ is defined as

$$\epsilon_{ij} = \frac{\Gamma[\Sigma \rightarrow S_i^* S_j^*] - \Gamma[\Sigma^* \rightarrow S_i S_j]}{\Gamma_{\text{tot}}[\Sigma]}. \quad (35)$$

Nonzero ϵ_{ij} can be generated from interference between a tree and a one-loop diagram which has an absorptive part. These diagrams are shown in Fig. 2. We compute these diagrams and find the leading order CP asymmetry as

$$\epsilon_{ij} = -\frac{1}{\pi} \frac{\text{BR}[\Sigma \rightarrow S_i^* S_j^*]}{|\eta_{ij}|^2} \times \left(\sum_{k \neq i} \frac{x_{i/k}}{1 - x_{i/k}} \text{Im}[\eta_{ij}^* \eta_{jk} \text{Tr}[Y^{S_i^\dagger} Y^{S_j}]] + i \leftrightarrow j \right), \quad (36)$$

where $x_{i/k} = M_{S_i}^2 / M_{S_k}^2$. Note that $\epsilon_{ij} = \epsilon_{ji}$ as $\eta_{ij} = \eta_{ji}$ in Eq. (34). The total CP asymmetry produced in decays of Σ is then given by

$$\epsilon = \epsilon_{11} + \epsilon_{22} + 2\epsilon_{12}, \quad (37)$$

for the present case. The asymmetry generated between the number densities of S_i and S_i^* due to the out-of-equilibrium decay of Σ gets further redistributed into the SM quarks through the decays of S_i . This happens at the temperatures

below the freeze-out of the asymmetry as $M_{S_i} \ll M_\Sigma$ and when S_i leaves thermal equilibrium.

The final baryon-to-entropy ratio is obtained as

$$Y_B = \frac{4}{3} \epsilon \frac{\kappa(K)}{g_*}, \quad (38)$$

where the factor of $4/3$ is the total $B - L$ quantum number of the final states in the decay of Σ . g_* is the effective number of the relativistic degrees of freedom at the time of decay. We have $g_* \simeq 125$ which include Σ , $S_{1,2}$ and the SM particles. $\kappa(K)$ is an efficiency factor which accounts for the washout of the asymmetries due to inverse decays and scattering processes listed above. K is a decay parameter which is a measure of out-of-equilibrium condition and it is defined as

$$K = \frac{\Gamma[\Sigma \rightarrow S_i^* S_j^*]}{2H(M_\Sigma)}, \quad (39)$$

with Hubble parameter

$$H(T) = 1.66 g_*^{1/2} \frac{T^2}{M_P}. \quad (40)$$

An exact value of κ is to be obtained by numerically solving the full Boltzmann equations as outlined in [20,53]. Nevertheless, an approximate analytical solution for κ exists which is suitable for the present setup. Assuming the initial thermal abundance for Σ , it is given by [54]

$$\kappa(x) = \frac{2}{xz_B(x)} \left(1 - \exp \left[-\frac{1}{2} x z_B(x) \right] \right), \quad (41)$$

with

$$z_B(x) = 2 + 4x^{0.13} \exp \left[-\frac{2.5}{x} \right]. \quad (42)$$

The above solution of κ takes into account the washout effects only by the inverse decay and it is valid for $K \leq 10^3$.

For $K > 10^3$, the scattering processes become more important and κ decreases exponentially. Note that for $K \leq 1$, $\kappa(K) \rightarrow 1$ implies no dilution in the baryon asymmetry due to washout.

We aim to show that there is enough CP violation available through Eq. (37) in the present framework such that it can account for the observed baryon to entropy ratio $Y_B^{\text{exp}} = (6.10 \pm 0.04) \times 10^{-10}$ [55]. For this, we first find the maximum possible value of ϵ and then evaluate the amount of damping permitted through Eq. (38) requiring that $Y_B \geq 6.0 \times 10^{-10}$. Assuming $\eta_{ij} = |\eta|e^{i\phi_{ij}}$ and $\text{BR}[\Sigma \rightarrow S_i^* S_j^*] \simeq 1/4$ for all i and j in Eqs. (36) and (37), we get (after some straightforward algebra),

$$\epsilon = \frac{1}{2\pi} \left(\frac{1 + x_{1/2}}{1 - x_{1/2}} \right) |\text{Tr}[Y^{S_{11}^\dagger} Y^{S_2}]| (\sin(\phi_{12} - \phi_{22} - \phi_Y) - \sin(\phi_{12} - \phi_{11} + \phi_Y)), \quad (43)$$

where $\phi_Y = \text{Arg}(\text{Tr}[Y^{S_{11}^\dagger} Y^{S_2}])$ is the phase that arise from the Yukawa couplings. For $x_{1/2} \ll 1$, maximization of ϵ leads to

$$\epsilon_{\text{max}} \simeq \frac{1}{\pi} |\text{Tr}[Y^{S_{11}^\dagger} Y^{S_2}]|. \quad (44)$$

