Exploring leptoquark effects in hyperon and kaon decays with missing energy

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Abstract

We entertain the possibility that scalar leptoquarks (LQs) generate consequential effects on the strangeness-changing decays of hyperons and kaons involving missing energy carried away by a pair of invisible fermions. Although such decays have suppressed rates in the standard model (SM), they could get much enhancement in the presence of the LQs. In order to respect the available data on the kaon modes $K \to \pi \nu \nu$ and at the same time enlarge the rates of the hyperon decays significantly, at least two different scalar LQs are needed. If the LQs have Yukawa couplings solely to SM fermions, we find that the hyperon rates cannot be sufficiently increased to levels within the reach of ongoing or near-future experiments because of the combined constraints from the measurements on kaon mixing and lepton-flavor-violating processes. However, if light right-handed neutrinos are included in the LQ interactions, their contributions can evade the leading restrictions and lead to hyperon rates which may be big enough to be probed by upcoming searches.
I. INTRODUCTION

The flavor-changing neutral current (FCNC) decays of strange hadrons involving missing energy are of great importance because they are known to be dominated by short-distance physics. Therefore these decays offer a good access not only to the underlying weak dynamics of the standard model (SM) but also to manifestations of possible new physics (NP) beyond it. Such processes arise mainly from the quark transition $s \rightarrow d \bar{E}$ which in the SM is induced by loop diagrams [1], with the missing energy ($\bar{E}$) being carried away by an undetected neutrino pair ($\nu \bar{\nu}$). This entails that in the SM the branching fractions of the corresponding hadron decays are very small, which makes them good places to look for NP signals. In the presence of NP there could be additional ingredients which modify the SM contribution and/or yield extra channels with one or more invisible nonstandard particles. These changes may bring about sizable effects discoverable by running or forthcoming searches.

To date the most intensive quests for $s \rightarrow d \bar{E}$ have focused on the kaon modes $K \rightarrow \pi \nu \bar{\nu}$ and produced upper limits on their branching fractions [2–5] which are fairly close to the SM expectations [6]. This implies that the room for possible NP in $K \rightarrow \pi \nu \bar{\nu}$ can no longer be considerable. Nevertheless, as shown in Refs. [7, 8], the impact of NP could still be substantial via operators having parity-odd $ds$ quark bilinears. This is because such operators do not contribute to $K \rightarrow \pi \nu \bar{\nu}$, hence not restrained by their data, and are instead subject to constraints from the channels with no or two pions, namely $K \rightarrow \bar{E}$ and $K \rightarrow \pi \pi' \bar{E}$, which currently have relatively weak empirical bounds [5, 9–11]. Thus, future measurements to improve upon the data on the latter two sets of kaon channels would be highly desirable.

The $s \rightarrow d \bar{E}$ operators influencing all of the aforementioned kaon modes affect their baryon counterparts as well and consequently may also be testable in strangeness-changing ($|\Delta S| = 1$) hyperon decays with missing energy [7, 8], on which no data exist. It is therefore of keen interest that there has recently been a proposal to look for them in the BESIII experiment [12]. As demonstrated in Refs. [7, 8], the existing constraints from $K \rightarrow \bar{E}$ and $K \rightarrow \pi \pi' \bar{E}$ are sufficiently loose to allow the branching fractions of these hyperon modes to rise far above their SM predictions to values within the proposed sensitivity reach of BESIII [12]. Its upcoming search for them may then be expected to come up with informative results regarding potential NP effects on the underlying quark transition.

The foregoing motivates us in this paper to consider these hyperon decays with missing energy in the contexts of relatively simple NP models. This also complements the earlier model-independent analyses in which the missing energy was assumed to be carried away by an invisible pair of spin-$1/2$ fermions [7] or spinless bosons [8]. Within a specific model the parameters determining the strength of the $s \rightarrow d \bar{E}$ operators often contribute to other observables, some of which may be well constrained by their respective data. In addressing this, one could thus learn how the different restrictions could probe various aspects of a model and what modifications to it, if still possible,
may need to be made to get around the constraints. Moreover, this exercise could offer rough
guidance about what kind of NP might be responsible if certain signatures beyond the SM appear
at some level in the quests for these hyperon modes.

The organization of the rest of the paper is the following. In Sec. II we introduce the new
particles in the theory, namely scalar LQs and right-handed SM-singlet fermions, and describe
their interactions with SM fermions. In Sec. III we derive the amplitudes for the hyperon and
kaon decay modes of interest and write down the expressions for their rates. In Sec. IV we deal
with the relevant constraints on the Yukawa couplings of the LQs and present our numerical
results. To illustrate the potential impact of the scalar LQs on these hyperon and kaon processes,
we discuss two simple scenarios. In the first one, only SM fermions participate in the Yukawa
interactions of the LQs. In the second scenario, we let the LQs couple directly to the new SM-
singlet fermions as well as to SM quarks, but not to SM leptons. In each of these two cases, two
different scalar LQs are necessary to respect the existing $K \rightarrow \pi \nu \nu$ data and simultaneously raise
the hyperon rates significantly.¹ In Sec. V we provide our conclusions.

II. INTERACTIONS

Among leptoquarks which can have renormalizable couplings to SM fermions without violating
baryon- and lepton-number conservation and SM gauge symmetries, there are three which have
spin 0 and can give rise at tree level to flavor-changing transitions with missing energy among down-
type quarks (as reviewed in Ref. [15]). We denote these LQs, with their $SU(3)_c \times SU(2)_L \times U(1)_Y$
assignments, as $\tilde{S}_1$ ($\bar{3}, 1, 1/3$), $\tilde{S}_2$ ($3, 2, 1/6$), and $S_3$ ($\bar{3}, 3, 1/3$). In this study we concentrate on the
contributions of the first two. In terms of their components,

$$\tilde{S}_1 = \tilde{S}_1^{1/3}, \quad \tilde{S}_2 = \left( \frac{\tilde{S}_2^{2/3}}{\tilde{S}_2^{-1/3}} \right),$$

where the superscripts indicate their electric charges. If we add to the theory three right-handed
fermions $N_{1,2,3}$ which are singlet under the SM gauge groups and hereafter referred to as right-
handed neutrinos, they can also couple to the LQs and SM quarks.

