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Baryon Number and Lepton Universality Violation in Leptoquark and Diquark Models

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We perform a systematic study of models involving leptoquarks and diquarks with masses well below the grand unification scale and demonstrate that a large class of them is excluded due to rapid proton decay. After singling out the few phenomenologically viable color triplet and sextet scenarios, we show that there exist only two leptoquark models which do not suffer from tree-level proton decay and which have the potential for explaining the recently discovered anomalies in B meson decays. Both of those models, however, contain dimension five operators contributing to proton decay and require a new symmetry forbidding them to emerge at a higher scale. This has a particularly nice realization for the model with the vector leptoquark $(3,1)_{2/3}$, which points to a specific extension of the Standard Model, namely the Pati-Salam unification model, where this leptoquark naturally arises as the new gauge boson. We explore this possibility in light of recent B physics measurements. Finally, we analyze also a vector diquark model, discussing its LHC phenomenology and showing that it has nontrivial predictions for neutron-antineutron oscillation experiments.

I. INTRODUCTION

Protons have never been observed to decay. Minimal grand unified theories (GUTs) [1, 2] predict proton decay at a rate which should have already been measured. The only four-dimensional GUTs constructed so far based on a single unifying gauge group with a stable proton require either imposing specific gauge conditions [3] or introducing new particle representations [4]. A detailed review of the subject can be found in [5]. Lack of experimental evidence for proton decay [6] imposes severe constraints on the form of new physics, especially on theories involving new bosons with masses well below the GUT scale. For phenomenologically viable models of physics beyond the Standard Model (SM) the new particle content cannot trigger fast proton decay, which seems like an obvious requirement, but is often ignored in the model building literature.

Simplified models with additional scalar leptoquarks and diquarks not triggering tree-level proton decay were discussed in detail in [7], where a complete list of color singlet, triplet and sextet scalars coupled to fermion bilinears was presented. An interesting point of that analysis is that there exists only one scalar leptoquark, namely $(3,2)_{\frac{7}{6}}$ (a color triplet electroweak doublet with hypercharge $7/6$) that does not cause tree-level proton decay. In this model dimension five operators that mediate proton decay can be forbidden by imposing an additional symmetry [8].

In this paper we collect the results of [7] and extend the analysis to vector particles. This scenario might be regarded as more appealing than the scalar case, since the new fields do not contribute to the hierarchy problem. We do not assume any global symmetries, but we do comment on how imposing a larger symmetry can remove proton decay that is introduced through nonrenormalizable operators, as in the scalar case.

Since many models for the recently discovered B meson decay anomalies [9, 10] rely on the existence of new scalar or vector leptoquarks, it is interesting to investigate which of the new particle explanations proposed in the literature do not trigger rapid proton decay. Surprisingly, the requirement of no

proton decay at tree level singles out only a few models, two of which involve the vector leptoquarks $(3,1)_{\frac{2}{3}}$ and $(3,3)_{\frac{2}{3}}$, respectively. Remarkably, these very same representations have been singled out as giving better fits to B meson decay anomalies data [11]. An interesting question we consider is whether there exists a UV complete extension of the SM containing such leptoquarks in its particle spectrum.

Finally, although the phenomenology of leptoquarks has been analyzed in great detail, there still remains a gap in the discussion of diquarks. In particular, neutron-antineutron $(n-\bar{n})$ oscillations have not been considered in the context of vector diquark models. We fill this gap by deriving an estimate for the $n-\bar{n}$ oscillation rate in a simple vector diquark model and discuss its implications for present and future experiments.

The paper is organized as follows. In Sec. II we study the order at which proton decay first appears in models including new color triplet and sextet representations and briefly comment on their experimental status. In Sec. III we focus on the unique vector leptoquark model which does not suffer from tree-level proton decay, can account for the anomalies in B meson decays, and has an appealing UV completion. In Sec. IV we analyze a model with a single vector color sextet, discussing its LHC phenomenology and implications for $n-\bar{n}$ oscillations. Section V contains conclusions.

Field	$SU(3)_c \times SU(2)_L \times U(1)_Y$ reps.
Scalar leptoquark	$(3,2)'_{\frac{7}{6}}$
Scalar diquark	$(3,1)_{\frac{2}{3}},\ (6,1)_{-\frac{2}{3}},\ (6,1)_{\frac{1}{3}},\ (6,1)_{\frac{4}{3}},\ (6,3)_{\frac{1}{3}}$
Vector leptoquark	$(3,1)'_{\frac{2}{3}},\ (3,1)_{\frac{5}{3}},\ (3,3)'_{\frac{2}{3}}$
Vector diquark	$(6,2)_{-\frac{1}{6}},\ (6,2)_{\frac{5}{6}}$

TABLE I. The only leptoquark and diquark models with a triplet or sextet color structure that do not suffer from tree-level proton decay. The primes indicate the existence of dim 5 proton decay channels.

II. VIABLE LEPTOQUARK AND DIQUARK MODELS

For clarity, we first summarize the combined results of [7] and this work in Table I, which shows the only color triplet and color sextet models that do not exhibit tree-level proton decay. The scalar case was investigated in [7], whereas in this paper we concentrate on vector particles. As explained below, the representations denoted by primes exhibit proton decay through dimension five operators (see also [8]). We note that although the renormalizable proton decay channels involving leptoquarks are well-known in the literature, to our knowledge the nonrenormalizable channels have not been considered anywhere apart from the scalar case in [8].

