

Composite leptoquarks and anomalies in B -meson decays

Ben Gripaios,^a M. Nardecchia^{a,b} and S.A. Renner^b

^a*Cavendish Laboratory, University of Cambridge,
J.J. Thomson Avenue, Cambridge, CB3 0HE, U.K.*

^b*DAMTP, University of Cambridge,
Wilberforce Road, Cambridge, CB3 0WA, U.K.*

E-mail: gripaios@hep.phy.cam.ac.uk, m.nardecchia@damtp.cam.ac.uk,
sar67@cam.ac.uk

ABSTRACT: We attempt to explain recent anomalies in semileptonic B decays at LHCb via a composite Higgs model, in which both the Higgs and an $SU(2)_L$ -triplet leptoquark arise as pseudo-Goldstone bosons of the strong dynamics. Fermion masses are assumed to be generated via the mechanism of partial compositeness, which largely determines the leptoquark couplings and implies non-universal lepton interactions. The latter are needed to accommodate tensions in the $b \rightarrow s\mu\mu$ dataset and to be consistent with a discrepancy measured at LHCb in the ratio of $B^+ \rightarrow K^+\mu^+\mu^-$ to $B^+ \rightarrow K^+e^+e^-$ branching ratios. The data imply that the leptoquark should have a mass of around a TeV. We find that the model is not in conflict with current flavour or direct production bounds, but we identify a few observables for which the new physics contributions are close to current limits and where the leptoquark is likely to show up in future measurements. The leptoquark will be pair-produced at the LHC and decay predominantly to third-generation quarks and leptons, and LHC13 searches will provide further strong bounds.

KEYWORDS: Beyond Standard Model, B-Physics, Technicolor and Composite Models

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1 Introduction

The first run of the LHC brought us the long-awaited discovery of the Higgs boson, but no firm evidence for the physics beyond the Standard Model (SM) needed to avoid fine-tuning of the weak scale. This was perhaps not unexpected, given the plethora of indirect constraints on new physics coming from, e.g., flavour physics and electroweak precision tests. Typically these point at scales of new physics way beyond a TeV; even when we invoke all the dynamical dirty tricks that we know of, the best we can do is to lower the possible scale of new physics to perhaps a few TeV. Therefore, there seems to be, nolens volens, at least a small tuning in the weak scale.

An unfortunate consequence of this is that, even if the electroweak scale is mostly natural, we may struggle to probe the associated dynamics at the LHC. At best, we might hope that one or two new states are anomalously light, such that we can either produce them on-shell, or see their effects indirectly in rare processes.¹ It is clear that discovery

¹Such light states might be present for a variety of reasons. For example, they might be desirable because they reduce fine-tuning (such as a light top squark and gluino in SUSY), or they might arise because of symmetries (such as additional pseudo-Goldstone boson states in composite Higgs models [1–3]).

of such states will require painstaking work, including careful scrutiny of all discrepancies between the data and SM predictions.

In this work, we ask, in this vein, whether anomalies recently observed in semileptonic decays of B -mesons at LHCb [4–6] can be explained by a model in which a scale of a few TeV arises naturally via strong dynamics. The necessary residual fine tuning required to generate the electroweak scale can be achieved by making the Higgs boson a pseudo-Goldstone boson (PGB) of global symmetries of the strong dynamics sector [7–10]. The Higgs potential (and thus the electroweak scale) arises due to the breaking of the global symmetries by the SM gauging and by couplings to fermions, and one can hope that there is an accidental cancellation in the various contributions, whence a somewhat lower electroweak scale emerges. The Yukawa couplings of the SM are assumed to arise via the mechanism of partial compositeness [11], which not only provides a rationale for the structure of masses and mixings observed in the quark sector, but also provides a paradigm for suppressing large flavour-violating effects in processes involving the light fermions, where the experimental constraints are strongest.

The general framework of partial compositeness is an obvious choice for explaining the anomalies, which appear in processes involving second and third generation quarks, and which appear to require new physics in muonic, but not electronic processes. To fit the detailed structure of the anomalies, we hypothesize that they are due to the presence in such a model of an anomalously light ($c.$ TeV, as it turns out) leptoquark. As pointed out in [12],² partial compositeness models necessarily feature a plethora of composite coloured fermion states, namely the composite quarks, and so it would be something of a surprise if they did not also feature composite coloured scalar states, which could couple as leptoquarks or diquarks [13]. Moreover, one can easily arrange for a leptoquark state to be rather lighter than the other resonances of the strong sector, by making it a PGB of the same symmetry breaking that gives rise to the Higgs boson.³

A disadvantage of such models is that, being strongly coupled, we cannot calculate ad libitum. But we can use naïve dimensional analysis (NDA) to compute and make predictions modulo $O(1)$ corrections. Using this framework, we find that the anomalies single out one among the possible SM irreps that allow leptoquark couplings, viz. a triplet under both $SU(3)_c$ and $SU(2)_L$. This leptoquark is one of those identified in a recent analysis [14] of LHCb B meson anomalies, in which just two, non-vanishing leptoquark couplings (to b quarks, s quarks and muons) were invoked in an ad hoc fashion to fit the anomalies. In contrast, the model considered here is underpinned by a complete (albeit presently uncalculable) framework for flavour physics, and all leptoquark couplings are non-vanishing, with magnitude fixed by the degrees of compositeness of each of the SM fermion multiplets, giving 15 mixing parameters. In the quark sector, all but one of these parameters is fixed by measurements of quark masses and the CKM matrix; there is more

²In fact, ref. [12], argued that evidence for such leptoquarks should first appear in $b \rightarrow s\mu\mu$ processes, precisely where the anomalies are now observed.

³The leptoquark is nevertheless expected to be somewhat heavier than the Higgs [12], both because it receives contributions to its potential from the QCD coupling and because we expect that the Higgs mass has been slightly tuned.

ambiguity in the lepton sector, but we find that everything can be fixed by assuming that the mixings of the left and right-handed lepton multiplets are comparable. This assumption is a plausible one, from the point of view of the UV flavour dynamics, and has the additional benefit that new physics (NP) corrections to the most severely constrained flavour-violating observable, $\mu \rightarrow e\gamma$, are minimized. As a result, we are left with just 3 free parameters in the model: the mass, M , of the leptoquark, the coupling strength, g_ρ , of the strong sector resonances, and the degree of compositeness, ϵ_3^q , of the third generation quark doublet. Furthermore, all processes to which the leptoquark contributes (with the exception of meson mixing) result in constraints on the single combination $x \equiv \sqrt{g_\rho}\epsilon_3^q/M$. Thus the model is extremely predictive. We find that the preferred range of x corresponds to plausible values of the 3 underlying parameters of the strongly coupled theory (in which the weak scale is slightly tuned), namely $g_\rho \sim 4\pi$, $M \sim \text{TeV}$, and $\epsilon_3^q \sim 1$. Thus, g_ρ and ϵ_3^q lie close to their maximal values, meaning that one cannot evade future direct searches at the LHC by scaling up M and g_ρ .

As for the existing bounds, we find that there is no obvious conflict, but that there is potential to see effects in $\mu \rightarrow e\gamma$, $K^+ \rightarrow \pi^+\nu\nu$, and $B^+ \rightarrow \pi^+\mu^+\mu^-$, in the near future. Moreover, the required mass range for the leptoquark is not far above that already excluded by LHC8, and so there is plenty of scope for discovery in direct production at LHC13.

The outline is as follows. In the next section, we describe the data anomalies and review fits thereto using higher-dimensional SM operators. We also show that they can be described by a leptoquark carrying the representation $(\bar{\mathbf{3}}, \mathbf{3}, \frac{1}{3})$ of the $SU(3) \times SU(2) \times U(1)$ gauge group. In section 3 we review the partial compositeness and strong dynamics paradigms. We show how the leptoquark can accompany the Higgs as a PGB of strong dynamics and exhibit symmetries that prevent proton decay, &c. In section 4, we discuss important constraints on the model and describe the prospects for direct searches for the leptoquark at LHC13 and indirect searches using flavour physics.

2 Status of $b \rightarrow s\ell\ell$ fits and leptoquark quantum numbers

The anomalies that we wish to explain were observed at LHCb in semileptonic B meson decays involving a $b \rightarrow s$ quark transition. These may be described via the low-energy, effective hamiltonian

$$\mathcal{H}_{\text{eff}} = -\frac{4G_F}{\sqrt{2}} (V_{ts}^* V_{tb}) \sum_i C_i^\ell(\mu) \mathcal{O}_i^\ell(\mu), \quad (2.1)$$

where \mathcal{O}_i^ℓ are a basis of $SU(3)_C \times U(1)_Q$ -invariant dimension-six operators giving rise to the flavour-changing transition. The superscript ℓ denotes the lepton flavour in the final state ($\ell \in \{e, \mu, \tau\}$), and the operators \mathcal{O}_i^ℓ are given in a standard basis by

$$\begin{aligned} \mathcal{O}_7^{(\prime)} &= \frac{e}{16\pi^2} m_b (\bar{s}\sigma_{\alpha\beta} P_{R(L)} b) F^{\alpha\beta}, \\ \mathcal{O}_9^{\ell(\prime)} &= \frac{\alpha_{\text{em}}}{4\pi} (\bar{s}\gamma_\alpha P_{L(R)} b) (\bar{\ell}\gamma^\alpha \ell), \\ \mathcal{O}_{10}^{\ell(\prime)} &= \frac{\alpha_{\text{em}}}{4\pi} (\bar{s}\gamma_\alpha P_{L(R)} b) (\bar{\ell}\gamma^\alpha \gamma_5 \ell). \end{aligned} \quad (2.2)$$

We neglect possible (pseudo-)scalar and tensor operators, since these have been shown [14, 15] to be constrained to be too small (in the absence of fine-tuning in the electron sector) to explain LHCb anomalies. In the SM, the operator coefficients are lepton universal and the operators that have non-negligible coefficients are \mathcal{O}_7 , \mathcal{O}_9^ℓ , and \mathcal{O}_{10}^ℓ , with

$$\begin{aligned} C_7^{\text{SM}} &= -0.319, \\ C_9^{\text{SM}} &= 4.23, \\ C_{10}^{\text{SM}} &= -4.41. \end{aligned} \tag{2.3}$$

at the scale m_b [16].

