$B - L$ violating nucleon decays as a probe of leptoquarks and implications for baryogenesis

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Abstract

We study the effective $B - L$ violating couplings for scalar and vector leptoquarks which can naturally induce dimension seven $B - L$ violating operators leading to very interesting nucleon decay modes such as $n \to e^- \pi^+, e^- K^+, \mu^- \pi^+, \mu^- K^+$ and $p \to \nu \pi^+$. This opens a new window to probe the nature and couplings of the scalar and vector leptoquarks in an ultraviolet model independent way which can provide an orthogonal probe for scalar and vector leptoquark solutions to the recent anomalous B-decay data. Furthermore, we also discuss how these $B - L$ violating interactions can also pave a new way to understand the origin of matter-antimatter asymmetry of the Universe.

1. Introduction

The renormalisable Standard Model (SM) Lagrangian interactions cannot violate baryon number, and consequently, baryon number violation can only arise through effective higher dimensional operators with dimension six or higher [1–3], making baryon number violation a very sensitive probe of new physics beyond the SM. The leading dimension six operators can natu-

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rally be generated in many grand unified theory (GUT) embedding of the SM and they conserve baryon number minus lepton number $(B - L)$ symmetry leading to proton decay modes such as $p \to e^+\pi^0$, $p \to \mu^+\pi^0$ and $p \to \bar{\nu}K^+$, which have been the primary focus of the recent experimental searches. On the other hand, $d = 7$ operators obeying the selection rule $\Delta(B - L) = -2$ [4,5] leading to interesting decay modes such as $n \to e^-\pi^+, e^-K^+, \mu^-\pi^+, \mu^-K^+$ and $p \to \nu\pi^+$ have received considerably less attention. These $B - L$ violating decay modes have been discussed in the context of $G_{\text{Pati-Salam}}$, and $SO(10)$ GUT theories in Ref. [6] and in the context of $SO(10)$ GUT theories in Refs. [7,8].[1] In these studies, driven by the GUT theory motivations, the mass scales of the mediating particles have been taken to be very heavy $\sim 10^{11} - 10^{13}$ GeV. In this work we point out that these $B - L$ violating nucleon decay modes provide a novel way to probe the nature and couplings of the scalar and vector leptoquarks in a model independent way.[2] This not only makes these $B - L$ violating nucleon decay modes extremely interesting for current and future experimental searches, but also gives way to a new type of mechanism to understand matter-antimatter asymmetry of the Universe where simultaneous baryon and lepton number violation gives rise to a $B - L$ violating asymmetry, which unlike $B - L$ conserving asymmetry, does not get washed out by $B + L$ violating sphaleron interactions near the electroweak scale [13]. Furthermore, this also motivates the experimental search for $B - L$ violating nucleon decay modes as an orthogonal probe for TeV scale scalar and vector leptoquark solutions to the B-decay anomalies, which are persistent with new data from the B-factories [14–26], and are currently one of the strongest hints of new physics beyond the standard model. To this end, the recent measurements of the ratio of branching fractions of $B \to K(K^*)\ell\ell$ decays into di-muons over di-electron modes $R_{K^{(*)}}$ and the ratio of branching fractions of $B \to D^{(*)}\ell^-\bar{\nu}$ decays into tau over other lepton modes $R_{D^{(*)}}$ are of particular interest because in these ratios the hadronic uncertainties cancel and consequently, these observables are sensitive to lepton flavour universality (LFU) violating new physics.

The latest measurement of R_K for the dilepton invariant mass squared bin $[1.1, 6]$ GeV^2 [22] shows a 2.5σ deviation falling short of the SM prediction [27,28]. Similarly, the measurement of R_{K^*} [23] display 2.3σ and 2.6σ deviations, below the SM predictions for the dilepton invariant mass squared bins $[0.045, 1.1]$ GeV^2 and $[1.1, 6]$ GeV^2, respectively [27,28]. Further, neutral current anomalies have also emerged in the observable $\Phi \equiv d\text{BR}(B_s \to \phi\mu\mu)/dm_{\mu\mu}^2$, in a similar kinematic regime ($m_{\mu\mu}^2 \in [1, 6]$ GeV^2) [25,29,30], amounting to a deviation of about 3σ. Deviations from the SM expectations have also been reported in the angular observable P_5' for the $B \to K^*\ell^+\ell^-$ decay. On the other hand, the current measured values of R_D [21,31] and R_{D^*} [19–21,31] exceed the expected SM values by about 1.4σ and 2.5σ respectively [32,33], and when combined they lead to a deviation of 3.1σ from the SM prediction [31,34,35].

These anomalies have resulted in a plethora of very interesting studies, among which two very popular scenarios which can address the above anomalies are the extensions of the SM with scalar leptoquarks and vector leptoquarks [36–97]. Vector leptoquark couplings have the potential to simultaneously explain both $R_{K^{(*)}}$ and $R_{D^{(*)}}$ [84], however a minimal $G_{\text{Pati-Salam}}$ GUT embedding of such vector leptoquark is inconsistent with very strong bounds from various flavour violating processes. Consequently, non-minimal extensions of such embedding is required to avoid such

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[1] Interestingly, these nucleon decay modes (which we refer to as $B - L$ violating throughout this work) conserve $B + L$ and in the literature also often have been referred to as $B + L$ conserving nucleon decay modes, e.g., [9]. For some recent works where such operators have been discussed in the context of some minimal SM extensions see for example [10,11].

[2] For discussion of $\Delta(B) = 2$ transitions involving trilinear di-quark interactions (forbidding leptoquark-like couplings) that can be probed in neutron-antineutron transitions and di-nucleon decays see Ref. [12].

Table 1

Transformation properties of scalar leptoquarks under the SM gauge group $G_{\text{SM}} \equiv SU(3)_c \times SU(2)_L \times U(1)_Y$ and the relevant $SU(5)$, $G_{\text{Pati-Salam}} \equiv SU(2)_L \times SU(2)_R \times SU(4)_c$, $SO(10)$ representations.

LQ	G_{SM}	$SU(5)$	$G_{\text{Pati-Salam}}$	$SO(10)$
S_1	$(\bar{\mathbf{3}}, \mathbf{1}, 1/3)$	$\mathbf{5}$, $\overline{\mathbf{45}}$, $\mathbf{50}$	$(\mathbf{1}, \mathbf{1}, \mathbf{6})$, $(\mathbf{1}, \mathbf{3}, \overline{\mathbf{10}})$	$\mathbf{10}$, $\mathbf{120}$, $\overline{\mathbf{126}}$
S_1'	$(\bar{\mathbf{3}}, \mathbf{1}, -2/3)$	$\mathbf{10}$	$(\mathbf{1}, \mathbf{3}, \mathbf{6})$	$\mathbf{120}\,\overline{\mathbf{126}}$
$\tilde{S}_1$	$(\bar{\mathbf{3}}, \mathbf{1}, 4/3)$	$\mathbf{45}$	$(\mathbf{1}, \mathbf{3}, \mathbf{6})$	$\mathbf{120}\,\overline{\mathbf{126}}$
$\tilde{R}_2$	$(\mathbf{3}, \mathbf{2}, 1/6)$	$\mathbf{10}$, $\mathbf{15}$	$(\mathbf{2}, \mathbf{2}, \mathbf{15})$	$\mathbf{120}$, $\overline{\mathbf{126}}$
R_2	$(\mathbf{3}, \mathbf{2}, 7/6)$	$\overline{\mathbf{45}}$, $\overline{\mathbf{50}}$	$(\mathbf{2}, \mathbf{2}, \mathbf{15})$	$\mathbf{120}$, $\overline{\mathbf{126}}$
S_3	$(\bar{\mathbf{3}}, \mathbf{3}, 1/3)$	$\overline{\mathbf{45}}$	$(\mathbf{3}, \mathbf{1}, \mathbf{6})$	$\mathbf{120}$, $\overline{\mathbf{126}}$

bounds [83–97]. On the other hand, single scalar leptoquark cannot simultaneously address $R_{K^{(*)}}$ and $R_{D^{(*)}}$ data without violating constraints from various flavour violating processes and consequently attempts have been made to use more than one leptoquark and to embed them in minimal GUT frameworks to such end [66,98]. In this work we will primarily focus on the scalar leptoquark states tabulated in Table 1 and the vector leptoquarks $U_1 : (\mathbf{3}, \mathbf{1}, 2/3)$ and $\tilde{V}_2 : (\bar{\mathbf{3}}, \mathbf{2}, -1/6)$.

The plan for rest of this paper is as follows. In section 2, we discuss the effective leptoquark couplings and $B - L$ violating operators with $d = 7$ and provide the relevant diagrams and estimation for the lifetime for nucleon decay modes. In section 3, we discuss possible UV completions for these interactions and the possible origins of the $B - L$ violating couplings in the context of these UV completions. In section 4, we discuss the possible connection between the B-decay anomalies and the $B - L$ violating nucleon decay rates. In section 5, we explore the implications of these $B - L$ violating interactions for baryogenesis and the possible correlation between the final baryon asymmetry and the nucleon decay lifetime. Finally, in section 5 we summarise and make our concluding remarks.