A. Proton decay in vector models

We first enumerate all possible dimension four interactions of the new vector color triplets and sextets with fermion bilinears respecting gauge and Lorentz invariance. A complete set of those operators is listed in Table II [12]. For the vector case there are two sources of proton decay. The first one comes from tree-level diagrams involving a vector color triplet exchange, as shown in Fig. 1. This excludes the representations $(3, 2)_{\frac{1}{6}}$ and $(3, 2)_{-\frac{5}{6}}$, since they would require unnaturally small couplings to SM fermions or very large masses to remain consistent with proton decay limits. The second source comes from dimension five operators involving the vector leptoquark representations $(3, 1)_{\frac{2}{3}}$ and $(3, 3)_{\frac{2}{3}}$:

$$\frac{1}{\Lambda} (\overline{Q}_L^c H^\dagger)\gamma^\mu d_R V_\mu , \quad \frac{1}{\Lambda} (\overline{Q}_L^c \tau^A H^\dagger)\gamma^\mu d_R V_\mu^A , \quad (1)$$

respectively. Those operators can be constructed if no additional global symmetry forbidding them is imposed and allow for the proton decay channel shown in Fig. 2, resulting in a lepton (rather than an antilepton) in the final state. The corresponding proton lifetime estimate is:

$$\tau_p \approx \left(2.5 \times 10^{32} \text{ years}\right) \left(\frac{M}{10^4 \text{ TeV}}\right)^4 \left(\frac{\Lambda}{M_{\text{Pl}}}\right)^2 , \quad (2)$$

where the leptoquark tree-level coupling and the coefficient of the dimension five operator were both set to unity. The numerical factor in front of Eq. (2) is the current limit on the proton lifetime from the search for $p \to K^+\pi^+e^-$ [13]. Even in the most optimistic scenario of the largest suppression of proton decay, i.e., when the new physics behind the dimension five operator does not appear below the Planck scale, those operators are still problematic for $M \lesssim 10^4$ TeV, which well includes the region of interest for the B meson decay anomalies.

The dimension five operators can be removed by embedding the vector leptoquarks into UV complete models. As argued in [8] for the scalar case, it is sufficient to impose a discrete subgroup $\mathcal{Z}_3$ of a global $\text{U}(1)_{B-L}$ to forbid the problematic dimension five operators. They are also naturally ab-

Operator	SU(3)$_c$	SU(2)$_L$	U(1)$_Y$	p decay
$\overline{Q}_L^c\gamma^\mu u_R V_\mu$	3	2	$-5/6$	tree-level
	$\bar{6}$	2	$-5/6$	–
$\overline{Q}_L^c\gamma^\mu d_R V_\mu$	3	2	$1/6$	tree-level
	$\bar{6}$	2	$1/6$	–
$\overline{Q}_L\gamma^\mu L_L V_\mu$	3	$1, 3$	$2/3$	dim 5
$\overline{Q}_L^c\gamma^\mu e_R V_\mu^*$	3	2	$-5/6$	tree-level
$\overline{L}_L^c\gamma^\mu u_R V_\mu^*$	3	2	$1/6$	tree-level
$\overline{L}_L^c\gamma^\mu d_R V_\mu^*$	3	2	$-5/6$	tree-level
$\overline{u}_R\gamma^\mu e_R V_\mu$	3	1	$5/3$	dim 7
$\bar{d}_R\gamma^\mu e_R V_\mu$	3	1	$2/3$	dim 5

TABLE II. Possible vector color triplet and sextet representations.
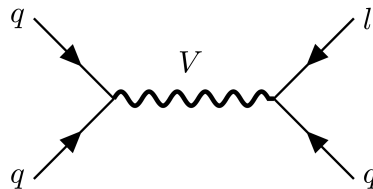

FIG. 1. Tree-level vector exchange triggering proton decay for $V = (3, 2)_{\frac{1}{6}}$ and $V = (3, 2)_{-\frac{5}{6}}$.
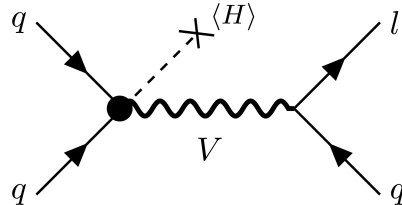

FIG. 2. Proton decay through a dimension five interaction for $V = (3, 1)_{\frac{2}{3}}$ and $V = (3, 3)_{\frac{2}{3}}$.

sent in models with gauged $\text{U}(1)_{B-L}$.[1]

Ultimately, as shown in Table I, there are only five color triplet or sextet vector representations that are free from tree-level proton decay, two of which produce dimension five proton decay operators. In the scalar case, as shown in [7], there are six possible representations with only one suffering from dimension five proton decay.

---

[1] We note that the dimension five operators (1) provide a baryon number violating channel which may be used to generate a cosmological baryon number asymmetry.

B. Leptoquark phenomenology

The phenomenology of scalar and vector leptoquarks has been extensively discussed in the literature [14–16] and we do not attempt to provide a complete list of all relevant papers here. For an excellent review and many references see [12], which is focused primarily on light leptoquarks.