Substituting $Y^{S_{1,2}}$ from Eq. (5) and using the fact that $\text{Tr}[F^\dagger G] = 0$ due to symmetric and antisymmetric properties of F and G respectively, we find

$$\epsilon_{\text{max}} \simeq \frac{2}{3\pi} \left| \sin 2\theta \text{Tr} \left[G^\dagger G - \frac{1}{10} F^\dagger F \right] \right|. \quad (45)$$

For $M_{S_i} \ll M_\Sigma$, one finds $\Gamma[\Sigma \rightarrow S_i^* S_j^*] \simeq \frac{|\eta_{ij}|^2 v_\sigma^2}{16\pi M_\Sigma}$. Substituting this in K and setting $|\eta_{ij}| = |\eta|$, we get

$$K = 1.07 \times \left(\frac{|\eta| v_\sigma}{10^{15} \text{ GeV}} \right)^2 \times \left(\frac{10^{15} \text{ GeV}}{M_\Sigma} \right)^3. \quad (46)$$

Using the above in $\kappa(K)$, the baryon to entropy ratio Y_B can be computed for a given value of ϵ_{max} . Demanding that $Y_B \geq 6.0 \times 10^{-10}$, the allowed regions are shown in Fig. 3. We observe that M_Σ cannot be much lighter than the scale of $B-L$ breaking as it leads to washout of the baryon asymmetry through very large K .

The possibility of generating baryon asymmetry discussed in this section utilizes the sextets from the GUT scalars of the Yukawa sectors only and requires the presence of both $\overline{\mathbf{126}}_H$ and $\mathbf{120}_H$. Alternatively, a similar mechanism also works if only one of them is present. However, this requires an additional copy of Σ -like sextet which can emerge from **54**-dimensional GUT scalar [19]. Note that baryon asymmetry can also be generated through thermal leptogenesis as the lepton number violation and right-handed neutrinos are inherently present in the $SO(10)$ GUTs [56]. The cases in

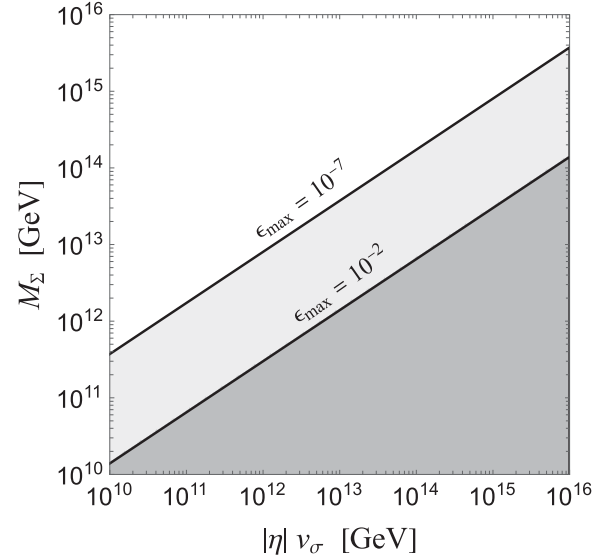


FIG. 3. Values of M_Σ and $|\eta|v_\sigma$ excluded by $Y_B > 6.0 \times 10^{-10}$ for $\epsilon_{\text{max}} = 10^{-7}$ (lighter gray) and $\epsilon_{\text{max}} = 10^{-2}$ (darker gray).

which the latter cannot account for the observed asymmetry due to the absence of required CP violation in the lepton sector and/or suitable mass spectrum and couplings of the right-handed neutrinos, sextets governed baryogenesis discussed above can provide a viable alternative.