We write the Lagrangian for the renormalizable interactions of the LQs with SM fermions plus
the right-handed neutrinos as

$$\mathcal{L}_{\text{LQ}} = \left( \gamma_{1,jy}^{\text{ll}} \overline{l}_y \tilde{S}_1 \tilde{l}_j + \tilde{\gamma}_{1,jy}^{\text{rr}} \overline{l}_j \tilde{S}_1 \tilde{l}_y \right) \tilde{S}_1 + \gamma_{2,jy}^{\text{ll}} \overline{d}_j \tilde{S}_2 \tilde{e}_y \tilde{l}_j + \tilde{\gamma}_{2,jy}^{\text{rr}} \overline{d}_j \tilde{S}_2 \tilde{l}_y \tilde{N}_y + \text{H.c.},$$

where the $\gamma_{jy}$ and $\tilde{\gamma}_{jy}$ are generally complex elements of the LQ Yukawa coupling matrices,
summation over family indices $j, y = 1, 2, 3$ is implicit, $q_j$ ($l_y$) and $d_j$ represent a left-handed

¹ Invoking two scalar LQs to suppress certain processes and enhance others has previously been applied in the
context of $B$-meson decays, such as by [13, 14].
quark (lepton) doublet and right-handed down-type quark singlet, respectively, and $\varepsilon = i\tau_2$ with $\tau_2$ being the second Pauli matrix. To treat decay amplitudes, we need to express $\mathcal{L}_{\text{LQ}}$ in terms of mass eigenstates. For the processes of interest, it is convenient to do so in the mass basis of the down-type fermions, in which case

$$
q_j = \begin{pmatrix} V^*_{ij} U_{iL}^j \\ D_{jL} \end{pmatrix}, \quad l_j = \begin{pmatrix} U_{jL} \nu_i \\ \ell_{jL} \end{pmatrix}, \quad d_j = D_{jR},
$$

where $i = 1, 2, 3$ is implicitly summed over, $V \equiv V^{\text{CKM}}$ ($U \equiv U^{\text{PMNS}}$) is the Cabibbo-Kobayashi-Maskawa quark (Pontecorvo-Maki-Nakagawa-Sakata neutrino) mixing matrix, $f_{L,R} = \frac{1}{2}(1 \pm \gamma_5)f$, and $U_{1,2,3}(= u, c, t)$, $D_{1,2,3}(= d, s, b)$, $\nu_{1,2,3}$, and $\ell_{1,2,3}(= e, \mu, \tau)$ refer to the mass eigenstates. Since the left-handed neutrinos’ masses are tiny, we can work with $\nu_j = U_{jL} \nu_i$ instead of $\nu_j$. Thus, Eq. (2) becomes

$$
\mathcal{L}_{\text{LQ}} = \left\{ Y_{1,jy}^L \left[ Y_{ij}^L (U_{iL})^c \ell_{jL} - (D_{jL})^c \nu_{jy} \right] + Y_{1,jy}^R (D_{jR})^c N_y \right\} \tilde{S}_1
$$

$$
+ Y_{2,jy}^R \tilde{D}_{jR} \left[ \ell_{jL} \tilde{s}_{S_2}^{2/3} - \nu_{jy} \tilde{s}_{S_2}^{-1/3} \right] + Y_{2,jy}^L \left[ \nu_{ij} \tilde{U}_{iL} \tilde{s}_{S_2}^{2/3} + \tilde{D}_{jL} \tilde{s}_{S_2}^{-1/3} \right] N_y
$$

+ H.c.

(4)

We take the left- and right-handed neutrinos to be all Dirac fermions. Furthermore, we suppose that the latter have masses which can be neglected for the processes of concern and are sufficiently long-lived that they do not decay inside detectors.

The scalar LQs can mediate $d_s f f'$ interactions at tree level, with $f = \nu$ or $N$. With both of the LQs being heavy, $\mathcal{L}_{\text{LQ}}$ leads to the low-energy effective Lagrangian

$$
-\mathcal{L}_{d_s f f'} = \overline{d}_y \gamma_5 S \overline{\nu}_y \left( C^0_{f_f} + \gamma_5 C^A_{f_f} \right) f' + \overline{d}_y \gamma_5 S \overline{\nu}_y \left( \tilde{C}^0_{f_f} + \gamma_5 \tilde{C}^A_{f_f} \right) f' + \text{H.c.},
$$

the pertinent coefficients being

$$
C^0_{f_f} = -C^A_{f_f} = \frac{-Y_{1,1x}^L Y_{1,2y}^L}{8m_2^2 S_1} + \frac{Y_{2,1x}^L Y_{2,2x}^L}{8m_2^2 S_2},
$$

$$
\tilde{C}^0_{f_f} = -\tilde{C}^A_{f_f} = \frac{Y_{1,1x}^L Y_{1,1y}^L}{8m_2^2 S_1} + \frac{Y_{2,1x}^L Y_{2,2x}^L}{8m_2^2 S_2},
$$

$$
C^0_{NN'} = C^A_{NN'} = \frac{-Y_{1,1x}^R Y_{1,2y}^R}{8m_2^2 S_1} + \frac{Y_{2,1x}^R Y_{2,2x}^R}{8m_2^2 S_2},
$$

$$
\tilde{C}^0_{NN'} = \tilde{C}^A_{NN'} = \frac{-Y_{1,1x}^L Y_{1,2y}^L}{8m_2^2 S_1} - \frac{Y_{2,1x}^L Y_{2,2x}^L}{8m_2^2 S_2},
$$

(6)

where $x$ and $y$ stand for lepton-family indices.