Low-scale leptoquarks have recently become a very active area of research due to their potential for explaining the experimental hints of new physics in B meson decays, in particular $B^+ \to K^+\ell^+\ell^-$ and $B^0 \to K^{*0}\ell^+\ell^-$, for which a deficit in the ratios $R_{K^{(*)}} = \mathrm{Br}(B \to K^{(*)}\mu^+\mu^-)/\mathrm{Br}(B \to K^{(*)}e^+e^-)$ with respect to the SM expectations has been reported [9, 10]. A detailed analysis of the anomalies can be found in [17–22]. Several leptoquark models have been proposed to alleviate this tension and are favored by a global fit to $R_{K^{(*)}}$, $R_{D^{(*)}}$ and other flavor observables. Surprisingly, not all of those models remain free from tree-level proton decay.

The leptoquark models providing the best fit to data with just a single new representation are: scalar $(3, 2)_{\frac{1}{6}}$ [23], vector $(3, 1)_{\frac{2}{3}}$ [11, 24] and vector $(3, 3)_{\frac{2}{3}}$ [25]. Among those, only the models with the vector leptoquarks $(3, 1)_{\frac{2}{3}}$ and $(3, 3)_{\frac{2}{3}}$ are naturally free from any tree-level proton decay, since for the scalar leptoquark $(3, 2)_{\frac{1}{6}}$ there exists a dangerous quartic coupling involving three leptoquarks and the SM Higgs [7] which triggers tree-level proton decay.

Interestingly, as indicated in Table II, both vector models $(3, 1)_{\frac{2}{3}}$ and $(3, 3)_{\frac{2}{3}}$ suffer from dimension five proton decay and require imposing an additional symmetry to eliminate it. An elegant way to do it would be to extend the SM symmetry by a gauged $\mathrm{U}(1)_{B-L}$. Actually, such an extended symmetry would eliminate also the tree-level proton decay in the model with the scalar leptoquark $(3, 2)_{\frac{1}{6}}$. However, as we will see in Sec. III, only in the case of the:

▶ vector leptoquark $(3, 1)_{\frac{2}{3}}$

there exists a very appealing SM extension which intrinsically contains such a state in its spectrum, simultaneously forbidding dimension five proton decay.

C. Diquark phenomenology

The literature on the phenomenology of diquarks is much more limited. It focuses on scalar diquarks [26] and predominantly looks at three aspects: LHC discovery reach for scalar diquarks [27–40], $n - \bar{n}$ oscillations mediated by scalar diquarks [7, 34, 41–46] and baryogenesis [7, 34, 42, 44–46]. Studies of vector diquarks investigate only their LHC phenomenology [35, 47–51], concentrating on their interactions with quarks.

In Sec. IV we close the gap in diquark phenomenology by discussing the implications of a vector diquark model for $n-\bar{n}$ oscillation experiments.

III. VECTOR LEPTOQUARK MODEL

As emphasized in Sec. II B, the SM extended by just the vector leptoquark $(3, 1)_{\frac{2}{3}}$ alone can explain the recently discovered anomalies in B meson decays, in agreement with all other experimental data [11, 24]. This model is unique, since apart from being free from tree-level proton decay, it has a very simple and attractive UV completion, which automatically forbids the dimension five proton decay operators.

A priori, any of the leptoquarks can originate from an extra GUT irrep, either scalar or vector. In particular, the vector leptoquark $(3, 1)_{\frac{2}{3}}$ could be a component of the vector 40 irrep of SU(5). Nevertheless, this generic explanation does not seem to be strongly motivated or predictive. The only other suggestion in the literature regarding the origin of the vector $(3, 1)_{\frac{2}{3}}$ state arises in composite models [52–54].

The third, perhaps most desirable option, that the vector leptoquark $(3, 1)_{\frac{2}{3}}$ is the gauge boson of a unified theory. Indeed, this scenario is realized if one considers partial unification based on the Pati-Salam gauge group:

▶ $\mathrm{SU}(4) \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$

at higher energies [55].[2] In this case the vector leptoquark $(3, 1)_{\frac{2}{3}}$ emerges naturally as the new gauge boson of the broken symmetry, and is completely independent of the symmetry breaking pattern. It is interesting that the Pati-Salam partial unification model can be fully unified into an SO(10) GUT.

The fermion irreps of the Pati-Salam model, along with their decomposition into SM fields, are

$$
\begin{aligned}
(4, 2, 1) &= (3, 2)_{\frac{1}{6}} \oplus (1, 2)_{-\frac{1}{2}} \\
(\bar{4}, 1, 2) &= (\bar{3}, 1)_{\frac{1}{3}} \oplus (\bar{3}, 1)_{-\frac{2}{3}} \oplus (1, 1)_1 \oplus (1, 1)_0 .
\end{aligned}
\tag{3}
$$

Interestingly, the theory is free from tree-level proton decay via gauge interactions. The explanation for this is straightforward. Since $\mathrm{SU}(4) \supset \mathrm{SU}(3)_c \times \mathrm{U}(1)_{B-L}$, this implies that $B-L$ is conserved. However, after the Pati-Salam group breaks down to the SM, the interactions of the leptoquark $(3, 1)_{\frac{2}{3}}$ with quarks and leptons have an accidental $B + L$ global symmetry. Those two symmetries combined result in both baryon and lepton number being conserved in gauge interactions, thus no proton decay can occur via a tree-level exchange of $(3, 1)_{\frac{2}{3}}$.

In addition, there are no gauge-invariant dimension five proton decay operators in the Pati-Salam model involving the vector leptoquark $(3, 1)_{\frac{2}{3}}$. This was actually expected from the fact that $\mathrm{SU}(4) \supset \mathrm{SU}(3)_c \times \mathrm{U}(1)_{B-L}$ and, as discussed in Sec. II B, a $\mathrm{U}(1)_{B-L}$ symmetry is sufficient to forbid such operators.