VII. RESULTS

We now discuss constraints on the mass scales of the various color sextet scalars from the observables quantified in the previous sections. Our emphasis is on the realistic $SO(10)$ models which are known to reproduce the observed fermion mass spectrum. It can be noted from the discussions so far that the various phenomena related to sextets involve two important parameters: (a) the Yukawa couplings with the quarks, i.e. the matrices F and G , and (b) the $B-L$ breaking VEV v_σ . In the $SO(10)$ models with minimal choices for the scalars in the Yukawa sector, both (a) and (b) are determined from the realistic fits to the quarks and lepton masses and mixing observables [5,56–62]. Using the results of the latest fit [56] performed for a nonsupersymmetric $SO(10)$ model with $\mathbf{10}_H$ and $\overline{\mathbf{126}}_H$ in the Yukawa sector, one typically finds

$$F \sim \frac{\lambda^4}{\alpha_2} \begin{pmatrix} \lambda^5 & \lambda^4 & \lambda^3 \\ \lambda^4 & \lambda^3 & \lambda^2 \\ \lambda^3 & \lambda^2 & \lambda \end{pmatrix}, \quad G \sim \frac{\lambda^4}{\alpha_3} \begin{pmatrix} 0 & \lambda^4 & \lambda^3 \\ -\lambda^4 & 0 & \lambda^2 \\ -\lambda^3 & -\lambda^2 & 0 \end{pmatrix}, \quad (47)$$

where $\lambda = 0.23$ is Cabibbo angle and we have suppressed $\mathcal{O}(1)$ coefficients of each of the elements of F and G . Note that the above form of F is taken directly from [56] while for

G we assume that its nonzero elements are of similar magnitude as those of F as observed in an earlier fit [57]. The parameters $\alpha_{2,3}$, with $|\alpha_{2,3}| \lesssim 1$ quantify the amount of Higgs doublet mixing as discussed in [56].

v_σ is determined by fitting the light neutrino masses and mixing assuming the dominance of type I seesaw mechanism and it implies

$$v_\sigma = \alpha_2 v'_S \simeq \alpha_2 \times 10^{15} \text{ GeV}, \quad (48)$$

where v'_S is explicitly defined in [56] and is found in the range $10^{14} - 10^{15}$ GeV. In addition to v_σ , F , and G , one also needs unitary matrices, U_f and U_{fc} , that diagonalize the quark mass matrices M_f for $f = u, d$. They are also determined from the fits in [56] and their generic forms are given by

$$U_u \sim U_d \sim U_{uc} \sim U_{dc} \sim \begin{pmatrix} 1 & \lambda & \lambda^3 \\ \lambda & 1 & \lambda^2 \\ \lambda^3 & \lambda^2 & 1 \end{pmatrix}. \quad (49)$$

Again, we have suppressed the coefficient of $\mathcal{O}(1)$ in writing the above.

For the subsequent analysis, we consider two example values for $\alpha_{2,3}$ as the following:

- (i) High-scale $B - L$ (HS): $\alpha_{2,3} = \lambda$. This implies

$$v_\sigma \simeq 10^{14} \text{ GeV}, \quad F \simeq \begin{pmatrix} \lambda^8 & \lambda^7 & \lambda^6 \\ \lambda^7 & \lambda^6 & \lambda^5 \\ \lambda^6 & \lambda^5 & \lambda^4 \end{pmatrix},$$

$$G \simeq \begin{pmatrix} 0 & \lambda^7 & \lambda^6 \\ -\lambda^7 & 0 & \lambda^5 \\ -\lambda^6 & -\lambda^5 & 0 \end{pmatrix}. \quad (50)$$

Since $\alpha_{2,3} = \lambda$, the lightest pair of the electroweak doublet Higgs contains a sizeable contribution from the doublets residing in $\overline{\mathbf{126}}_H$ and $\mathbf{120}_H$. F and G are required to be relatively small in this case. Note that $\alpha_{2,3}$ cannot be taken much greater than λ as in that case contribution of $\mathbf{10}_H$ to the fermion masses become negligible and it is known that $\overline{\mathbf{126}}_H$ and $\mathbf{120}_H$ alone cannot reproduce the observed fermion mass spectrum [57].

- (ii) Intermediate-scale $B - L$ (IS): $\alpha_{2,3} = \lambda^5$. This leads to

$$v_\sigma \simeq 10^{11} \text{ GeV}, \quad F \simeq \begin{pmatrix} \lambda^4 & \lambda^3 & \lambda^2 \\ \lambda^3 & \lambda^2 & \lambda \\ \lambda^2 & \lambda & 1 \end{pmatrix},$$

$$G \simeq \begin{pmatrix} 0 & \lambda^3 & \lambda^2 \\ -\lambda^3 & 0 & \lambda \\ -\lambda^2 & -\lambda & 0 \end{pmatrix}. \quad (51)$$

In this case, F and G can have relatively stronger couplings with the fermions as the lightest Higgs have suppressed contributions from the doublets residing in $\overline{\mathbf{126}}_H$ and $\mathbf{120}_H$.

