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Phenomenology of TeV-scale scalar leptoquarks in the EFT

Shaouly Bar-Shalom,[1,*] Jonathan Cohen,[1,†] Amarjit Soni,[2,‡] and Jose Wudka[3,§]
[1]*Physics Department, Technion-Institute of Technology, Haifa 32000, Israel*
[2]*Physics Department, Brookhaven National Laboratory, Upton, New York 11973, USA*
[3]*Physics Department, University of California, Riverside, California 92521, USA*

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We examine new aspects of leptoquark (LQ) phenomenology using effective field theory (EFT). We construct a complete set of leading effective operators involving SU(2) singlets scalar LQ and the Standard Model fields up to dimension six. We show that, while the renormalizable LQ-lepton-quark interaction Lagrangian can address the persistent hints for physics beyond the Standard Model in the B-decays $\bar{B} \to D^{(*)}\tau\bar{\nu}$, $\bar{B} \to \bar{K}\ell^+\ell^-$ and in the measured anomalous magnetic moment of the muon, the LQ higher dimensional effective operators may lead to new interesting effects associated with lepton number violation. These include the generation of one-loop and two-loops sub-eV Majorana neutrino masses, mediation of neutrinoless double-β decay and novel LQ collider signals. For the latter, we focus on third generation LQ (ϕ_3) in a framework with an approximate Z_3 generation symmetry and show that one class of the dimension five LQ operators may give rise to a striking asymmetric same-charge $\phi_3\phi_3$ pair-production signal, which leads to low background same-sign leptons signals at the LHC. For example, with $M_{\phi_3} \sim 1$ TeV and a new physics scale of $\Lambda \sim 5$ TeV, we expect at the 13 TeV LHC with an integrated luminosity of 300 fb^{-1}, about 5000 positively charged $\tau^+\tau^+$ events via $\boldsymbol{pp \to \phi_3\phi_3 \to \tau^+\tau^+ + 2 \cdot j_b}$ ($j_b = b$-jet), about 500 negatively charged $\tau^-\tau^-$ events with a signature $\boldsymbol{pp \to \phi_3\phi_3 \to \tau^-\tau^- + 4 \cdot j + 2 \cdot j_b}$ ($j =$ light jet) and about 50 positively charged $\ell^+\ell^+$ events via $\boldsymbol{pp \to \ell^+\ell^+ + 2 \cdot j_b + \not{E}_T}$ for any of the three charged leptons, $\ell^+\ell^+ = e^+e^+, \mu^+\mu^+, \tau^+\tau^+$. It is interesting to note that, in the LQ EFT framework, the expected same-sign lepton signals have a rate which is several times larger than the QCD LQ-mediated opposite-sign leptons signals, $gg, q\bar{q} \to \phi_3\phi_3^* \to \ell^+\ell^- + X$. We also consider the same-sign charged lepton signals in the LQ EFT framework at higher energy hadron colliders such as a 27 TeV HE-LHC and a 100 TeV FCC-hh.

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I. INTRODUCTION

The electroweak (EW) and strong interactions of the Standard Model (SM) have been very successfully tested at the low-energy (GeV-scale) and high-energy (EW-scale) frontiers as well as in precision measurements [1]. However, despite the impressive success of the SM at sub-TeV energies, it is widely believed that it is an effective low-energy framework of a more complete UV theory that should address the experimental and theoretical indications for new physics beyond the SM (BSM), such as the indirect detection of dark matter and dark energy, the measurements of neutrino masses, the flavor and hierarchy problems residing in the SM's scalar sector and the long sought higher symmetry which unifies the fundamental forces.

The scale of the new physics (NP) that may shed light on these fundamental questions in particle physics and address the deficiencies of the SM might be beyond the reach of present and future high-energy colliders. Nonetheless, the underlying UV theory may contain new particles with masses spanning over many orders of magnitudes, similar to the hierarchical mass pattern observed in nature and embedded in the SM. Indeed, although direct searches at high-energy colliders have not yet led to a discovery of new heavy particles, there have been intriguing and persistent hints in the past several years in favor of new TeV-scale degrees of freedom from measured anomalies associated with possible violations of lepton universality in B-decays: $\bar{B} \to D^{(*)}\tau\bar{\nu}$ [2–4] and $\bar{B} \to \bar{K}\ell^+\ell^-$ [5], as well as in the anomalous magnetic moment of the muon [6].

Out of these three anomalies, the most striking and least expected is the anomalous enhanced $\bar{B} \to D^{(*)}\tau\bar{\nu}$ rate measured by *BABAR* [2], Belle [3] and LHCb [4] (a $\sim 4\sigma$ effect). In the SM this decay occurs at tree-level and is

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[*] shaouly@physics.technion.ac.il
[†] jcohen@tx.technion.ac.il
[‡] adlersoni@gmail.com
[§] jose.wudka@ucr.edu

mediated by the Wcb charged current coupling, so that the measured deviation requires a relatively large tree-level NP contribution near the TeV scale to compete with the "classic" SM tree-level diagram. Promising candidates that address this large effect in $\bar{B} \to D^{(*)}\tau\bar{\nu}$ are TeV-scale leptoquarks (LQ's); in addition to this phenomenological role, these particles also appear naturally in theories that address some of the most fundamental questions in particle physics (see [7] and references therein) such as grand unification [8] and compositeness [9], where they can also arise as pseudo-Nambu Goldstone bosons [10] and lead to interesting collider signals [11,12]. They are also involved in models for neutrino masses [13]. In some cases, the effects of scalar LQ are similar to that of the scalar partners of the quarks in R-parity violating supersymmetry models [14,15], which can have similar couplings to quark-lepton pairs.

Given their theoretical appeal, and their potential role in addressing the B anomalies, it is of interest to study LQ phenomenology within the context of BSM physics. That is, allowing for the presence of excitations heavier than the LQs. This we shall do using an effective field theory, which will include the LQs as (relatively) low-energy excitations, and the effective interactions generated by heavier physics of scale Λ. Indeed, the mere presence of the TeV-scale renormalizable LQ framework (e.g., its Yukawa-like couplings to a quark-lepton pair which is being used in order to address the B-anomalies) suggests that this EFT higher-dimensional expansion is well defined and can be constructed in principle to any order (see Sec. IV). If the NP scale Λ is much higher than the multi TeV-scale, i.e., $\Lambda \gg 10$ TeV, then the effects of these higher-dimensional LQ effective interactions will be negligible. On the other hand, the purpose of this paper is to investigate the extent to which the LHC can probe physics beyond the LQ mass; we will see that, if $\Lambda \sim 5$–15 TeV, then the higher dimensional LQ effective interactions can produce unique collider signatures that may be observable at the LHC, and, in some cases, at rates that are *higher* than for the usual channels. We will also see that the physics at scale Λ, responsible for the effective LQ interactions, is also intimately connected with various possible mechanism of neutrino mass generation, so that a study of LQ phenomenology at the LHC can provide also information about the neutrino sector.

In this work we will concentrate on the study of the interactions and phenomenology of TeV-scale scalar LQs, which are SU(2) singlets and transform either as a right-handed down-type quark,[1] $\phi(3,1,-\frac{1}{3})$, or as a right-handed up-type quark, $\phi(3,1,\frac{2}{3})$, under the SM gauge group; since the BSM effects of both types of LQ have similar characteristics, in the bulk of the paper we will explore the effects and underlying physics of the down-type LQ, and towards the end of the paper we will shortly address the underlying physics and effects that are expected for an up-type LQ.

We construct the complete set of effective operators up to dimension six that involve the LQs and SM fields, and use this LQ EFT framework to demonstrate the impact of heavy physics on ϕ collider phenomenology, and on low-energy lepton number violating (LNV) phenomena such as Majorana neutrino masses and neutrinoless double beta decay. This model-independent formalism provides a broader and a more reliable view of the expected physics associated with TeV-scale LQs, and lays the ground for further investigations of ϕ-related phenomenology at high-energy colliders. For example, we find that the higher dimensional LQ interactions in the EFT framework may lead to very interesting, essentially background free, same-sign lepton signals at the LHC and/or at future colliders.

Lastly, we want to stress that while our starting motivation for this work was the B-anomalies, the confirmation of the anomalies is not needed for our work to have merit. Indeed, as was mentioned above, leptoquarks dynamics may be linked to well motivated extensions of the SM, such as composite theories and R-parity violating supersymmetry and, in particular, they play an important role in grand unified theories.

The paper is organized as follows: in the following section we summarize the effects of the renormalizable LQ interaction Lagrangian $\mathcal{L}_{\phi\mathrm{SM}}$; in Sec. III we review the LHC phenomenology of the scalar LQ in the ϕSM framework, and in Sec. IV we construct the effective theory beyond $\mathcal{L}_{\phi\mathrm{SM}}$, listing all the higher-dimensional effective operators involving the down-type LQ $\phi(3,1,-\frac{1}{3})$ up to dimension six. In Sec. V we study the $\Delta L = 2$ low-energy effects associated with the dimension five operators, and in Sec. VI we explore the leading signals of the down-type and up-type LQ, $\phi(3,1,-\frac{1}{3})$ and $\phi(3,1,\frac{2}{3})$, in the EFT framework at the 13 TeV LHC as well as at higher energy (27 and 100 TeV) hadron colliders. In Sec. VII we summarize and in the Appendix we list all dimension six operators for the down-type LQ.