III. AMPLITUDES AND RATES

In the hyperon sector the interactions in Eq. (3) induce the decay modes $\mathfrak{B} \to \mathfrak{B}' f f'$ involving the pairs of baryons $\mathfrak{B} \mathfrak{B'} = \Lambda n, \Sigma^+ p, \Xi^0 \Lambda, \Xi^0 \Sigma^0, \Xi^- \Sigma^-$, all having spin 1/2, and $\Omega^- \to \Xi^- f f'$,
where Ω− has spin 3/2. To deal with the amplitudes for these decays, we need the baryonic matrix elements of the quark bilinears in Eq. (5) which can be estimated with the aid of chiral perturbation theory at leading order [7, 16]. For \( \mathcal{B} \to \mathcal{B}' f \bar{f}' \) the results are [7]

\[
\langle \mathcal{B}' | \bar{d} \gamma^\eta s | \mathcal{B} \rangle = \mathcal{V}_{\mathcal{B} \mathcal{B}'} \bar{u}_{\mathcal{B}'} \gamma^\eta u_{\mathcal{B}}, \quad \langle \mathcal{B}' | \bar{d} \gamma^\eta \gamma_5 s | \mathcal{B} \rangle = \mathcal{A}_{\mathcal{B} \mathcal{B}'} \bar{u}_{\mathcal{B}'} \left( \gamma^\eta - \frac{m_{\mathcal{B}'}}{m_K - Q^2} \gamma_5 \right) \gamma_5 u_{\mathcal{B}},
\]

where \( \mathcal{V}_{\mathcal{B} \mathcal{B}'} = -3/\sqrt{6}, -1, 3/\sqrt{6}, -1/\sqrt{2} \) and \( \mathcal{A}_{\mathcal{B} \mathcal{B}'} = -(D + 3F)/\sqrt{6}, D - F, (3F - D)/\sqrt{6}, -(D + F)/\sqrt{2}, D + F \) for \( \mathcal{B} \mathcal{B}' = \Lambda n, \Sigma^+ p, \Xi_0^0 \Lambda, \Xi_0^0 \Sigma^0, \Xi^- \Sigma^- \), respectively, \( \bar{u}_{\mathcal{B}'} \) and \( u_{\mathcal{B}} \) represent the Dirac spinors of the baryons, \( Q = p_{\mathcal{B}'} - p_{\mathcal{B}} \), with \( p_{\mathcal{B}} \) and \( p_{\mathcal{B}'} \) being their four-momenta, and \( m_{\mathcal{B}} \) and \( m_{\mathcal{B}'} \) in Eq. (7) refer to the isospin-averaged masses of \( \mathcal{B}^{(i)} \) and the kaons. For \( \Omega^- \to \Xi^- f \bar{f}' \) we have [7]

\[
\langle \Xi^- | \bar{d} \gamma^\eta \gamma_5 s | \Omega^- \rangle = C \bar{u}_\Xi \left( \frac{\tilde{Q}^\eta \tilde{Q}_\epsilon}{m_K^2 - Q^2} u_\Xi \right),
\]

where \( \tilde{Q} = p_{\Omega^-} - p_{\Xi^-} \) and \( u_\Xi^\eta \) is a Rarita-Schwinger spinor. The constants \( D, F \), and \( C \) above occur in the lowest-order chiral Lagrangian and can be fixed from baryon data. In numerical work, we will adopt the same values of these and other input parameters as those given in Ref. [7].

With Eqs. (7) and (8), we derive the amplitudes for \( \mathcal{B} \to \mathcal{B}' f \bar{f}' \) and \( \Omega^- \to \Xi^- f \bar{f}' \). Since in the scenarios under consideration the \( f \) and \( f' \) masses are taken to be small compared to the pion mass, we then arrive at the differential rates [7]

\[
\frac{d\Gamma_{\mathcal{B} \to \mathcal{B}' f \bar{f}'}}{ds} = \frac{\lambda_{\mathcal{B} \mathcal{B}'}}{32\pi^3 m_{\mathcal{B}}^3} \left\{ \frac{\lambda_{\mathcal{B} \mathcal{B}'}}{3} + (m_{\mathcal{B}} - m_{\mathcal{B}'})^2 \left[ c_{f f'}^\eta \right]^2 \mathcal{V}_{\mathcal{B} \mathcal{B}'}^2 + \left[ \lambda_{\mathcal{B} \mathcal{B}'} \right] + (m_{\mathcal{B}} + m_{\mathcal{B}'})^2 \left[ c_{f f'}^\eta \right]^2 A_{\mathcal{B} \mathcal{B}'}^2 \right\},
\]

\[
\frac{d\Gamma_{\Omega^- \to \Xi^- f \bar{f}'}}{ds} = \frac{\lambda_{\Omega^- \Xi^-}}{384\pi^3 m_{\Omega^-}^3} \frac{C_{\Xi^-}}{3m_{\Omega^-}^2} \left( \frac{\lambda_{\Omega^- \Xi^-}}{3} + 4 \left[ (m_{\Omega^-} + m_{\Xi^-})^2 - \hat{s} \right] \left[ c_{f f'}^\eta \right]^2 \right),
\]