2. Effective leptoquark couplings and $B - L$ violating operators with $d = 7$

The dimension seven B–violating effective operators invariant under the SM gauge group which follow the selection rule $\Delta(B - L) = +2$ and mediate $(B - L)$ violating nucleon decays are given by [4,5]

$$
\begin{aligned}
&\mathcal{O}_1 = (Q_i Q_j)(d^c L_l)^* H_l^* \epsilon_{ij} \epsilon_{kl}, &\quad &\mathcal{O}_2 = (Q_i Q_j)(d^c L_j)^* H_i^*, \\
&\mathcal{O}_3 = (d^c d^c)^* (Q_i e^c) H_i^*, &\quad &\mathcal{O}_4 = (d^c d^c)^* (u^c L_i)^* H_j^* \epsilon_{ij} \\
&\mathcal{O}_5 = (d^c u^c)^* (d^c L_i)^* H_j^* \epsilon_{ij}, &\quad &\mathcal{O}_6 = (d^c d^c)^* (d^c L_i)^* H_i, \\
&\mathcal{O}_7 = (d^c D_\mu d^c)^* (\overline{L}_i \gamma^\mu Q_i), &\quad &\mathcal{O}_8 = (d^c D_\mu L_i)^* (\overline{d^c} \gamma^\mu Q_i), \\
&\mathcal{O}_9 = (d^c D_\mu d^c)^* (\overline{d^c} \gamma^\mu e^c) ,
\end{aligned}
\tag{1}
$$

where i, j, k, l are $SU(2)_L$ indices, and all the fermion fields are written in terms of left handed spinors. We will first show that some of these operators can arise naturally in the presence of effective trilinear couplings involving two scalar (vector) leptoquarks and the SM Higgs doublet. Interestingly, at the effective level the coupling coefficient of these trilinear interactions have mass dimension unity and is proportional to the $B - L$ breaking scale, which can in principle be decoupled from the gauge coupling unification scale in a given UV complete model. We will first treat the problem at an effective level where we will only fix the effective mass scales corresponding to the coupling coefficient of this trilinear interactions and the mass scales of the leptoquark

states to show that these operators can give rise to nucleon decay rates which can be probed at the current and future experimental nucleon decay search facilities. The relevant effective trilinear scalar interactions invariant under the SM gauge group $G_{\text{SM}} \equiv SU(3)_c \times SU(2)_L \times U(1)_Y$ are given by

$$\mathcal{L}_{\text{scalar}} \supset -\lambda_1 \tilde{R}_2^\dagger H S_1^\dagger - \tilde{\lambda}_2 \tilde{R}_2^\dagger H^\dagger S_1'^\dagger - \lambda_2 R_2^\dagger H S_1'^\dagger - \lambda_3 \tilde{R}_2^{\dagger i} (\vec{\tau}.S_3)^{\dagger \, ij} H^k + \text{h.c.}, \qquad (2)$$

where i, j, k are $SU(2)_L$ indices. For a relevant discussion of some similar relevant interactions in the context of $R-$ parity violating supersymmetric interactions see Ref. [9] and in the context of some minimal SM extensions see Refs. [10,11]. To see the $B-L$ assignments of the scalar leptoquark states and the above couplings let us write down the Yukawa interactions for the relevant scalar leptoquarks at an effective level—

$$\mathcal{L}_{\text{Y}_{S_1}} \supset -(y_{S_1}^1 Q_L^i L_L^j \epsilon^{ij} + y_{S_1}^2 \epsilon^c u_L^c d_L^c) S_1 - (y_{S_1}^3 \epsilon^c Q_L^i Q_L^j \epsilon^{ij} + y_{S_1}^4 u_L^c e_L^c + y_{S_1}^5 d_L^c \nu_L^c) S_1^\dagger \qquad (3)$$

$$\mathcal{L}_{\text{Y}_{S_1', \tilde{S}_1}} \supset -y_{S_1'}^1 u_L^c \nu_L^c S_1'^\dagger - y_{S_1'}^2 \epsilon^c d_L^c d_L^c S_1' - y_{\tilde{S}_1}^1 d_L^c e_L^c \tilde{S}_1 - y_{\tilde{S}_1}^2 \epsilon^c u_L^c u_L^c \tilde{S}_1^\dagger \qquad (4)$$

$$\mathcal{L}_{\text{Y}_{\tilde{R}_2, R_2}} \supset -y_{\tilde{R}_2}^1 d_L^c L_L^i \tilde{R}_2^{\dagger \, j} \epsilon^{ij} - y_{\tilde{R}_2}^2 \nu_L^c Q_L^i \tilde{R}_2^{\dagger \, j} \epsilon^{ij} - y_{R_2}^1 u_L^c L_L^i R_2^j \epsilon^{ij} - y_{R_2}^2 e_L^c Q_L^i R_2^{\dagger \, j} \epsilon^{ij} \qquad (5)$$

$$\mathcal{L}_{\text{Y}_{S_3}} \supset -y_{S_3}^1 Q_L^i L_L^k \epsilon^{ij} (\vec{\tau}.\vec{S}_3)^{jk} + y_{S_3}^2 \epsilon^c Q_L^i Q_L^k \epsilon^{ij} (\vec{\tau}.\vec{S}_3)^{\dagger \, jk} \qquad (6)$$

where ϵ^c corresponds to the $SU(3)_C$ tensor; i, j, k are $SU(2)_L$ indices. From Eqs. (3), (4), (5) and (6) it follows that B and L can not uniquely be defined for S_1, S_1', $\tilde{S}_1$ and S_3 in the presence of all allowed interactions; however, the quantum number $B-L$ can be uniquely defined for all the leptoquarks and $(B-L)_{S_1, S_1', S_3} = +2/3$, $(B-L)_{\tilde{S}_1} = -2/3$, $(B-L)_{\tilde{R}_2, R_2} = +4/3$. From a symmetry point of view, in the presence of these scalar leptoquarks as new physics beyond the SM this simple observation hints towards $B-L$ as a fundamental local or global symmetry, and naturally motivates UV embedding of scalar leptoquarks in frameworks like $SO(10)$ or $G_{\text{Pati-Salam}}$, which can manifest $B-L$ as a local symmetry at a high scale. Using the $B-L$ assignments for the leptoquark states it can be readily verified from Eq. (2) that each of the terms carry a $B-L$ charge -2 and consequently together with the Yukawa interaction given in Eqs. (3), (4), (5) and (6) they can give rise to $B-L$ violating proton decay operators given in Eq. (1) with $\Delta(B-L) = +2$.[3] The vector leptoquarks $\tilde{V}_2 : (\bar{3}, 2, -1/6)$ and $U_1 : (3, 1, 2/3)$ can also induce the $d = 7$ operators of Eq. (1) through the effective coupling

$$\mathcal{L}_{\text{vector}} \supset -\chi_1 \tilde{V}_2 U_1 H^\dagger + \text{h.c.}, \qquad (7)$$

where the vector leptoquarks couple to the SM fields through the interactions

$$\mathcal{L}_{U_1, \tilde{V}_2} \supset +x_1 \bar{Q}_L \gamma^\mu U_{1,\mu} L_L + x_2 \bar{d}_L^c \gamma^\mu U_{1,\mu}^\dagger e_L^c + x_3 \bar{u}_L^c \gamma^\mu U_{1,\mu}^\dagger \nu_L^c + + \tilde{x}_1 \bar{u}_L^{c\,i} \gamma^\mu \tilde{V}_{2,\mu}^j \epsilon^{ij} L_L^i$$
$$+ \tilde{x}_2 \bar{Q}_L^i \gamma^\mu \epsilon^{ij} \tilde{V}_{2,\mu}^{\dagger \, j} \nu_L^c + \tilde{x}_3 \bar{d}_L^c \gamma^\mu \tilde{V}_{2,\mu}^\dagger Q_L + \text{h.c.} \qquad (8)$$

The couplings in Eq. (8) leads to the $B-L$ assignments $(B-L) = (4/3, 2/3)$ for $(U_1, \tilde{V}_2)$. Consequently, Eq. (7) together with the Yukawa interaction given in Eq. (8)) lead to $\Delta(B-L) = 2$ nucleon decay. In Fig. 1 we show some representative diagrams inducing $d = 7$ $B-L$ violating

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[3] Note that one can also have $d = 12$ nucleon decays in the presence of trilinear couplings involving three scalar leptoquarks [99] and trilinear couplings involving a triplet Higgs scalar, a singlet and a triplet leptoquark [100].
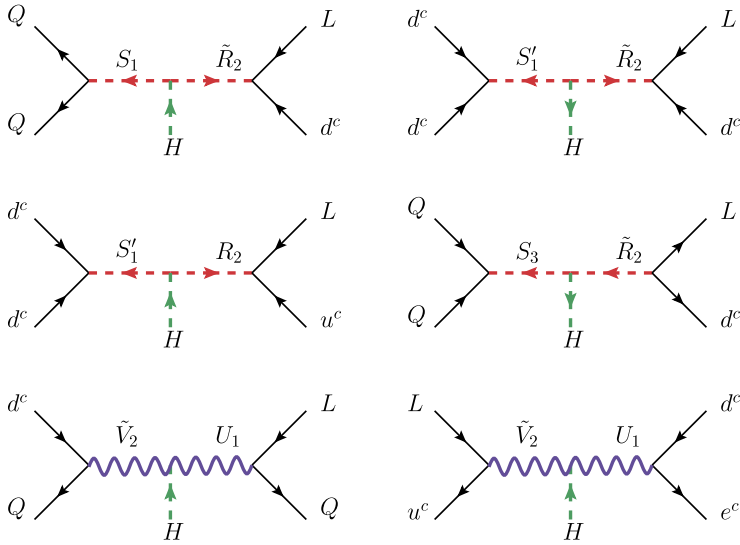
Fig. 1. Some representative diagrams realising the effective $B - L$ violating $d = 7$ operators induced by the trilinear leptoquark couplings.

nucleon decay operators in the presence of the scalar and vector leptoquark couplings discussed above. For illustration purposes we will consider the simple case of partial lifetime for the decay mode $n \to e^- \pi^+$. The lifetime for this decay mode induced by the trilinear coupling involving scalar leptoquarks $P_{1,2}$ and H, can be estimated by [101]

$$\Gamma_s(n \to e^- \pi^+) \approx \frac{|Y^*_{P_1} Y_{P_2}|^2}{64\pi} G^2 \frac{\beta_H^2 m_p}{f_\pi^2} \frac{\lambda^2 v^2}{M^4_{P_1} M^4_{P_2}}, \tag{9}$$

where λ correspond to the dimension-ful $B - L$ violating trilinear coupling of $P_{1,2}$ and H; Y_{P_1} and Y_{P_2} are the Yukawa couplings for the scalar leptoquarks P_1 and P_2. $v \equiv \langle H^0 \rangle = 174$ GeV, $\beta_H \simeq 0.012$ GeV^3 is the nucleon decay matrix element and $G \equiv (1 + D + F) \simeq 1.3$ where the D and F correspond to contributions from chiral Lagrangian factors. The decay rate induced by the effective coupling involving vector leptoquarks U_1, $\tilde{V}_2$ and H given in Eq. (10) is given by

$$\Gamma_v(n \to e^- \pi^+) \approx \frac{|x^*_{U_1} x_{\tilde{V}_2}|^2}{64\pi} G^2 A \frac{|\bar{\alpha}_H|^2 m_p}{f_\pi^2} \frac{\chi_1^2 v^2}{M^4_{T_1} M^4_{T_2}}. \tag{10}$$

Here x_{U_1} and $x_{\tilde{V}_2}$ are the couplings of U_1 and $\tilde{V}_2$ with SM fields, $A = |A_L|^2 R$ with $R = [(A_{SR} + A_{SL})(1 + |V_{ud}|^2)^2]$ for $SO(10)$ where $A_L \sim 1.25$ and $A_L A_{SL} = A_L A_{SR} \sim 3.2$ [102,103]. $\alpha_H = \bar{\alpha}_H G \simeq 0.012$ GeV^3.