The first tension with the SM was observed last year in angular observables in the semileptonic decay $B \rightarrow K^* \mu^+ \mu^-$ [4, 5]. The rôle of theoretical hadronic uncertainties in the discrepancy is not yet clear [16–23]. Nevertheless, several model-independent analyses [19, 24–27] have been performed on the $B \rightarrow K^* \mu^+ \mu^-$ decay data, as well as on other, relevant, semileptonic and leptonic processes, allowing for the possibility of new physics contributions to the effective operators in eq. (2.2). There seems to be a consensus that, if only a single Wilson coefficient is allowed to be non-vanishing, then NP contributions to the effective operator \mathcal{O}_9^μ are preferred, with the NP coefficient C_9^{NP} of this operator being negative. A number of models of NP were proposed to explain this effect [28–33].

Earlier this year LHCb measured another discrepancy in B decays. To wit, it was found that a certain ratio, R_K , of branching ratios of $B \rightarrow K \mu^+ \mu^-$ to $B \rightarrow K e^+ e^-$ lay 2.6σ below the SM prediction [6]. Specifically, the observable is defined as

$$R_K = \frac{\int_1^6 dq^2 \frac{d\Gamma(B^+ \rightarrow K^+ \mu^+ \mu^-)}{dq^2}}{\int_1^6 dq^2 \frac{d\Gamma(B^+ \rightarrow K^+ e^+ e^-)}{dq^2}}, \tag{2.4}$$

where q^2 is the invariant mass of the di-lepton pair and the integral is performed over the interval $q^2 \in [1, 6] \text{ GeV}^2$. Like the $B \rightarrow K^* \mu^+ \mu^-$ decay, these processes proceed via a $b \rightarrow s \ell \ell$ transition. The observable R_K has the advantage of being theoretically well-understood, predicted to be almost exactly 1 in the SM [34] (specifically, 1.0003 ± 0.0001 when mass effects are taken into account [35]). A discrepancy in R_K cannot be explained by lepton-flavour-universal NP, nor by any of the sources of theoretical uncertainty that might underlie the $B \rightarrow K^* \mu^+ \mu^-$ anomalies. Analyses and fits including the R_K data and other recent measurements were performed in [14, 22, 36, 37]. Due to the lepton non-universality required by the R_K data, these analyses allowed the electronic and muonic Wilson coefficients to differ. They found that a negative contribution to C_9^μ remains favoured, while contributions to electronic Wilson coefficients C_i^e were found to be consistent with zero, but could have large deviations therefrom, due to larger experimental uncertainties in electronic measurements.

One could argue that the ‘axial-vector’ basis, whilst convenient for studying physics below the weak scale, is not the most natural choice in the context of models of NP above the weak scale, which must respect the chiral gauge symmetries of the SM. In the absence of multiple couplings or particles that have been somehow tuned (perhaps by

additional symmetries), NP is likely to generate operators that are coupled to a specific lepton chirality, and thus aligned with a ‘chiral basis’ in which $C_9 = -C_{10}$, $C_9 = C_{10}$, $C'_9 = -C'_{10}$, $C'_9 = C'_{10}$. Given this, the recent analyses have also made use of this basis [14, 22, 36, 37]. They find that, when looking at NP contributions in a single Wilson coefficient at a time, the best fit in this basis is achieved by a negative contribution to $C_9^\mu = -C_{10}^\mu$.

Therefore, of the possible scalar leptoquarks,⁴ which always generate contributions to one Wilson coefficient in the chiral basis, the obvious choice to explain the anomalies appears to be that with quantum numbers $(\bar{\mathbf{3}}, \mathbf{3}, 1/3)$, which contributes to the combination $C_9^\mu = -C_{10}^\mu$ at tree level.⁵ This leptoquark was already considered to explain R_K in [14], in a scenario in which its only non-zero couplings were to $b\mu$ and to $s\mu$.⁶

3 Details of the composite model

3.1 Flavour structure and leptoquark couplings

With the required quantum numbers of the leptoquark in hand, we now embed the leptoquark in a composite Higgs model.⁷ We assume, then, the presence of a new strong sector and of an elementary sector. The strong sector is characterised by a mass scale m_ρ and by a single coupling among the resonances, which is denoted by g_ρ . We expect the strong sector in isolation to have a global symmetry \mathcal{G} which is spontaneously broken by the strong dynamics to a subgroup \mathcal{H} . The SM gauge interactions are introduced in the strong sector by gauging a subgroup of \mathcal{H} . We identify the Goldstone bosons coming from the breaking \mathcal{G}/\mathcal{H} with the Higgs boson H and the leptoquark state Π . So as to avoid large contributions to other flavour observables, we seek a model in which the coset space contains only H and Π .

We postulate that the SM fermion Yukawa couplings are generated via the paradigm of partial compositeness [11]. The basic assumption is that elementary states f_i^a (where $a \in \{q, u, d, \ell, e\}$ and i is the family index) couple linearly to fermionic operators $\bar{\mathcal{O}}_i^a$ of the strong sector. For example, the relevant lagrangian required to generate the masses of the up quarks is, schematically,

$$\mathcal{L} \supset g_\rho \epsilon^q \bar{\mathcal{O}}^q q + g_\rho \epsilon^u \bar{\mathcal{O}}^u u + m_\rho (\bar{\mathcal{O}}^q \mathcal{O}^q + \bar{\mathcal{O}}^u \mathcal{O}^u) + g_\rho \bar{\mathcal{O}}^q H \mathcal{O}^u. \quad (3.1)$$

After electroweak symmetry breaking (EWSB), the resulting light mass eigenstates correspond to the SM fields and are given by linear combinations of the form

$$f_{\text{SM}}^a = \cos \theta^a f^a + \sin \theta^a \mathcal{O}^a, \quad (3.2)$$

with $\sin \theta^a = O(\epsilon^a)$. Thus, the parameters ϵ_i^a have a physical meaning: they measure the degree of compositeness of the SM fields. If $\epsilon_i^a \lesssim 1$, we have that (at leading order in ϵ) $f_{\text{SM}} \approx f$ and the projections of the composite operators onto the SM fields are given

⁴For a review see [38].

⁵Vector leptoquarks with charges $(\mathbf{3}, \mathbf{1}, 2/3)$ or $(\mathbf{3}, \mathbf{3}, 2/3)$ also generate the required structure, but cannot be directly realized as Goldstone bosons.

⁶Another interpretation of both the R_K anomaly and a deviation seen at CMS, in the context of R -parity violating supersymmetry, was given in [39].

⁷For a recent review see [40].

by $(\mathcal{O}^a)_{\text{SM}} \sim \epsilon^a f_{\text{SM}}$. In this way, projecting operators such as $g_\rho \overline{\mathcal{O}}^q H \mathcal{O}^u$ along the SM components, we can read off the strength of the Yukawa interactions. In particular, for the the up and down quarks, we have

$$(Y_u)_{ij} \sim g_\rho \epsilon_i^q \epsilon_j^u, \quad (Y_d)_{ij} \sim g_\rho \epsilon_i^q \epsilon_j^d. \quad (3.3)$$

Throughout this section, we use the symbol \sim to mean a relation that holds up to an unknown $O(1)$ coefficient whose value is fixed by the uncalculable strong sector dynamics. With an appropriate choice of the values of ϵ_i^q , ϵ_i^u , and ϵ_i^d , it is possible to reproduce the hierarchy of the quark masses and the mixing angles of the CKM matrix. We find

$$\begin{aligned} g_\rho v \epsilon_i^q \epsilon_i^u &\sim m_i^u, & g_\rho v \epsilon_i^q \epsilon_i^d &\sim m_i^d, & (3.4) \\ \frac{\epsilon_1^q}{\epsilon_2^q} &\sim \lambda, & \frac{\epsilon_2^q}{\epsilon_3^q} &\sim \lambda^2, & \frac{\epsilon_1^q}{\epsilon_3^q} &\sim \lambda^3, \end{aligned}$$

where v is the Higgs VEV, $\lambda = 0.23$ is the Cabibbo angle and m_i^u and m_i^d are the masses of the up- and down-type quarks, respectively. In our framework, then, the Yukawa sector is described by 10 parameters $(g_\rho, \epsilon_i^q, \epsilon_i^u, \epsilon_i^d)$. The phenomenological relations (3.4) can be used to reduce the number of free parameters that we can use to fit the anomalies. Indeed, there are 8 independent relations in (3.4) and we choose to parametrize everything in terms of g_ρ and ϵ_3^q . In the lepton sector, there is more arbitrariness in the values of ϵ_i^ℓ and ϵ_i^e . This is due to the fact that there are several mechanisms that can be envisaged for introducing mass terms in the neutrino sector. In order to make progress, we shall assume the left and right mixing parameters to be of the same order, $\epsilon_i^\ell \approx \epsilon_i^e$. This assumption about the unknown flavour dynamics at high scales is a plausible one, but it also has the phenomenological advantage that it mitigates constraints on NP coming from lepton flavour violating (LFV) observables, such as $\mu \rightarrow e\gamma$, which are the most problematic flavour-violating observables for partial compositeness models [41, 42]. Indeed, physics at the scale m_ρ generates a contribution to the radiative LFV decays of the form $\Gamma(\ell^i \rightarrow \ell^j \gamma) \sim \left| \epsilon_i^\ell \epsilon_j^e \right|^2 + \left| \epsilon_j^\ell \epsilon_i^e \right|^2$. Considering the mass constraints $\epsilon_i^\ell \epsilon_i^e = \frac{m_i^e}{g_\rho v} \delta_{ij}$, it is easy to show that $\left| \epsilon_i^\ell \epsilon_j^e \right|^2 + \left| \epsilon_j^\ell \epsilon_i^e \right|^2$ is minimized when

$$\frac{\epsilon_i^\ell}{\epsilon_j^\ell} \sim \frac{\epsilon_i^e}{\epsilon_j^e} \sim \sqrt{\frac{m_i^e}{m_j^e}}. \quad (3.5)$$

Evidently, this condition is implied by (but does not imply) our assumption that the left and right leptonic mixings are equal.