A low-scale $B - L$ breaking VEV, corresponding to $v_\sigma \ll 10^{11}$, would require $\alpha_2 \ll \lambda^5$ and it makes some of the couplings in F nonperturbative within this class of realistic models. Further small α_2 with perturbative values of couplings in F implies that the charged fermion masses arise dominantly from $\mathbf{10}_H$ and $\mathbf{120}_H$. This has been disfavored by the fits [57]. Alternatively, the above correlation between v_σ and the Yukawa couplings can also be understood as follows. The right-handed neutrino mass matrix is given by $M_R = v_\sigma F$ in this realistic model. The order of magnitude of the elements of M_R is more or less determined by the light neutrino masses induced through the type I seesaw mechanism. This, therefore, implies smaller F for the near-GUT scale v_σ and relatively large F for intermediate values of v_σ .

Before we proceed to estimate the constraints on the sextet scalars for HS and IS cases discussed above, let us outline a model-independent limit on their masses from the direct search experiments. The color sextets can be pair-produced at the LHC from gluon fusion [63–65]. Unlike all the observables considered in this paper, this process does not depend on the couplings with quarks and, therefore, provides a robust limit on the masses of the sextets. Nonobservation of deviation from the SM results so far implies [65]

$$M_\Phi \geq 1 \text{ TeV}. \quad (52)$$

The other direct search methods, such as resonant production and single-top production, depend on the couplings of the sextet with quarks. They are known to provide relatively milder limits for small values of the couplings [20,66]. Therefore, we consider the above lower limit and study the other constraints in the mass range $10^3 - 10^{16}$ GeV of the sextet scalars.

A. Light Σ, S_1

First, we consider $1 \text{ TeV} \leq M_\Sigma, M_{S_i} < M_{\text{GUT}}$ while the remaining sextet fields stay close to the GUT scale. Using $|\eta_{11}| = 1$ and $\theta = \pi/4$ in c_1 given in Eq. (22), we compute the neutron-antineutron transition time from Eq. (27) for the high- and intermediate-scale $B - L$ symmetry as described above. The relevant Yukawa couplings are evaluated using Eqs. (49)–(51) which give

$$|(U_{dc}^\dagger F^* U_{dc}^*)_{11}| \sim |(U_{uc}^\dagger F^* U_{dc}^*)_{11}| \sim |(U_{uc}^\dagger G^* U_{dc}^*)_{11}| \simeq \lambda^8,$$

$$|(U_{dc}^\dagger F^* U_{dc}^*)_{11}| \sim |(U_{uc}^\dagger F^* U_{dc}^*)_{11}| \sim |(U_{uc}^\dagger G^* U_{dc}^*)_{11}| \simeq \lambda^4, \quad (53)$$

for the HS and IS cases, respectively. We also compute contributions of Σ and S_1 to the meson-antimeson mixing using the derived expressions, Eqs. (12)–(14), and impose the constraints from Eq. (15). For the matching scale, we use $M_\Phi = (M_\Sigma + M_{S_1})/2$. The limits on M_Σ and M_{S_1} arising from meson-antimeson and neutron-antineutron oscillations are displayed in Fig. 4. We also indicate in, Fig. 4, a region which disfavors simultaneously light Σ and S_1 due to nonperturbativity of the effective quartic coupling as discussed in Sec. V.

As it can be seen from Fig. 4, either Σ or S_1 can be substantially lighter than the GUT scale in the case of high-scale v_σ . The perturbativity of quartic coupling forbids both of them from being lighter than 10^{13} GeV, simultaneously. Feeble couplings with quarks allow TeV scale Σ or S_1 to remain practically unconstrained from the $|\Delta F| = 2$ or $|\Delta B| = 2$ processes. For $v_\sigma = 10^{11}$ GeV, M_{S_1} (M_Σ) can be as light as 1 (10) TeV provided the other sextet is heavier than v_σ . Light M_{S_1} , in this case, predict relatively faster neutron-antineutron transition time in comparison to M_Σ as can be seen from the right panel in Fig. 4. Overall, the constraint imposed by quartic coupling's perturbativity almost rules out the possibility of observing $n-\bar{n}$ in near-future experiments for both HS and IS cases.