3. UV completion and origin of $B - L$ violating couplings

At an effective level the computation of the decay rates given in Eqs. (9) and (10) depends on the effective mass scales corresponding to the coupling coefficient of the $B - L$ violating effective interaction involving two leptoquarks and H, and the leptoquark mass scales. In this section we will show that one can have observable $B - L$ violating nucleon decay rates for some of the

leptoquarks with mass as low as TeV scale, which is very interesting because the explanation the B-decay anomalies with scalar (vector) leptoquarks require the masses of such leptoquarks to be at around the TeV scale, and such mass scales of the leptoquarks will be probed by the direct searches at the LHC in the near future. Moreover, if the collider searches find leptoquark new physics at around TeV scale then the search for $B - L$ violating nucleon decays also provides an unique opportunity to indirectly probe $B - L$ breaking scale far beyond the reach of the current collider searches.

To discuss the realisations of the effective leptoquark couplings in the UV completion $SO(10)$, $G_{\text{Pati-Salam}} \equiv SU(2)_L \times SU(2)_R \times SU(4)_c$ and $SU(5)$, it is convenient to write down the decompositions of $SO(10)$ multiplets under the smaller subgroup representations. Under $G(2,2,4) \equiv SU(2)_L \times SU(2)_R \times SU(4)_C$ the $SO(10)$ multiplets decompose as

$$
\begin{aligned}
10 &= (2,2,1) + (1,1,6) \\
16 &= (2,1,4) + (1,2,\overline{4}) \\
120 &= (2,2,1) + (1,1,10) + (1,1,\overline{10}) + (3,1,6) + (1,3,6) + (2,2,15) \\
126 &= (1,1,6) + (3,1,10) + (1,3,\overline{10}) + (2,2,15) \\
210 &= (1,1,1) + (1,1,15) + (3,1,15) + (1,3,15) + (2,2,6) + (2,2,10) + (2,2,\overline{10}) \, .
\end{aligned}
\tag{11}
$$

Under the $G(5,1) \equiv SU(5) \times U(1)$ subgroup various multiplets decompose as follows:

$$
\begin{aligned}
10 &= 5(2) + \overline{5}(-2) \\
16 &= 1(-5) + \overline{5}(3) + 10(-1) \\
120 &= 5(2) + \overline{5}(-2) + 10(-6) + \overline{10}(6) + 45(2) + \overline{45}(-2) \\
126 &= 1(-10) + \overline{5}(-2) + 10(-6) + \overline{15}(6) + 45(2) + \overline{50}(-2) \\
210 &= (1,1,1) + (1,1,15) + (3,1,15) + (1,3,15) + (2,2,6) + (2,2,10) + (2,2,\overline{10}) \, .
\end{aligned}
\tag{12}
$$

The effective $B - L$ violating couplings can naturally arise in GUT theories like $SO(10)$ and $G_{\text{Pati-Salam}} \equiv SU(2)_L \times SU(2)_R \times SU(4)_c$ where $B - L$ is a part of the local gauge symmetry. In such theories the effective $B - L$ violating couplings given in Eqs. (2) and (7) can arise when the local $B - L$ symmetry is broken by giving a vacuum expectation value to the SM singlet field Δ^c carrying $B - L = -2$. Δ^c is present in the $\overline{126}_H$ multiplet of $SO(10)$ and corresponds to $(1,3,\overline{10})$ under $G_{\text{Pati-Salam}}$ [7,8]. Note that in SO(10) GUT Δ^c can also generate large Majorana masses for the right handed neutrinos through the couplings of the form $\nu^c \nu^c \Delta^c$.

In Table 2 we summarise the possible realisations of the relevant effective trilinear couplings given in Eq. (2) in the context of $SO(10)$, $G_{\text{Pati-Salam}}$, and $SU(5)$ UV completion. The effective trilinear coupling $\tilde{R}_2^\dagger H S_1^\dagger$ can arise from the quartic coupling $(2,2,15) \cdot (2,2,15) \cdot (1,1,6) \cdot (1,3,\overline{10})$ under $G_{\text{Pati-Salam}}$ where $\tilde{R}_2^\dagger \subset (2,2,15)$, $S_1^\dagger \subset (1,1,6)$, $H \subset (2,2,15)$ and $\Delta^c \subset (1,3,\overline{10})$. In $SO(10)$ such quartic coupling can be realised as $(126)^4$, $(126)^2(\overline{126})^2$ and $(126)^2(\overline{126}\,10)$. The effective trilinear couplings $\tilde{R}_2^\dagger H^\dagger S_1'^\dagger$ and $R_2^\dagger H S_1'^\dagger$ can arise from the $G_{\text{Pati-Salam}}$ quartic term $(2,2,15) \cdot (2,2,15) \cdot (1,3,6) \cdot (1,3,10)$, where $S_1'^\dagger \subset (1,3,6)$ and $R_2 \subset (2,2,15)$. In $SO(10)$ such quartic coupling can be realised by $(126)^2(120)^2$ and $(126\,\overline{126})(120)^2$. The quartic term $(126)^2(120)^2$ also contains the piece $(2,2,15) \cdot (2,2,15) \cdot (3,1,6) \cdot (1,3,10)$ under $G_{\text{Pati-Salam}}$ which can give the effective trilinear coupling $\tilde{R}_2 S_3 H^\dagger$ with

Table 2
Transformation properties of scalar leptoquarks under the SM gauge group $G_{\text{SM}} \equiv SU(3)_c \times SU(2)_L \times U(1)_Y$ and the relevant $SU(5)$, $G_{\text{Pati-Salam}} \equiv SU(2)_L \times SU(2)_R \times SU(4)_c$, $SO(10)$ representations.

G_{SM} effective coupling	$SO(10)$	$G_{\text{Pati-Salam}}$	$SU(5)$
$\tilde{R}_2^\dagger H S_1^\dagger$	$(126)^4$; $(126)^2(\overline{126})^2$; $(126)^2(\overline{126}\,10)$	$(2,2,15) \cdot (2,2,15) \cdot$ $(1,1,6) \cdot (1,3,\overline{10})$	$\overline{15}\,45\,45(24)$; $5\,\overline{10}\,45(24)$; $5\,5\,15(24)$; $\overline{5}\,\overline{5}\,10\,24$; $5\,\overline{15}\,45\,24$
$\tilde{R}_2^\dagger H^\dagger S'^\dagger_1$; $R_2^\dagger H S'^\dagger_1$	$(126)^2(120)^2$; $(126\,\overline{126})(120)^2$	$(2,2,15) \cdot (2,2,15) \cdot$ $(1,3,6) \cdot (1,3,10)$	$5\,10\,10(24)$; $5\,10\,15\,24$; $\overline{5}\,10\,\overline{45}(24)$; $\overline{5}\,10\,\overline{50}\,24$
$\tilde{R}_2 S_3 H^\dagger$	$(126)^2(120)^2$	$(2,2,15) \cdot (2,2,15) \cdot$ $(3,1,6) \cdot (1,3,10)$	$\overline{15}\,45\,45(24)$; $5\,\overline{10}\,45(24)$; $5\,\overline{15}\,45\,24$

$S_3 \supset (3,1,6)$. The effective Yukawa couplings given in Eqs. (3), (4), (5) and (6) can arise from the Yukawa couplings [104,105] $16_i\,16_j\,10_H$, $16_i\,16_j\,\overline{126}_H$ and $16_i\,16_j\,120_H$ in $SO(10)$ GUT theory. Note that, the Higgs multiplets which can couple to the fermion bi-linears of interest $16_i\,16_j$ are 10_H, 126_H and 120_H, with the couplings of the 10_H and 126_H being symmetric in flavour indices (i, j) and those of the 120_H being antisymmetric. In particular, the SM Higgs doublet $(1, 2, 1/2)$ is assumed to be a linear combination of h: $(1, 2, 1/2)$ and $h^\dagger$: $(1, 2, -1/2)$ which belongs to 10_H and $\overline{126}_H$ multiplets of $SO(10)$. The 16 multiplets are not associated with scalars which will lead to an incorrect $B - L$ assignment for the SM Higgs doublet. The terms relevant to the generation of the $d = 7$ operators can be obtained decomposing the relevant $SO(10)$ invariants

$$\mathcal{L}(16_i\,16_j\,10_H) = h_{ij}^{(10)}\Big[(u_i^c Q_j + v_i^c L_j)\,h - (d_i^c Q_j + e_i^c L_j)\,h^\dagger$$
$$+ \Big(\frac{\epsilon}{2}Q_i Q_j + u_i^c e_j^c - d_i^c v_j^c\Big) S_1^\dagger + \Big(\epsilon u_i^c d_j^c + Q_i L_j\Big) S_1\Big], \tag{13}$$