————

[2] Similar unification models have recently been constructed based on the gauge group $\mathrm{SU}(4) \times \mathrm{SU}(2)_L \times \mathrm{U}(1)$ [56, 57], in which the new gauge bosons also coincide with the vector leptoquark $(3, 1)_{2/3}$.

The quark and lepton mass eigenstates are related to the gauge eigenstates through $n_f \times n_f$ unitary matrices, with $n_f = 3$ the number of families of quarks and leptons. Expressing the interactions that couple the $(3, 1)_{\frac{2}{3}}$ vector leptoquark to the quark and the lepton in each irrep of Eqs. (3) in terms of mass eigenstates, one must include unitary matrices, similar to the Cabibbo-Kobayashi-Maskawa (CKM) matrix for the quarks in the SM, that measure the misalignment of the lepton and quark mass eigenstates:

$$\mathcal{L} \supset \frac{g_4}{\sqrt{2}} V_\mu \Big[L_{ij}^u \left(\bar{u}^i \gamma^\mu P_L \nu^j \right) + L_{ij}^d \left(\bar{d}^i \gamma^\mu P_L e^j \right) + R_{ij} \left(\bar{d}^i \gamma^\mu P_R e^j \right) \Big] + \text{h.c.} \, . \quad (4)$$

The SU(4) gauge coupling constant, g_4, is not an independent paramter but fixed by the QCD coupling constant at the scale M of the masses of the vector bosons of SU(4); to leading order $g_4(M) = \sqrt{4\pi\alpha_s(M)} \approx 1.03$ at $M = 2$ TeV. The unitary matrices L^u and L^d, and CKM matrix V and the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix U satisfy $L^u = V L^d U$.

In the SM flavor-changing neutral currents with $\Delta B = -\Delta S = 1$ are described by the effective Lagrangian [58, 59]:

$$\mathcal{L} = -\frac{4 G_F}{\sqrt{2}} V_{tb} V_{ts}^* \left(\sum_{k=7}^{10} C_k \mathcal{O}_k + \sum_{i,j} C_\nu^{ij} \mathcal{O}_\nu^{ij} + ... \right), \quad (5)$$

where the ellipsis denote four-quark operators, $\mathcal{O}_7$ and $\mathcal{O}_8$ are electro- and chromo-magnetic-moment-transition operators, and $\mathcal{O}_9$, $\mathcal{O}_{10}$ and $\mathcal{O}_\nu$ are semi-leptonic operators involving either charged leptons or neutrinos:[3]

$$\mathcal{O}_{9(10)} = \frac{e^2}{(4\pi)^2} \left[\bar{s} \gamma_\mu P_L b \right] \left[\bar{\mu} \gamma^\mu (\gamma_5) \mu \right] , \quad (6)$$

$$\mathcal{O}_\nu^{ij} = \frac{2 e^2}{(4\pi)^2} \left[\bar{s} \gamma^\mu P_L b \right] \left[\bar{\nu}^i \gamma^\mu P_L \nu^j \right] . \quad (7)$$

Chirally-flipped $(b_{L(R)} \to b_{R(L)})$ versions of all these operators are denoted by primes and are negligible in the SM. New physics (NP) can generate modifications to the Wilson coefficients of the above operators, and, moreover, it can generate additional terms in the effective Lagrangian,

$$\Delta\mathcal{L} = -\frac{4 G_F}{\sqrt{2}} V_{tb} V_{ts}^* \left(C_S \mathcal{O}_S + C_P \mathcal{O}_P + C_S' \mathcal{O}_S' + C_P' \mathcal{O}_P' \right) \quad (8)$$

in the form of scalar operators:

$$\mathcal{O}_S^{(\prime)} = \frac{e^2}{(4\pi)^2} \left[\bar{s} P_{R(L)} b \right] \left[\bar{\mu}\mu \right] , \quad (9)$$

$$\mathcal{O}_P^{(\prime)} = \frac{e^2}{(4\pi)^2} \left[\bar{s} P_{R(L)} b \right] \left[\bar{\mu}\gamma_5\mu \right] . \quad (10)$$

Tensor operators cannot arise from short distance NP with the SM linearly realized and, moreover, under this assumption $C_P = -C_S$ and $C_P' = C_S'$ [60].

Exchange of the $(3, 1)_{\frac{2}{3}}$ vector leptoquark gives tree level contributions to the Wilson coefficients at its mass scale, M:

$$\Delta C_9(M) = -\Delta C_{10}(M) = -\frac{2\pi^2}{\sqrt{2}\, G_F M^2} \frac{g_4^2}{e^2} \frac{L_{b\mu}^{d*} L_{s\mu}^d}{V_{tb} V_{ts}^*}, \quad (11)$$

$$\Delta C_9'(M) = \Delta C_{10}'(M) = -\frac{2\pi^2}{\sqrt{2}\, G_F M^2} \frac{g_4^2}{e^2} \frac{R_{b\mu}^* R_{s\mu}}{V_{tb} V_{ts}^*} , \quad (12)$$

$$\Delta C_S(M) = -\frac{4\pi^2}{\sqrt{2}\, G_F M^2} \frac{g_4^2}{e^2} \frac{L_{b\mu}^{d*} R_{s\mu}}{V_{tb} V_{ts}^*} , \quad (13)$$