where \( \hat{s} = (p_f + p_{f'})^2 \) and \( \lambda_{XY} = m_X^2 - 2(m_Y^2 + \hat{s})m_X + (m_Y^2 - \hat{s})^2 \). Here the coefficients \( c_{f f'}^\eta \) are the LQ-mediated ones in Eq. (3). Following Ref. [7], in our numerical treatment of the rates we will include form-factor effects not yet taken into account in Eqs. (7) and (8). Particularly, in Eq. (9) we apply the changes \( \mathcal{V}_{\mathcal{B} \mathcal{B}'} \to (1 + 2\hat{s}/M_V^2)\mathcal{V}_{\mathcal{B} \mathcal{B}'} \) and \( \mathcal{A}_{\mathcal{B} \mathcal{B}'} \to (1 + 2\hat{s}/M_A^2)\mathcal{A}_{\mathcal{B} \mathcal{B}'} \) with \( M_V = 0.97(4) \) GeV and \( M_A = 1.25(15) \) GeV, in line with the parametrization commonly employed in experimental analyses of hyperon semileptonic decays [17, 21]. Furthermore, in Eq. (10) we modify \( C \) to \( C/(1 - \hat{s}/M_A^2) \).

In the kaon sector, \( \mathcal{L}_{\text{d} f f'} \) in Eq. (5) is consequential for \( K \to \pi f f' \) and \( K \to \pi \pi f f' \) but not for \( K \to f f' \) which is helicity suppressed because we assume \( m_{f f'} \simeq 0 \) in our cases of concern.\(^2\)

\(^2\) Many other LQ scenarios with consequences for various kaon processes have been explored before [22, 32].
It follows that the relevant mesonic matrix elements are

\[
\langle \pi^-(p_\pi)|\bar{d}\gamma^\eta s|K^-(p_K)\rangle = p_K^0 + p_\pi^\eta,
\]
\[
\langle \pi^0(p_0)\pi^-(p_-)|\bar{d}\gamma^\eta\gamma_5 s|K^0\rangle = \frac{i\sqrt{2}}{f_K} \left((p_0^\eta - p_-^\eta) + \frac{(p_0^\mu - p_-^\mu)\bar{q}_\mu\bar{q}^\eta}{m_K^2 - \bar{q}^2}\right),
\]
\[
\langle \pi^0(p_1)\pi^0(p_2)|\bar{d}\gamma^\eta\gamma_5 s|K^0\rangle = \frac{i}{f_K} \left((p_1^\eta + p_2^\eta) + \frac{(p_1^\mu + p_2^\mu)\bar{q}_\mu\bar{q}^\eta}{m_K^2 - \bar{q}^2}\right),
\]
(11)

where \( f_K = 155.6(4) \) MeV \[^3\] is the kaon decay constant, \( \bar{q} = p_K^- - p_0^- - p_- = p_K^0 - p_1^- - p_2^- \), and we have ignored form-factor effects in these matrix elements. Assuming isospin symmetry and making use of charge conjugation, we also have \( \langle \pi^0|d\gamma^\eta s|K^0\rangle = -\langle \pi^-|d\gamma^\eta s|K^-\rangle/\sqrt{2} = -\langle \pi^0|\bar{s}\gamma^\eta d|K^0\rangle \) and \( \langle \pi^0\pi^-|\bar{s}\gamma^\eta\gamma_5 d|K^0\rangle = (\pi^0\pi^-|\bar{d}\gamma^\eta\gamma_5 s|K^0) \). Hence the amplitudes for \( K^- \to \pi^- (\pi^0) f \bar{f}' \) and \( K_L \to \pi^0 (\pi^0) f \bar{f}' \) are \([3, 31]\)

\[
\mathcal{M}_{K^-\to\pi^- f \bar{f}'} = 2 \bar{u}_f \gamma_K [C^{\pi, A}_{f f'}(\gamma_5 C^{\pi, A}_{f f'})] v_{f'}, \quad \mathcal{M}_{K_L\to\pi^0 f \bar{f}'} = \bar{u}_f \gamma_K [C^{\pi, A}_{f f'} - \gamma_5 (C^{\pi, A}_{f f'} - C^{\pi, A}_{f f'})] v_{f'},
\]
\[
\mathcal{M}_{K^-\to\pi^0 f \bar{f}'} = \frac{i\sqrt{2}}{f_K} \bar{u}_f [\bar{q}_0^\eta - \bar{q}_-^\eta] (\tilde{c}_f^{\pi, A} + \gamma_5 \tilde{c}_f^{\pi, A}) v_{f'},
\]
\[
\mathcal{M}_{K_L\to\pi^0 f \bar{f}'} = \frac{i}{\sqrt{2}f_K} \bar{u}_f [\bar{q}_1^\eta + \bar{q}_2^\eta] (\tilde{c}_f^{\pi, A} + \gamma_5 \tilde{c}_f^{\pi, A}) v_{f'},
\]
(12)

in the \( m_{f, f'} = 0 \) limit. If \( f' \neq f \), we additionally have \( \mathcal{M}_{K_L\to\pi^0 f \tilde{f}'} \) and \( \mathcal{M}_{K_L\to\pi^0 f \bar{f}'} \) which equal \( \mathcal{M}_{K_L\to\pi^0 f \bar{f}'} \) and \( \mathcal{M}_{K_L\to\pi^0 f \tilde{f}'} \), respectively, but with \( f \) and \( f' \) interchanged. If \( C^{\pi, A}_{f f'} \) and \( C^{\pi, A}_{f f'} \) with \( f' \neq f \) are nonzero, they are generally independent of \( C^{\pi, A}_{f f'} \) and \( C^{\pi, A}_{f f'} \) and induce \( K^- \to \pi^- (\pi^0) f \bar{f}' \) as well as \( \Omega^- \to \Xi^+ f \bar{f}' \) and \( \Omega^- \to \Xi^- f \bar{f}' \).