$$\mathcal{L}(16_i\,16_j\,120_H) \supset g_{ij}^{(120)}\Big[(d_i Q^j + e_i^c L_j)\,h^\dagger - (u_i^c Q_j + v_i^c L_j)\,h - \sqrt{2}Q_i L_j\,S_1$$
$$- \sqrt{2}(u_i^c e_j^c - d_i^c v_j^c)\,S_1^\dagger - \frac{i}{\sqrt{3}}(d_i^c Q_j - 3e_i^c L_j)\,h^{(2)\dagger}$$
$$+ \frac{i}{\sqrt{3}}(u_i^c Q_j - 3v_i^c L_j)\,h^{(2)} - 2e_i^c Q_j\,R_2^\dagger + 2v_i^c Q_j\,\tilde{R}_2^\dagger - 2d_i^c L_j\,\tilde{R}_2$$
$$+ 2u_i^c L_j\,R_2 - i\,\epsilon\,d_i^c d_j^c\,S'_1 + 2\,i\,u_i^c v_j^c\,S'^\dagger_1 + \sqrt{2}\,i\,\epsilon\,d_i^c u_j^c\,S_1^{(2)}$$
$$+ \sqrt{2}\,i\,(d_i^c v_j^c - e_i^c u_j^c)\,S_1^{(2)\dagger} - \frac{\epsilon}{\sqrt{2}}Q_i Q_j\,S_3^\dagger - \sqrt{2}\,Q_i L_j\,S_3$$
$$- 2\,i\,d_i^c e_j^c\,\tilde{S}_1^\dagger + i\,\epsilon\,u_i^c u_j^c\,\tilde{S}_1\Big], \tag{14}$$

$$\mathcal{L}(16_i\,16_j\,\overline{126}_H) \supset f_{ij}^{(126)}\Big[(u_i^c Q_j - 3v_i^c L_j)\,h - (d_i^c Q_j - 3e_i^c L_j)\,h^\dagger$$
$$+ \sqrt{3}i\,\Big(\frac{\epsilon}{2}Q_i Q_j - u_i^c e_j^c + v_i^c d_j^c\Big) S_1^\dagger + \sqrt{3}i\,(Q_i L_j - \epsilon u_i^c d_j^c)\,S_1$$
$$+ \sqrt{6}(d_i^c v_j^c + u_i^c e_j^c)\,S_1^{(2)\dagger} + 2\sqrt{3}i\,d_i^c L_j\,\tilde{R}_2 - 2\sqrt{3}i\,v_i^c Q_j\,\tilde{R}_2^\dagger$$
$$+ 2\sqrt{3}\,u_i^c v_j^c\,S'^\dagger_1 - 2\sqrt{3}i\,u_i^c L_j\,R_2 + 2\sqrt{3}i\,e_i^c Q_j\,R_2^\dagger$$
$$- 2\sqrt{3}\,d_i^c e_j^c\,\tilde{S}_1^\dagger + \sqrt{6}i\,Q_i L_j\,S_3\Big], \tag{15}$$

where parenthesised superscripts have been used to distinguish more than one fields with the same quantum numbers appearing in some decompositions. In principle, after the GUT symmetry breaking various such subfields with the same SM quantum number would mix among themselves. The $h(1, 2, 1/2)$ and $h^\dagger(1, 2, -1/2)$ fields from 10_H and $\overline{126}_H$ as well as any other Higgs multiplet with the same quantum number would mix to give a SM Higgs doublet which is a linear combination of all such fields. In $SU(5)$ GUT theory, the effective trilinear $B - L$ violating coupling e.g. $\tilde{R}_2^\dagger H S_1^\dagger$ can correspond to a number of possibilities as shown in Table 2. Considering one example of such possibility $\overline{15}\,45\,45$ where $\tilde{R}_2^\dagger \subset \overline{15}$, $S_1^\dagger \subset 45$, $H \subset 45$; one can embed such coupling as a quartic coupling $(126)^4$ in $SO(10)$, which corresponds to $\overline{15}(6) \cdot 45(2) \cdot 45(2) \cdot 1(-10)$ under the intermediate symmetry subgroup $SU(5) \times U(1)$. The vector leptoquarks $\tilde{V}_2$ and U_1 can also be embedded in $SO(10)$ framework and the covariant derivative of the 126_H multiplet would contain the quartic term $\tilde{V}_2 U_1 H^\dagger \Delta^c$ which leads to the effective term in Eq. (7). It is important to note that the relevant vector leptoquarks should be contained in the fundamental representation 45 of $SO(10)$ in order to be the fundamental gauge bosons associated with the $SO(10)$ symmetry breaking.

4. Connection to B-decay anomalies and observable $B - L$ violating nucleon decay rates

A particularly interesting solution to both $R_{K^{(*)}}$ and $R_{D^{(*)}}$ anomalies is to have R_2 and S_3 leptoquarks at around TeV scale. The relevant effective Lagrangian for the charged current anomalies $R_{D^{(*)}}$ is given by

$$\mathcal{L}_{\text{eff}}^{d \to u\ell\bar{\nu}} = -\frac{4\,G_F}{\sqrt{2}} V_{ud} \big[(1 + g_{V_L})(\bar{u}_L \gamma_\mu d_L)(\bar{\ell}_L \gamma^\mu \nu_L) + g_{S_L}(\mu)\,(\bar{u}_R d_L)(\bar{\ell}_R \nu_L) \\ + g_T(\mu)\,(\bar{u}_R \sigma_{\mu\nu} d_L)(\bar{\ell}_R \sigma^{\mu\nu} \nu_L) \big]. \tag{16}$$

Interestingly, the relevant effective Wilson coefficients can get significant contributions from R_2 to resolve $R_{D^{(*)}}$ anomalies for $m_{R_2} \sim 1$ TeV after taking into account the constraints from the measurements of $R_{D^{(*)}}^{\mu/e} = \mathcal{B}(B \to D^{(*)}\mu\bar{\nu})/\mathcal{B}(B \to D^{(*)}e\bar{\nu})$, $\mathcal{B}(B \to \tau\bar{\nu})$ and the kaon LFU ratio $R_{e/\mu}^K = \Gamma(K^- \to e^-\bar{\nu})/\Gamma(K^- \to \mu^-\bar{\nu})$. On the other hand, S_3 can contribute to $b \to s$ (semi-)leptonic transition to explain the anomalous data, but can give a negligibly small contribution to $R_{D^{(*)}}$ once the constraints from various favour violating processes are taken into account. The standard left-handed effective Hamiltonian for the $b \to s$ (semi-)leptonic transition is given by

$$\mathcal{H}_{\text{eff}}^{b \to sll} = -\frac{4G_F \lambda_t}{\sqrt{2}} \sum_{i=7,9,10} C_i(\mu)\mathcal{O}_i(\mu), \tag{17}$$

where $\lambda_t = V_{tb}V_{ts}^*$. The most relevant operators are

$$\mathcal{O}_{9(10)} = \frac{e^2}{(4\pi)^2} \big(\bar{s}_L \gamma_\mu b_L\big)\big(\bar{l}\gamma^\mu(\gamma^5)l\big). \tag{18}$$

S_3 can contribute at the tree-level inducing $\delta C_9^{\mu\mu} = -\delta C_{10}^{\mu\mu}$ [66]. For a TeV scale S_3 one can address the $R_{K^{(*)}}$ anomalies while being consistent with various flavour violating constraints such as $\mathcal{B}(B_s \to \mu\mu)$, $B \to K^{(*)}\nu\nu$, $B_s - \bar{B}_s$ mixing amplitude etc. Using various available experimental measurements of $R_{K^{(*)}}$ in different high and low q^2 bins and the available experimental measurements on the angular observables for $b \to s\mu\mu$ and $b \to see$, and the latest average for $\text{BR}(B_s \to \mu\mu)$ one finds the global fit ranges [97]

$$-0.50\,(-0.58) \geq C_{9,sb}^{\mu\mu} = -C_{10,sb}^{\mu\mu} \geq (-0.83)-0.91\,,$$
$$0.00\,(-0.12) \geq C_{9,sb}^{ee} = -C_{10,sb}^{ee} \geq (-0.45)-0.55\,, \qquad (19)$$

at the 2σ (1σ) level.[4]

The possibility of accommodating such a scenario for S_3 leptoquark in the context of $SU(5)$ GUT completion has been discussed in Ref. [106,107] and it has been explicitly pointed out that the term $10_i\,10_j\,45_H$ gives rise to the di-quark coupling for S_3 and consequently, such a coupling must be absent to avoid constraints from $d=6$ proton decay if S_3 leptoquark mass is around TeV scale. Here we will be interested in the $SO(10)$ breaking to the SM gauge group through the intermediate subgroup $G_{\text{Pati-Salam}} \equiv SU(2)_L \times SU(2)_R \times SU(4)_c$. From the $SO(10)$ GUT completion perspective, from Eq. (14) we note that the coupling $16_i\,16_j\,120_H$ would lead to $d=6$ rapid proton decay due to the presence of both di-quark and leptoquark couplings of S_3 leptoquark and therefore is inconsistent with a TeV scale S_3 explaining $R_{K^{(*)}}$ anomalies. More interestingly, we note that in the coupling $16_i\,16_j\,\overline{126}_H$ the possibility of such a rapid proton decay due to $d=6$ operator induced by S_3 is absent and furthermore, we also notice that R_2 (which can possibly explain resolve $R_{D^{(*)}}$ anomalies for $m_{R_2} \sim 1$ TeV) can be embedded in such a coupling without inducing any rapid nucleon decay modes. Therefore, we will focus on this scenario where the pair $S_3 - \tilde{R}_2$ belongs to the $\overline{126}_H$ multiplet and have masses at around TeV scale not inducing both $d=6$ and $d=7$ nucleon decays in the absence of the di-quark coupling of S_3. However, from Eq. (15) it follows that the coupling $16_i\,16_j\,\overline{126}_H$ induce S_1 mediated $d=6$ nucleon decay and consequently the experimental searches for $d=6$ nucleon decay constrain their mass scales to be relatively large $M_{S_1} > 10^{11}$ GeV for a benchmark Yukawa coupling $Y \sim 10^{-3}$. On the other hand, another submultiplet of $\overline{126}_H$, $\tilde{R}_2$ is not subject to such constrains and can have TeV scale mass. Therefore, if the interactions of the pair $S_3 - R_2$ resolving $R_{K^{(*)}}$ and $R_{D^{(*)}}$ anomalies simultaneously, are traced back to the $SO(10)$ invariant coupling $16_i\,16_j\,\overline{126}_H$, then one can have observable $B-L$ violating $d=7$ nucleon decays due to the $\overline{126}_H$ submultiplet pair $S_1 - \tilde{R}_2$.