$$\Delta C_S'(M) = -\frac{4\pi^2}{\sqrt{2}\, G_F M^2} \frac{g_4^2}{e^2} \frac{R_{b\mu}^* L_{s\mu}^d}{V_{tb} V_{ts}^*} , \quad (14)$$

$$\Delta C_\nu^{ij}(M) = 0 . \quad (15)$$

Experimental bounds on $R_{K^{(*)}\nu} = \text{Br}(B \to K^{(*)}\nu\bar{\nu}) / \text{Br}(B \to K^{(*)}\nu\bar{\nu})^{\text{SM}} = \frac{1}{3} \sum_{ij} |\delta^{ij} + C_\nu^{ij}/C_\nu^{\text{SM}}|^2$, where $C_\nu^{\text{SM}} \simeq -6.35$ [61], severely constrain models of B anomalies. As seen above, the $(3, 1)_{\frac{2}{3}}$ vector leptoquark evades this constraint by giving no correction at all to C_ν; the result holds generally for this type of NP mediator at tree level [11, 62]. It has been pointed out that generally the condition $\Delta C_\nu(M) = 0$ is not preserved by renormalization group running of the Wilson coefficients [63]. Because of the flavor structure of the interaction in Eq. (4) there are no "penguin" or wave function renormalization contributions to the running of ΔC_ν down to the electroweak scale. The only contribution comes from the renormalization by exchange of SU(2) gauge bosons that mixes the singlet operator $(\bar{q}\gamma^\mu P_L e)(\bar{e}\gamma_\mu P_L q)$ into the triplet, $(\bar{q}\tau^a\gamma^\mu P_L e)(\bar{e}\tau^a\gamma_\mu P_L q)$, resulting in

$$\Delta C_\nu^{ij}(M_W) = -\frac{3}{4\sqrt{2}\, G_F M^2} \frac{g_4^2}{\sin^2\theta_w} \ln\left(\frac{M}{M_W}\right) \frac{S_{sj} S_{bi}^*}{V_{tb} V_{ts}^*} , \quad (16)$$

where $S = V^\dagger L^u = L^d U$. The vector contribution to the rate does not interfere with the SM, which implies $R_{K^{(*)}\nu} - 1 = \frac{1}{3} \sum_{ij} |C_\nu^{ij}/C_\nu^{\text{SM}}|^2$; using $R_{K^{(*)}\nu} < 4.3$ [64], we obtain the condition $M > 0.8$ TeV. Since $\ln(M/M_W)$ is not large for $M \approx 2$ TeV, the leading log term is subject to sizable ($\approx 100\%$) corrections. However, a complete one-loop calculation is beyond the scope of this work.

We pause to comment on the remarkable cancellation of the interference term between the SM and NP contributions to the rate for $B \to K^{(*)}\nu\bar{\nu}$ and the absence of a sum over generations in the pure NP contribution to the rate. These observations hold generally for any vector leptoquark model that couples universally to quark and lepton generations. This can be easily seen by not rotating to the neutrino mass eigenstate basis, a good approximation for the nearly massless neutrinos. Vector leptoquark exchange leads to an effective interaction with flavor structure $(\bar{s}_L \gamma^\mu \nu_L^2)(\bar{\nu}_L^3 \gamma^\mu b_L)$, while the SM always involves a sum over same neutrino flavors $\sim \sum_j \bar{\nu}_L^j \gamma^\mu \nu_L^j$. There are never common final states to the

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[3] We have assumed the B anomalies in $R_{K^{(*)}}$ arise from a suppression in the μ channel relative to the SM. The Pati-Salam $(3, 1)_{2/3}$ vector leptoquark model can equally well accommodate an enhancement in the electron channel, but a suppression in the muon channel is preferred by global fits to data that include angular moments in $B \to K^*\mu\mu$ and various branching fractions in addition to $R_{K^{(*)}}$.

SM and the NP mediated interactions and therefore no interference. Moreover, there is a single flavor configuration in the final state of the NP mediated interaction ($\bar{\nu}^3 \nu^2$) while there are three configurations in the SM case ($\bar{\nu}^j \nu^j$, $j = 1, 2, 3$).

Bounds on $\Delta C'_{9,10}$ and $\Delta C_S^{(\prime)}$ can be accommodated by adjusting $R_{s\mu}$ and $R_{b\mu}$. B-anomalies are best fit by $\Delta C_9 = -\Delta C_{10} \approx -0.6$ which requires $(g_4^2/M^2) L_{b\mu}^{d*} L_{s\mu}^d \approx 1.8 \times 10^{-3}$ TeV^{-2}, or $L_{b\mu}^{d*} L_{s\mu}^d \approx 7.2 \times 10^{-3}$ for a leptoquark mass of $M = 2$ TeV.