Since, as mentioned in Sec. [3], possible NP in \( C^{\pi, A}_{f f'} \) is not expected to be considerable,\(^3\) hereafter we suppose that it has no influence on these coefficients altogether and instead enters via \( \tilde{c}_{f f'}^{\pi, A} \) alone. Accordingly, among the kaon modes we need to examine the NP impact on only the four-body ones. From the last two amplitudes in Eq. (12), we arrive at their double-differential decay rates

\[
\frac{d^2\Gamma_{K^-\to\pi^0 f \bar{f}'}(s, \bar{\zeta})}{ds d\bar{\zeta}} = \frac{\beta^3 \beta_\zeta^{1/2}}{9(4\pi)^5 f_K^2 m_K^3} \left(3\beta_\zeta + 12\bar{s}\bar{\zeta}\right) |\tilde{c}_f^{\pi, A}|^2,
\]
\[
\frac{d^2\Gamma_{K_L\to\pi^0 f \bar{f}'}(s, \bar{\zeta})}{ds d\bar{\zeta}} = \frac{\beta^3 \beta_\zeta^{1/2}}{24(4\pi)^5 f_K^2 m_K^3} |\tilde{c}_f^{\pi, A} + \tilde{c}_{\bar{f}'}^{\pi, A}|^2 = \frac{d^2\Gamma_{K_L\to\pi^0 f \bar{f}'}(s, \bar{\zeta})}{ds d\bar{\zeta}},
\]
(13)

where we have set \( m_{f, f'} = 0 \) and taken Eq. (6) into account,

\[
\beta_\zeta = \sqrt{1 - \frac{4m_\pi^2}{s}}, \quad \bar{\zeta} = (p_0 + p_-)^2 = (p_1 + p_2)^2, \quad \beta_\bar{\zeta} = \left(m_K^2 - s - \bar{\zeta}\right)^2 - 4\bar{s}\bar{\zeta}. \quad (14)
\]

\(^3\) Studies on the implications of NP affecting exclusively \( C_{nf, nf}^{\pi, A} \) for the hyperon modes can be found in \([33, 34]\).
The $\hat{s}$ and $\hat{\zeta}$ integration ranges are $0 \leq \hat{s} \leq (m_{K^-} - m_0 - 2m_\pi)^2$ and $4m_\pi^2 \leq \hat{\zeta} \leq (m_{K^-} - m_0 - \hat{s}^{1/2})^2$ for the $K^-$ and $K_L$ channels, respectively.

IV. CONSTRAINTS AND NUMERICAL RESULTS

For definiteness and simplicity, we look at a couple of distinct cases as examples. In the first one (A) the right-handed neutrinos, $N_{1,2,3}$, are absent. In the second scenario (B) they are present but the LQs do not couple directly to SM leptons, $Y_1^{\mu\mu} = Y_2^{\nu\nu} = 0$. We will treat the nonzero Yukawa couplings phenomenologically, allowing one or more of them to vanish or differ substantially from the others, in order to avoid various constraints.

A. SM fermions only

Choosing $c_{\nu\nu'}^V = 0$ in Eq. (6) implies

\[
\hat{c}_{\nu\nu'}^V = -c_{\nu\nu'}^A = \frac{Y_{1,1,2}^{\mu\mu} Y_{1,2,1}^{\nu\nu}}{4m_{S_1}^2} = \frac{Y_{2,2,1}^{\mu\mu} Y_{2,1,2}^{\nu\nu}}{4m_{S_2}^2},
\]

Consequently, as Eqs. (9), (10), and (12) indicate, the $K \rightarrow \pi E$ decays are now due to the SM alone, whereas $K \rightarrow \pi\pi' E$ and the hyperon modes still receive the LQ contributions. In the absence of $N_{1,2,3}$, we have set $c_{\nu\nu'}^V = c_{\nu\nu'}^A = 0$.

The interactions in Eq. (4) can affect the mixing of neutral kaons $K^0$ and $\bar{K}^0$ via four-quark operators arising from box diagrams, with LQs and SM leptons running around the loops, and given in the effective Hamiltonian $[6, 13, 32]$

\[
H_{\Delta S=2}^{LQ} = \left(\sum_x \frac{Y_{1,1,2}^{\mu\mu} Y_{1,2,1}^{\nu\nu}}{128\pi^2 m_{S_1}^2}\right)^2 \bar{s}_L \gamma_5 d_L \bar{d}_L \gamma_5 \gamma_5 d_L + \left(\sum_x \frac{Y_{2,2,1}^{\mu\mu} Y_{2,1,2}^{\nu\nu}}{64\pi^2 m_{S_2}^2}\right)^2 \bar{s}_R \gamma_5 d_R \bar{d}_R \gamma_5 \gamma_5 d_R + \text{H.c.}
\]

(16)

Its matrix element between the $K^0$ and $\bar{K}^0$ states contributes to the kaon-mixing parameters $\Delta m_K \propto \text{Re} M_{KK}$ and $\epsilon_K \propto \text{Im} M_{KK}$, where $M_{KK} = \langle \bar{K}^0 | H_{\Delta S=2} | K^0 \rangle / (2m_{K^0})$. Given that $\Delta m_K$ and $\epsilon_K$ have been well measured, the constraints from them on the LQ contributions are stringent [32]. Nevertheless, we see from Eq. (16) that these restrictions can be evaded by assigning the nonzero elements of the first and second rows of each contributing Yukawa matrix to separate columns.