It is important to stress that S_3 and R_2 don't appear together in the relevant interactions in Eq. (2) and it is necessary to have the pairings $S_3 - \tilde{R}_2$ and $S_1' - R_2$ instead, to have the relevant nucleon decay modes. Both of these relevant pairings however, originate from the $SO(10)$ interaction $(126)^2(120)^2$ as noted in Table 2, with one of the 126 corresponding to Δ^c carrying $B-L = -2$. Therefore, the crucial point to note here is that we need both 126_H and 120_H scalar multiplets to be present in a $SO(10)$ embedding of S_3 (essential for explaining $R_K^{(*)}$ data) and R_2 (essential for explaining $R_D^{(*)}$ data) to facilitate the relevant pairings. For example, assuming that S_3 or R_2 corresponds to 126_H: the multiplet 120_H must contain $\tilde{R}_2$ and S_1' to realise the interactions in Eq. (2).

As a benchmark point taking $M_{\tilde{R}_2} \sim 1$ TeV, $M_{S_1} \sim 10^{16}$ GeV, $y_{S_1,\tilde{R}_2} \sim 10^{-3}$ we find from Eq. (9) that the lifetime $\tau \approx 3 \times 10^{33}$ yrs for $\lambda_2 = 10^{11}$ GeV, which is in the observable range of current and future experiments. Such a choice for the mass spectrum can be motivated by an intermediate symmetry $G_{\text{Pati-Salam}} \equiv SU(2)_L \times SU(2)_R \times SU(4)_c$, which fixes the scale of $B-L$ symmetry breaking and consequently fixes the scale of λ_2. Interestingly, when the intermediate symmetry is $G_{\text{Pati-Salam}} \equiv SU(2)_L \times SU(2)_R \times SU(4)_c$ in the $SO(10)$ embedding of the scalar

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[4] Note that due to RG-running a large leptoquark coupling to third generation lepton can potentially induce a non-negligible lepton-universal contribution to $b \to s\ell\ell$ transitions via a log-enhanced anapole photon penguin contribution, as noted in Ref. [108].
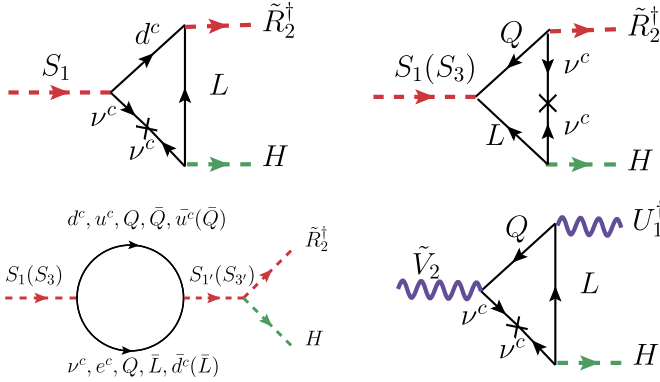
Fig. 2. One loop corrections responsible for generating $(B - L)$ asymmetry in $S_{1,3}$ and $\tilde{V}_2$ decays.

leptoquarks, one obtains gauge coupling unification at around 10^{16} GeV for the intermediate symmetry scale $M_I = M_{G_{\text{Pati-Salam}}} \approx 10^{11}$ GeV [109–116].

For vector leptoquarks we note that a TeV scale U_1 can address the anomalies in $b \to s\ell\ell$ and $b \to c\ell\bar{\nu}$ data simultaneously. However, the strong bound from charged lepton flavour violating decay processes can only be satisfied if U_1 couples to the SM fields through a mixing with vector-like counterparts of the SM fields and moreover the first generation mixing is further suppressed to account for the bounds from $K_L \to \mu e$ and $K \to \pi \mu e$ data (e.g. by some additional flavour symmetry) [85–87,97]. To account for such a suppression we parametrise $x_{U_1, \tilde{V}_2} = \beta g_U$ where $\beta \sim 10^{-3}$. Taking $M_{U_1} \sim 1$ TeV and $M_{\tilde{V}_2} \sim 10^{16}$ GeV (as it can also mediate $d = 6$ nucleon decay) we obtain from Eq. (10) the lifetime $\tau \approx 1.8 \times 10^{34}$ yrs for the effective $B - L$ violating scale $\chi_1 = 10^{11}$ GeV, which is again in the observable range of current and future experiments searching for nucleon decays.

5. Implications for baryogenesis and comments on neutrino masses

The $(B - L)$ violating effective couplings in Eqs. (2) and (7) open up a new possibility of realising high scale or resonant baryogenesis.[5] Some of the decay modes which can give rise to high scale baryogenesis are $S_1 \to \tilde{R}_2^\dagger H$, $S_1' \to \tilde{R}_2^\dagger H^\dagger$, $S_1' \to R_2^\dagger H$, with S_1 and S_1' having GUT scale masses $\sim 10^{16}$ GeV and $M_{\tilde{R}_2, R_2} \sim$ TeV. On the other hand, if $M_{S_3} > M_{\tilde{R}_2}$, with both of them having masses at around TeV scale then one can have resonant baryogenesis through the decay mode $S_3 \to \tilde{R}_2^\dagger H$ in the presence of at least two generations of S_3 with nearly degenerate masses. For vector leptoquarks one can have high scale baryogenesis through the decay mode $\tilde{V}_2 \to U_1^\dagger H$. In Fig. 2 we show the one loop diagrams which can interfere with the tree level decay modes of $S_{1,3}$ and $\tilde{V}_2$ to provide the CP violation. The computation of the final asymmetry is UV model dependent and must also take into account the potential gauge and Yukawa washout processes [118].

We will consider here the simplest case of the $(B - L)$–violating decay of the scalar S_1, which we shall assume to be very heavy with mass $\sim 10^{16}$ GeV. In a UV complete theory such as $SO(10)$ the leptoquark S_1 is in general a linear combination of fields from the decomposition

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[5] A complete list of trilinear couplings which can contribute to baryogenesis can be found for example in Ref. [117].

of multiplets 10 and $\overline{126}$ and 120 fields. Firstly, we note that the decay $S_1 \to \tilde{R}_2^\dagger H$ would violate $(B - L)$ by 2 units, and because H carries zero $(B - L)$ charge the final state has $(B - L) = -4/3$. On the other hand S_1 can also decay via the $B - L$ conserving fermionic modes $S_1 \to ff'$ with $(B - L) = 2/3$ in the final state.[6]

In the absence of more than one generation of S_1 only the top two diagrams in Fig. 2 contributes and the contribution coming from these diagrams are proportional to the Majorana masses for ν^c. Assuming the Majorana mass matrix of the ν^c fields to be diagonal and real, the asymmetry contribution coming from these two diagrams is given by

$$
\begin{aligned}
\epsilon_{B-L}^{\mathrm{v}} &= \epsilon_{B-L}^{\mathrm{I}} + \epsilon_{B-L}^{\mathrm{II}} \\
&= -\frac{f_r}{\pi} \mathrm{Im} \left[\frac{\mathrm{Tr}\{y_{S_1}^1 \, y_{\tilde{R}_2}^{2\dagger} \, M_{\nu^c} \, F_{S_1,\nu^c,\tilde{R}_2} \, y_{\nu^c LH}\} \lambda_1}{|\lambda_1|^2} \right] \\
&\quad + \frac{f_r}{\pi} \mathrm{Im} \left[\frac{\mathrm{Tr}\{y_{S_1}^5 \, y_{\tilde{R}_2}^{1\dagger} \, y_{\nu^c LH} \, M_{\nu^c} \, F_{\tilde{R}_2,\nu^c,S_1}\} \lambda_1}{|\lambda_1|^2} \right]
\end{aligned}
\tag{20}
$$

Here we define the Yukawa coupling matrix corresponding to the coupling $\nu^c LH$, as $y_{\nu^c LH}$ and M_{ν^c} is the real diagonalised mass matrix of ν^c fields. f_r corresponds to the branching ratio $\mathrm{Br}(S_1^\dagger \to \tilde{R}_2 H^\dagger)$. The function $F_{i,j,k} \equiv F(M_i, M_j, M_k)$ is defined as

$$
F(M_i, M_j, M_k) = \ln\left(1 + \frac{M_i^2}{M_j^2}\right) + \Theta\left(1 - \frac{M_j^2}{M_k^2}\right)\left(1 - \frac{M_j^2}{M_k^2}\right),
\tag{21}
$$

where Θ stands for the step function, accounting for additional ways of cutting the diagram for the case $M_{\nu^c} < M_{\tilde{R}_2}$.

In addition to the above contributions, another additional contribution can arise in any realistic UV complete model where there are more than one S_1 fields. Denoting the heavier S_1 as $S_{1'}$ the contribution coming from the bottom-left diagram of Fig. 2 is given by[7]

$$
\epsilon_{B-L}^{s} = \frac{f_r}{\pi} \mathrm{Im} \left[\frac{\mathrm{Tr}\{y_{S_{1'}}^{5\dagger} \, y_{S_1}^5 \, G(S_1, S_{1'}, \nu^c)\}(\lambda_{1'})^*(\lambda_1)}{|\lambda_1|^2} \right],
\tag{22}
$$

where the loop function $G(S_1, S_{1'}, \nu^c) \equiv G(M_{S_1}, M_{S_{1'}}, \nu^c)$ is defined as

$$
G(M_i, M_j, M_k) = \left(\frac{M_i^2 - M_k^2}{M_i^2 - M_j^2}\right) \Theta\left(1 - \frac{M_k^2}{M_i^2}\right)\left(1 - \frac{M_j^2}{M_i^2}\right).
\tag{23}
$$

For the case where S_1 and $S_{1'}$ are degenerate in mass, a resonant enhancement is possible. For a prescription of such a case see for example Ref. [119].

The branching ratio $f_r \equiv \Gamma(S_1^\dagger \to \tilde{R}_2 H^*)/(\Gamma(S_1^\dagger \to \tilde{R}_2 H^*) + \Gamma(S_1^\dagger \to ff))$ in the asymmetry can be estimated using the partial decay widths

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[6] The interplay between the $(B - L)$–conserving decays and the $(B - L)$–preserving decays of leptoquarks are crucial for inducing generating the asymmetry as first pointed out in Refs. [7,8].