It has been suggested that the vector leptoquark may also account for the anomaly in semileptonic decays to τ-leptons [11, 65]. Defining, as is customary, $R_{D^{(*)}} = \mathrm{Br}(B \to D^{(*)}\tau\nu)/\mathrm{Br}(B \to D^{(*)}\ell\nu)$, the SM predicts [66] (see also [67–70]) $R_D = 0.299(3)$ and $R_{D^*} = 0.257(3)$. These branching fractions have been measured by Belle [71–73], BaBar [74, 75] and LHCb [76], and the average gives [77] $R_D = 0.403(47)$ and $R_{D^*} = 0.310(17)$. The effect of leptoquarks on B semileptonic decays to τ is described by the following terms of the effective Lagrangian for charged current interactions [78, 79]:

$$\mathcal{L} \supset -\frac{4G_F}{\sqrt{2}} V_{cb} \Big[(U_{\tau j} + \epsilon_L^j)(\bar{c}\gamma^\mu P_L b)(\bar{\tau}\gamma_\mu P_L \nu^j)$$
$$+ \epsilon_{s_R}^j (\bar{c} P_R b)(\bar{\tau} P_L \nu^j) \Big] + \mathrm{h.c.} \quad (17)$$

with

$$\epsilon_L^j = \frac{g_4^2}{4\sqrt{2}\,G_F M^2} \frac{L_{cj}^u L_{b\tau}^{d*}}{V_{cb}} \,, \quad \epsilon_{s_R}^j = -\frac{g_4^2}{2\sqrt{2}\,G_F M^2} \frac{L_{cj}^u R_{b\tau}^*}{V_{cb}}. \quad (18)$$

The B_c lifetime [80] and $B_c \to \tau\nu$ branching fraction [81] impose severe constraints on ϵ_{s_R}; these are accommodated by abating $R_{b\tau}$. Hence,

$$\frac{R_{D^{(*)}}}{R_{D^{(*)}}^{\mathrm{SM}}} \approx 1 + 2\,\mathrm{Re}\Big[\sum_j \epsilon\, L_{cj}^u L_{b\tau}^{d*} \Big] \approx 1 + 2\,\mathrm{Re}\big[\epsilon\, (V L^d)_{c\tau} L_{b\tau}^{d*} \big]$$
$$\approx 1 + 2\,\mathrm{Re}\big[\epsilon\, L_{s\tau}^d L_{b\tau}^{d*} \big] \leq 1 + \epsilon \,, \quad (19)$$

where,

$$\epsilon = \frac{g_4^2}{4\sqrt{2}\,G_F V_{cb} M^2} \approx 0.1 \left(\frac{2\,\mathrm{TeV}}{M} \right)^2 . \quad (20)$$

IV. VECTOR DIQUARK MODEL

In this section we discuss the properties of a model with just one additional representation – the vector color sextet:

$$V_\mu = \begin{pmatrix} V_u \\ V_d \end{pmatrix}_\mu^{\alpha\beta} = (6, 2)_{-\frac{1}{6}} \,, \quad (21)$$

which is obviously free from proton decay. Although in the SM all fundamental vector particles are gauge bosons, we can still imagine that such a vector diquark arises from a vector
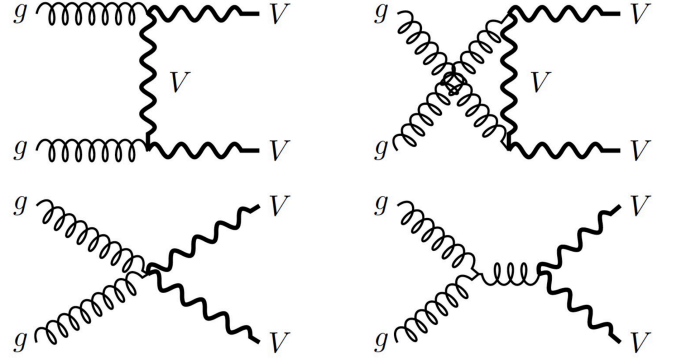

FIG. 3. Processes contributing to pair production of vector diquarks.

GUT representation, for instance from a vector 40 irrep of SU(5) [82]. The Lagrangian for the model is given by:

$$\mathcal{L}_V = -\tfrac{1}{4}(D_{[\mu} V_{\nu]})^\dagger D^{[\mu} V^{\nu]} + M^2 V_\mu^\dagger V^\mu$$
$$- \Big[\lambda_{ij}\, (\overline{Q}_L^c)_\alpha^i \gamma^\mu (d_R)_\beta^j (V^\dagger)_\mu^{\alpha\beta} + \mathrm{h.c.} \Big] , \quad (22)$$

where $\alpha, \beta = 1, 2, 3$ are SU(3)$_c$ indices, $i, j = 1, 2, 3$ are family indices and there is an implicit contraction of the SU(2)$_L$ indices. We assume that the mass term arises from a consistent UV completion.

Among the allowed higher dimensional operators, $n - \bar{n}$ oscillations, as we discuss in Sec. IV B, are mediated by the dimension five terms:

$$\mathcal{O}_1 = \frac{c_1}{\Lambda} V_\mu^{\alpha\alpha'} \epsilon\, V_\nu^{\beta\beta'} (\bar{u}_R^c)^\delta \sigma^{\mu\nu} d_R^{\delta'} \epsilon_{\alpha\beta\delta}\, \epsilon_{\alpha'\beta'\delta'} \,,$$
$$\mathcal{O}_2 = \frac{c_2}{\Lambda} \big[\partial_\mu (V^\mu)^{\alpha\alpha'} \epsilon\, V_\nu^{\beta\beta'} \big] \big[(V^\nu)^{\delta\delta'} \epsilon\, H \big] \epsilon_{\alpha\beta\delta}\, \epsilon_{\alpha'\beta'\delta'} \,. \quad (23)$$

A. LHC phenomenology

Several studies of constraints and prospects for discovering vector diquarks at the LHC can be found in the literature. Most of the analysis have focused on the case of a sizable diquark coupling to quarks [47–49], although an LHC four-jet search that is essentially independent of the strength of the diquark coupling to quarks has also been considered [35].