II. RENORMALIZABLE LQ INTERACTIONS

We define the renormalizable extension of the SM which contains the LQ as

$$\mathcal{L}_{\phi\mathrm{SM}} = \mathcal{L}_{\mathrm{SM}} + \mathcal{L}_{Y,\phi} + \mathcal{L}_{H,\phi}, \tag{1}$$

where, for the down-type LQ $\phi(3,1,-\frac{1}{3})$, the Yukawa-like and scalar interaction pieces are

$$\mathcal{L}_{Y,\phi} = y^L_{q\ell}\bar{q}^c i\tau_2\ell\phi^* + y^R_{ue}\bar{u}^c e\phi^* + y^L_{qq}\bar{q}^c i\tau_2 q\phi + y^R_{ud}\bar{u}^c d\phi + \mathrm{H.c.}, \tag{2}$$

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[1] In our notation $X(c,w,y)$, indicates that particle X transforms under SU(3) representation c, SU(2) dimension w and carries hypercharge y.

$$\mathcal{L}_{H,\phi} = |D_\mu \phi|^2 - M_\phi^2 |\phi|^2 + \lambda_\phi |\phi|^4 + \lambda_{\phi H} |\phi|^2 |H|^2, \quad (3)$$

with q and ℓ the SU(2) left-handed quark and lepton doublets, respectively, while u, d, e are the right-handed SU(2) singlets; also, $\psi^c = C\bar{\psi}^T$.

A few comments are in order regarding the ϕSM Lagrangian defined in Eqs. (1)–(3),

(i) The last two Yukawa-like ϕ-quark-quark terms of $\mathcal{L}_{Y,\phi}$ in Eq. (2) violate the baryon number and can potentially mediate proton decay (see e.g., [16]). The Yukawa-like LQ couplings involving the first and second generations are then either vanishingly small [i.e., $(y^L_{qq})_{ij}, (y^R_{ud})_{ij} \to 0$ for $i, j \neq 3$] or are forbidden, e.g., by means of a symmetry.

(ii) The first two Yukawa-like ϕ-quark-lepton terms of $\mathcal{L}_{Y,\phi}$ in Eq. (2) (i.e., $\propto y^L_{q\ell}, y^R_{ue}$) can address the enhanced rate measured in the tree-level $\bar{B} \to D^{(*)}\tau\bar{\nu}$ decay as well as the 1-loop anomalies observed in $\bar{B} \to \bar{K}\ell^+\ell^-$ and the muon magnetic moment [14,17–21], when $M_\phi \sim \mathcal{O}(1)$ TeV and couplings $y^L_{q\ell}, y^R_{ue} \sim \mathcal{O}(0.1-1)$. It should be noted, though, that these down-type LQ ϕ-quark-lepton interactions are not sufficient for a simultaneous explanation of all these anomalies [22–28].

(iii) The LQ—Higgs interaction term $\propto \lambda_{\phi H}$ in Eq. (3) may play an important role in stabilizing the EW vacuum [29].

(iv) As will be discussed below, within the renormalizable ϕSM framework, $\mathcal{L}_{\phi\mathrm{SM}}$, LQ phenomenology and leading signals at the LHC are completely determined by the two Yukawa-like parameters $y^L_{q\ell}, y^R_{ue}$ and the LQ mass M_ϕ (ignoring the baryon number violating couplings).

III. PHENOMENOLOGY OF SCALAR LEPTOQUARKS IN THE ϕSM FRAMEWORK

In the limit $y^L_{q\ell}, y^R_{ue} \to 0$ the only production channels of a scalar LQ at the LHC are the tree-level QCD $\phi\phi^*$ pair-production via $gg \to \phi\phi^*$ and the s-channel gluon exchange in $q\bar{q}$-fusion $q\bar{q} \to \phi\phi^*$; see e.g., [30–38]. The corresponding typical $\phi\phi^*$ pair-production cross section at the 13 TeV LHC is $\sigma_{\phi\phi^*} \sim 5(0.01)$ fb for $M_\phi \sim 1(2)$ TeV [36]. Turning on the Yukawa-like ϕ-quark-lepton interactions in Eq. (2) adds another tree-level t-channel lepton exchange diagram to $q\bar{q} \to \phi\phi^*$, which, however, is subdominant. Thus, LQ pair-production at the LHC is essentially independent of its Yukawa-like couplings to a quark-lepton pair.

On the other hand, with sizable $y^L_{q\ell}, y^R_{ue}$ Yukawa terms, the LQ ϕ can also be singly produced at tree-level by the quark-gluon fusion processes $qg \to \phi\ell$; for $\phi = (3, 1, -\frac{1}{3})$ there are two production channels $ug \to \phi\ell_i$ and $dg \to \phi\nu_i$, where $i = 1, 2, 3$ is a generation index and both channels include two diagrams: an s-channel q-exchange and t-channel ϕ-exchange. The single LQ production channel is in fact dominant if ϕ has $\mathcal{O}(1)$ Yukawa-like couplings to the first generation quarks: $\sigma_\phi^{\mathrm{single}} = \sigma(qg \to \phi\ell) \propto y^2_{q\ell}$ (here $q = u$, d and $\ell = e, \nu_e$), and with $y_{q\ell} \sim \mathcal{O}(1)$ one obtains $\sigma_\phi^{\mathrm{single}}(pp_{(ug)} \to \phi e) \sim 100(2)$ fb and $\sigma_\phi^{\mathrm{single}}(pp_{(dg)} \to \phi\nu_e) \sim 50(0.5)$ fb for $M_\phi = 1(2)$ TeV; see e.g., [36].

The search for LQ is then performed assuming two distinct LQ decay channels that correspond to its two Yukawa-like interactions in the ϕSM: $\phi \to e_i j$ and $\phi \to \nu j$, with $\Gamma(\phi \to e_i j/\nu j) \sim |y|^2 m_\phi/16\pi$, where y is the corresponding ϕ-lepton-quark coupling, and the quark and lepton masses are neglected. Thus, the overall LQ signatures at the LHC contain either two leptons and two jets with large transverse momentum, $e_i^+ e_j^- jj$ and/or $e_i jj +$ missingE_T, when the LQ are pair-produced [39–44,46], or two leptons and a jet with large transverse momentum, $e_i^+ e_j^- j$ and $e_i j +$ missingE_T, when the LQ is singly produced.

Indeed, searches for first and second generations LQ pair-production (i.e. for LQ with couplings only to quark-lepton pairs of the first and second generations) yield stronger bounds than the ones for third generation LQ, since the detector sensitivity to the different flavors of high-p_T leptons and quarks varies. In addition, these bounds strongly depend on the LQ decay pattern, i.e., branching ratios to the different quark-lepton pairs. For example, the current bounds on the mass of a first (second) generation LQ assuming $pp \to \phi\phi^* \to e^+e^-/\mu^+\mu^- + jj$ and $\mathrm{BR}(\phi \to e/\mu + j) \sim 1$ is $M_\phi \gtrsim 1.5$ TeV [41,44].

Third generation LQ are particularly motivated, due to their potential role in explaining the observed anomalies in B-physics discussed above, but also on more general aspects concerning the underlying UV physics, e.g., the dynamical generation of fermion masses in composite scenarios [45]. Recent searches for a pair-produced third generation scalar LQ, decaying via $\phi \to t\tau, b\nu_\tau$ and/or $\phi \to b\tau$, have yielded weaker bounds: $M_\phi \gtrsim 1$ TeV [39,40,42,43,46,47]. On the other hand, the bound on the mass of a $\phi(3, 1, -\frac{1}{3})$ that couples exclusively to a top-muon pair (and can, therefore, address the anomalous muon magnetic moment and the anomaly measured in $\bar{B} \to \bar{K}\ell^+\ell^-$), obtained in the search for $pp \to \phi\phi^* \to t\bar{t}\mu^+\mu^-$, is $M_\phi \gtrsim 1.4$ TeV [43], i.e., comparable to the lower limit on the mass of a first and second generation LQ. Furthermore, a search for a singly produced third generation scalar LQ which decays exclusively via $\phi \to b\tau$ has also been performed recently by CMS; they exclude such a LQ up to a mass of 740 GeV [48].

Finally, another important LQ-mediated signal is the t-channel LQ exchange in the Drell-Yan lepton pair-production process $q\bar{q} \to \ell^+\ell^-$. In particular, this channel becomes important in the large LQ-lepton-quark coupling regime, since the corresponding cross section scales as $\sigma(q\bar{q} \to \ell^+\ell^-) \propto y_{q\ell}^4$, thus providing a complimentary sensitivity to the LQ dynamics as the LHC [7,37,38,49–52]; in particular yielding better access to larger LQ masses where the QCD on shell LQ pair production channel is suppressed.

IV. EFT BEYOND THE ϕSM FRAMEWORK

In this section we focus on the EFT extension of the renormalizable Lagrangian in Eqs. (1)–(3), for the down-type LQ $\phi(3,1,-\frac{1}{3})$. The effects of the NP which underlies the ϕSM framework in Eqs. (1)–(3) can be parametrized by a series of effective operators O_i, which are constructed using the ϕSM fields and whose coefficients are suppressed by inverse powers of the NP scale Λ,

$$\mathcal{L} = \mathcal{L}_{\phi\mathrm{SM}} + \sum_{n=5}^{\infty} \frac{1}{\Lambda^{n-4}} \sum_i f_i O_i^{(n)}, \tag{4}$$

where n is the mass dimension of $O_i^{(n)}$ and we assume decoupling and weakly coupled heavy NP, so that n equals the canonical dimension. The dominating NP effects are then expected to be generated by contributing operators with the lowest dimension (n value) that can be generated at tree-level in the underlying theory.