[7] We assume $M_{S_1} - M_{S_{1'}} \gg \Gamma_{S_1}$, so that there is no resonant enhancement for the decay.

$$\Gamma(S_1^\dagger \to \tilde{R}_2 H^*) = \frac{|\lambda_1|^2}{8\pi M_{S_1}} \left(1 - \frac{M_{\tilde{R}_2}^2}{M_{S_1}^2} \right),$$

$$\Gamma(S_1^\dagger \to ff) = \frac{\mathrm{Tr}(y_{S_1}^{f\dagger} y_{S_1}^f)}{4\pi} M_{S_1}. \tag{24}$$

The baryon asymmetry to entropy ratio Y_B is given by

$$Y_B \equiv \frac{n_B - n_{\overline{B}}}{s} = \frac{\epsilon_{B-L}}{g_*} d, \tag{25}$$

where g_* is the total number of relativistic degrees of freedom and d in Eq. (25) is the dilution factor taking into account the partial wash out processes, which can be obtained exactly by solving the Boltzmann equation. The dilution factor can be approximately estimated by

$$d \simeq \begin{cases} 1 & (K < 1) \\ \frac{0.3}{K (\ln K)^{0.6}} & (K \gg 1), \end{cases} \tag{26}$$

where

$$K = \left. \frac{\Gamma(S_1^\dagger \to \tilde{R}_2 H^*)}{2\mathrm{H}} \right|_{T = M_{S_1}}, \tag{27}$$

with H denoting the Hubble expansion rate given by

$$\mathrm{H} = 1.66 \, g_*^{1/2} \frac{T^2}{M_{\mathrm{Pl}}}. \tag{28}$$

In Fig. 3, we show the dependence of the final baryon asymmetry on the couplings for $y_{S_1} \sim y_{\tilde{R}_2}$, for different values of M_{ν^c} and $y_{\nu^c LH}$, while choosing the benchmark values for rest of the parameters: $M_{\tilde{R}_2} = 1$ TeV, $M_{S_1} = 10^{16}$ GeV, $\lambda = 10^{11}$ GeV and $y_{S_1}^f = 10^{-3}$. A direct correlation between the final baryon asymmetry and the lifetime for nucleon decay mode can also be obtained under the simplifying assumption that the leptoquark dominantly couples to first generation SM couplings.[8]

In Fig. 4, we show the correlation between the final baryon asymmetry and the partial lifetime for the decay mode $n \to e^- \pi^+$ for different values of M_{ν^c} and $y_{\nu^c LH}$ with the same benchmark values for the rest of the parameters as in Fig. 3. It is important to note that various partial wash out processes can occur by means of various $(B - L)$ violating interactions of the right-handed neutrinos, after the ν^c fields acquire Majorana masses after $(B - L)$ breaking. Depending on the region of the parameter space, it may be desirable to have partial wash-out. On the other hand, any wash out effects can be inhibited by decoupling the ν^c fields at the same temperature as the S_1 field ($M_{\nu^c} \sim M_{S_1}$).

Furthermore, the generation of baryon asymmetry can also be directly linked to the neutrino masses in the presence of Δ^c which can induce the effective $B - L$ violating couplings of the leptoquarks and generate large Majorana masses for the right handed neutrinos through a term $\nu^c \nu^c \Delta^c$. Alternatively, neutrino masses can also be generated through one loop corrections involving the $B - L$ violating vertices and up (down) type quarks in the loop [98]. Thus depending

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[8] This assumption can always be relaxed to introduce non-trivial couplings of the leptoquark with second and third generation fermion couplings, therefore introducing a number of new parameters generalising y_{S_1} and $y_{\tilde{R}_2}$ to matrices.
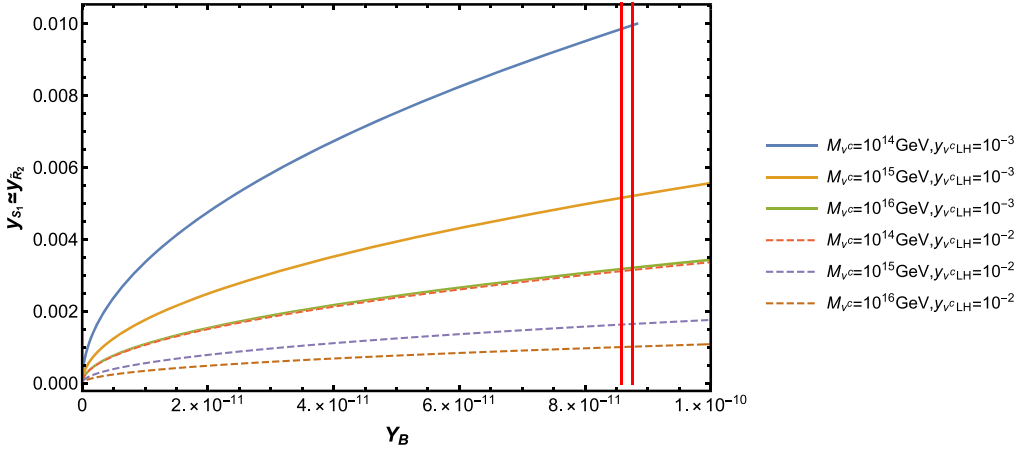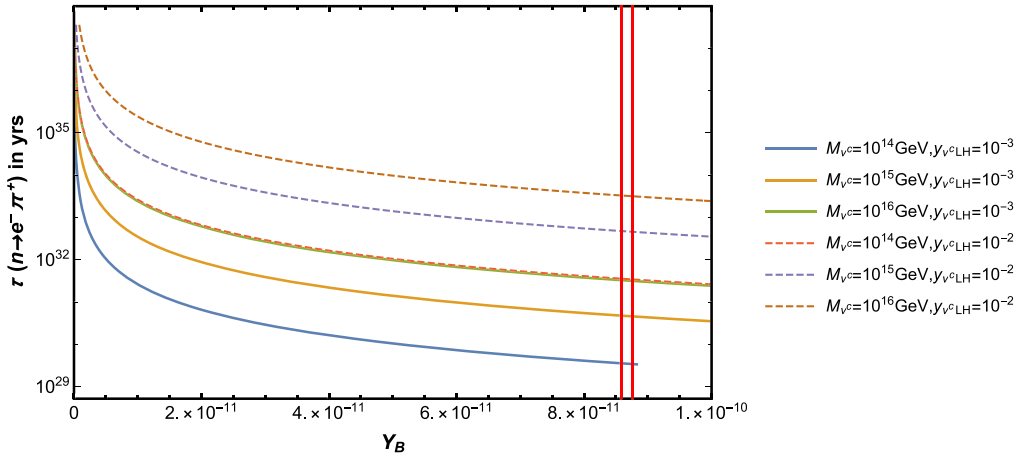
Fig. 3. The dependence of the final baryon asymmetry on the leptoquark Yukawa couplings $y_{S_1} \sim y_{\tilde{R}_2}$, for different values of M_{ν^c} and $y_{\nu^c LH}$. The vertical red lines correspond to the current best fit for the baryon asymmetry to entropy ratio provided by the PLANK collaboration. We have chosen the benchmark values $M_{\tilde{R}_2} = 1$ TeV, $M_{S_1} = 10^{16}$ GeV, $\lambda = 10^{11}$ GeV and $y_{S_1}^f = 10^{-3}$ for this plot. (For interpretation of the colours in the figure(s), the reader is referred to the web version of this article.)

Fig. 4. Correlation between the final baryon asymmetry and the partial lifetime for the decay mode $n \to e^- \pi^+$ for different values of M_{ν^c} and $y_{\nu^c LH}$ with the benchmark values $M_{\tilde{R}_2} = 1$ TeV, $M_{S_1} = 10^{16}$ GeV, $\lambda = 10^{11}$ GeV and $y_{S_1}^f = 10^{-3}$.

on the mass scales and UV completed model these two contributions will compete to generate the correct neutrino masses, nevertheless, both these possibilities have the $B - L$ violating scale embedded in them which can be probed by the $B - L$ violating nucleon decay modes.

6. Conclusion

In conclusion, we have shown that the effective $B - L$ violating couplings of scalar and vector leptoquarks can naturally induce dimension seven $B - L$ violating nucleon decay modes. The

decay lifetime of such $B - L$ violating nucleon decay modes are sensitive to the ratio of effective coupling proportional to the $B - L$ breaking scale and the mass scales of the leptoquarks. Consequently, a discovery of a $d = 7$ $B - L$ violating nucleon decay mode will certainly provide a strong evidence in the favour of leptoquark new physics and complimented by the upcoming results from the B-factories and the collider searches such observation can potentially probe the $B - L$ breaking scale as high as 10^{11} GeV, paving way for a new baryogenesis mechanism and giving interesting hints towards our understanding of neutrino masses. We have discussed some examples of possible UV completion for the relevant leptoquark interactions (in the context of GUT theories such as $SU(5)$, $G_{\text{Pati-Salam}} \equiv SU(2)_L \times SU(2)_R \times SU(4)_c$ and $SO(10)$) and have explored the possibility of having observable $B - L$ violating nucleon decay rates for some of the leptoquarks with mass as low as TeV scale. This is particularly very interesting because the explanation of the B-decay anomalies with scalar (vector) leptoquarks require the masses of such leptoquarks to be at around the TeV scale. Furthermore, if the upcoming collider searches find leptoquark new physics at around TeV scale then the search for $B - L$ violating nucleon decays will provide a unique opportunity to indirectly probe $B - L$ breaking scale far beyond the reach of the current collider searches and can have further implications for baryogenesis as we have discussed in the case of the simple example of scalar leptoquarks. We have also explored the correlation between the final baryon asymmetry and the lifetime for nucleon decay mode to this end.