There are severe limits on vector diquark masses arising from LHC searches for non-SM dijet signals [83, 84]. For a coupling $\lambda_{ij} \approx 1$ ($i, j = 1, 2$) those searches result in a bound on the vector diquark mass

$$M_{\lambda \approx 1} \gtrsim 8\ \mathrm{TeV} \,. \quad (24)$$

Lowering the value of the coupling to $\lambda_{ij} \approx 0.01$ completely removes the LHC constraints from dijet searches and, at the same time, does not affect the strength of the four-jet signal arising from gluon fusion (see Fig. 3). Using the results of the analysis of four-jet events at the LHC presented in [35], the currently collected 37 fb^{-1} of data by the ATLAS experiment [84] with no evident excess above the SM background constrains the vector diquark mass to be

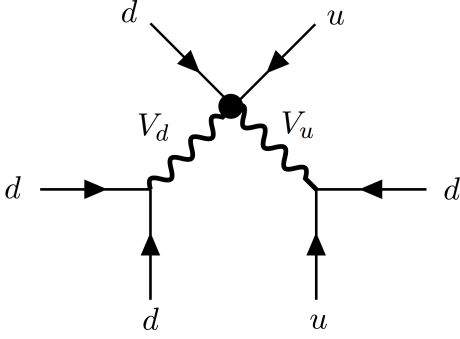$$M_{\lambda \ll 1} \gtrsim 2.5\ \mathrm{TeV} \,. \quad (25)$$

FIG. 4. Process mediating $n - \bar{n}$ oscillations via the vector diquark $V = (6, 2)_{-\frac{1}{6}}$.

B. Neutron-antineutron oscillations

In the light of null results from proton decay searches [6], the possibility of discovering $n - \bar{n}$ oscillations has recently gained increased interest [85]. The most important reason for this is that the matter-antimatter asymmetry of the Universe requires baryon number to be violated at some point during its evolution. If processes with $\Delta B = 1$ are indeed suppressed or do not occur in Nature at all, the next simplest case involves $\Delta B = 2$, a baryon number breaking pattern that may result in $n - \bar{n}$ oscillations without proton decay. Moreover, the SM augmented only by right-handed neutrinos can have an additional $U(1)_{B-L}$ gauge symmetry without introducing any gauge anomalies. Through this symmetry, processes with $\Delta B = 2$ would be accompanied by $\Delta L = 2$ lepton number violating ones, which in turn are intrinsically connected to the seesaw mechanism generating naturally small neutrino masses [86].

Models constructed so far propose $n - \bar{n}$ oscillations mediated by scalar diquarks, as mentioned in Sec. II C. Those oscillations proceed through a triple scalar vertex, so that the process is described at low energies by a local operator of dimension nine. Below we show that $n - \bar{n}$ oscillations can also be mediated by vector diquarks, in particular within the model with just one new representation discussed in this section.

The least suppressed channel is via a dimension five gauge invariant quartic interaction $\mathcal{O}_1$ in Eq. (22) involving two vector diquarks $(6, 2)_{-\frac{1}{6}}$ and the SM up and down quarks, as shown in Fig. 4, ultimately leading to $n - \bar{n}$ oscillations through a low energy effective interaction local operator of dimension nine, as in the case of scalar diquarks.

With the simplifying assumption that $c_1 \approx 1$ and neglecting other operators contributing to the signal we now estimate the rate of $n - \bar{n}$ oscillations in this model. The effective Hamiltonian corresponding to the operator $\mathcal{O}_1$ is,

$$
\mathcal{H}_{\text{eff}} \approx -\frac{\lambda_{11}^2}{M^4 \Lambda} (\overline{d_L^c})^\alpha \gamma_\mu d_R^{\alpha'} (\overline{u_L^c})^\beta \gamma_\nu d_R^{\beta'} (\overline{u_L^c})^\delta \sigma^{\mu\nu} d_R^{\delta'}
$$
$$
\times (\epsilon_{\alpha\beta\delta} \, \epsilon_{\alpha'\beta'\delta'} + \epsilon_{\alpha'\beta\delta} \, \epsilon_{\alpha\beta'\delta'} + \epsilon_{\alpha\beta'\delta} \, \epsilon_{\alpha'\beta\delta'} + \epsilon_{\alpha\beta\delta'} \, \epsilon_{\alpha'\beta'\delta})
$$
$$
+ \text{h.c.} . \tag{26}
$$

Combining this with the results of [7, 87], we obtain an esti-

mate for the $n - \bar{n}$ transition matrix element,[4]

$$
\left| \langle \bar{n} | \mathcal{H}_{\text{eff}} | n \rangle \right| \approx \frac{10^{-4} \, |\lambda_{11}|^2}{M^4 \Lambda} \, \text{GeV}^6 , \tag{28}
$$

where $|n\rangle$ is the neutron state at zero momentum. Current experimental limit [88] implies, assuming $\lambda_{11} \approx 0.01$ (which is well below the LHC bound from dijet searches, as discussed earlier), that

$$
M \gtrsim 2.5 \, \text{TeV} \left(\frac{10^8 \, \text{TeV}}{\Lambda} \right)^{1/4} . \tag{29}
$$

An interesting limit on the vector diquark mass is derived if we assume that the physics behind the triple diquark interaction with the SM Higgs is related to the physics responsible for providing the diquark its mass, i.e., for $\Lambda \approx M$. In such case the $n - \bar{n}$ oscillation search provides the bound:

$$
M \gtrsim 90 \, \text{TeV} , \tag{30}
$$

much stronger than the LHC limit. If there is new physics around that energy scale, it should be discovered by future $n - \bar{n}$ oscillation experiments with increased sensitivity [85], which are going to probe the vector diquark mass scale up to ~ 175 TeV. This is especially interesting since models with TeV-scale diquarks tend to improve gauge coupling unification [45, 46].