Before listing the specific form of the higher dimension operators, $O_i^{(n)}$, it is useful to denote their generic structure in the form,

$$O_i^{(n)} \in \phi^a H^b \psi^c D^d, \tag{5}$$

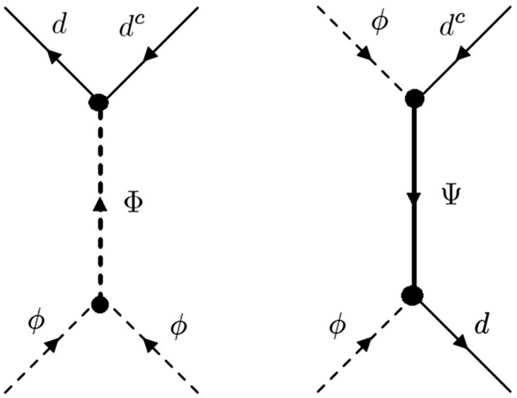
where a, b, c, d are integers representing the multiplicity of the corresponding factors: $O_i^{(n)}$ contains a LQ fields ϕ or ϕ^*, b Higgs fields H or $\tilde{H}$, c fermionic fields ψ and d covariant derivatives D. Group contractions and which fields are acted on by the derivatives are not specified.

We find that there are only two possible dimension-five operators involving the LQ $\Phi(3,1,-\frac{1}{3})$ and the SM fields —both violating lepton number by two units. To see that, note that the dimension-five operators with $c = 0$ in Eq. (5) are all absent because of gauge invariance. Furthermore, operators of the form $\phi^2\psi^2$ must contain the fermion bilinear $\bar{\psi}_L\psi_R$, so that only a single gauge invariant dimension five operator of this form survives [with two possible SU(3) color contractions which are not specified],

$$O_{d^2\phi^2}^{(5)} = \bar{d}d^c\phi^2, \tag{6}$$

which violates lepton number by two units.

The diagrams that can generate the dimension five operator $\bar{d}d^c\phi^2$ at tree-level in the underlying heavy theory are depicted in Fig. 1; the corresponding heavy NP must contain a heavy scalar $\Phi(6,1,-\frac{2}{3})$ and/or the heavy fermions $\Psi(1,1,0)$, $\Psi(8,1,0)$.



(a) (b)

FIG. 1. Tree-level graphs in the underlying heavy theory that generate the dimension five effective operator $\bar{d}d^c\phi^2$. Φ and Ψ stand for a heavy scalar and heavy fermion, respectively, with quantum numbers $\Phi(6,1,-\frac{2}{3})$ and $\Psi(1,1,0)$ or $\Psi(8,1,0)$ (see text).

Dimension five operators of the class $\phi\psi^2 D$ can be shown to be equivalent to operators without a derivative using integration by parts and, therefore, can be ignored. Thus, the remaining class of dimension five operators is of the form $\phi\psi^2 H$ and, therefore, must also contain the fermion bilinear $\bar{\psi}_L\psi_R$. The only gauge invariant operator of this form, which also violates lepton number by two units is

$$O_{\ell d\phi H}^{(5)} = \bar{\ell}d\tilde{H}\phi^*. \tag{7}$$

The heavy physics generating this operator at tree-level must contain a heavy scalar $\Phi(3,2,\frac{1}{6})$ and/or the heavy fermions $\Psi(1,1,0)$, $\psi(1,3,0)$ or $\Psi(3,2,-\frac{5}{6})$; see Fig. 2.

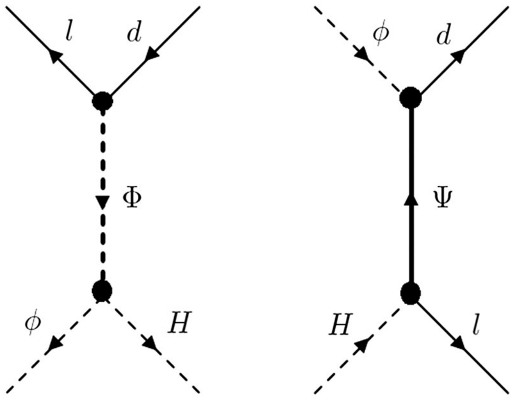
We recall that there is also a unique dimension five operator that can be constructed using the SM fields only; the so-called Weinberg operator [53],

$$O_W^{(5)} = \bar{\ell}^c\tilde{H}^\star\tilde{H}^\dagger\ell, \tag{8}$$

that can be generated in the underlying theory at tree-level by an exchange of a heavy scalar $\Phi(1,3,0)$ and/or the heavy fermions $\Psi(1,1,0)$, $\Psi(1,3,0)$.

Therefore, the overall dimension five effective operator extension of $\mathcal{L}_{\phi\mathrm{SM}}$ is

$$\Delta\mathcal{L}_{\phi\mathrm{SM}}^{(5)} = \frac{f_W}{\Lambda_W}\bar{\ell}^c\tilde{H}^\star\tilde{H}^\dagger\ell + \frac{f_{\ell d\phi H}}{\Lambda_{\ell d\phi H}}\bar{\ell}d\tilde{H}\phi^* + \frac{f_{d^2\phi^2}}{\Lambda_{d^2\phi^2}}\bar{d}d^c\phi^2 + \mathrm{H.c.}, \tag{9}$$



(a) (b)

FIG. 2. Tree-level graphs in the underlying heavy theory that generate the dimension five effective operator $\bar{\ell} d \tilde{H} \phi^*$. Φ and Ψ stand for a heavy scalar and heavy fermion, respectively, with quantum numbers $\Phi(3,2,\frac{1}{6})$ and $\Psi(1,1,0)$, $\Psi(1,3,0)$ or $\Psi(3,2,-\frac{5}{6})$ (see text).

where we have kept a general notation assigning each of these operators their own effective scale. Note, for example, that the heavy fermionic state $\Psi(1,1,0)$ can generate all three dimension five operators in Eq. (9), in which case they will have a common scale. On the other hand, as we will see below, the effective scale, $\sim f/\Lambda$, of the Weinberg operator $\bar{\ell}^c \tilde{H}^\star \tilde{H}^\dagger \ell$ and the operator $\bar{\ell} d \tilde{H} \phi^*$ must be considerably suppressed in order to obtain sub-eV Majorana neutrino masses. This leaves us with a single viable dimension five operator, $\bar{d} d^c \phi^2$, which can generate a sub-eV neutrino mass at two-loops (see next section) with a scale low enough for it to be relevant for collider LQ phenomenology.

In the Appendix we construct the complete set of the dimension six operators involving the down-type scalar LQ $\phi(3,1,-\frac{1}{3})$ and the SM fields.[2]

V. THE DIMENSION FIVE OPERATORS AND LOW ENERGY $\Delta L = 2$ EFFECTS

As mentioned earlier, while the ϕSM renormalizable interaction Lagrangian, $\mathcal{L}_{\phi\mathrm{SM}}$, can address the BSM effects associated with the current B-physics anomalies, other aspects of NP associated with LNV require new higher-dimensional effective interactions of the LQ with the SM fields. In particular, the dimension five operators in Eq. (9) violate lepton number by two units and can, therefore, generate Majorana neutrino masses, mediate neutrinoless double beta decay and also give rise to interesting same-sign lepton signals at the LHC.

In this section we investigate in more detail the low energy $\Delta L = 2$ effects associated with these operators, while in the next section we discuss the potential $\Delta L = 2$ collider signals.

A. Majorana neutrino masses

As is well known, the dimension five Weinberg operator $\bar{\ell}^c \tilde{H}^\star \tilde{H}^\dagger \ell$ can generate a tree-level Majorana neutrino mass through the type I [if it is generated by the exchange of the heavy fermion $\Psi(1,1,0)$] and/or type III [if it is generated by $\Psi(1,3,0)$] seesaw mechanisms. In either case, the resulting Majorana neutrino mass is

$$m_\nu(\Lambda) \sim f_W \cdot \frac{v^2}{\Lambda_W}, \tag{10}$$

where v is the Higgs vacuum expectation value (VEV) and f_W and Λ_W are the Wilson coefficient and NP scale of the Weinberg operator (see Eq. (9).

Therefore, there are two extreme cases for generating $m_\nu \lesssim 1$ eV from $O_W^{(5)}$: either $\Lambda_W \sim \mathcal{O}(10^{14})$ GeV and $f_W \sim \mathcal{O}(1)$ or, if the NP scale is at the TeV range, i.e., $\Lambda_W \sim \mathcal{O}(1)$ TeV, then $f_W \sim \mathcal{O}(10^{-11})$. In both cases the effect of the Weinberg operator at TeV-scale energies is negligible.

The operators $\bar{\ell} d \tilde{H} \phi^*$ and $\bar{d} d^c \phi^2$ can also generate a Majorana neutrino mass term at 1-loop and 2-loops order, respectively, via the diagrams depicted in Fig. 3. In particular, this involves insertions of the dimension five coupling strengths $f_{\ell d \phi H}$ and $f_{d^2 \phi^2}$ as well as the Yukawa-like LQ-quark-lepton renormalizable interaction $\propto y_{q\ell}^L$ of the ϕSM Lagrangian in Eq. (2). For the $\bar{\ell} d \tilde{H} \phi^*$ case, the resulting 1-loop Majorana mass is[3]

$$m_\nu(\Lambda) \sim \frac{3 m_d}{16\pi^2} \frac{f \cdot y_{q\ell}^L}{\sqrt{2}} \frac{v}{\Lambda} \ln\left(\frac{\Lambda^2}{M_\phi^2}\right), \tag{11}$$

where $\Lambda = \Lambda_{\ell d \phi H}$ and $f = f_{\ell d \phi H}$ are the NP scale and Wilson coefficient of the dimension five operator $\bar{\ell} d \tilde{H} \phi^*$; m_d is the mass of the down-quark in the loop and M_ϕ is the leptoquark mass. Thus, setting e.g., $\Lambda = 5$ TeV and $M_\phi = 1$ TeV, we obtain

$$\frac{m_\nu(\Lambda = 5\ \mathrm{TeV})}{f \cdot y_{q\ell}^L} \sim 10^{-3} \cdot m_d, \tag{12}$$

so that, for $f \cdot y_{q\ell}^L \sim \mathcal{O}(1)$, the resulting Majorana mass is $m_\nu \sim \mathcal{O}(\mathrm{KeV})$ for $m_d \sim \mathcal{O}(\mathrm{MeV})$ (i.e., the d-quark) and $m_\nu \sim \mathcal{O}(\mathrm{MeV})$ for $m_d \sim \mathcal{O}(\mathrm{GeV})$ (i.e., the b-quark). Thus,

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[2] We have used the Mathematicanotebook of [54] to validate the EFT extension of $\mathcal{L}_{\phi\mathrm{SM}}$ which is presented in this work.