Declaration of competing interest

Acknowledgements

References

[1] S. Weinberg, Phys. Rev. Lett. 43 (1979) 1566, https://doi.org/10.1103/PhysRevLett.43.1566.
[2] F. Wilczek, A. Zee, Phys. Rev. Lett. 43 (1979) 1571, https://doi.org/10.1103/PhysRevLett.43.1571.
[3] L.F. Abbott, M.B. Wise, Phys. Rev. D 22 (1980) 2208, https://doi.org/10.1103/PhysRevD.22.2208.
[4] S. Weinberg, Phys. Rev. D 22 (1980) 1694, https://doi.org/10.1103/PhysRevD.22.1694.
[5] H.A. Weldon, A. Zee, Nucl. Phys. B 173 (1980) 269, https://doi.org/10.1016/0550-3213(80)90218-7.
[6] J.C. Pati, A. Salam, U. Sarkar, Phys. Lett. B 133 (1983) 330, https://doi.org/10.1016/0370-2693(83)90157-0.
[7] K.S. Babu, R.N. Mohapatra, Phys. Rev. D 86 (2012) 035018, https://doi.org/10.1103/PhysRevD.86.035018, arXiv: 1203.5544 [hep-ph].
[8] K.S. Babu, R.N. Mohapatra, Phys. Rev. Lett. 109 (2012) 091803, https://doi.org/10.1103/PhysRevLett.109. 091803, arXiv:1207.5771 [hep-ph].
[9] F. Vissani, Phys. Rev. D 52 (1995) 4245, https://doi.org/10.1103/PhysRevD.52.4245, arXiv:hep-ph/9503227.
[10] P.H. Gu, E. Ma, U. Sarkar, Phys. Rev. D 94 (11) (2016) 111701, https://doi.org/10.1103/PhysRevD.94.111701, arXiv:1608.02118 [hep-ph].
[11] E. Ma, K. Tsumura, Phys. Rev. D 98 (3) (2018) 035037, https://doi.org/10.1103/PhysRevD.98.035037, arXiv: 1805.05056 [hep-ph].

[12] I. Baldes, N.F. Bell, R.R. Volkas, Phys. Rev. D 84 (2011) 115019, https://doi.org/10.1103/PhysRevD.84.115019, arXiv:1110.4450 [hep-ph].

[13] V.A. Kuzmin, V.A. Rubakov, M.E. Shaposhnikov, Phys. Lett. B 155 (1985) 36, https://doi.org/10.1016/0370-2693(85)91028-7.

[14] J.P. Lees, et al., BaBar Collaboration, Phys. Rev. Lett. 109 (2012) 101802, arXiv:1205.5442 [hep-ex].

[15] J.P. Lees, et al., BaBar Collaboration, Phys. Rev. D 88 (7) (2013) 072012, arXiv:1303.0571 [hep-ex].

[16] M. Huschle, et al., Belle Collaboration, Phys. Rev. D 92 (7) (2015) 072014, arXiv:1507.03233 [hep-ex].

[17] I. Adachi, et al., Belle Collaboration, arXiv:0910.4301 [hep-ex].

[18] A. Bozek, et al., Belle Collaboration, Phys. Rev. D 82 (2010) 072005, arXiv:1005.2302 [hep-ex].

[19] R. Aaij, et al., LHCb Collaboration, Phys. Rev. Lett. 115 (11) (2015) 111803; Phys. Rev. Lett. 115 (15) (2015) 159901 (Erratum), arXiv:1506.08614 [hep-ex].

[20] S. Hirose, et al., Belle Collaboration, Phys. Rev. Lett. 118 (21) (2017) 211801, arXiv:1612.00529 [hep-ex].

[21] A. Abdesselam, et al., Belle Collaboration, arXiv:1904.08794 [hep-ex].

[22] R. Aaij, et al., LHCb Collaboration, Phys. Rev. Lett. 122 (19) (2019) 191801, arXiv:1903.09252 [hep-ex].

[23] R. Aaij, et al., LHCb Collaboration, J. High Energy Phys. 1708 (2017) 055, arXiv:1705.05802 [hep-ex].

[24] A. Abdesselam, et al., Belle Collaboration, arXiv:1904.02440 [hep-ex].

[25] R. Aaij, et al., LHCb Collaboration, J. High Energy Phys. 1509 (2015) 179, arXiv:1506.08777 [hep-ex].

[26] S. Wehle, et al., Belle Collaboration, Phys. Rev. Lett. 118 (11) (2017) 111801, arXiv:1612.05014 [hep-ex].

[27] M. Bordone, G. Isidori, A. Pattori, Eur. Phys. J. C 76 (8) (2016) 440, arXiv:1605.07633 [hep-ph].

[28] B. Capdevila, A. Crivellin, S. Descotes-Genon, J. Matias, J. Virto, J. High Energy Phys. 1801 (2018) 093, arXiv:1704.05340 [hep-ph].

[29] W. Altmannshofer, D.M. Straub, Eur. Phys. J. C 75 (8) (2015) 382, arXiv:1411.3161 [hep-ph].

[30] A. Bharucha, D.M. Straub, R. Zwicky, J. High Energy Phys. 1608 (2016) 098, arXiv:1503.05534 [hep-ph].

[31] Y. Amhis, et al., HFLAV Collaboration, Eur. Phys. J. C 77 (12) (2017) 895, arXiv:1612.07233 [hep-ex], Results Spring 2019: https://hflav.web.cern.ch/, 2019.

[32] D. Bigi, P. Gambino, Phys. Rev. D 94 (9) (2016) 094008, arXiv:1606.08030 [hep-ph].

[33] D. Bigi, P. Gambino, S. Schacht, J. High Energy Phys. 1711 (2017) 061, arXiv:1707.09509 [hep-ph].

[34] Z. Ligeti, M. Papucci, D.J. Robinson, J. High Energy Phys. 1701 (2017) 083, arXiv:1610.02045 [hep-ph].

[35] A. Crivellin, J. Fuentes-Martin, A. Greljo, G. Isidori, Phys. Lett. B 766 (2017) 77, arXiv:1611.02703 [hep-ph].

[36] I. Doršner, S. Fajfer, A. Greljo, J.F. Kamenik, N. Košnik, Phys. Rep. 641 (2016) 1, https://doi.org/10.1016/j.physrep.2016.06.001, arXiv:1603.04993 [hep-ph].

[37] M. Algueró, B. Capdevila, A. Crivellin, S. Descotes-Genon, P. Masjuan, J. Matias, J. Virto, arXiv:1903.09578 [hep-ph].

[38] J. Aebischer, W. Altmannshofer, D. Guadagnoli, M. Reboud, P. Stangl, D.M. Straub, arXiv:1903.10434 [hep-ph].

[39] M. Ciuchini, A.M. Coutinho, M. Fedele, E. Franco, A. Paul, L. Silvestrini, M. Valli, arXiv:1903.09632 [hep-ph].

[40] A. Datta, J. Kumar, D. London, arXiv:1903.10086 [hep-ph].

[41] A. Arbey, T. Hurth, F. Mahmoudi, D.M. Santos, S. Neshatpour, arXiv:1904.08399 [hep-ph].

[42] R.X. Shi, L.S. Geng, B. Grinstein, S. Jäger, J. Martin Camalich, arXiv:1905.08498 [hep-ph].

[43] D. Bardhan, D. Ghosh, arXiv:1904.10432 [hep-ph].

[44] A.K. Alok, A. Dighe, S. Gangal, D. Kumar, J. High Energy Phys. 1906 (2019) 089, arXiv:1903.09617 [hep-ph].

[45] A.K. Alok, D. Kumar, J. Kumar, S. Kumbhakar, S.U. Sankar, J. High Energy Phys. 1809 (2018) 152, arXiv:1710.04127 [hep-ph].

[46] D. Ghosh, M. Nardecchia, S.A. Renner, J. High Energy Phys. 1412 (2014) 131, arXiv:1408.4097 [hep-ph].

[47] S.L. Glashow, D. Guadagnoli, K. Lane, Phys. Rev. Lett. 114 (2015) 091801, arXiv:1411.0565 [hep-ph].

[48] B. Bhattacharya, A. Datta, D. London, S. Shivashankara, Phys. Lett. B 742 (2015) 370, arXiv:1412.7164 [hep-ph].

[49] M. Freytsis, Z. Ligeti, J.T. Ruderman, Phys. Rev. D 92 (5) (2015) 054018, arXiv:1506.08896 [hep-ph].

[50] M. Ciuchini, A.M. Coutinho, M. Fedele, E. Franco, A. Paul, L. Silvestrini, M. Valli, Eur. Phys. J. C 77 (10) (2017) 688, arXiv:1704.05447 [hep-ph].

[51] G. Hiller, M. Schmaltz, Phys. Rev. D 90 (2014) 054014, arXiv:1408.1627 [hep-ph].

[52] B. Gripaios, M. Nardecchia, S.A. Renner, J. High Energy Phys. 1505 (2015) 006, arXiv:1412.1791 [hep-ph].

[53] S. Sahoo, R. Mohanta, Phys. Rev. D 91 (9) (2015) 094019, arXiv:1501.05193 [hep-ph].

[54] I. de Medeiros Varzielas, G. Hiller, J. High Energy Phys. 1506 (2015) 072, arXiv:1503.01084 [hep-ph].

[55] R. Alonso, B. Grinstein, J. Martin Camalich, J. High Energy Phys. 1510 (2015) 184, arXiv:1505.05164 [hep-ph].

[56] M. Bauer, M. Neubert, Phys. Rev. Lett. 116 (14) (2016) 141802, arXiv:1511.01900 [hep-ph].

[57] C. Hati, G. Kumar, N. Mahajan, J. High Energy Phys. 1601 (2016) 117, arXiv:1511.03290 [hep-ph].

[58] S. Fajfer, N. Kosnik, Phys. Lett. B 755 (2016) 270, arXiv:1511.06024 [hep-ph].

[59] D. Das, C. Hati, G. Kumar, N. Mahajan, Phys. Rev. D 94 (2016) 055034, arXiv:1605.06313 [hep-ph].

[60] D. Bečirević, S. Fajfer, N. Kosnik, O. Sumensari, Phys. Rev. D 94 (11) (2016) 115021, arXiv:1608.08501 [hep-ph].

[61] S. Sahoo, R. Mohanta, A.K. Giri, Phys. Rev. D 95 (3) (2017) 035027, arXiv:1609.04367 [hep-ph].