Sub-GUT scale Σ and S_i can also account for the baryon asymmetry of the Universe through the thermal baryogenesis as discussed in Sec. VI. The maximum CP asymmetry obtained using Eq. (45) for the two cases discussed above is found to be

$$\epsilon_{\max} = \begin{cases} \mathcal{O}(10^{-7}) & \text{for HS} \\ \mathcal{O}(10^{-2}) & \text{for IS.} \end{cases} \quad (54)$$

As it can be read from Fig. 3, the above implies

$$M_\Sigma \gtrsim \begin{cases} 2 \times 10^{14} \text{ GeV} & \text{for HS} \\ 5 \times 10^{10} \text{ GeV} & \text{for IS} \end{cases} \quad (55)$$

such that $Y_B \geq 6.0 \times 10^{-10}$. This region favored by sextet-generated baryogenesis is also shown in Fig. 4. For the very light S_1 and $v_\sigma \simeq 10^{11}$ GeV, this region can be probed through improved measurements of $n-\bar{n}$ oscillations.

B. Light Σ , S

Assuming $|\eta_2| = 1$ in Eq. (17), we now assess the constraints on light Σ and S assuming the remaining sextets at the GUT scale. Unlike $S_{1,2}$, S couples to the only up-type quarks and mediates $D^0-\bar{D}^0$ oscillations at the tree level. This puts severe constraints on the strongly coupled TeV scale S . The constraints derived from various considerations are displayed in Fig. 5. It can be seen that the limits from the meson-antimeson oscillations and perturbativity of the effective quartic coupling imply no observable $n-\bar{n}$ transition rate in the near-future experiments in the realistic renormalizable $SO(10)$ based models.

C. Light Σ , S

Next, we consider light Σ and S with the remaining sextets decoupled at the GUT scale. The results are shown

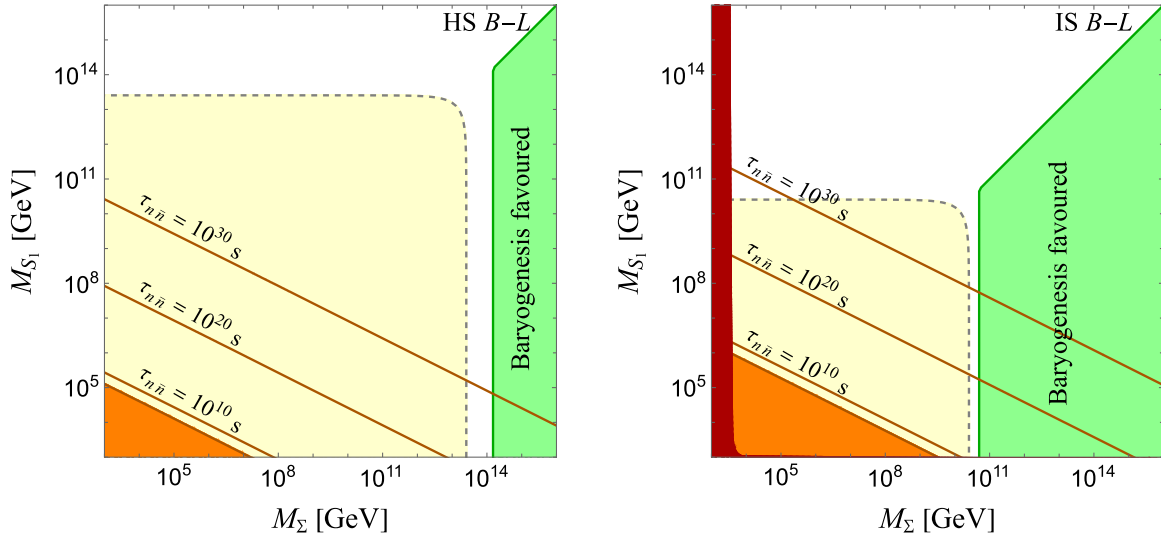


FIG. 4. Constraints on the masses of Σ and S_1 for the high (left panel) and intermediate (right panel) $B-L$ breaking scale. The red regions are excluded by the most dominant constraint from the neutral meson-antimeson oscillations. The region shaded in orange is excluded by the current limit $\tau_{n\bar{n}} > 4.7 \times 10^8$ seconds [33]. The diagonal red lines, from bottom to top, correspond to $\tau_{n\bar{n}} = 10^{10}$, 10^{20} , and 10^{30} seconds, respectively. The yellow region bounded by the dashed contour is disfavored by the perturbativity of the effective quartic couplings while the region shaded in green is favored by the baryogenesis constraints.