[3] See also Eq. 26 in [15] for an analogous down-quark—down-squark 1-loop Majorana mass term in R-parity violating supersymmetry.

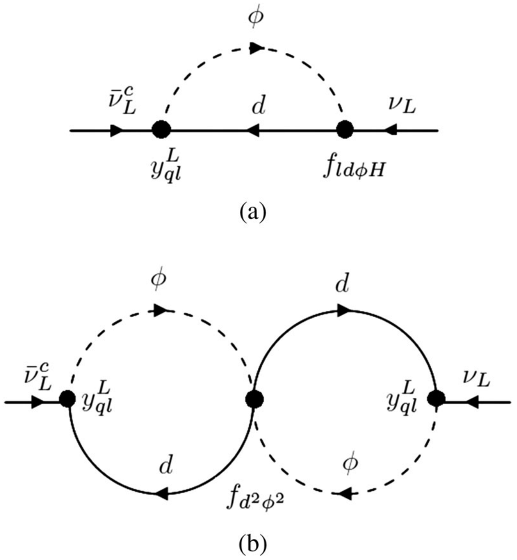

FIG. 3. The one-loop and two-loops diagrams (a) and (b) which generates a Majorana mass term with the Yukawa-like LQ-quark-lepton interaction ($\propto y^L_{q\ell}$) and the dimension five operators $\bar{\ell}d\tilde{H}\phi^*$ and $\bar{d}d^c\phi^2$ (with the coupling strength $f_{\ell d\phi H}$ and $f_{d^2\phi^2}$, respectively). See also text.

in order to obtain sub-eV Majorana neutrino masses when $\Lambda = \mathcal{O}(\mathrm{TeV})$ we should have $f\cdot y^L_{q\ell} \lesssim \mathcal{O}(10^{-3})$ for the d-quark loop and $f\cdot y^L_{q\ell} \lesssim \mathcal{O}(10^{-6})$ for the b-quark loop. In particular, if ϕ is a third generation LQ [i.e., having $\mathcal{O}(1)$ couplings only to the third generation SM fermions; see next section], then $y^L_{b\nu} \sim \mathcal{O}(1)$ and, therefore, the corresponding dimension five coupling strength should be suppressed to the level $f_{\ell d\phi H} \lesssim \mathcal{O}(10^{-6})$ if $\Lambda_{\ell d\phi H} \sim$ 5 TeV, in order to obtain e.g., $m_{\nu_\tau} \lesssim 1$ eV (ignoring off diagonal generation couplings). We note that other interesting mechanisms for generating light Majorana neutrino masses from 1-loop LQ exchanges that are intimately related to the down-quark mass matrix have been discussed in [55–61]. These studies, however, were based on renormalizable LQ extensions of the SM.

The 2-loop Majorana mass generated by the $\bar{d}d^c\phi^2$ class of dimension 5 operators is (see Fig. 3)

$$m_\nu(\Lambda) \sim \frac{f\cdot (y^L_{q\ell})^2}{(16\pi^2)^2}\frac{3m_d^2}{\Lambda}\cdot \ln^2\left(\frac{\Lambda^2}{M_\phi^2}\right), \qquad (13)$$

where here $\Lambda = \Lambda_{d^2\phi^2}$ and $f = f_{d^2\phi^2}$ are the NP scale and Wilson coefficient of the dimension five operator $\bar{d}d^c\phi^2$. Thus, setting again $\Lambda = 5$ TeV and $M_\phi = 1$ TeV, we obtain in the 2-loop case,

$$\frac{m_\nu(\Lambda = 5\ \mathrm{TeV})}{f\cdot (y^L_{q\ell})^2} \sim 10^{-4}\cdot \frac{m_d^2}{\mathrm{TeV}}, \qquad (14)$$

which, as in the 1-loop case, depends on the down-quark mass in the loops or, equivalently, on the LQ generation (defined through its renormalizable couplings to the quark-lepton pairs; see discussion above). In particular, here also, it is useful to distinguish between the three cases where ϕ couples to first, second or third generation quarks:

d-quark case ($y^L_{q\ell} = y^L_{d\nu}$ and $f = f_{d^2\phi^2}$):

In this case the 2-loop neutrino mass is too small, $m_\nu \sim 10^{-4}$ eV, when $f_{d^2\phi^2}\cdot (y^L_{d\nu})^2 \sim \mathcal{O}(1)$, so that no useful bound can be set on the scale of the dimension 5 operator involving the first generation down-quarks $\bar{d}d^c\phi^2$. Indeed, the collider effects of this operator, with a scale $\Lambda_{d^2\phi^2} \sim$ 5–15 TeV and $f_{d^2\phi^2} \sim \mathcal{O}(1)$, will be studied in the next sections.

s-quark case ($y^L_{q\ell} = y^L_{s\nu}$ and $f = f_{s^2\phi^2}$):

The resulting neutrino mass in this case is consistent with oscillation data, $m_\nu \sim$ eV, for a NP scale of several TeV and $\mathcal{O}(1)$ couplings, i.e., $f_{s^2\phi^2}\cdot (y^L_{s\nu})^2 \sim \mathcal{O}(1)$. Therefore, here also, no useful bound can be put on the corresponding dimension 5 operator $\bar{s}s^c\phi^2$.

b-quark case ($y^L_{q\ell} = y^L_{b\nu}$ and $f = f_{b^2\phi^2}$):

This corresponds to the third generation LQ case, for which we obtain $m_\nu \sim$ KeV with $f_{b^2\phi^2}\cdot (y^L_{b\nu})^2 \sim \mathcal{O}(1)$ and a NP scale of several TeV. Thus, in this case, the neutrino mass bound constrains the dimension 5 operator $\bar{b}b^c\phi^2$ or the corresponding LQ couplings: either $\Lambda_{b^2\phi^2} \sim \mathcal{O}(1000)$TeV or $f_{b^2\phi^2}\cdot (y^L_{b\nu})^2 \sim \mathcal{O}(10^{-3})$.

Finally, we wish to further comment on the link between neutrino masses and the underlying heavy physics. As noted in the previous section, the heavy fermionic states $\Psi(1,1,0)$ and $\Psi(1,3,0)$ can generate at tree-level both the Weinberg operator $\bar{\ell}^c\tilde{H}^\star\tilde{H}^\dagger\ell$ and the operator $\bar{\ell}d\tilde{H}\phi^*$, while $\Psi(1,1,0)$ can generate all three types of dimension five operators $\bar{\ell}^c\tilde{H}^\star\tilde{H}^\dagger\ell$, $\bar{\ell}d\tilde{H}\phi^*$ and $\bar{d}d^c\phi^2$. Therefore, in this setup there are several scenarios that do not require small coupling constants:

(1) The heavy fermionic state $\Psi(1,1,0)$ is responsible for generating all dimension five operators $\bar{\ell}^c\tilde{H}^\star\tilde{H}^\dagger\ell$, $\bar{\ell}d\tilde{H}\phi^*$ and $\bar{d}d^c\phi^2$, with a typical mass scale of $M_\Psi \sim \mathcal{O}(10^{14})$ GeV. In this case, the Majorana neutrino mass term will be generated at tree-level through the type I seesaw mechanisms by the Weinberg operator $\bar{\ell}^c\tilde{H}^\star\tilde{H}^\dagger\ell$, whereas the 1-loop and 2-loops contribution from the operators $\bar{\ell}d\tilde{H}\phi^*$ and $\bar{d}d^c\phi^2$ will be negligible.

(2) The heavy fermionic state $\Psi(1,3,0)$ is responsible for generating both operators $\bar{\ell}^c\tilde{H}^\star\tilde{H}^\dagger\ell$ and $\bar{\ell}d\tilde{H}\phi^*$, with a typical mass scale of $M_\Psi \sim \mathcal{O}(10^{14})$ GeV, while the operator $\bar{d}d^c\phi^2$ is generated by another heavy mediator. In this case, the Majorana neutrino mass term can be generated again at tree-level through the type I or type III seesaw mechanisms by the Weinberg operator $\bar{\ell}^c\tilde{H}^\star\tilde{H}^\dagger\ell$ and the 1-loop contribution from the operator $\bar{\ell}d\tilde{H}\phi^*$ will be

subdominant. This holds also in the case that the Weinberg operator is generated by the heavy scalar $\Phi(1,3,0)$ if $M_\Phi \sim \mathcal{O}(10^{14})$ GeV and a corresponding $\mathcal{O}(1)$ Wilson coefficient. Note that, in this case, a 2-loop Majorana mass term can be generated as well by the operator $\bar{d}d^c\phi^2$, depending on the couplings involved (see discussion above).

(3) The Weinberg operator is not relevant to neutrino masses; i.e., there are no heavy $\Phi(1,3,0)$, $\Psi(1,1,0)$ and $\Psi(1,3,0)$ states in the underlying theory. In this case, neutrino masses are not generated through the seesaw mechanism, but they may be still generated at 1-loop or at 2-loops by the dimension five operators $\bar{\ell}d\tilde{H}\phi^*$ and $\bar{d}d^c\phi^2$ as described above, if these operators are generated at tree-level in the underlying theory by other heavy states (see previous section).