In this way, we are able to fix all parameters in the lepton sector in terms of g_ρ , and so all the NP effects of the model are parameterized by M , g_ρ , and ϵ_q^3 . The phenomenological inputs and the expressions of the various mixing parameters are summarised in tables 1 and 2.

We may now determine the leptoquark couplings, as follows. Similarly to [44], below the scale of the strongly-coupled resonances we can describe the low energy physics by an effective field theory (EFT) of the form

$$\mathcal{L} = \frac{m_\rho^4}{g_\rho^2} \mathcal{L}^{(0)} \left(\frac{g_\rho \epsilon_i^a f_i^a}{m_\rho^{3/2}}, \frac{D^\mu}{m_\rho}, \frac{g_\rho H}{m_\rho}, \frac{g_\rho \Pi}{m_\rho} \right). \quad (3.6)$$

Fermion	Mass
e	0.487 MeV
μ	103 MeV
τ	1.78 GeV
d	$2.50^{+1.08}_{-1.03}$ MeV
s	47^{+14}_{-13} MeV
b	2.43 ± 0.08 GeV
u	$1.10^{+0.43}_{-0.37}$ MeV
c	0.53 ± 0.07 GeV
t	150.7 ± 3.4 GeV

Table 1. Values of running fermion masses at the scale $\mu = 1$ TeV [43].

Mixing Parameter	Value
$\epsilon_1^q = \lambda^3 \epsilon_3^q$	$1.15 \times 10^{-2} \epsilon_3^q$
$\epsilon_2^q = \lambda^2 \epsilon_3^q$	$5.11 \times 10^{-2} \epsilon_3^q$
$\epsilon_1^u = \frac{m_u}{vg_\rho} \frac{1}{\lambda^3 \epsilon_3^q}$	$5.48 \times 10^{-4} / (g_\rho \epsilon_3^q)$
$\epsilon_2^u = \frac{m_c}{vg_\rho} \frac{1}{\lambda^2 \epsilon_3^q}$	$5.96 \times 10^{-2} / (g_\rho \epsilon_3^q)$
$\epsilon_3^u = \frac{m_t}{vg_\rho} \frac{1}{\epsilon_3^q}$	$0.866 / (g_\rho \epsilon_3^q)$
$\epsilon_1^d = \frac{m_d}{vg_\rho} \frac{1}{\lambda^3 \epsilon_3^q}$	$1.24 \times 10^{-3} / (g_\rho \epsilon_3^q)$
$\epsilon_2^d = \frac{m_s}{vg_\rho} \frac{1}{\lambda^2 \epsilon_3^q}$	$5.29 \times 10^{-3} / (g_\rho \epsilon_3^q)$
$\epsilon_3^d = \frac{m_b}{vg_\rho} \frac{1}{\epsilon_3^q}$	$1.40 \times 10^{-2} / (g_\rho \epsilon_3^q)$
$\epsilon_1^\ell = \epsilon_1^e = \left(\frac{m_e}{g_\rho v}\right)^{1/2}$	$1.67 \times 10^{-3} / g_\rho^{1/2}$
$\epsilon_2^\ell = \epsilon_2^e = \left(\frac{m_\mu}{g_\rho v}\right)^{1/2}$	$2.43 \times 10^{-2} / g_\rho^{1/2}$
$\epsilon_3^\ell = \epsilon_3^e = \left(\frac{m_\tau}{g_\rho v}\right)^{1/2}$	$0.101 / g_\rho^{1/2}$

Table 2. Partial compositeness mixing parameters and values.

In the strongly-coupled, UV theory we expect the presence of an operator of the form $g_\rho \Pi \bar{\mathcal{O}}^L \mathcal{O}^Q$, where \mathcal{O}^Q (or \mathcal{O}^L) is a composite operator with the same quantum numbers as a SM quark (or lepton). Below the scale m_ρ , this operator generates a contribution to \mathcal{L} of the form $\sim g_\rho \epsilon_i^\ell \epsilon_j^q \Pi \ell_i q_j$. At low energies, the renormalizable lagrangian of the model is

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + (D^\mu \Pi)^\dagger D_\mu \Pi - M^2 \Pi^\dagger \Pi + \lambda_{ij} \bar{q}_{Lj}^c i \tau_2 \tau_a \ell_{Li} \Pi + \text{h.c.}, \quad (3.7)$$

with $\lambda_{ij} = g_\rho c_{ij} \epsilon_i^\ell \epsilon_j^q$, where we have omitted quartic terms involving H and Π that are not relevant to our discussion. Note that we have explicitly re-introduced the c_{ij} parameters that are expected to be of $O(1)$, but are otherwise unknown. We summarise the values of the leptoquark couplings in table 3.

$\lambda_{ij}/(c_{ij}g_\rho^{1/2}\epsilon_3^q)$	$j = 1$	$j = 2$	$j = 3$
$i = 1$	1.92×10^{-5}	8.53×10^{-5}	1.67×10^{-3}
$i = 2$	2.80×10^{-4}	1.24×10^{-3}	2.43×10^{-2}
$i = 3$	1.16×10^{-3}	5.16×10^{-3}	0.101

Table 3. Values of leptoquark couplings, λ_{ij} , where i denotes the lepton generation label and j the quark generation label.

3.2 Coset structure

Here we supply a coset space construction that gives rise to the required SM quantum numbers for the Higgs and leptoquark fields. First we describe the pattern of spontaneous breaking of the symmetry of the strong sector \mathcal{G}/\mathcal{H} , and the embedding of the SM gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$ therein. We then discuss additional symmetry structure required to avoid constraints from nucleon decay and neutron-antineutron oscillations.

To build a coset, we start from the minimal composite Higgs model [10], in which a single SM Higgs doublet arises from the spontaneous breaking of $SO(5)$ to $SU(2)_H \times SU(2)_R$, with H transforming as a $(\mathbf{2}, \mathbf{2})$ of the unbroken subgroup. We must now enlarge the coset space somehow to include the leptoquark Π and its conjugate Π^\dagger . To see how this may be achieved, consider first a model with just the leptoquark and no Higgs boson. This can be achieved using $SO(9)$ broken to $SU(4) \times SU(2)_\Pi$. The 6 Goldstone bosons, (Π, Π^\dagger) , transform as $(\mathbf{6}, \mathbf{3})$.

Now form the direct product of $SO(5)$ and $SO(9)$ and consider the coset space

$$\frac{SO(9) \times SO(5)}{SU(4) \times SU(2)_\Pi \times SU(2)_H \times SU(2)_R}. \quad (3.8)$$

This has, of course, the same Goldstone boson content as the two models above. The trick is to somehow embed the SM gauge group in \mathcal{H} so as to get the right charges for H and Π . To do so we embed $SU(3)_C$ into $SU(4)$. Explicitly, $SU(4)$ contains a maximal subgroup $SU(3)_C \times U(1)_\psi$, and the decomposition of the 6-d irrep of $SU(4)$ under this group gives $\mathbf{6} = \mathbf{3}_{2/3} + \bar{\mathbf{3}}_{-2/3}$. We then embed $SU(2)_L$ as the diagonal subgroup of $SU(2)_H \times SU(2)_\Pi$, while the hypercharge gauge group $U(1)_Y$ is embedded as $T_Y = -\frac{1}{2}T_\psi + T_{3R} + T_X$, where T_ψ generates $U(1)_\psi$, T_{3R} belongs to the $SU(2)_R$ algebra, and $U(1)_X$ is an additional symmetry (under which the Higgs and the leptoquark are uncharged) which may be required to reproduce the correct SM hypercharge assignments. It is now straightforward to show that the SM quantum numbers of H and $\Pi + \Pi^\dagger$ are respectively $(\mathbf{1}, \mathbf{2}, 1/2)$ and $(\bar{\mathbf{3}}, \mathbf{3}, 1/3) + (\mathbf{3}, \bar{\mathbf{3}}, -1/3)$, as required.