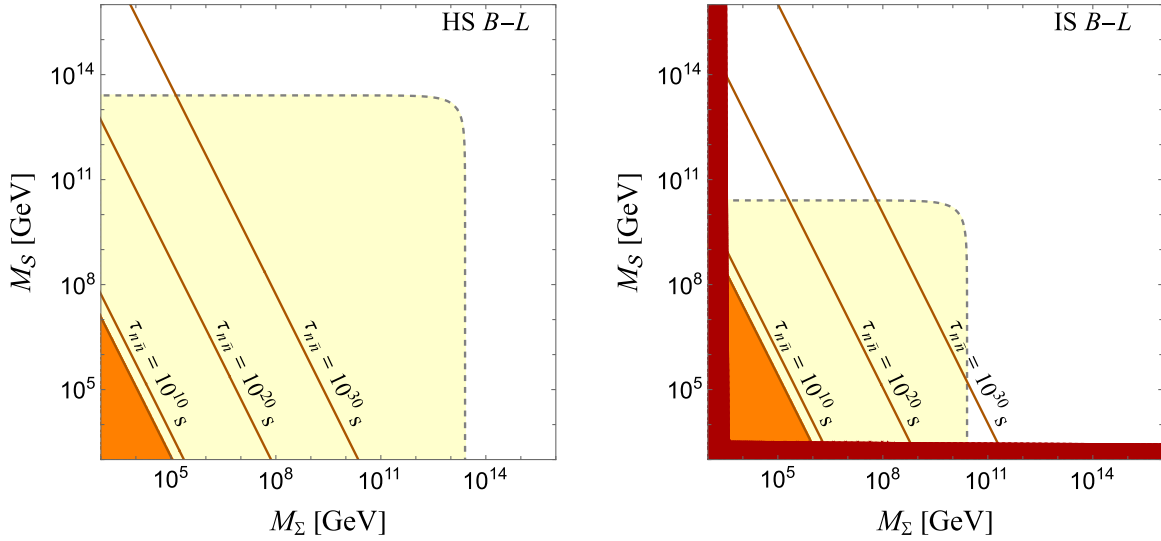


FIG. 5. Constraints on the masses of Σ and \mathcal{S} for the high (left panel) and intermediate (right panel) $B-L$ breaking scale. The other details are the same as the caption of Fig. 4.

in Fig. 6. We set $|\eta_3| = 1$ for this analysis and consider two cases for the values of F and v_σ as discussed earlier. Unlike S_i , the electroweak triplet \mathcal{S} mediate quark flavor-violating interactions at tree level in the both up- and down-type quark sector. This results in a relatively large upper bound on $M_{\mathcal{S}}$ in the case of the strong Yukawa coupling. The constraints from $n-\bar{n}$ oscillations and perturbativity of the effective quartic coupling are similar to the ones obtained in the case of light Σ and S_1 . Sub-GUT scale Σ and \mathcal{S} alone cannot generate baryon asymmetry and an additional copy of Σ -like scalar would be required if viable baryogenesis is to be realized in this case as discussed in Sec. VI.

D. Light $\Sigma, \bar{\mathcal{S}}$

Finally, we consider a case for $1 \text{ TeV} \leq M_\Sigma, M_{\bar{\mathcal{S}}} < M_{\text{GUT}}$ and the GUT scale masses for the remaining sextets. Setting $|\eta_4| = 1$, the obtained results are shown in Fig. 7. $\bar{\mathcal{S}}$ originates solely from $\mathbf{120}_H$ and it has flavor antisymmetric couplings with left-chiral up- and down-type quarks. As it can be seen from $Y^{\bar{\mathcal{S}}}$ in Eq. (5), its diagonal couplings vanish for $U_u = U_d$. The latter, however, is not supported by the nontrivial quark mixings. Hence, the effective diagonal couplings are generated by the quark mixing leading to a nonvanishing rate for $n-\bar{n}$ oscillations at the leading order. Like $S_{1,2}$, $\bar{\mathcal{S}}$ also contributes to the $|\Delta F| = 2$