B. Neutrinoless double beta decay

The dimension five operator $\bar{d}d^c\phi^2$ can mediate neutrinoless double beta decay ($0\nu\beta\beta$) via the diagram depicted in Fig. 4. This requires both the dimension five operator $\bar{d}d^c\phi^2$ and the Yukawa-like renormalizable coupling of ϕ to the right-handed first generation u-quark and electron, i.e., the term $\propto y^R_{ue}$ in $\mathcal{L}_{Y,\phi}$ (see Eq. (2). If ϕ is a third generation leptoquark, we expect $y^R_{ue} \ll 1$ (see discussion in the next section) in which case the $0\nu\beta\beta$ decay rate will be significantly suppressed.

The limit on $0\nu\beta\beta$ decay is usually expressed in terms of the electron-electron element of the neutrino mass matrix. The current bound is $|(m_\nu)_{ee}| < 0.1$–0.5 eV, depending on the $0\nu\beta\beta$ experiment; see e.g., [62]. This translates into a bound on the corresponding parton-level amplitude for $0\nu\beta\beta$ [63],

$$\frac{p_{\rm eff}}{G_F^2}|\mathcal{A}_{0\nu\beta\beta}| \simeq \frac{|(m_\nu)_{ee}|}{p_{\rm eff}} < 5\times 10^{-9}, \tag{15}$$

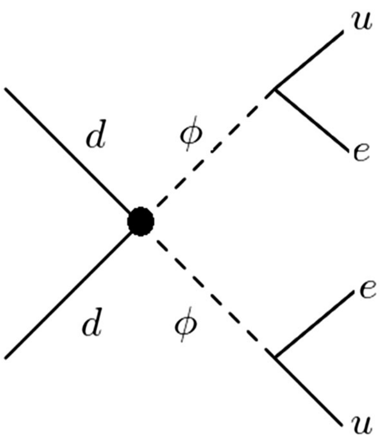

FIG. 4. Tree-level graph that generates neutrinoless double beta decay via the dimension five operator $\bar{d}d^c\phi^2$. See also text.

where $p_{\rm eff} \sim 100$ MeV is the neutrino effective momentum obtained by averaging the corresponding nuclear matrix element contribution.

In our case, the $0\nu\beta\beta$ amplitude corresponding to the diagram in Fig. 4 can be estimated as

$$\mathcal{A}_{0\nu\beta\beta} \sim \frac{f\cdot|y^R_{ue}|^2}{\Lambda M_\phi^4}, \tag{16}$$

where $f = f_{d^2\phi^2}$ and $\Lambda = \Lambda_{d^2\phi^2}$. Therefore, using Eq. (15) we obtain

$$\frac{\Lambda}{\rm TeV} \gtrsim 150\cdot\frac{f\cdot|y^R_{ue}|^2}{(M_\phi/{\rm TeV})^4}. \tag{17}$$

In particular, we find that no useful bound can be imposed on the scale of the dimension five operator $\bar{d}d^c\phi^2$, assuming $f_{d^2\phi^2} \sim \mathcal{O}(1)$ and a TeV-scale LQ mass, $M_\phi \sim \mathcal{O}(1\ {\rm TeV})$, if the LQ ϕ is a third generation LQ (as assumed below), i.e., having a suppressed Yukawa-like coupling to the first generation right-handed fermions: $y^R_{ue} < 0.1$.

VI. COLLIDER PHENOMENOLOGY OF A THIRD GENERATION SCALAR LEPTOQUARK IN THE EFT

We next discuss the expected NP signals of the down-type $\phi(3,1,-\frac{1}{3})$ and up-type $\phi(3,1,\frac{2}{3})$ LQs at the 13 TeV LHC and also at future higher energy hadron colliders such as a 27 TeV high-energy LHC (HE-LHC) and a 100 TeV future circular proton-proton collider (FCC-hh) [64].

All cross sections presented in this section were calculated using MadGraph5 [65] at LO parton-level, for which a dedicated universal FeynRules output (UFO) model for the LQ-SM EFT framework defined in Eq. (4) was produced for the MadGraph5 sessions using FeynRules [66]. The LO nnpdf3 PDF set (NNPDF30-lo-as-0130 [67]) was used in all the calculations presented below. Also, all cross sections were calculated with a dynamical scale choice for the central value of the factorization (μ_F) and renormalization (μ_R) scales corresponding to the sum of the transverse mass in the hard-process, and, for consistency with the EFT framework, a cut on the center of mass energy of $\sqrt{\hat{s}} < \Lambda$ was placed using Mad-Analysis55 [68], where several values of Λ (the scale of NP) were used for the processes considered below.[4]

Furthermore, we will assume throughout the rest of the paper that $\phi(3,1,-\frac{1}{3})$ and $\phi(3,1,\frac{2}{3})$, under consideration in this section, are third generation leptoquarks and denote them generically by ϕ_3. In particular, we assume that the LQ-lepton-quark Yukawa-like couplings of ϕ_3 to the first and second generations SM fermions in the corresponding

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[4] The UFO model files are available upon request.

renormalizable ϕSM Lagrangian are much smaller than its couplings to the third generation quark-lepton pair, e.g., to a $t\tau$ and/or $b\nu_\tau$ pairs in the case of the down-type LQ $\phi(3,1,-\frac{1}{3})$ (see Eqs. (2).

This scenario can be realized by imposing an approximate Z_3 generation symmetry under which the physical states of the SM fermions (i.e., mass eigenstates) transform as

$$\psi^k \to e^{i\alpha(\psi^k)\tau_3}\psi^k, \qquad \tau_3 \equiv 2\pi/3, \tag{18}$$

where k is the generation index and $\alpha(\psi^k)$ are the Z_3 charges of ψ^k.

Consider for example the down-type LQ $\phi(3,1,-\frac{1}{3})$: if the Z_3 charges equal the generation index, i.e., $\alpha(\psi^k) = k$, and $\alpha(\phi) = 3$, then only terms in $\mathcal{L}_{\phi\mathrm{SM}}$ involving the third generation are allowed. In particular, assuming the baryon number conservation and thus ignoring the Z_3-allowed LQ interactions with the third generation quarks (i.e., $\phi\bar t^c_R b_R$ and $\phi\bar t^c_L b_L$) that would in general allow for proton decay, we have

$$\mathcal{L}_{Y,\phi_3} \approx y^L_{q_3\ell_3}(\bar t^c_L\tau_L + \bar b^c_L\nu_{\tau L})\phi^* + y^R_{u_3e_3}\bar t^c_R\tau_R\phi^* + \mathrm{H.c.}, \tag{19}$$

where we will assume that the above Yukawa-like LQ-quark-lepton third generation couplings are $\mathcal{O}(1)$.

The Z_3 generation symmetry is exact in the limit where the quark mixing CKM matrix V is diagonal, so that Z_3-breaking effects will in general be proportional to the square of the small off diagonal CKM elements $|V_{cb}|^2$, $|V_{ub}|^2$, $|V_{ts}|^2$, $|V_{td}|^2$, and will, therefore, be suppressed (see also [28,69,70]). In particular, the Z_3 generation symmetry is assumed to be broken in the underlying heavy theory and can, therefore, be traced to the higher dimensional operators. For example, the off diagonal SM Yukawa couplings may be generated by the dimension six operators,

$$\Delta\mathcal{L}^{(6)}_{Y,H} = (f_{uH}\bar q_L\tilde H u_R + f_{dH}\bar q_L H d_R)\frac{H^\dagger H}{\Lambda^2} + \mathrm{H.c.}, \tag{20}$$

where, if e.g., $\Lambda \sim 1.5$, 3 or 5 TeV and $f_{uH}, f_{dH} \sim \mathcal{O}(1)$, then the resulting effective Yukawa couplings, $y_{\mathrm{eff}} = f_{uH,dH}\cdot v^2/\Lambda^2$, are $y_{\mathrm{eff}} \sim \mathcal{O}(y_b^{\mathrm{SM}})$, $y_{\mathrm{eff}} \sim \mathcal{O}(y_c^{\mathrm{SM}})$ or $y_{\mathrm{eff}} \sim \mathcal{O}(y_s^{\mathrm{SM}})$, respectively, where y_q^{SM} are the corresponding Yukawa couplings in the SM (see [71]).

The Z_3 breaking terms in the LQ sector will also be generated in the effective theory through higher dimensional operators. To demonstrate that consider for example the dimension five operator $\bar d d^c\phi^2$ in Eq. (6). As was shown in Sec. IV, this operator can be generated at tree-level in the UV theory by exchanging e.g., a heavy scalar $\Phi(6,1,-\frac{2}{3})$ [see diagram (a) in Fig. 1]. Thus, if $\Phi(6,1,-\frac{2}{3})$ couples to the first and/or second generation down-quarks, then the Z_3 generation symmetry is broken and the scale of generation breaking is the mass of $\Phi(6,1,-\frac{2}{3})$, M_Φ. In particular, the Z_3 generation breaking effects in this case will be proportional to $g_{\Phi dd}\cdot g_{\Phi\phi\phi}/M_\Phi$, where $g_{\Phi dd}$ and $g_{\Phi\phi\phi}$ are the couplings of the heavy $\Phi(6,1,-\frac{2}{3})$ to a dd-pair and a $\phi\phi$-pair, respectively. The matching to the EFT framework of Eq. (9) can be done by replacing $M_\Phi \to \Lambda_{d^2\phi^2}$ and $g_{\Phi dd}\cdot g_{\Phi\phi\phi} \to f_{d^2\phi^2}$.