[62] P. Cox, A. Kusenko, O. Sumensari, T.T. Yanagida, J. High Energy Phys. 1703 (2017) 035, arXiv:1612.03923 [hep-ph].

[63] A. Crivellin, D. Müller, T. Ota, J. High Energy Phys. 1709 (2017) 040, arXiv:1703.09226 [hep-ph].

[64] D. Bečirević, O. Sumensari, J. High Energy Phys. 1708 (2017) 104, arXiv:1704.05835 [hep-ph].

[65] Y. Cai, J. Gargalionis, M.A. Schmidt, R.R. Volkas, J. High Energy Phys. 1710 (2017) 047, arXiv:1704.05849 [hep-ph].

[66] I. Doršner, S. Fajfer, D.A. Faroughy, N. Košnik, J. High Energy Phys. 1710 (2017) 188, arXiv:1706.07779 [hep-ph].

[67] A. Greljo, B.A. Stefanek, Phys. Lett. B 782 (2018) 131, arXiv:1802.04274 [hep-ph].

[68] S. Sahoo, R. Mohanta, J. Phys. G 45 (8) (2018) 085003, arXiv:1806.01048 [hep-ph].

[69] D. Bečirević, I. Doršner, S. Fajfer, N. Košnik, D.A. Faroughy, O. Sumensari, Phys. Rev. D 98 (5) (2018) 055003, arXiv:1806.05689 [hep-ph].

[70] C. Hati, G. Kumar, J. Orloff, A.M. Teixeira, J. High Energy Phys. 1811 (2018) 011, arXiv:1806.10146 [hep-ph].

[71] I. de Medeiros Varzielas, S.F. King, J. High Energy Phys. 1811 (2018) 100, arXiv:1807.06023 [hep-ph].

[72] J. Aebischer, A. Crivellin, C. Greub, Phys. Rev. D 99 (5) (2019) 055002, arXiv:1811.08907 [hep-ph].

[73] I. De Medeiros Varzielas, S.F. King, Phys. Rev. D 99 (9) (2019) 095029, arXiv:1902.09266 [hep-ph].

[74] H. Yan, Y.D. Yang, X.B. Yuan, arXiv:1905.01795 [hep-ph].

[75] I. Bigaran, J. Gargalionis, R.R. Volkas, arXiv:1906.01870 [hep-ph].

[76] O. Popov, M.A. Schmidt, G. White, arXiv:1905.06339 [hep-ph].

[77] N.G. Deshpande, X.G. He, Eur. Phys. J. C 77 (2) (2017) 134, arXiv:1608.04817 [hep-ph].

[78] W. Altmannshofer, P.S. Bhupal Dev, A. Soni, Phys. Rev. D 96 (9) (2017) 095010, arXiv:1704.06659 [hep-ph].

[79] D. Das, C. Hati, G. Kumar, N. Mahajan, Phys. Rev. D 96 (9) (2017) 095033, arXiv:1705.09188 [hep-ph].

[80] K. Earl, T. Gregoire, arXiv:1806.01343 [hep-ph].

[81] S. Trifinopoulos, Eur. Phys. J. C 78 (10) (2018) 803, arXiv:1807.01638 [hep-ph].

[82] S. Trifinopoulos, arXiv:1904.12940 [hep-ph].

[83] N. Assad, B. Fornal, B. Grinstein, Phys. Lett. B 777 (2018) 324, arXiv:1708.06350 [hep-ph].

[84] D. Buttazzo, A. Greljo, G. Isidori, D. Marzocca, J. High Energy Phys. 1711 (2017) 044, arXiv:1706.07808 [hep-ph].

[85] L. Calibbi, A. Crivellin, T. Li, Phys. Rev. D 98 (11) (2018) 115002, arXiv:1709.00692 [hep-ph].

[86] M. Bordone, C. Cornella, J. Fuentes-Martin, G. Isidori, Phys. Lett. B 779 (2018) 317, arXiv:1712.01368 [hep-ph].

[87] M. Blanke, A. Crivellin, Phys. Rev. Lett. 121 (1) (2018) 011801, arXiv:1801.07256 [hep-ph].

[88] M. Bordone, C. Cornella, J. Fuentes-Martín, G. Isidori, J. High Energy Phys. 1810 (2018) 148, arXiv:1805.09328 [hep-ph].

[89] J. Kumar, D. London, R. Watanabe, Phys. Rev. D 99 (1) (2019) 015007, arXiv:1806.07403 [hep-ph].

[90] A. Angelescu, D. Bečirević, D.A. Faroughy, O. Sumensari, J. High Energy Phys. 1810 (2018) 183, arXiv:1808.08179 [hep-ph].

[91] S. Balaji, R. Foot, M.A. Schmidt, Phys. Rev. D 99 (1) (2019) 015029, arXiv:1809.07562 [hep-ph].

[92] B. Fornal, S.A. Gadam, B. Grinstein, Phys. Rev. D 99 (5) (2019) 055025, arXiv:1812.01603 [hep-ph].

[93] M.J. Baker, J. Fuentes-Martín, G. Isidori, M. König, Eur. Phys. J. C 79 (4) (2019) 334, arXiv:1901.10480 [hep-ph].

[94] C. Cornella, J. Fuentes-Martin, G. Isidori, arXiv:1903.11517 [hep-ph].

[95] L. Da Rold, F. Lamagna, arXiv:1906.11666 [hep-ph].

[96] R. Barbieri, C.W. Murphy, F. Senia, Eur. Phys. J. C 77 (1) (2017) 8, arXiv:1611.04930 [hep-ph].

[97] C. Hati, J. Kriewald, J. Orloff, A.M. Teixeira, J. High Energy Phys. 1912 (2019) 006, https://doi.org/10.1007/JHEP12(2019)006, arXiv:1907.05511 [hep-ph].

[98] I. Doršner, S. Fajfer, N. Košnik, Eur. Phys. J. C 77 (6) (2017) 417, https://doi.org/10.1140/epjc/s10052-017-4987-2, arXiv:1701.08322 [hep-ph].

[99] T. Hambye, J. Heeck, Phys. Rev. Lett. 120 (17) (2018) 171801, https://doi.org/10.1103/PhysRevLett.120.171801, arXiv:1712.04871 [hep-ph].

[100] P.H. Gu, U. Sarkar, Mod. Phys. Lett. A 28 (2013) 1350159, https://doi.org/10.1142/S0217732313501599, arXiv:1110.4581 [hep-ph].

[101] P. Nath, P. Fileviez Perez, Phys. Rep. 441 (2007) 191, https://doi.org/10.1016/j.physrep.2007.02.010, arXiv:hep-ph/0601023.

[102] M.K. Parida, R.L. Awasthi, P.K. Sahu, J. High Energy Phys. 1501 (2015) 045, https://doi.org/10.1007/JHEP01(2015)045, arXiv:1401.1412 [hep-ph].

[103] S. Bertolini, L. Di Luzio, M. Malinsky, Phys. Rev. D 87 (8) (2013) 085020, https://doi.org/10.1103/PhysRevD.87.085020, arXiv:1302.3401 [hep-ph].

[104] C.S. Aulakh, A. Girdhar, Nucl. Phys. B 711 (2005) 275, https://doi.org/10.1016/j.nuclphysb.2005.01.008, arXiv:hep-ph/0405074.

[105] C.S. Aulakh, S.K. Garg, arXiv:hep-ph/0612021.

[106] I. Dorsner, S. Fajfer, N. Kosnik, Phys. Rev. D 86 (2012) 015013, https://doi.org/10.1103/PhysRevD.86.015013, arXiv:1204.0674 [hep-ph].

[107] D. Bečirević, I. Doršner, S. Fajfer, D.A. Faroughy, N. Košnik, O. Sumensari, arXiv:1806.05689 [hep-ph].

[108] A. Crivellin, C. Greub, D. Müller, F. Saturnino, Phys. Rev. Lett. 122 (1) (2019) 011805, arXiv:1807.02068 [hep-ph].

[109] K.S. Babu, R.N. Mohapatra, Phys. Rev. Lett. 70 (1993) 2845, https://doi.org/10.1103/PhysRevLett.70.2845, arXiv:hep-ph/9209215.

[110] B. Bajc, G. Senjanovic, F. Vissani, PoS HEP 2001 (2001) 198, arXiv:hep-ph/0110310.

[111] B. Bajc, G. Senjanovic, F. Vissani, Phys. Rev. Lett. 90 (2003) 051802, https://doi.org/10.1103/PhysRevLett.90.051802, arXiv:hep-ph/0210207.

[112] T. Fukuyama, N. Okada, J. High Energy Phys. 0211 (2002) 011, https://doi.org/10.1088/1126-6708/2002/11/011, arXiv:hep-ph/0205066.

[113] H.S. Goh, R.N. Mohapatra, S.P. Ng, Phys. Lett. B 570 (2003) 215, https://doi.org/10.1016/j.physletb.2003.08.011, arXiv:hep-ph/0303055.

[114] K.S. Babu, C. Macesanu, Phys. Rev. D 72 (2005) 115003, https://doi.org/10.1103/PhysRevD.72.115003, arXiv:hep-ph/0505200.

[115] S. Bertolini, T. Schwetz, M. Malinsky, Phys. Rev. D 73 (2006) 115012, https://doi.org/10.1103/PhysRevD.73.115012, arXiv:hep-ph/0605006.

[116] A.S. Joshipura, K.M. Patel, J. High Energy Phys. 1109 (2011) 137, https://doi.org/10.1007/JHEP09(2011)137, arXiv:1105.5943 [hep-ph].

[117] E. Ma, M. Raidal, U. Sarkar, Phys. Rev. D 60 (1999) 076005, https://doi.org/10.1103/PhysRevD.60.076005, arXiv:hep-ph/9811240.

[118] P.H. Gu, U. Sarkar, Phys. Lett. B 663 (2008) 80, https://doi.org/10.1016/j.physletb.2008.04.001, arXiv:0711.2727 [hep-ph].

[119] A. Pilaftsis, T.E.J. Underwood, Nucl. Phys. B 692 (2004) 303, https://doi.org/10.1016/j.nuclphysb.2004.05.029, arXiv:hep-ph/0309342.