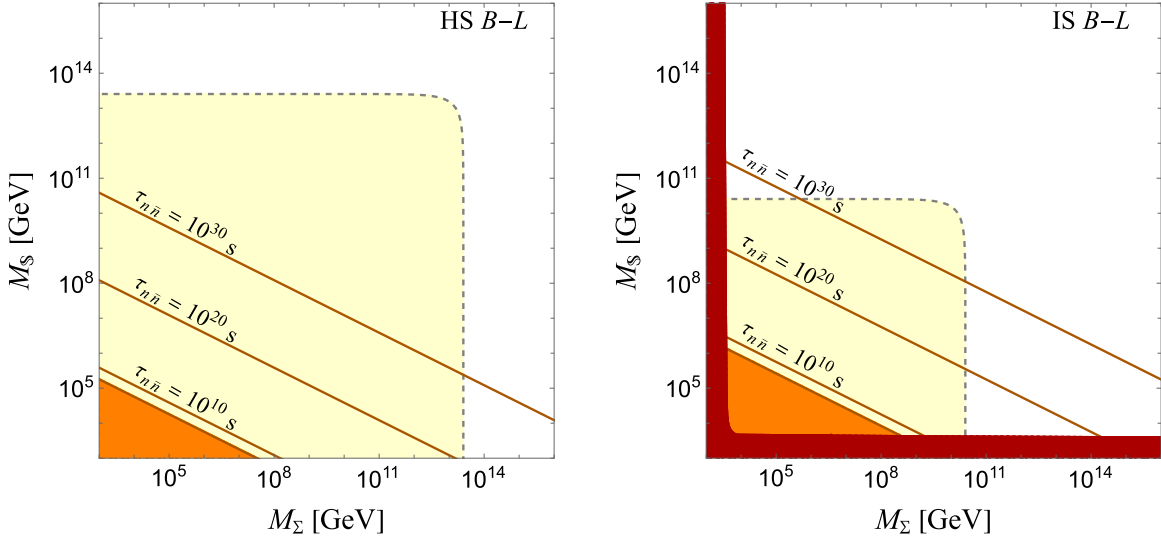


FIG. 6. Constraints on the masses of Σ and \mathcal{S} for the high (left panel) and intermediate (right panel) $B-L$ breaking scale. The other details are the same as the caption of Fig. 4.

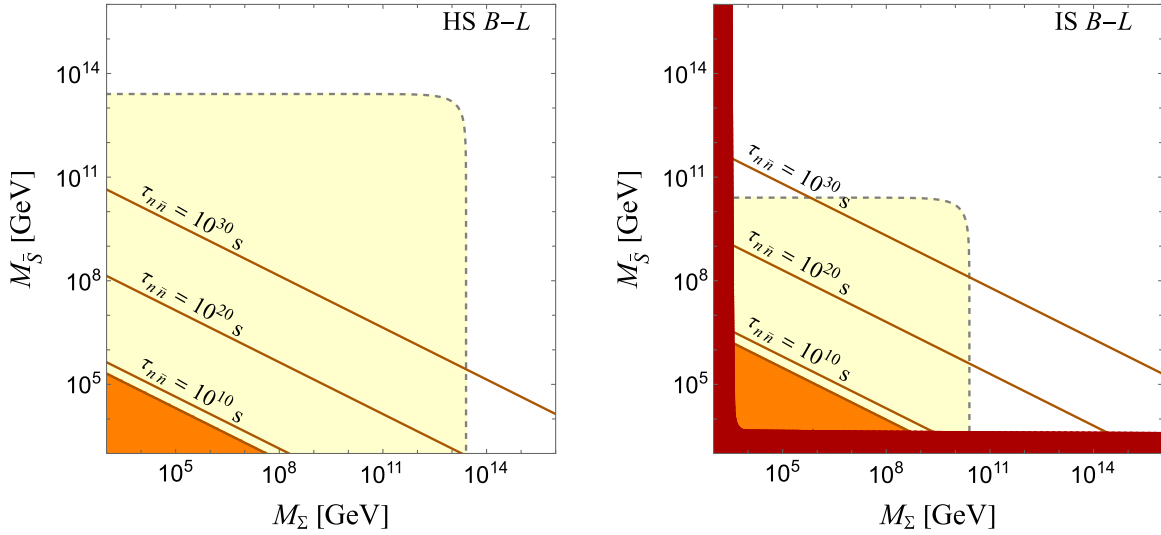


FIG. 7. Constraints on the masses of Σ and \bar{S} for the high (left panel) and intermediate (right panel) $B-L$ breaking scale. The other details are the same as the caption of Fig. 4.

processes at one-loop level. However, these processes put stronger constraints on the light \bar{S} in comparison to $S_{1,2}$ as can be seen from Figs. 4 and 7. This is attributed to the relatively large Yukawa coupling of \bar{S} due to different Clebsch-Gordan factors. Overall, the constraints on the masses of sub-GUT scale S and \bar{S} are similar.