We next need to show that the necessary Yukawa and leptoquark couplings can be generated by linear mixing of the elementary fermions of the SM with resonances of the strong sector carrying suitable representations of the group \mathcal{H} . In fact, a number of representations are available. One suitable assignment is

$$\begin{aligned} \mathcal{O}_q &\sim (4, \mathbf{1}, \mathbf{2}, \mathbf{2})_{+1/2}, & \mathcal{O}_u &\sim (4, \mathbf{1}, \mathbf{1}, \mathbf{1})_{+1/2}, & \mathcal{O}_d &\sim (4, \mathbf{1}, \mathbf{1}, \mathbf{3})_{+1/2}, \\ \mathcal{O}_L &\sim (4, \mathbf{3}, \mathbf{2}, \mathbf{1})_{-1/2}, & \mathcal{O}_e &\sim (4, \mathbf{3}, \mathbf{3}, \mathbf{1})_{-1/2} \end{aligned} \quad (3.9)$$

where the subscript denotes the charge under the $U(1)_X$ symmetry. It is straightforward to check that this assignment permits tri-linear couplings between the fermionic resonances and H and Π that yield the desired Yukawa and leptoquark couplings after mixing with the elementary fermions.

An advantage of this assignment is that we can use it to protect $\Gamma(Z \rightarrow b\bar{b})$. This is desirable since, with $g_\rho \sim 4\pi$, there would otherwise be sizable corrections to $\Gamma(Z \rightarrow b\bar{b})$, even with $m_\rho \sim 10 \text{ TeV}$. The protection cannot be achieved in exactly the same way as in [45], because the semi-direct product $(SU(2)_L \times SU(2)_R) \rtimes \mathbb{Z}_2$ used there is not a subgroup of \mathcal{G} . But we can instead use the symmetry $(SU(2)_H \times SU(2)_R) \rtimes \mathbb{Z}_2$, with much the same result. In a nutshell (for more details, see [46]), the group⁸ $SO(5) \subset G$ contains not just $SU(2)_H \times SU(2)_R$, but also the larger subgroup $\mathcal{K} \equiv (SU(2)_H \times SU(2)_R) \rtimes \mathbb{Z}_2$. We require: (i) that this larger group be contained in \mathcal{H} ; (ii) that b_L couple to a resonance of the strong sector transforming as a $(2, 2)$ under $SU(2)_H \times SU(2)_R \subset \mathcal{K}$ and as either the trivial irrep or the sign irrep under $\mathbb{Z}_2 \subset \mathcal{K}$; and (iii) that the coupling of b_L to the strong sector respect the subgroup $(U(1)_H \times U(1)_R) \rtimes \mathbb{Z}_2 \subset \mathcal{K}$. With these three requirements, a straightforward modification of the arguments given in [45, 46] shows that there can be no corrections to $\Gamma(Z \rightarrow b\bar{b})$.

There is, however, a disadvantage with this assignment, in that the linear mixing between \mathcal{O}_q and q breaks the $SU(2)_L \times SU(2)_R$ custodial symmetry, which is often invoked to protect $\frac{m_W}{m_Z}$. Since $\epsilon_3^q = 1$, these corrections are unsuppressed. Happily, we find thanks to $m_\rho \sim 10 \text{ TeV}$ and to the presence of light custodians [47, 48], we are consistent with the bounds coming from EWPT observables.⁹

The global symmetry \mathcal{G} is broken explicitly by the gauging of the SM group, as well as by the linear couplings between the elementary and composite sector. As a result of these breakings, the PGBs get a mass term. NDA suggests that the main contribution to the effective potential of the Higgs comes from the top Yukawa coupling. This implies a negative contribution to the Higgs mass parameter [10], which can trigger EWSB, and the resulting Higgs mass is expected to be of order $m_H^2 \sim \frac{y_t^2}{16\pi^2} m_\rho^2$. In contrast to the Higgs boson, the composite leptoquark gets its dominant mass term contribution from QCD. The resulting leptoquark mass is of order $m_\Pi^2 \sim \frac{g_s^2}{16\pi^2} m_\rho^2$ and is positive-definite, avoiding the danger of colour- and charge-breaking vacua.

We now move on to discuss constraints from nucleon decay, &c. In models with TeV scale strong dynamics, we cannot expect the accidentally symmetries of the SM that lead to conservation of baryon and lepton number to be preserved. This problem is exacerbated in our model with a light leptoquark state, since the SM gauge symmetry allows a $(\bar{\mathbf{3}}, \mathbf{3}, \frac{1}{3})$ leptoquark to couple to both qq and $q\ell$, and thus mediate proton decay.

We now assess whether additional global symmetries can be imposed to prevent such decays. Our objective is to allow the coupling to $q\ell$, but not that to qq . Evidently, then,

⁸Or rather, strictly speaking, its universal cover $Sp(2)$.

⁹Note that with the alternative assignment $\mathcal{O}_q \sim (4, 2, 1, 1)$, $\mathcal{O}_u \sim (4, 2, 2, 2)$, $\mathcal{O}_d \sim (4, 2, 2, 2)$, $\mathcal{O}_L \sim (4, 2, 1, 1)$, $\mathcal{O}_e \sim (4, 2, 2, 2)$, (for which an additional $U(1)_X$ is not necessary), the linear mixing between \mathcal{O}_q and q is $SU(2)_L \times SU(2)_R$ invariant, and corrections to $\frac{m_W}{m_Z}$ are suppressed by powers of $(\epsilon_3^u)^4 \ll 1$. But then one must relinquish custodial protection of $\Gamma(Z \rightarrow b\bar{b})$.

q and ℓ must carry different charges, e_q and e_ℓ , say, under such a symmetry.¹⁰ We must now decide whether the leptoquark itself should carry charge or not.

The easiest option to realise is for the leptoquark not to carry a charge. Then the corresponding symmetry can lie outside of the $SO(9)$ group of which the Π is a Goldstone boson. Then the leptoquark coupling is allowed if $e_q + e_\ell = 0$. A problem with any such symmetry is that it cannot forbid decays of 3 quarks to 3 anti-leptons. So, while the usual suspects, like $p \rightarrow e^+\pi^0$ are forbidden, decays such as $p \rightarrow e^+2\bar{\nu}$ and $n \rightarrow 3\bar{\nu}$ are not. In our framework the most stringent bound comes from searches for $pp \rightarrow \mu^+2\bar{\nu}$ decays, where $\Gamma < 1.0 \times 10^{-63}$ GeV [49]. The leading contribution to this processes is generated by the dimension-9 operator $(qqd^{c\dagger})(\ell\ell e^{c\dagger})$, with τ neutrinos. A NDA estimate gives

$$\Gamma(p \rightarrow \mu^+\bar{\nu}_\tau\bar{\nu}_\tau)_{NDA} = \frac{m_p^{11}}{(4\pi)^3} \left(g_\rho^4 \frac{\epsilon_1^d(\epsilon_1^q)^2 \epsilon_2^\ell (\epsilon_3^\ell)^2}{M^5} \right)^2 = 4.7 \times 10^{-53} \text{ GeV}^{-1}. \quad (3.10)$$

It is then clear that the searches for such decays suffice to rule out a model with compositeness at multi-TeV scales. In comparing with the bound, we have used the values

$$M = 1 \text{ TeV}, \quad g_\rho = 4\pi, \quad \text{and} \quad \epsilon_3^q = 1, \quad (3.11)$$

and we shall continue to do so henceforth.

We need, therefore, to explore the alternative option, which is to look for a symmetry that lies (at least partly) within $SO(9)$, such that the leptoquark is charged. A simple expedient is to use the $\mathbb{Z}_2 \subset SO(9)$ symmetry whose non-trivial element in the defining representation of $SO(9)$ is the matrix $\begin{pmatrix} -I_6 & 0 \\ 0 & I_3 \end{pmatrix}$, where I_n is the $n \times n$ identity matrix. This element commutes with $SO(6) \times SO(3)$ (and therefore is unbroken by the gauging of the SM subgroup) but anti-commutes with the broken generators in $SO(9)/SO(6) \times SO(3)$, meaning that the leptoquarks transform under \mathbb{Z}_2 as $\Pi \rightarrow -\Pi$. Now, by insisting that the \mathbb{Z}_2 be unbroken by the strong dynamics and the couplings to elementary fermions, the diquark coupling Πqq is forbidden. Provided, moreover, that the elementary q and ℓ are assigned opposite charges, the leptoquark coupling $\Pi \ell q$ is allowed. Yukawa couplings can be retained by assigning the elementary (u^c, d^c) and e^c to have the same charges as q and ℓ , respectively.

Such a symmetry (which may be thought of as either a baryon or lepton parity) stabilizes nucleons completely, and so also solves potential problems from generic operators generated by the heavier resonances of the strong dynamics. Its drawback is that it cannot forbid neutron-antineutron oscillations,¹¹ for which there are again strong experimental constraints. There are two dimension 9 operators in the EFT that could give a contribution to this process, namely $qqqq(d^c d^c)^\dagger$ and $u^c d^c d^c u^c d^c d^c$. The low-energy effects of these operators are subject to large hadronic uncertainties; we estimate a rough bound on the necessary scale as $\Lambda \gtrsim 100$ TeV.

¹⁰We assume that all particles come in 1-d representations of the symmetry, so as not to have to introduce additional states.

¹¹For a review see [50].

In our leptoquark model, we expect to generate the operator

$$\frac{g_\rho^4(\epsilon_1^q)^4(\epsilon_1^d)^2}{M^5} qqqq(d^c d^c)^\dagger. \quad (3.12)$$

Using the nominal values in (3.11) and matching with the previous expression, we find $\Lambda = 188$ TeV. Given the high dimension of the operator, this scale comes with a large uncertainty, but it would seem that we are safe.

Finally, we remark that we could, of course, invoke both symmetries discussed above, in order to forbid both nucleon decay and oscillations absolutely.