We thus, allow for higher dimensional interactions of ϕ_3 with the lighter SM fermion generations, keeping in mind that these are *a priori* suppressed in the EFT by inverse powers of the NP scale (e.g., by $1/\Lambda$ if it originates from the dimension five operators) and that, in this case, Λ represents the scale of breaking the Z_3 generation symmetry.[5]

A. The down-type scalar LQ $\phi(3,1,-\frac{1}{3})$

We now consider the LHC signals of the down-type third generation LQ $\phi_3 = \phi_3(3,1,-\frac{1}{3})$ under investigation. Following our above setup where ϕ_3 is expected to have suppressed couplings to first and second generation fermions, single ϕ_3 production will occur through the channel $gb \to \phi_3\nu_\tau$, with a cross section $\sigma(pp_{(gb)} \to \phi_3\nu_\tau) \sim 3.5(0.025)$ fb for $M_{\phi_3} = 1(2)$ TeV and $y^L_{b\nu_\tau} = 1$ [36]. Also, with subleading couplings to the first and second generation fermions, the main channels for ϕ_3 pair-production will be gluon and $q-\bar q$ fusion, where the typical cross sections are $\sigma(pp_{(gg,q\bar q)} \to \phi_3\phi_3^*) \sim 5.5(0.01)$ fb for $M_{\phi_3} = 1(2)$ TeV [36] (with no cut on the $\phi_3\phi_3^*$ invariant mass) and do not depend on the ϕ_3-quark-lepton couplings. Thus, assuming that ϕ_3 decays via $\phi_3 \to t\tau^-$ and/or $\phi_3 \to b\nu_\tau$ with a 50% branching ratio into each channel, we find e.g., $\sigma(pp_{(gg,q\bar q)} \to \phi_3\phi_3^* \to t\bar t\tau^-\tau^+) \sim 1.4$ fb at a 13 TeV LHC if $M_{\phi_3} \sim 1$ TeV. A dedicated search in this channel was carried by CMS in [46], where no evidence for this signal was found, setting a limit on the LQ mass of $M_{\phi_3} \gtrsim 900$ GeV at 95% confidence level for $\mathrm{BR}(\phi_3 \to t\tau^-) = 1$.

As mentioned above, LQ phenomenology changes in the presence of the higher dimensional effective operators. In particular, additional potentially interesting ϕ_3 production channels are opened at the LHC. However, most of them will have a too small cross section at the 13 TeV LHC, due to the $1/\Lambda^n$ suppression in the EFT expansion, so that the leading effects are produced by the dimension five operators involving ϕ_3 in Eq. (9). Recall, however, that the operator $\bar\ell d\tilde H\phi^*$ is expected to have suppressed effects because of a large effective scale, as required for consistency with sub-eV neutrino masses (cf. the previous section).

We are therefore left with only one dimension five operator, $\bar d d^c\phi^2$, that can potentially mediate interesting

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[5] Note that the couplings of ϕ_3 to the first and second generations fermions can also be loop generated by the renormalizable LQ-quark-lepton couplings. In this case they are suppressed by the corresponding loop factor and CKM elements and are, therefore, subdominant.

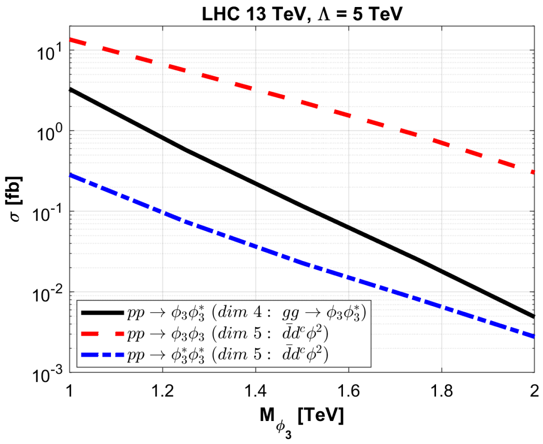

FIG. 5. Pair-production cross sections of the down-type LQ ϕ_3 at the 13 TeV LHC with $\Lambda_{d^2\phi^2} = 5$ TeV: $pp \to \phi_3\phi_3$ (dashed line), $pp \to \phi_3\phi_3^*$ (solid line) and $pp \to \phi_3^*\phi_3^*$ (dashed-dot line) (see also text).

ϕ_3 pair-production signals at the LHC. In particular, we find that this operator may yield a strikingly large *asymmetric* same-sign(charge) $\phi_3\phi_3$ signal at the LHC via $dd \to \phi_3\phi_3$, which is more than an order of magnitude larger than the charged conjugate channel $\bar{d}\,\bar{d} \to \phi_3^*\phi_3^*$, due to the different fractions of d and $\bar{d}$ in the incoming protons; see Fig. 5. The hard cross section for $dd \to \phi_3\phi_3$ (which equals that of the charged conjugate one $\bar{d}\,\bar{d} \to \phi_3^*\phi_3^*$) is

$$\hat{\sigma}(dd \to \phi_3\phi_3) = \frac{\beta f^2}{12\pi\Lambda^2}, \tag{21}$$

where (cf. Eq. (9) $\Lambda = \Lambda_{d^2\phi^2}$, $f = f_{d^2\phi^2}$, $\beta^2 = 1\text{–}4M_{\phi_3}^2/\hat{s}$, and $\sqrt{\hat{s}}$ is the center of mass energy of the hard process. For example, if $\Lambda_{d^2\phi^2} = 5$ TeV (and with a cut on the $\phi_3\phi_3$ invariant mass, $M_{\phi_3\phi_3} < 5$ TeV), we find[6]

$$\sigma(pp \to \phi_3\phi_3)_{M_{\phi_3}\sim 1\ \mathrm{TeV}} \sim 14\ \mathrm{fb},$$
$$\sigma(pp \to \phi_3\phi_3)_{M_{\phi_3}\sim 2\ \mathrm{TeV}} \sim 0.3\ \mathrm{fb}. \tag{22}$$

This can be compared to the gluon-fusion cross section of the opposite-charge $\phi_3\phi_3^*$ pair-production signal, $pp_{(gg)} \to \phi_3\phi_3^*$, for which the hard cross section (see e.g., [30,31]),

$$\hat{\sigma}(gg \to \phi_3\phi_3^*) = \frac{\pi\alpha_s^2}{96\hat{s}} \cdot \left\{\beta(41 - 31\beta^2) - (17 - 18\beta^2 + \beta^4) \cdot log\left(\frac{1+\beta}{1-\beta}\right)\right\}, \tag{23}$$

drops with the energy as $1/\hat{s}$ and yields a cross section of (again with $M_{\phi_3\phi_3^*} < 5$ TeV),

$$\sigma(pp \to \phi_3\phi_3^*)_{M_{\phi_3}\sim 1\ \mathrm{TeV}} \sim 3\ \mathrm{fb},$$
$$\sigma(pp \to \phi_3\phi_3^*)_{M_{\phi_3}\sim 2\ \mathrm{TeV}} \sim 0.005\ \mathrm{fb}. \tag{24}$$

We thus see that the same-sign $\phi_3\phi_3$ rate is expected to be larger than the opposite-sign $\phi_3\phi_3^*$ rate at the 13 TeV LHC, in particular, $\sigma(pp \to \phi_3\phi_3)/\sigma(pp \to \phi_3\phi_3^*) \sim 5(60)$ for $M_{\phi_3} = 1(2)$ TeV.

Taking into account the leading ϕ_3 decays $\phi_3 \to t\tau^-$ and $\phi_3 \to b\nu_\tau$, this signal will in turn give rise to the new asymmetric signatures ($j_b = b$-jet):

(i) $pp \to \phi_3\phi_3 \to 2 \cdot j_b + \not{E}_T$
(ii) $pp \to \phi_3\phi_3 \to tt\tau^-\tau^-$
(iii) $pp \to \phi_3\phi_3 \to t\tau^- + j_b + \not{E}_T$

with a cross section which is more than an order of magnitude larger than the charged conjugate channels.

While $pp \to 2 \cdot j_b + \not{E}_T$ may not be unique to ϕ_3 pair-production and may be more challenging due to the larger background expected in this channel, the signal of same-sign top-quark pair in association with a pair of same-sign negatively charged τ-leptons, $pp \to tt\tau^-\tau^-$, and the single top—single τ signature, $pp \to t\tau^- + j_b + \not{E}_T$, may give striking new asymmetric $\phi_3\phi_3$ signals.

For example, if the scale of the NP underlying $\mathcal{L}_{\phi\mathrm{SM}}$ is $\Lambda = 5$ TeV, the LQ mass is $M_{\phi_3} \sim 1$ TeV and its leading branching ratios are $\mathrm{BR}(\phi_3 \to t\tau^-) = \mathrm{BR}(\phi_3 \to b\nu_\tau) = 0.5$, then we expect $\sigma(pp \to tt\tau^-\tau^-) \sim 3.4$ fb; while $\sigma(pp \to \bar{t}\bar{t}\tau^+\tau^+) \sim 0.07$ fb; see Fig. 5. The former is about 5 times larger than the rate for the gluon-fusion $\phi_3\phi_3^*$ signal $pp \to t\bar{t}\tau^+\tau^-$, for which a dedicated search has already been performed by CMS [46] with null results.