VIII. CONCLUSIONS

Scalars which transform as two index symmetric representations of the $SU(3)_C$ of the SM gauge symmetry have been considered actively for their various phenomenological applications in bottom-up approaches. We evaluate the possibility of these sextets originating from the realistic renormalizable models of gauge and quark-lepton unification based on the $SO(10)$ symmetry. Five distinct color sextets are naturally accommodated in this class of models: $\Sigma \sim (6, 1, -\frac{2}{3})$, $S \sim (6, 1, \frac{1}{3})$, $\bar{S} \sim (\bar{6}, 1, -\frac{1}{3})$, $\mathcal{S} \sim (6, 1, \frac{4}{3})$, and $\mathbb{S} \sim (\bar{6}, 3, -\frac{1}{3})$. Deriving their couplings with the quarks, we compute effective operators contributing to the electrically neutral meson-antimeson and baryon-antibaryon oscillations. The latter arises from $B-L$ breaking induced by VEV, v_σ , of an SM singlet field σ . We also evaluate the effective quartic coupling of the sextet scalars which is prone to receive large contributions from $B-L$ breaking effects.

The noteworthy points from our present study are the following:

- (i) Four pairs of the sextets, i.e. Σ - S , Σ - \mathcal{S} , Σ - \mathbb{S} and Σ - \bar{S} , can give rise to the lifetime of neutron-antineutron transitions observable in near-future experiments provided that both the sextets in a pair are lighter than 10^5 – 10^8 GeV. This possibility is almost entirely excluded by the perturbativity of the quartic couplings in the models with $v_\sigma \geq 10^{11}$ GeV.

- (ii) Observable n - \bar{n} oscillation along with perturbative effective quartic couplings can be achieved if $v_\sigma < 10^8$ GeV and couplings of the sextets with quarks are of $\mathcal{O}(1)$. However, this generically leads to relatively light right-handed neutrino masses inconsistent with the type I seesaw mechanism in the realistic $SO(10)$ models.
- (iii) For $M_\Sigma > 10^{11}$ GeV, S can be as light as of $\mathcal{O}(\text{TeV})$ while the masses of S , \bar{S} , and \mathcal{S} can be $\gtrsim 10$ TeV. Similarly, for S , S , \bar{S} , and \mathcal{S} heavier than 10^{11} GeV, M_Σ can be of $\mathcal{O}(10)$ TeV or heavier. The lower limits on the masses of these light sextets come almost entirely from meson-antimeson oscillations and/or from direct searches.
- (iv) $SO(10)$ models with both $\overline{\mathbf{126}}_H$ and $\mathbf{120}_H$ in the Yukawa sector leads to the existence of a pair of sextets, $S_{1,2}$, with quantum numbers identical to that of S . Σ heavier than v_σ and $S_{1,2} \ll M_\Sigma$, in this case, provides a novel and viable possibility of generating baryon asymmetry of the Universe.

Many of the above observations follow from the fact that the couplings of sextets with quarks and the $B-L$ scale are strongly correlated in the renormalizable class of $SO(10)$ GUTs and they cannot take arbitrary values as typically assumed in the bottom-up approaches. On the other hand, a positive signal of n - \bar{n} oscillation in near-future experiments will rule out this simplest and predictive framework of grand unification. This study provides an interesting example of how a well-defined model in the ultraviolet can lead to a restrictive class of new physics at low energies making the former a falsifiable theory.

In the present work, our aim has been to study the phenomenological constraints on the color sextet scalars which have direct couplings with the quarks governed by the

$SO(10)$ grand unification. We do not discuss the constraints on the scalar mass spectra which may arise from the full scalar potential and requirement of consistent symmetry breaking. A careful treatment of this requires the specification of a full model beyond the Yukawa sector and it is a highly model-dependent exercise. Nevertheless, availing the freedom to choose a suitable set of GUT scalars and with an appropriate choice and tuning of parameters in the scalar potential, it is expected that the desired scalar mass spectra can be realized in concrete models.

ACKNOWLEDGMENTS

We acknowledge fruitful discussions with Namit Mahajan. S.K.S. also thanks Dayanand Mishra and Gurucharan Mohanta for useful discussions. This work is partially supported under the Mathematical Research Impact Centric Support (MATRICS) Project No. MTR/2021/000049 funded by the Science and Engineering Research Board (SERB), Department of Science and Technology (DST), Government of India.

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