C. Flavor constraints

Apart from LHC bounds on the diquark coupling to first and second generation quarks λ_{ij} $(i, j = 1, 2)$ coming from dijet searches, as discussed in Sec. IV A, other stringent constraints arise from the absence of experimental evidence for flavor changing neutral currents. In particular, for the vector diquark model studied here, this includes constraints from neutral meson mixing and radiative B meson decays. Similar bounds were calculated for scalar diquark models in [7, 89, 90]. In the vector diquark case the electric dipole moment constraints do not arise, since there is only one term in the Lagragian coupling the vector diquark to quarks.

To derive the bounds from neutral meson mixing, it is useful to integrate out the vector diquark and analyze the relevant effective Hamiltonian terms:

$$
\mathcal{H}_{\text{eff}} \supset \frac{1}{M^2} \Big[\lambda_{11} \lambda_{22}^* \big(\overline{s_R} \, \gamma^\mu d_R \big)^2 + \lambda_{11} \lambda_{33}^* \big(\overline{b_R} \, \gamma^\mu d_R \big)^2
$$
$$
+ \lambda_{22} \lambda_{33}^* \big(\overline{b_R} \, \gamma^\mu s_R \big)^2 \Big] + \text{h.c.} , \tag{31}
$$

where a Fierz transformation has been performed. Comparing this with the current $K^0 - \overline{K^0}$, $B_d - \overline{B_d}$ and $B_s - \overline{B_s}$ mixing

---

[4] The single particle state normalization adopted here is:

$$
\langle n(\vec{p}) | n(\vec{p}') \rangle = \frac{E}{m} (2\pi)^3 \delta^{(3)}(\vec{p} - \vec{p}') . \tag{27}
$$

constraints [91], we obtain for $M \approx 2.5$ TeV:

$$\left|\text{Re}\left[\lambda_{11}\lambda_{22}^*\right]\right| \lesssim 5.6 \times 10^{-6} \,, \quad \left|\text{Im}\left[\lambda_{11}\lambda_{22}^*\right]\right| \lesssim 2.1 \times 10^{-8} \,,$$
$$\left|\text{Re}\left[\lambda_{11}\lambda_{33}^*\right]\right| \lesssim 2.1 \times 10^{-5} \,, \quad \left|\text{Im}\left[\lambda_{11}\lambda_{33}^*\right]\right| \lesssim 6.3 \times 10^{-6} \,,$$
$$\left|\text{Re}\left[\lambda_{22}\lambda_{33}^*\right]\right| \lesssim 4.8 \times 10^{-4} \,, \quad \left|\text{Im}\left[\lambda_{22}\lambda_{33}^*\right]\right| \lesssim 4.8 \times 10^{-4} \,,$$
$$(32)$$

with all the numbers in Eq. (32) scaling like M^2.

The radiative B meson decay bounds come mainly from $B \to K^* \gamma$ measurements and apply to the product of couplings $\lambda_{33}\lambda_{23}^*$. A detailed analysis for scalar diquarks was performed in [90] and it is very similar in the vector diquark case. However, if one assumes a hierarchical structure of the (λ_{ij}) matrix, the bound on λ_{33} from Eq. (32) is strong enough such that λ_{23} is hardly constrained at all. A careful determination of B decay constraints on other λ_{ij} couplings requires analyzing various decay channels, as described in [92], and is beyond the scope of this paper.

V. CONCLUSIONS

We have shown that lack of experimental evidence for proton decay singles out only a handful of phenomenologically viable leptoquark models. In addition, even leptoquark models with tree-level proton stability contain dangerous dimension five proton decay mediating operators and require an appropriate UV completion to remain consistent with experiments. This is especially relevant for the Standard Model extension involving the vector leptoquarks $(3,1)_{\frac{2}{3}}$ or $(3,3)_{\frac{2}{3}}$, since those are the only two models with a single new representation that do not suffer from tree-level proton decay and can explain the recently discovered anomalies in B meson decays.

The property which makes the vector leptoquark $(3,1)_{\frac{2}{3}}$ even more appealing is that it fits perfectly into the simplest Pati-Salam unification model, where it can be identified as the new gauge boson. If such an exciting scenario is indeed realized in nature, the B physics experiments can be used to actually probe the scale and various properties of grand unification!

In the second part of the paper we focused on a model with a vector diquark $(6,2)_{-\frac{1}{6}}$ and showed that neutron-antineutron oscillations can be mediated by such a vector particle. The model is somewhat constrained by LHC dijet searches; however, it can still yield a sizable neutron-antineutron oscillation signal, that can be probed in current and upcoming experiments. It can also give rise to significant four-jet event rates testable at the LHC.

It would be interesting to explore whether a vector diquark with a mass at the TeV scale can improve gauge coupling unification in non-supersymmetric grand unified theories, similarly to the scalar case [93], providing even more motivation for upgrading the neutron-antineutron oscillation experimental sensitivity.

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