4 Phenomenological analysis

At tree level, the effects of the leptoquark on flavour physics observables can be studied using the effective lagrangian

$$\mathcal{L}_{LQ}^{\text{eff}} = \sum_{ijkl} \frac{\lambda_{ij}(\lambda_{lk})^*}{4M^2} [(\bar{q}_j \tau^a \gamma^\mu P_L q_k)(\bar{\ell}_i \tau^a \gamma_\mu P_L \ell_\ell) + 3(\bar{q}_j \gamma^\mu P_L q_k)(\bar{\ell}_i \gamma_\mu P_L \ell_\ell)], \quad (4.1)$$

where $i, j, k, \ell \in \{1, 2, 3\}$ are generation indices. We work in a basis where the CKM matrix acts on the up sector such that q_j is the quark doublet, $q_j = \left(V_{CKM}^{\dagger jk} u_L^k, d_L^j \right)^T$, and ℓ_i is the lepton doublet, $\ell_i = (\nu^i, e_L^i)^T$. We assume that the mass differences between the components of the leptoquark triplet are small compared to the masses themselves, so that the components can be assumed to have a common mass, M . Therefore we may write

$$\begin{aligned} \mathcal{L}_{LQ}^{\text{eff}} = \sum_{ijkl} \frac{\lambda_{ij}(\lambda_{lk})^*}{2M^2} & \left[2(\bar{d}_L \gamma^\mu d_L)_{kj} (\bar{e}_L \gamma_\mu e_L)_{\ell i} + 2(\bar{u}'_L \gamma^\mu u'_L)_{kj} (\bar{\nu}_L \gamma_\mu \nu_L)_{\ell i} \right. \\ & + (\bar{d}_L \gamma^\mu d_L)_{kj} (\bar{\nu}_L \gamma_\mu \nu_L)_{\ell i} + (\bar{u}'_L \gamma^\mu u'_L)_{kj} (\bar{e}_L \gamma_\mu e_L)_{\ell i} \\ & \left. + (\bar{u}'_L \gamma^\mu d_L)_{kj} (\bar{e}_L \gamma_\mu \nu_L)_{\ell i} + (\bar{d}_L \gamma^\mu u'_L)_{kj} (\bar{\nu}_L \gamma_\mu e_L)_{\ell i} \right], \end{aligned} \quad (4.2)$$

where $u_L^{\prime j} = V_{CKM}^{\dagger jk} u_L^k$. All unprimed fields are mass eigenstates.¹²

We now comment briefly on the qualitative consequences of the various operators that appear above.

(i) Flavour changing neutral currents (FCNC) in the down quark sector

These are generated by the operators $(\bar{d}_L \gamma^\mu d_L)_{kj} (\bar{e}_L \gamma_\mu e_L)_{\ell i}$ and $(\bar{d}_L \gamma^\mu d_L)_{kj} (\bar{\nu}_L \gamma_\mu \nu_L)_{\ell i}$. They can mediate meson decays via the transitions $b \rightarrow s\ell\ell$, $b \rightarrow s\nu\nu$, $s \rightarrow d\ell\ell$, $s \rightarrow d\nu\nu$, $b \rightarrow d\ell\ell$ and $b \rightarrow d\nu\nu$.

The $b \rightarrow s\ell\ell$ transition is the main motivation for this work and will be discussed in more detail below. The decays involving neutrinos can have large NP contributions, because couplings to tau neutrinos are large in the partial compositeness framework considered here. We provide a quantitative analysis of the decays $B \rightarrow K^{(*)}\nu\nu$ and

¹²We neglect neutrino masses.

$K \rightarrow \pi\nu\nu$ below. Constraints on leptoquark couplings from measurements of (lepton-flavour-conserving) K and B decays are summarized in table 4 below, excluding $b \rightarrow s\ell\ell$ and $b \rightarrow s\nu\nu$ processes, which will be discussed in more detail in the text. Lepton-flavour-violating (LFV) processes, recently investigated in [51], are also possible in our set-up, but current bounds on these are weak. We will comment more on LFV processes in section 4.2.5.

(ii) FCNC in the up quark sector

These are generated by the operators $(\bar{u}'_L \gamma^\mu u'_L)_{kj} (\bar{\nu}_L \gamma_\mu \nu_L)_{\ell i}$ and $(\bar{u}'_L \gamma^\mu u'_L)_{kj} (\bar{e}_L \gamma_\mu e_L)_{\ell i}$. They can mediate decays of charmed mesons via $c \rightarrow u\ell\ell$ and $c \rightarrow u\nu\nu$ transitions. Bounds on these processes are weak, and we know of no bounds for decays with τ leptons or neutrinos in the final state, which would receive the largest NP contributions. These operators can also generate top decays into u or c quarks plus a pair of charged leptons or of neutrinos. The rates of these decays will be very small relative to current limits on FCNC top quark decays [49] (which in any case search specifically for $t \rightarrow Zq$, meaning they cannot be directly applied to leptoquarks). Since current measurements of FCNC in the up sector do not provide strong constraints on our model, we will not discuss them further.

(iii) Charged currents

These are generated by the operators $(\bar{u}'_L \gamma^\mu d_L)_{kj} (\bar{\nu}_L \gamma_\mu \nu_L)_{\ell i}$ and $(\bar{d}_L \gamma^\mu u'_L)_{kj} (\bar{\nu}_L \gamma_\mu \nu_L)_{\ell i}$. Processes generated by these operators are also present at tree level in the SM, so NP contributions are not expected to be large relative to the SM predictions. The largest NP rates will occur in processes with τ or ν_τ in the final state.

With these considerations in mind, in the remainder of this section we discuss the values of the model parameters that are needed to fit recent B -decay anomalies and then list important constraints on the model and predictions for its effects in other processes.

4.1 Anomalies in B decays

4.1.1 Fit to muonic $\Delta B = \Delta S = 1$ processes

We consider recent results of [22], in which a fit to all available data on muonic (or lepton-universal) $\Delta B = \Delta S = 1$ processes is described. A part of that work involved allowing one Wilson Coefficient (or chiral combination thereof) to vary while assuming all other coefficients are set to their SM values (for details of the fit please see [22]). The best fit value found in this way for the chiral combination relevant to our leptoquark is $C_9^{NP\mu} = -C_{10}^{NP\mu} = -0.55$, with 1σ and 2σ ranges

$$C_9^{NP\mu} = -C_{10}^{NP\mu} \in [-0.74, -0.36] \quad (\text{at } 1\sigma), \tag{4.3}$$

$$C_9^{NP\mu} = -C_{10}^{NP\mu} \in [-0.95, -0.19] \quad (\text{at } 2\sigma). \tag{4.4}$$

It can be seen, by comparing the effective leptoquark lagrangian in (4.2) with the effective hamiltonian in (2.1), that, for our model,

$$C_9^{\mu NP} = -C_{10}^{\mu NP} = \left[\frac{4G_F e^2 (V_{ts}^* V_{tb})}{16\sqrt{2}\pi^2} \right]^{-1} \frac{\lambda_{22}^* \lambda_{23}}{2M^2} = -0.24 c_{22}^* c_{23} (\epsilon_3^q)^2 \left(\frac{M}{\text{TeV}} \right)^{-2} \left(\frac{g_\rho}{4\pi} \right), \quad (4.5)$$

and the requirements on the parameters are

$$\text{Re}(c_{22}^* c_{23}) = 2.30 \left(\frac{4\pi}{g_\rho} \right) \left(\frac{1}{\epsilon_3^q} \right)^2 \left(\frac{M}{\text{TeV}} \right)^2 \quad (\text{Best fit}), \quad (4.6)$$

$$\text{Re}(c_{22}^* c_{23}) \in [1.50, 3.08] \left(\frac{4\pi}{g_\rho} \right) \left(\frac{1}{\epsilon_3^q} \right)^2 \left(\frac{M}{\text{TeV}} \right)^2 \quad (\text{at } 1\sigma), \quad (4.7)$$

$$\text{Re}(c_{22}^* c_{23}) \in [0.79, 3.96] \left(\frac{4\pi}{g_\rho} \right) \left(\frac{1}{\epsilon_3^q} \right)^2 \left(\frac{M}{\text{TeV}} \right)^2 \quad (\text{at } 2\sigma). \quad (4.8)$$

Thus, if this anomaly is to be explained, there are 3 immediate implications for the parameters of our model:

1. the mass of the leptoquark states should be low enough, $M \lesssim 1 \text{ TeV}$, to be within the reach of the second run of the LHC;
2. the left-handed doublet of the third quark family should be largely composite, $\epsilon_3^q \sim 1$;
3. the composite sector must be genuinely strongly interacting, $g_\rho \sim 4\pi$.

Indeed, if any one of these does not hold then we are forced to set $\text{Re}(c_{22}^* c_{23}) \gg 1$, implying an inconsistency with the EFT paradigm described in the previous section.