With an integrated luminosity of ~300 fb^{-1}, $\Lambda = 5$ TeV and $M_{\phi_3} \sim 1$ TeV, about 1000 $tt\tau^-\tau^-$ events with an invariant mass smaller than 5 TeV are expected. After the top-quarks decay hadronically via $t \to W^+b \to 2 \cdot j + b$ ($j =$ light jet) with a $\mathrm{BR}(t \to W^+b \to 2 \cdot j + b) \sim 2/3$, we

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[6] There are no SM contributions to the processes studied here, and also none of the tree-level generated dimension six operators that we list in the Appendix contribute to them. Furthermore, other dimension six operators which do not involve the LQ fields and which can, in principle, be generated by the heavy mediators in Figs. 1 and 2 [e.g., four-fermion operators such as $d\bar{d}\ell^+\ell^-$ and $(d\bar{d})^2$], do not affect the same-sign lepton signals considered in this work. Thus, the dimension five operators that we consider generate the leading contributions to these processes. In particular, potential corrections to the leading-order cross sections presented in this section can be generated either by loop-generated dimension six operators and/or by dimension seven operators. The former are suppressed by a factor of $E/(16\pi^2\Lambda)$ (E is the typical energy of the process) and can, therefore, be neglected here, while the latter are suppressed typically by $(E/\Lambda)^2$, and, therefore, their size depends on the relevant energy scale of the process. In particular, for the s-channel process [see Fig. 1(a)] the corrections can reach 50%, while for t or u channel processes [see Fig. 1(b)] the relevant energy scale is much smaller and the corrections are again negligible.

expect about 450 same-sign $\tau^-\tau^-$ events with a high jet-multiplicity signature: $pp \to \tau^-\tau^- + 4\cdot j + 2\cdot j_b$ and with a statistical error of $\sim\sqrt{450} \sim 20$ events and no irreducible background (see also discussion below).[7] Note also that roughly the same number of events are expected for the $t\tau^-$ production signal $pp \to t\tau^- + j_b + \not{E}_T$, which leads to $pp \to \tau^- + 2\cdot j + 2\cdot j_b + \not{E}_T$, when the top-quark decays hadronically via $t \to W^+ b \to 2\cdot j + b$. This single-$\tau$ signal lack a unique characterization akin to the same-sign lepton signature in pair LQ production and might, therefore, be harder to trace.

It is also useful to define the inclusive same-charge $\tau\tau$ asymmetry,

$$\mathcal{A}_{\tau\tau} \equiv \frac{\sigma(pp \to \tau^-\tau^- + X_j) - \sigma(pp \to \tau^+\tau^+ + X_j)}{\sigma(pp \to \tau^-\tau^- + X_j) + \sigma(pp \to \tau^+\tau^+ + X_j)}, \qquad (25)$$

where we have assumed again that the top-quark decays hadronically via $t \to W^+ b \to 2\cdot j + b$ and X_j stands for any accompanying jets in the final state, i.e., for events with prompt same-sign $\tau\tau$ and no missing transverse energy (MET). When the ϕ_3 mass is in the range $1\ \text{TeV} \lesssim M_{\phi_3} \lesssim 2$ TeV, we expect $\mathcal{A}_{\tau\tau} \to 1$ since this asymmetry receives its most significant contribution from the $\phi_3\phi_3$ and $\phi_3^*\phi_3^*$ channels (see Fig. 5). The SM background for the same-sign $\tau^-\tau^-$ events with no MET, from processes that can mimic this final state, is expected to be significantly suppressed, in particular, after imposing the appropriate kinematical and selection cuts (see also comment below). We, therefore, expect the above same-charge asymmetry $\mathcal{A}_{\tau\tau}$ to be close to a 100%.

The statistical significance, N_{SD}, with which this asymmetry can be detected at the LHC is

$$N_{SD} \sim \sqrt{\sigma_{\tau\tau}\cdot L}\cdot \mathcal{A}_{\tau\tau}\cdot\sqrt{\epsilon}, \qquad (26)$$

where $\sigma_{\tau\tau}$ is the inclusive cross section $\sigma(pp \to \tau^-\tau^- + X_j)$ and ϵ is the corresponding combined efficiency for the simultaneous measurement of this final state. Thus, with an integrated luminosity of 300 inverse fb [recall that $\sigma(pp \to \tau^-\tau^- + X_j) \sim 1.5$ fb for $\Lambda = 5$ TeV and $M_{\phi_3} = 1$ TeV] and a combined efficiency of $\epsilon \sim 0.01$, this asymmetry can be detected with about a $\sim 2\sigma$ significance. At the high-luminosity LHC with 3000 inverse fb this asymmetry should be accessible with a statistical significance of $N_{SD} \sim 7$.

Finally, we wish to further comment on the potential background to the LNV same-sign lepton signals considered here and in the following section. Although these signals have formally no irreducible SM background (since lepton number is conserved in the SM), they can be "contaminated" by reducible background that can mimic these signatures due to higher-order effects (e.g., initial and final state radiation), particle/jets miss-identification, $\tau^\pm$ reconstruction limitations, heavy flavor decays and alike. However, due to the distinct characteristics of our same-sign (isolated) lepton-pair signals, such a background can in principle be reduced to the desired level with the appropriate kinematical and selection cuts as well as veto requirements, e.g., on the MET and the energy distribution of the jets in the process; see e.g., the recent SUSY searches in same-sign lepton events at the LHC, performed by the CMS [72] and ATLAS [73] Collaborations. An example of such a potential background is the SM process $pp \to t(\to bW^+ \to bjj)\bar{t}(\to \bar{b}W^- \to \bar{b}\tau^-\bar{\nu}_\tau)W^-(\to \tau^-\nu_\tau)$ and the charged conjugate channel (considered also in [72,73]), which lead to same-sign $\tau^\pm\tau^\pm$ events that can mimic our LNV LQ mediated signals, e.g., from $pp \to t(\to bW^+ \to bjj)t(\to bW^+ \to bjj)\tau^-\tau^-$. This SM process can, therefore, also "contaminate" the asymmetry $\mathcal{A}_{\tau\tau}$ in Eq. (25), since $\sigma(pp \to t\bar{t}W^+) \sim 2\sigma(pp \to t\bar{t}W^-)$. However, not only that this background has a cross section of the same order of the LNV signal considered above, i.e., $\sigma(pp \to t\bar{t}W^\pm \to \tau^\pm\tau^\pm + X_j + \not{E}_T) \sim \sigma(pp \to tt\tau^-\tau^- \to \tau^-\tau^- + X_j) \sim \mathcal{O}(1)$ fb, it also contains a different energy distribution of the MET and jets in the process and can, therefore, be significantly reduced with the proper selection cuts and veto requirements.

B. The up-type scalar LQ $\phi(3,1,\frac{2}{3})$

We wish to briefly comment here on the phenomenology and LHC signals expected for an up-type LQ $\phi(3,1,\frac{2}{3})$ in the EFT framework. The renormalizable Yukawa-like interactions of this LQ contain only the term $y^R_{d^i d^j}\bar{d}^{ci}_R d^j_R\phi$, where $y^R_{d^i d^j}$ is antisymmetric due to SU(3) (color) gauge invariance. Note, however, that this diquark LQ coupling violates baryon number and, in the presence of the higher dimensional LQ couplings to quark-lepton pairs (see below), may mediate proton decay. We therefore, assume that it is either negligibly small or forbidden due to a symmetry.

In the up-type $\phi(3,1,\frac{2}{3})$ case, we find that there are four dimension five operators (in addition to the Weinberg operator of Eq. (8),[8]

$$\Delta\mathcal{L}^{(5)}_{\phi\text{SM}} = \frac{f_{\ell u\phi H}}{\Lambda_{\ell u\phi H}}\bar{\ell}u\tilde{H}\phi^* + \frac{f_{\ell d\phi H}}{\Lambda_{\ell d\phi H}}\bar{\ell}dH\phi^* + \frac{f_{qe\phi H}}{\Lambda_{qe\phi H}}\bar{q}eH\phi + \frac{f_{u^2\phi^2}}{\Lambda_{u^2\phi^2}}\bar{u}u^c\phi^2 + \text{H.c.}. \qquad (27)$$

---

[7] This estimate does not include the τ-decay branching ratio into a specific final state.

[8] We note that if both the down-type and up-type LQ are included as light degrees of freedom in the low-energy framework, then four more dimension five operators can be constructed in the EFT extension: $\bar{q}\ell^c\phi_d^*\phi_u^*$, $\bar{u}e^c\phi_d^*\phi_u^*$, $\bar{q}q^c\phi_d\phi_u$ and $\bar{d}u^c\phi_d\phi_u$, where we have used here the subscripts d and u to distinguish between them.

The fourth operator in Eq. (27), $\bar{u}u^c\phi^2$, will give rise to a similar same-sign asymmetric $\phi_3\phi_3$ signals via $uu \to \phi_3\phi_3$ (and the much smaller charged conjugate one $\bar{u}\,\bar{u} \to \phi_3^*\phi_3^*$), with a considerably larger cross section than the same-sign down-type LQ pair-production one, due to the larger u-quark content/PDF in the protons. For example, with $\Lambda_{u^2\phi^2} = 5$ TeV and the invariant mass cut $M_{\phi_3\phi_3} < 5$ TeV, we find for the up-type LQ case,

$$\sigma(pp \to \phi_3\phi_3)_{M_{\phi_3}\sim 1\ \mathrm{TeV}} \sim 77\ \mathrm{fb},$$
$$\sigma(pp \to \phi_3\phi_3)_{M_{\phi_3}\sim 2\ \mathrm{TeV}} \sim 3\ \mathrm{fb}, \tag{28}$$

which is about 25(600) times larger than the expected opposite-charged $\phi_3\phi_3^*$ signal for $M_{\phi_3} = 1(2)$ TeV; see Eq. (24).