The interactions in Eq. (4) also give rise to lepton-flavor-violating $ds\ell\ell'$ operators in

\[
-L_{ds\ell\ell'} = \bar{d} \gamma_5 \gamma_5 (V_{\ell\ell'} + \gamma_5 A_{\ell\ell'}) \ell' + \bar{d} \gamma_5 \gamma_5 \gamma_5 (\bar{V}_{\ell\ell'} + \gamma_5 \bar{A}_{\ell\ell'}) \ell' + \text{H.c.},
\]

(17)

where from tree-level LQ-exchange diagrams $[13, 31, 32]$

\[
V_{\ell\ell'} = -A_{\ell\ell'} = \bar{V}_{\ell\ell'} = -\bar{A}_{\ell\ell'} = \frac{Y_{2,2,1}^{\mu\mu} Y_{2,1,2}^{\nu\nu}}{8m_{S_2}^2}.
\]

(18)
Accordingly, if the pair of the lepton-family indices of $\mathbf{y}_{1,2}^{\nu L} \mathbf{y}_{1,2}^{\nu L}$ has the values $xy = 12, 21$, corresponding to $\ell\ell' = e\mu, \mu\mu$, the $\mathcal{A}^\gamma s$ and $\mathcal{A}^\gamma s$ terms in Eq. (17), plus their H.c., will induce $K \rightarrow \pi^+ \mu^+$ and $K_L \rightarrow e^+ \mu^+$, respectively, both of which have severe experimental restraints. They can be eluded by ensuring that $xy$ is neither 12 nor 21. These choices also imply that one-loop LQ contributions to the $\mu \rightarrow e\gamma, 3e$ decays and $\mu \rightarrow e$ transitions in nuclei also escape their search limits.

We can write down a sample set of Yukawa matrices that fulfill these requisites:

$$
\mathbf{Y}_1^{LL} = \begin{pmatrix}
0 & 0 & y_{1,13} \\
0 & y_{1,22} & 0 \\
0 & 0 & 0 \\
\end{pmatrix}, \quad \mathbf{Y}_2^{RL} = \begin{pmatrix}
y_{2,12} & 0 & 0 \\
0 & 0 & y_{2,23} \\
0 & 0 & 0 \\
\end{pmatrix}.
$$

They lead to

$$
\mathbf{V}_{\nu\nu'} = -\mathbf{V}_{\nu'\nu} = \frac{y_{1,13}^* y_{1,22}}{4m_1^2} = \frac{y_{2,23}^* y_{2,12}}{4m_2^2},
$$

$$
\mathbf{V}_{\tau\mu} = -\mathbf{V}_{\tau'\mu} = -\mathbf{A}_{\tau\mu} = \frac{y_{2,23}^* y_{2,12}}{8m_2^2},
$$

the other coefficients in $\mathcal{L}_{d\ell\ell'}$ and $\mathcal{L}_{d\ell\ell'}$ in Eqs. (5) and (17) being zero. Alternatively, we could pick one of three other sets, with the corresponding products of nonzero matrix elements being given by $(y_{1,12}^* y_{1,23}, y_{2,22}^* y_{2,13})$, $(y_{1,11}^* y_{1,23}, y_{2,21}^* y_{2,13})$, and $(y_{1,13}^* y_{1,21}, y_{2,23}^* y_{2,11})$, respectively.

If $xy = 13, 31, 23, 32$ in Eq. (18), the operators in Eq. (17) generate the lepton-flavor-violating $|\Delta S| = 1$ tau decays $\tau \rightarrow \ell K_s, \ell K_s^*, \ell \pi^K$ with $\ell = e, \mu$, and so their empirical bounds cannot be avoided by the selections of the Yukawa couplings made in the previous paragraph but are much milder than the aforementioned restrictions on $e\mu$ violation. Numerically, among the $\tau$ decay constraints, the one on $\tau^- \rightarrow \mu^- K^{0*}$ turns out to be the least demanding on the coefficients [33] and applies to the combinations already displayed in Eq. (21). Using the results of Ref. [35], we find that this translates into

$$
\frac{|y_{1,13}^* y_{1,22}|}{m_1^2} = \frac{|y_{2,23}^* y_{2,12}|}{m_2^2} < \frac{0.036}{\text{TeV}^2}.
$$

We can now put this together with the branching fractions of the kaon decays $K \rightarrow \pi^0 \bar{f}^f$ and the hyperon ones, $\mathcal{B} \rightarrow \mathcal{B}^f \bar{f}^f$ and $\Omega^- \rightarrow \Xi^- \bar{f}^f$. Thus, employing the rate formulas from Sec. III we arrive at

\begin{align}
B(K^- \rightarrow \pi^0 \bar{\nu}_\tau \bar{\nu}_\mu) & < 1.0 \times 10^{-10}, & B(K_L \rightarrow \pi^0 \nu_\tau \bar{\nu}_\mu) & < 6.9 \times 10^{-10}, \\
B(\Lambda \rightarrow n \nu_\tau \bar{\nu}_\mu) & < 1.1 \times 10^{-8}, & B(\Sigma^+ \rightarrow p \nu_\tau \bar{\nu}_\mu) & < 2.9 \times 10^{-9}, \\
B(\Xi^0 \rightarrow \Lambda \nu_\tau \bar{\nu}_\mu) & < 1.6 \times 10^{-10}, & B(\Xi^0 \rightarrow \Sigma^0 \nu_\tau \bar{\nu}_\mu) & < 2.2 \times 10^{-9}, \\
B(\Xi^0 \rightarrow \Sigma^0 \nu_\tau \bar{\nu}_\mu) & < 2.7 \times 10^{-9}, & B(\Omega^- \rightarrow \Xi^- \nu_\tau \bar{\nu}_\mu) & < 1.3 \times 10^{-7},
\end{align}
where \( \mathcal{B}(K_L \to \pi^0 \pi^0 \nu_\tau \nu_\mu) \) is due to \( K_L \to \pi^0 \pi^0 \nu_\mu \bar{\nu}_\tau \) and \( K_L \to \pi^0 \pi^0 \nu_\tau \bar{\nu}_\mu \) which have the same rate according to the second line of Eq. (13).