4.1.2 R_K

R_K , as defined in eq. (2.4), has been recently measured by LHCb to be $R_K = 0.745_{-0.074}^{+0.090} \pm 0.036$ [6]. Roughly, adding errors in quadrature, we therefore take the measured value at the 1σ level to be within the range [0.66, 0.84]. The model prediction, including the effect of the leptoquark, is given approximately by

$$R_K \approx \frac{\left| C_{10}^{\text{SM}} + C_{10}^{\mu NP} + C_{10}^{\mu'} \right|^2 + \left| C_9^{\text{SM}} + C_9^{\mu NP} + C_9^{\mu'} \right|^2}{\left| C_{10}^{\text{SM}} + C_{10}^{e NP} + C_{10}^{e'} \right|^2 + \left| C_9^{\text{SM}} + C_9^{e NP} + C_9^{e'} \right|^2}, \quad (4.9)$$

which can be found from the full expression by neglecting the coefficient of the dipole operator, C_7 . (In the SM C_7 has a magnitude less than 10% that of C_9 or C_{10} , and NP contributions to it are constrained small by the measured branching ratio of $B \rightarrow X_s \gamma$). The NP values of the Wilson coefficients are

$$C_9^{\mu NP} = -C_{10}^{\mu NP} = -0.24 c_{22}^* c_{23} (\epsilon_3^q)^2 \left(\frac{M}{\text{TeV}} \right)^{-2} \left(\frac{g_\rho}{4\pi} \right), \quad (4.10)$$

$$C_9^{e NP} = -C_{10}^{e NP} = -1.2 \times 10^{-3} c_{11}^* c_{23} (\epsilon_3^q)^2 \left(\frac{M}{\text{TeV}} \right)^{-2} \left(\frac{g_\rho}{4\pi} \right). \quad (4.11)$$

The values of C_9^{SM} and C_{10}^{SM} are given in eq. (2.3).

We see that, due to the structure of partial compositeness, NP contributions in the decay $B^+ \rightarrow K^+ e^+ e^-$ are negligible. Neglecting these and the quadratic terms in $C_{9,10}^{\mu NP}$, we obtain

$$\text{Re}(c_{22}^* c_{23}) \in [1.42, 2.98] \left(\frac{4\pi}{g_\rho}\right) \left(\frac{1}{\epsilon_3^q}\right)^2 \left(\frac{M}{\text{TeV}}\right)^2 \quad (\text{at } 1\sigma). \quad (4.12)$$

The allowed region thus has reasonable overlap with the 1σ region found above using a fit to muonic $\Delta B = \Delta S = 1$ observables. Therefore our model is able to fit the muonic data and R_K with no tension between the two. Of course, this is hardly surprising as several works [15, 22, 36] have pointed out the compatibility of the $b \rightarrow s\mu\mu$ data with R_K if the NP is predominantly in the muon sector, rather than the electron sector. This feature is automatic in models with partial compositeness.

4.2 Important constraints and predictions

The largest couplings of the composite leptoquark are to third generation quarks and leptons. Therefore, generically, the most important constraints and predictions will be in processes involving third generation quarks and fermions in initial or final states and also processes with third-generation fermions in a loop.¹³ This section will look at some of these processes, discussing implications of current measurements on our model, as well as highlighting promising channels for probing our scenario with future measurements.

4.2.1 $b \rightarrow s\nu\nu$

Due to the $SU(2)_L$ structure of the leptoquark, it will couple to neutrinos as well as charged leptons and thus induce $b \rightarrow s\nu\nu$ transitions. The importance of this channel in general for pinning down NP has been recently emphasised in [52]. These $B \rightarrow K^* \nu\bar{\nu}$ and $B \rightarrow K \nu\bar{\nu}$ decays are good channels to look for large effects from the composite leptoquark we consider. Indeed, since the identity of the neutrino cannot be determined in these experiments, large contributions from the processes involving tau neutrinos are expected in our model. Thus our model predicts a much larger rate than that expected in models where NP couples only to the second generation lepton doublet.

Current NP bounds from these decays can be found in [52], which are quoted in terms of ratios to Standard Model predictions. With a slight alteration of the notation of [52], so as not to cause confusion with the notation used here, the relevant quantities, and the limits thereon, are

$$R_K^{*\nu\nu} \equiv \frac{\mathcal{B}(B \rightarrow K^* \nu\bar{\nu})}{\mathcal{B}(B \rightarrow K^* \nu\bar{\nu})_{\text{SM}}} < 3.7, \quad (4.13)$$

and

$$R_K^{\nu\nu} \equiv \frac{\mathcal{B}(B \rightarrow K \nu\bar{\nu})}{\mathcal{B}(B \rightarrow K \nu\bar{\nu})_{\text{SM}}} < 4.0. \quad (4.14)$$

The leptoquark can in principle induce transitions involving any combination of neutrino flavours, since it couples to all generations and also has flavour-violating couplings. There

¹³Of course, this is only true generically, since the sensitivity depends not only on the size of the NP contribution, but also on the experimental feasibility and also the size and nature of the competing SM contributions.

will be interference between NP and SM processes only in flavour-conserving transitions. The NP contributions to the $\bar{\nu}_\tau\nu_\tau$ and $\bar{\nu}_\mu\nu_\mu$ processes will induce a shift from unity in $R_K^{\nu\nu}$ and $R_K^{(*)\nu\nu}$ given by

$$\Delta(R_K^{(*)\nu\nu})^{\tau\tau} = \left[0.220 \operatorname{Re}(c_{32}^*c_{33}) + 0.0363 |c_{32}^*c_{33}|^2 (\epsilon_3^q)^2 \left(\frac{M}{\text{TeV}}\right)^{-2} \left(\frac{g_\rho}{4\pi}\right) \right] (\epsilon_3^q)^2 \left(\frac{M}{\text{TeV}}\right)^{-2} \left(\frac{g_\rho}{4\pi}\right),$$

$$\Delta(R_K^{(*)\nu\nu})^{\mu\mu} \approx 1.27 \times 10^{-2} \operatorname{Re}(c_{32}^*c_{33}) (\epsilon_3^q)^2 \left(\frac{M}{\text{TeV}}\right)^{-2} \left(\frac{g_\rho}{4\pi}\right).$$

(The expression for $\Delta(R_K^{(*)\nu\nu})^{\mu\mu}$ is approximate, because we have kept only the interference term with the Standard Model, which is large compared to the term from purely NP contributions.) The next biggest contribution comes from $\bar{\nu}_\mu\nu_\tau$ and $\bar{\nu}_\tau\nu_\mu$ final states. In these cases, there is no interference with the SM and the contribution is

$$\Delta(R_K^{(*)\nu\nu})^{\mu\tau} + \Delta(R_K^{(*)\nu\nu})^{\tau\mu} = 2.10 \times 10^{-3} (|c_{22}^*c_{33}|^2 + |c_{32}^*c_{23}|^2) (\epsilon_3^q)^4 \left(\frac{M}{\text{TeV}}\right)^{-4} \left(\frac{g_\rho}{4\pi}\right)^2. \tag{4.15}$$

As is clear from these equations, the most important contribution comes from the $\bar{\nu}_\tau\nu_\tau$ process. It is possible to pass the bound $\Delta(R_K^{(*)\nu\nu})^{\tau\tau} < 2.7$ in a large fraction of the parameter space. Furthermore, large deviations in $R_K^{\nu\nu}$ and $R_K^{*\nu\nu}$ ($\sim 25\%$ of the SM contribution) represent an interesting prediction of our composite leptoquarks scenario, which will be testable at the upcoming Belle II experiment [52, 53]. Our prediction can be compared with the case in which the leptoquark has only muonic couplings, in which the contributions to $\Delta(R_K^{(*)\nu\nu})$ are $\lesssim 5\%$ (see section 4.5 of [52]).

4.2.2 $K^+ \rightarrow \pi^+\nu\nu$

Given that measurements involving neutrinos have the ability to probe some of the largest couplings in our model — those involving third generation leptons — it is necessary to check other rare meson decays with final state neutrinos.

Following [54], (but rescaling the bound given there to match the slightly more recent measurement in [49]), the measurement of $\mathcal{B}(K^+ \rightarrow \pi^+\nu\nu)$ produces a bound (at 95% confidence level) on the real NP coefficient $\delta C_{\nu\bar{\nu}}$ (defined in [54]) of

$$\delta C_{\nu\bar{\nu}} \in [-6.3, 2.3]. \tag{4.16}$$

The branching ratio is given in terms of $\delta C_{\nu\bar{\nu}}$ by

$$\mathcal{B}(K^+ \rightarrow \pi^+\nu\nu) = 8.6(9) \times 10^{-11} [1 + 0.96\delta C_{\nu\bar{\nu}} + 0.24(\delta C_{\nu\bar{\nu}})^2]. \tag{4.17}$$

Our leptoquark contributes to $\delta C_{\nu\bar{\nu}}$ as

$$\delta C_{\nu\bar{\nu}} = 0.62 \operatorname{Re}(c_{31}c_{32}^*) \left(\frac{g_\rho}{4\pi}\right) (\epsilon_3^q)^2 \left(\frac{M}{\text{TeV}}\right)^{-2}, \tag{4.18}$$

via the dominant process involving a pair of tau neutrinos. So with $c_{31} \sim c_{32} \sim O(1)$, and $M \sim \text{TeV}$, our scenario passes current bounds.

However the NA62 experiment, due to begin data-taking in 2015, will measure $\mathcal{B}(K^+ \rightarrow \pi^+ \nu \nu)$ to an accuracy of 10% of the SM prediction [55]. This means it will be able to shrink the bounds on $\delta C_{\nu\bar{\nu}}$ to

$$\delta C_{\nu\bar{\nu}} \in [-0.2, 0.2] \tag{4.19}$$

at 95%. Thus, if $c_{31} \sim c_{32} \sim O(1)$ and $M \sim \text{TeV}$, measurements at NA62 will be sensitive to our leptoquark.