The numbers in Eq. (23) are far bigger than their SM counterparts, \( \mathcal{B}(K^- \to \pi^0 \pi^- \nu \bar{\nu})_{\text{SM}} \sim 10^{-14} \) and \( \mathcal{B}(K_L \to \pi^0 \pi^0 \nu \bar{\nu})_{\text{SM}} \sim 10^{-13} \) \cite{36,38}, but still lie significantly below the existing measured bounds \( \mathcal{B}(K^- \to \pi^0 \pi^- \nu \bar{\nu})_{\exp} < 4.3 \times 10^{-5} \) \cite{9} and \( \mathcal{B}(K_L \to \pi^0 \pi^0 \nu \bar{\nu})_{\exp} < 8.1 \times 10^{-7} \) \cite{10}, both at 90% confidence level. Likewise, the hyperon results in Eq. (24) greatly exceed their SM counterparts, within the \( 10^{-13} \text{-} 10^{-11} \) range \cite{7}, but are not yet close to the corresponding sensitivity levels of BESIII estimated in Ref. \cite{12} for the branching fractions of \( \Lambda \to n \nu \bar{\nu}, \Sigma^+ \to p \nu \bar{\nu}, \Xi^0 \to \Lambda \nu \bar{\nu}, \Xi^0 \to \Sigma^0 \nu \bar{\nu}, \) and \( \Omega^- \to \Xi^- \nu \bar{\nu}, \) which are \( 3 \times 10^{-7}, 4 \times 10^{-7}, 8 \times 10^{-7}, 9 \times 10^{-7}, \) and \( 2.6 \times 10^{-5}, \) respectively.

One could try to get around instead the \( \tau \) decay constraints by arranging the nonvanishing elements of the Yukawa matrices to be on, say, the third columns and let the elements be subject to the kaon-mixing restrictions. This option turns out to restrain the couplings more strongly than Eq. (22) and hence produces comparatively smaller kaon and hyperon rates. Another possibility would be to enlist the third scalar LQ mentioned earlier, \( S_3 \) \((3,3,1/3)\), but we have found that it would not improve on the situation treated in the last paragraph.

The above exercise then suggests that with only SM fermions participating in the Yukawa interactions of the scalar LQs it would be unlikely for the \(|\Delta S| = 1\) hyperon decays with missing energy to have rates sizable enough to be within the reach of BESIII. It may therefore require future facilities such as super charm-tau factories to test this particular LQ scenario.

B. Including right-handed neutrinos

If \( N_{1,2,3} \) are present and have the couplings specified in Eq. (1), then upon setting \( c_{N_3} = 0 \) in Eq. (6) we get

\[
\tilde{c}_{\nu \nu'} = \tilde{c}_{\nu' \nu} = \frac{-\hat{Y}_{1,1x}^{LR} \hat{Y}_{1,2y}^{RR}}{4m_S^2} = \frac{-\hat{Y}_{1,2x}^{LR} \hat{Y}_{2,1y}^{RR}}{4m_S^2}.
\] (25)

If, in addition, the LQs do not interact at tree level with SM leptons, then \( Y_{1,1}^{LL} = Y_{2,2}^{LL} = 0 \), which leads to \( \tilde{c}_{\nu \nu'} = \tilde{c}_{\nu' \nu} = 0 \).

Similarly to the preceding case, the Yukawa couplings in Eq. (25) affect \( K^0 \bar{K}^0 \) mixing through box diagrams, with the LQs and right-handed neutrinos going around the loops, giving rise to the effective Hamiltonian

\[
\tilde{H}_{|\Delta S| = 2}^{LQ} = \frac{\left(\sum_{x} \bar{Y}_{2,2x}^{LL} \hat{Y}_{2,1x}^{LR*}\right)^2}{128\pi^2 m_S^2} \bar{s}_L \gamma^0 \gamma_5 d_L \bar{s}_L \gamma_\eta d_L + \frac{\left(\sum_{x} \bar{Y}_{1,1x}^{LR} \hat{Y}_{1,2x}^{RR*}\right)^2}{128\pi^2 m_S^2} \bar{s}_R \gamma^0 \gamma_5 d_R \bar{s}_R \gamma_\eta d_R + \text{H.c.}
\] (26)

Evidently, the limitations from kaon-mixing data can again be eluded by assigning the nonzero elements of the first and second rows of the \( \hat{Y} \) matrices to separate columns. As for processes which do not conserve charged-lepton flavor, their data no longer yield pertinent constraints on the \( \hat{Y} \) matrix elements because the LQs now do not interact directly with SM leptons.
The main restrictions would then come from the empirical bounds on $K \to \pi\pi'\bar{E}$. To illustrate their impact, we consider this sample set of Yukawa matrices that satisfy the above requirements:

$$\tilde{Y}_1^{RR} = \begin{pmatrix} 0 & \tilde{y}_{1,12} & 0 \\ 0 & 0 & \tilde{y}_{1,23} \\ 0 & 0 & 0 \end{pmatrix}, \quad \tilde{Y}_2^{LR} = \begin{pmatrix} 0 & 0 & \tilde{y}_{2,13} \\ 0 & \tilde{y}_{2,22} & 0 \\ 0 & 0 & 0 \end{pmatrix}. \quad (27)$$