4.2.3 Meson mixing

The leptoquark we consider can mediate mixing between neutral mesons via box diagrams. This effect will be largest in B_s mesons. From [38], the bound produced on the leptoquark couplings when both leptons exchanged in the box are taus (the dominant contribution in our scenario) is

$$|\lambda_{33}\lambda_{32}^*|^2 < \frac{196\pi^2 M^2 \Delta m_{B_s^0}^{NP}}{f_{B_s^0}^2 m_{B_s^0}}. \tag{4.20}$$

From [56], $f_{B_s^0} = 0.231 \text{ GeV}$, and

$$\Delta m_{B_s^0}^{\text{SM}} = (17.3 \pm 2.6) \times 10^{12} \hbar s^{-1} = (1.14 \pm 0.17) \times 10^{-8} \text{ MeV}, \tag{4.21}$$

while from [49], the measured value of the mass splitting is

$$\Delta m_{B_s^0} = 17.69 \times 10^{12} \hbar s^{-1} = 1.2 \times 10^{-8} \text{ MeV}. \tag{4.22}$$

Taking the uncertainty in the prediction to be roughly the size of the NP contribution, $|\Delta m_{B_s^0}^{NP} / \Delta m_{B_s^0}^{\text{SM}}| < 0.15$ (as in [14]), then

$$|\lambda_{33}\lambda_{32}^*|^2 < 0.017 \left(\frac{M}{\text{TeV}} \right)^2. \tag{4.23}$$

In terms of the parameters of our model this becomes

$$|c_{33}c_{23}^*| < 4.2 \left(\frac{4\pi}{g_\rho} \right)^2 \left(\frac{M}{\text{TeV}} \right)^2 \left(\frac{1}{\epsilon_3^q} \right)^4. \tag{4.24}$$

We are able to pass this bound taking $O(1)$ values for c_{33} and c_{23} and taking the other parameters at values necessary to fit the anomalies as discussed above. The leptoquark will also contribute to mixing of other neutral mesons. However bounds from the measurement of mixing observables are generally weaker than bounds from meson decays (see eg. [57]).

4.2.4 $\mu \rightarrow e\gamma$ and other radiative processes

The leptoquark has only left handed couplings, meaning that we will not get chiral enhancements to the branching ratio of $\mu \rightarrow e\gamma$. Nevertheless, the bound on $\mathcal{B}(\mu \rightarrow e\gamma)$ is tight enough to be relevant for the model. The largest contributions come from diagrams with a loop containing either a top or a bottom quark, together with the leptoquark. The most recent measurement was performed by the MEG collaboration [58], who found a bound at

90% confidence level of $\mathcal{B}(\mu^+ \rightarrow e^+\gamma) < 5.7 \times 10^{-13}$. Using the formula for the rate given in [38], and neglecting all but the processes involving 3rd generation quarks in the loop,

$$|\lambda_{23}^* \lambda_{13}| < 7.3 \times 10^{-4} \left(\frac{M}{\text{TeV}} \right)^2, \quad (4.25)$$

which amounts to a bound on $c_{23}^* c_{13}$ of

$$|c_{23}^* c_{13}| < 1.4 \left(\frac{4\pi}{g_\rho} \right) \left(\frac{M}{\text{TeV}} \right)^2 \left(\frac{1}{\epsilon_3^q} \right)^2. \quad (4.26)$$

This turns out to be a strong constraint for our model. Given that our EFT paradigm assumes $c_{ij} \sim O(1)$, the bound is, roughly, saturated.

Given our flavour structure we expect an even larger contribution to $\tau \rightarrow \mu\gamma$ than to $\mu \rightarrow e\gamma$. However the current bound on the branching ratio of this process is $\mathcal{B}(\tau \rightarrow \mu\gamma) < 4.4 \times 10^{-8}$ [49], which is several orders of magnitude larger than the model prediction.

The process $b \rightarrow s\gamma$ can be generated via similar diagrams. Current bounds on this process, which leave room for NP contributions up to about 30% of the SM prediction, lead to a bound on the combination $|c_{33}^* c_{32}|$ of roughly $|c_{33}^* c_{32}| \lesssim 100 \left(\frac{4\pi}{g_\rho} \right) \left(\frac{M}{\text{TeV}} \right)^2 \left(\frac{1}{\epsilon_3^q} \right)^2$.

4.2.5 Comments on other constraints and predictions

Despite the fact that contributions from leptoquark diagrams will be largest for processes containing taus (or tau neutrinos) in the final state, we have not yet mentioned any bounds from meson decays with τ leptons in the final state. This is because existing bounds are very weak due to the relative difficulty of tau measurements. The current bound [59] on the decay $B \rightarrow K\tau^+\tau^-$ from BaBar, $\mathcal{B}(B \rightarrow K\tau^+\tau^-) < 3.3 \times 10^{-3}$, is several orders of magnitude larger than the NP prediction. Likewise the recent Belle measurement of $\mathcal{B}(B^+ \rightarrow \tau^+\nu)$ [60] has error bars much larger than the NP contribution (as does the SM prediction).

We have discussed $b \rightarrow s\ell\ell$ processes and anomalies in previous subsections. Bounds from meson decays mediated by other FCNC processes in the down sector are summarised in table 4. The most constraining of these measurements is from the branching ratio of $B^+ \rightarrow \pi^+\mu^+\mu^-$, for which the bound is approximately saturated.

Our leptoquark can appear in diagrams which contribute to the muon anomalous magnetic moment, an observable which currently has a 2.2–2.7 σ discrepancy with SM calculations [61]. However, as was pointed out in [62, 63], if a leptoquark couples only to one chirality of muon, as is the case for us, the couplings would need to be very large to explain the measurement. Our scenario produces a prediction several orders of magnitude too small (for a mass of $O(1 \text{ TeV})$), and so does nothing to alleviate the current tension between the SM and experiment.

One hallmark of our model is that there should be only very small NP effects in the electron sector. So decay measurements involving electrons should see no significant deviations from the Standard Model in our scenario. A recent paper [22] contains a table with predictions of ratios of observables with muons in the final state versus those with electrons for $b \rightarrow s\ell\ell$ processes. The predictions of our leptoquark model will, to a good

Decay	(ij)(kl)*	$ \lambda_{ij}\lambda_{kl}^* / (\frac{M}{\text{TeV}})^2$	$ c_{ij}c_{kl}^* (\frac{g_\rho}{4\pi}) (\epsilon_3^g)^2 / (\frac{M}{\text{TeV}})^2$
$K_S \rightarrow e^+e^-$	(12)(11)*	< 1.0	$< 4.9 \times 10^7$
$K_L \rightarrow e^+e^-$	(12)(11)*	$< 2.7 \times 10^{-3}$	$< 1.3 \times 10^5$
† $K_S \rightarrow \mu^+\mu^-$	(22)(21)*	$< 5.1 \times 10^{-3}$	$< 1.2 \times 10^3$
$K_L \rightarrow \mu^+\mu^-$	(22)(21)*	$< 3.6 \times 10^{-5}$	< 8.3
$K^+ \rightarrow \pi^+e^+e^-$	(11)(12)*	$< 6.7 \times 10^{-4}$	$< 3.3 \times 10^4$
$K_L \rightarrow \pi^0e^+e^-$	(11)(12)*	$< 1.6 \times 10^{-4}$	$< 7.8 \times 10^3$
$K^+ \rightarrow \pi^+\mu^+\mu^-$	(21)(22)*	$< 5.3 \times 10^{-3}$	$< 1.2 \times 10^3$
$K_L \rightarrow \pi^0\nu\bar{\nu}$	(31)(32)*	$< 3.2 \times 10^{-3}$	< 42.5
† $B_d \rightarrow \mu^+\mu^-$	(21)(23)*	$< 3.9 \times 10^{-3}$	< 46.0
$B_d \rightarrow \tau^+\tau^-$	(31)(33)*	< 0.67	$< 4.6 \times 10^2$
† $B^+ \rightarrow \pi^+e^+e^-$	(11)(13)*	$< 2.8 \times 10^{-4}$	$< 6.9 \times 10^2$
† $B^+ \rightarrow \pi^+\mu^+\mu^-$	(21)(23)*	$< 2.3 \times 10^{-4}$	< 2.7

Table 4. 90% confidence level bounds [57] on leptoquark couplings from branching ratios of (semi-)leptonic meson decays involving $b \rightarrow d$ and $s \rightarrow d$, rescaled to $M = 1$ TeV. A dagger denotes bounds that have been rescaled to newer measurements [49]. The final column gives bounds on partial compositeness parameters in units of the nominal values in (3.11).

approximation, coincide with those of the third column of their table, which contains the predictions for a scenario with NP only in $C_9^{\mu NP} = -C_{10}^{\mu NP} = -0.5$.

The leptoquark we consider will mediate lepton flavour violating processes. However we find that all current bounds are well above rates predicted for the leptoquark contribution. Lepton flavour violation in the context of B decays was recently discussed in detail in [51]. There, the authors consider a model in which, similarly to our case, the NP contributions to $b \rightarrow s\ell\ell$ decays arise in a $V - A$ structure (ie. $C_9^\ell = -C_{10}^\ell$) and the largest effects are in the third generation of quarks and leptons. Interestingly, a special case of our model can be made to fit into their framework, if we take all the $O(1)$ coefficients c_{ij} to be equal (and for simplicity, equal to 1). Then the coupling denoted G in [51] is given by $G = (g_\rho/M^2) (\epsilon_3^g)^2 \frac{m_\tau}{v}$, and the mixing matrices U_{L3i}^ℓ and U_{L3i}^d therein are given by $U_{L3}^d = (\lambda^3, \lambda^2, 1)$ and $U_{L3i}^\ell = \sqrt{m_i/m_\tau}$. With these choices, we find that all bounds quoted in [51] are comfortably satisfied by the composite leptoquark model. More precise bounds on LFV processes will certainly provide an interesting test of our model and other lepton non-universal scenarios.

Another recent paper [64] proposes double ratios of branching ratios as clean probes of NP that is not lepton universal and couples to right-handed quarks. Since the leptoquark we consider has no couplings to right-handed quarks, measurements of these would be a useful test of the model if the B anomalies persist.

4.3 Direct searches at the LHC

If the leptoquark is light enough, as the arguments in section 4.1 suggest it should be, it will be pair-produced at the LHC with sizable cross-section via QCD interactions. The leptoquark field comprises 3 charge eigenstates, $\Pi_{4/3}$, $\Pi_{1/3}$ and $\Pi_{-2/3}$, with charges $4/3$, $1/3$ and $-2/3$ respectively. Since we expect them to be rather heavier than the top, if we assume all c_{ij} coefficients in the couplings to have a modulus equal to 1, their branching ratio to third generation quarks and leptons is around 94% or greater. So they predominantly decay as follows:

$$\begin{aligned}\Pi_{4/3} &\rightarrow \bar{\tau} \bar{b}, \\ \Pi_{1/3} &\rightarrow \bar{\tau} \bar{t} \text{ or } \Pi_{1/3} \rightarrow \bar{\nu}_{\tau} \bar{b}, \\ \Pi_{-2/3} &\rightarrow \bar{\nu}_{\tau} \bar{t}.\end{aligned}$$

The branching ratios are quite sensitive to the c_{ij} coefficients, however, so other decay modes (e.g. involving second generation leptons) may be important for different values of c_{ij} , even if they are all still $O(1)$. The bounds and branching ratios in this section have been derived under the assumption that the modulus of all c_{ij} coefficients be equal to 1, but we will comment on the impact of lifting this assumption towards the end of the section.

There will be electroweak mass splittings between the three leptoquark states, allowing the heavier ones to decay to the lighter ones, but these decays will be subdominant to those through the leptoquark couplings, if the mass splittings are small. Of the LHC leptoquark searches, dedicated searches for third generation leptoquarks will put the strongest limits on our leptoquarks [65]. The $\Pi_{-2/3}$ leptoquark will decay to tops and missing energy, so stop searches, which look for the same signature, will apply. Likewise sbottom searches will apply to $\Pi_{1/3}$. A recent CMS search [66] ruled out leptoquarks decaying wholly to τ and b up to a mass of 740 GeV. This bound roughly applies to the leptoquark $\Pi_{4/3}$. This leptoquark's branching ratio to τ and b is 0.94 (over the mass range of the search, the variation is only in higher decimal places), so the bound on it from [66] is roughly 720 GeV. Another CMS search [67] puts bounds on leptoquarks decaying to either top and tau or bottom and neutrino with a combined branching ratio of 100%. Since the $\Pi_{1/3}$ state has a combined branching ratio of 97% to these final states, to a good approximation the results of this search should apply. This search implies a bound of 570 GeV on the mass of the $\Pi_{1/3}$, which at this mass has a branching ratio of 0.40 to top and tau. A bound from an ATLAS stop search [68] can be applied to the remaining leptoquark state, $\Pi_{-2/3}$. In one scenario considered in the search, the stop is presumed to decay wholly to a top and the lightest neutralino, and a 640 GeV bound on the mass of the stop is quoted, assuming that the neutralino is massless. The production mechanism for the $\Pi_{-2/3}$ leptoquark is identical to that for the stop, which is assumed in the search to be directly pair produced. Furthermore at a mass of 640 GeV, the branching ratio of $\Pi_{-2/3}$ to top and neutrino is greater than 99.5%. We can hence take the 640 GeV bound to apply directly to the mass of the leptoquark $\Pi_{-2/3}$. Since we are assuming small mass splittings between the charge eigenstates in the leptoquark multiplet, a bound on the mass of one eigenstate roughly

corresponds to a bound on them all. So we can apply the strongest of the bounds given above to the mass M ; we therefore conclude that $M > 720$ GeV.

However all these bounds are found by assuming that the $O(1)$ coefficients c_{ij} in the leptoquark couplings all have a modulus equal to 1. The limits can change quite a lot if this is not the case. In particular, the ratio of λ_{3j} to λ_{2j} is $(c_{3j}\sqrt{m_\tau})/(c_{2j}\sqrt{m_\mu})$, ie. $\sim 4c_{3j}/c_{2j}$. So the branching ratio to e.g. top and muon can be larger than that to top and tau if $c_{33}/c_{23} \lesssim 0.25$. By contrast, the difference between the third and second generations of quarks is harder to overcome by changes in the c_{ij} coefficients, since the hierarchy in the mixing parameters is larger. A measured bound on leptoquarks decaying to third generation quarks and second generation leptons would be a useful measurement to cover a case where c_{33} were accidentally small.

5 Conclusions

We have argued that current anomalies in semileptonic B decays seen at LHCb are consistent with a composite Higgs model featuring an additional, light leptoquark. This leptoquark has quantum numbers $(\bar{\mathbf{3}}, \mathbf{3}, 1/3)$ under the SM gauge group and couplings to the SM fermions that are largely fixed by the partial compositeness paradigm. We have identified a possible coset structure that contains both the SM Higgs and the leptoquark as pseudo-Goldstone bosons of the strong sector, which allows them to be rather lighter than the resonances of the strong sector, and with a natural explanation of the sizes and signs of their squared-mass parameters (bar a small, unavoidable, residual fine-tuning in the electroweak scale).

The partial compositeness framework automatically implies lepton non-universality in the leptoquark couplings. In this way, the departure from unity of the value of the ratio of the branching ratio of $B \rightarrow K\mu^+\mu^-$ to that of $B \rightarrow Ke^+e^-$ measured at LHCb this year can be accommodated, as can earlier anomalies in measurements of angular observables in $B \rightarrow K^*\mu^+\mu^-$ decays. The framework predicts large couplings to third-generation leptons, hence large deviations in observables involving tau leptons or neutrinos in the final state. These processes therefore provide a good check for the model, but the predictions are not in conflict with current bounds. In our scenario, with parameters chosen to fit the LHCb $b \rightarrow sll$ anomalies, we predict deviations of $\sim 25\%$ from the SM value in $\mathcal{B}(B \rightarrow K\nu\bar{\nu})$, which will be testable at the Belle II experiment. And the NA62 experiment, starting in 2015, will measure $\mathcal{B}(K^+ \rightarrow \pi^+\nu\bar{\nu})$ to sufficient accuracy that the effects of the model should be visible there.

We find that the model is consistent with all other flavour constraints, however we have identified a few processes for which the NP contributions are at or close to the current bounds. These are $\mathcal{B}(\mu \rightarrow e\gamma)$, the mass splitting Δm_{B_s} in B_s meson mixing, and $\mathcal{B}(B^+ \rightarrow \pi^+\mu^+\mu^-)$. More precise measurements of these will also test the model. It should be remembered, however, that the leptoquark couplings, λ_{ij} , within partial compositeness are each only predicted up to an $O(1)$ factor, c_{ij} . Thus predictions can only be made to an accuracy of order one or so, and even tight constraints could be evaded if, for a particular process, the combination of c_{ij} factors involved is accidentally small. It should be noted

that, in particular, none of the processes listed in this paragraph get their dominant contributions from the same combination of couplings that are involved in the LHCb B decay anomalies. Thus, the various $O(1)$ factors are not determined by fitting the anomalies. However, the fact that the framework can make predictions for a wide range of processes, due to non-zero couplings to all SM fermions, means that it is, nevertheless, falsifiable.

If the composite leptoquark is the cause of the measured discrepancies in B decays, there are three implications for the model. Firstly, the composite sector must be strongly interacting, $g_\rho \sim 4\pi$. Secondly, the left handed doublet of the third quark generation must be highly composite $\epsilon_3^q \sim 1$. Thirdly, the leptoquark should have a mass of around a TeV, meaning that there is scope for its discovery at LHC13. Current LHC8 bounds, analysed in section 4.3, exclude the leptoquark up to masses of $M > 720$ GeV (under an assumption on the coefficients involved in the couplings). Large third generation couplings ensure that the three charge eigenstates of the leptoquark triplet decay mostly to third generation quarks and leptons, so searches for third generation leptoquarks are effective for constraining their mass.

We finally comment on the physics associated with the strong sector at the scale m_ρ . First of all let us estimate the value of that scale. According to the discussion in section 3 we expect that $M^2 \sim \frac{\alpha_s}{4\pi} m_\rho^2$. Given that we need $M \sim 1$ TeV to explain the LHCb anomalies, we obtain $m_\rho \sim 10$ TeV. With such a scale for the composite sector it has been shown [69] that the structure of partial compositeness is enough to suppress dangerous contributions to indirect search observables with the exception of the electron EDM and the radiative LFV decay $\mu \rightarrow e\gamma$.¹⁴ The further suppression required in these channels might be obtained by departing from the hypothesis of lepton flavour anarchy in the strong sector.¹⁵ Very roughly, the amount of tuning needed to accommodate the right values of the EW scale and of the Higgs mass [77–81] is expected to be, at best, at or below the per cent level, which is not much worse than the amount already required in generic supersymmetric extensions of the SM, given current bounds. We feel that this is a not unreasonable price to pay, given the additional benefits of a motivated flavour paradigm and the power to explain the LHCb anomalies.

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¹⁴Other observables close to the current experimental sensitivity are the neutron EDM, $\epsilon_K, \epsilon'/\epsilon$ and $B \rightarrow X_s \gamma$.

¹⁵Various non-anarchic flavour structures have been considered in the context of partial compositeness, see for example [42, 70–76]. In our framework a possible solution is to invoke an unbroken $U(1)_e$ symmetry protecting the electron family number. This will lead to vanishing leptoquark couplings with the electron field, but our phenomenological analysis is unaffected.

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