They give

$$\tilde{c}_{\bar{N}_2 N_3}^\nu = \tilde{c}_{\bar{N}_2 N_3}^A = \frac{-\tilde{y}_{1,12}^* \tilde{y}_{1,23}}{4m_{\tilde{S}_1}^2} = \frac{-\tilde{y}_{2,22}^* \tilde{y}_{2,13}}{4m_{\tilde{S}_2}^2}, \quad (28)$$

the other $\tilde{c}_{ff'}^\nu$ vanishing.\(^4\) Since the $\mathcal{B}(K_L \to \pi^0\pi^0\nu\bar{\nu})_{\text{exp}}$ bound mentioned in the last subsection is stronger than the $\mathcal{B}(K^- \to \pi^0\pi^-\nu\bar{\nu})_{\text{exp}}$ one, we can impose $\mathcal{B}(K_L \to \pi^0\pi^0\bar{E})_{\text{LQ}} < 8 \times 10^{-7}$, where $\mathcal{B}(K_L \to \pi^0\pi^0\bar{E})_{\text{LQ}} = \mathcal{B}(K_L \to \pi^0\pi^0N_2\bar{N}_3) + \mathcal{B}(K_L \to \pi^0\pi^0N_3\bar{N}_2) = 2\mathcal{B}(K_L \to \pi^0\pi^0N_2\bar{N}_3)$. This translates into $|\tilde{c}_{\bar{N}_2 N_3}^\nu|^2 < 9.4 \times 10^{-14}$ GeV$^{-4}$ and consequently

$$\left|\frac{\tilde{y}_{1,12}^* \tilde{y}_{1,23}}{m_{\tilde{S}_1}^2}\right| = \left|\frac{\tilde{y}_{2,22}^* \tilde{y}_{2,13}}{m_{\tilde{S}_2}^2}\right| < \frac{1.2}{\text{TeV}^2}. \quad (29)$$

Incorporating this into the hyperon decay rates from Sec. \[III\] we obtain the maximal branching fractions

$$\mathcal{B}(\Lambda \to nN_2\bar{N}_3) < 1.3 \times 10^{-5}, \quad \mathcal{B}(\Sigma^+ \to pN_2\bar{N}_3) < 3.5 \times 10^{-6},$$

$$\mathcal{B}(\Xi^0 \to \Lambda N_2\bar{N}_3) < 1.9 \times 10^{-6}, \quad \mathcal{B}(\Xi^0 \to \Sigma^0N_2\bar{N}_3) < 2.6 \times 10^{-6},$$

$$\mathcal{B}(\Xi^- \to \Sigma^-N_2\bar{N}_3) < 3.2 \times 10^{-6}, \quad \mathcal{B}(\Omega^- \to \Xi^-N_2\bar{N}_3) < 1.5 \times 10^{-4}. \quad (30)$$

Their upper values well exceed the corresponding estimated BESIII sensitivity levels quoted in the previous subsection. It follows that BESIII might uncover NP clues in these processes or, if not, come up with improved restrictions on the Yukawa interactions of the LQs with the right-handed neutrinos.

This example and the preceding one indicate that experiments on these hyperon transitions can serve as a valuable tool to discriminate NP models. Furthermore, the acquired data could supply information on potential NP contributions which is complementary to that gained from the kaon sector, including restraints on them which may be stricter than those provided by kaon measurements.

V. CONCLUSIONS

We have explored the effects of scalar leptoquarks on $dsff'$ interactions that involve light invisible spin-1/2 fermions, $\mathbf{f}$ and $\mathbf{f}'$, and induce the FCNC decays of strange hadrons with

\(^4\) As before, there are alternatives, such as those with the corresponding products of nonzero matrix elements being given by $(\tilde{y}_{1,13} \tilde{y}_{1,22}, \tilde{y}_{2,23} \tilde{y}_{2,12}), (\tilde{y}_{1,11} \tilde{y}_{1,23}, \tilde{y}_{2,21} \tilde{y}_{2,13})$, and $(\tilde{y}_{1,13} \tilde{y}_{1,21}, \tilde{y}_{2,23} \tilde{y}_{2,11})$. 

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missing energy, $\mathcal{E}$, carried away by the fermions. We concentrate on the case in which the LQ impact is insignificant on the kaon mode $K \rightarrow \pi \mathcal{E}$ but can be substantial on the hyperon decays $\mathcal{B} \rightarrow \mathcal{B}' \mathcal{E}$ and $\Omega^- \rightarrow \Xi^- \mathcal{E}$, as well as on $K \rightarrow \pi \pi' \mathcal{E}$. This can occur because $K \rightarrow \pi \mathcal{E}$ are sensitive exclusively to $dsff'$ operators with parity-even quark bilinears, while the hyperon modes are affected not only by this kind of operators but also by those with parity-odd quark bilinears.

We show that this possibility can be realized in a model containing scalar LQs by employing two of them which have different chiral couplings to the $d$ and $s$ quarks. If only SM fermions participate in the Yukawa interactions of the LQs, we find that the combination of restrictions from kaon mixing and lepton-flavor-violating processes does not permit the hyperon rates to reach values likely to be detectable in the near future. In contrast, with the addition of light SM-singlet right-handed neutrinos which also couple directly to the LQs, extra $s \rightarrow d \mathcal{E}$ channels can arise, with the right-handed neutrinos being emitted invisibly, and moreover the constraints from lepton-flavor-violation data can be evaded. As a consequence, the resulting hyperon rates are allowed to increase to levels which may be discoverable by the ongoing BESIII or future efforts such as at super charm-tau factories. Our analysis based on simple NP scenarios illustrates the importance of experimental quests for these hyperon modes, as the outcomes would complement the results of measurements on their kaon counterparts and could help distinguish NP models.

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5 It is worth noting that considerable branching fractions of the hyperon modes $\mathcal{B} \rightarrow \mathcal{B}' \mathcal{E}$ and $\Omega^- \rightarrow \Xi^- \mathcal{E}$ are also attainable if the missing energy is carried away by a single massless dark photon [39].
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