In contrast to the case of the down-type LQ (which decays via its renormalizable couplings to quark-lepton pairs), the decay pattern of the up-type LQ considered here will be controlled by its dimension five interactions with the SM fields in Eq. (27). In particular, it will decay via either $\phi \to de^+$ and/or $\phi \to u\nu$, where d, u, e, ν stand here for a down-quark, up-quark, charged lepton and neutrino of any generation, with a corresponding coupling which is suppressed by $\sim v/\Lambda$; e.g., for the decay $\phi \to u\nu$ the coupling is $f_{lu\phi H} \cdot (v/\Lambda_{lu\phi H})$. Thus, assuming as an example that its dominant dimension five couplings are to the third generation SM fermions, then here also, when it decays via either $\phi_3 \to b\tau^+$ and/or $\phi_3 \to t\nu_\tau$, we expect the new asymmetric signals:

(i) $pp \to \phi_3\phi_3 \to tt + \not{E}_T$
(ii) $pp \to \phi_3\phi_3 \to \tau^+\tau^+ + 2\cdot j_b$
(iii) $pp \to \phi_3\phi_3 \to t\tau^+ + j_b + \not{E}_T$

each having a cross section which is several orders of magnitude larger than the charged conjugate channels.

Despite obvious parallels, there are important differences between the above signals and the ones expected for the down-type LQ:

1. The same-sign $\tau^+\tau^+$ signal $pp \to \phi_3\phi_3 \to \tau^+\tau^+ + 2\cdot j_b$ for the up-type LQ has opposite lepton charges than the corresponding signal for the down-type LQ. Therefore, the asymmetry $\mathcal{A}_{\tau\tau}$ flips signs in the up-type LQ case.
2. Similarly, in single LQ production, the final τ lepton is positive for the up-type LQ and negative for the down-type.
3. The same-sign $\tau\tau$ signal has a lower jet multiplicity than in the case of the down-type LQ.

(iv) The same-charge top-quark pair signal $pp \to \phi_3\phi_3 \to tt + \not{E}_T$ can also yield a same-sign lepton signal $pp \to \ell^+\ell^+ + 2\cdot j_b + \not{E}_T$, involving any of the charged leptons, i.e., $\ell^+\ell^+ = e^+e^+, \mu^+\mu^+, \tau^+\tau^+$, if the top-quark decays leptonically via $t \to W^+b \to \ell^+\nu_\ell b$.

Thus, the most promising signals in up-type $\phi_3\phi_3$ pair-production are $pp \to \tau^+\tau^+ + 2\cdot j_b$ and $pp \to tt + \not{E}_T \to \ell^+\ell^+ + 2\cdot j_b + \not{E}_T$, containing two positive charged leptons (for which the background is low) and two high-p_T tagged b-jets. For $\Lambda = 5$ TeV, $M_{\phi_3} = 1$ TeV and assuming $\mathrm{BR}(\phi_3 \to b\tau^+) = \mathrm{BR}(\phi_3 \to t\nu_\tau) = 0.5$, the overall cross sections for these signals (with an invariant mass smaller than 5 TeV) are expected to be

$$\sigma(pp \to \tau^+\tau^+ + 2\cdot j_b) \sim 20\ \mathrm{fb},$$
$$\sigma(pp \to \ell^+\ell^+ + 2\cdot j_b + \not{E}_T) \sim 0.2\ \mathrm{fb}, \tag{29}$$

where, as mentioned above, for the same-charged top-quark pair signal, $pp \to tt + \not{E}_T \to \ell^+\ell^+ + 2\cdot j_b + \not{E}_T$, this cross section applies to any one of the same-charged leptons,

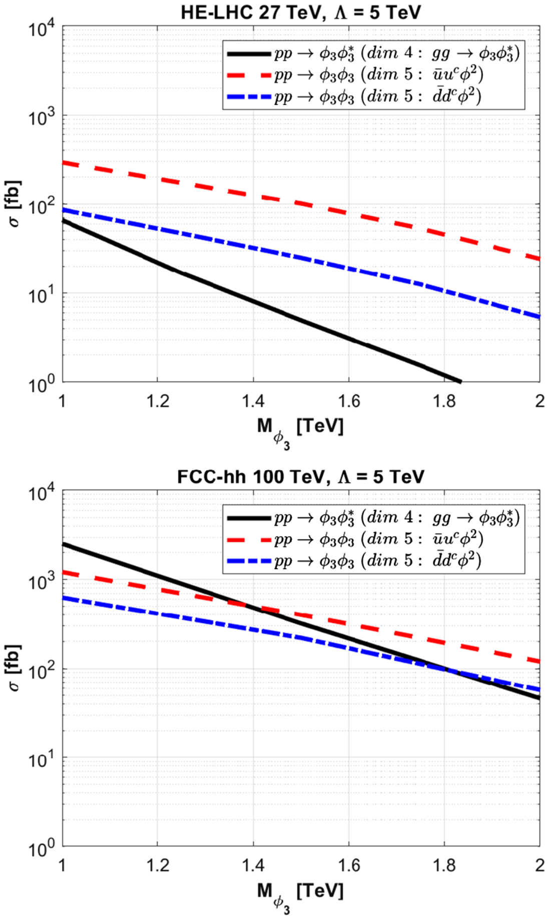

FIG. 6. Pair-production cross sections of the down-type and up-type LQ, as a function of the LQ mass, for a NP scale $\Lambda = 5$ TeV, at a 27 TeV HE-LHC (upper plot) and a 100 TeV FCC-hh (lower plot): the QCD cross section via $gg, q\bar{q} \to \phi_3\phi_3^*$ (solid line), the same-charge up-type LQ pair-production cross section via $uu \to \phi_3\phi_3$ (dashed line) and the same-charge down-type LQ pair-production cross section via $dd \to \phi_3\phi_3$ (dashed-dotted line). See also text.

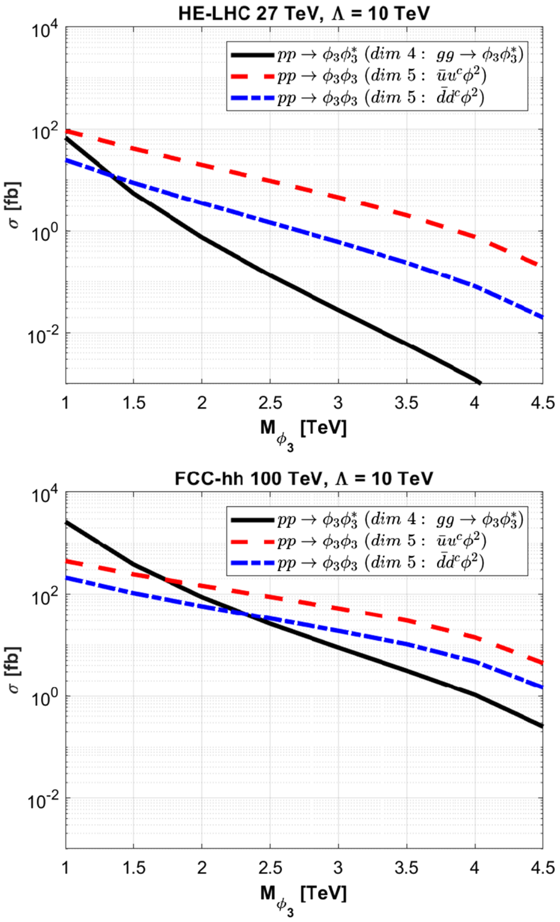

FIG. 7. Same as Fig. 6 for Λ = 10 TeV.

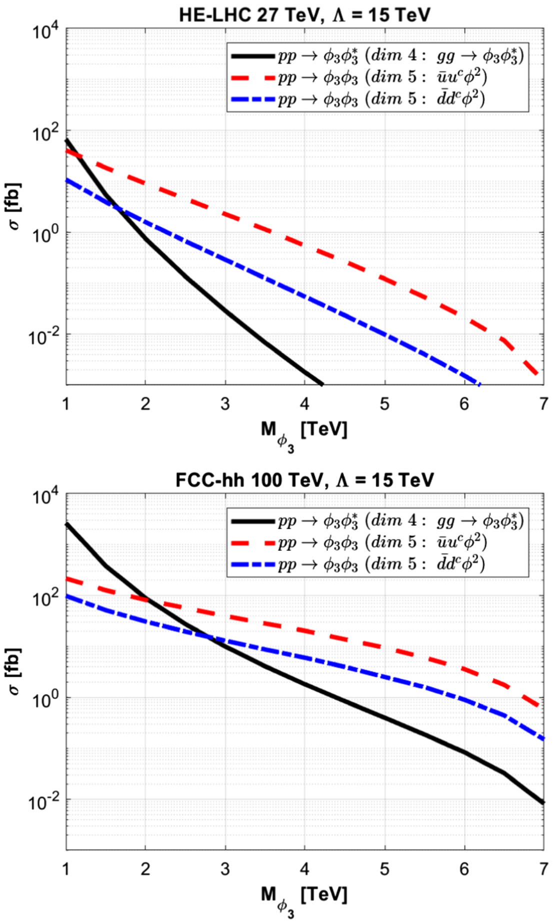

FIG. 8. Same as Fig. 6 for Λ = 15 TeV.

i.e., $\ell^+\ell^+ = e^+e^+, \mu^+\mu^+$ or $\tau^+\tau^+$, when the top-quarks decay leptonically with $\mathrm{BR}(t \to W^+ b \to \ell^+\nu_\ell b) \sim 0.1$.

Considering the same-sign dilepton asymmetry defined in Eq. (25), in the up-type LQ case we find that $\mathcal{A}_{\tau\tau}$ may be detected with a statistical significance of $N_{SD} \sim 8$, with an integrated luminosity of 300 inverse fb and a combined efficiency of $\epsilon \sim 0.01$ (see Eq. (26). On the other hand, a statistically significant signal of the asymmetries $\mathcal{A}_{ee/\mu\mu}$ will require the 13 TeV HL-LHC with an integrated luminosity of 3000 inverse fb.

C. Expectations at higher energy hadron colliders

As can be seen from Fig. 5, the LQ production cross sections sharply drop with the LQ mass at the 13 TeV LHC for LQ masses $M_\phi > 1$ TeV. This is due to the limited phase space at the 13 TeV LHC for producing TeV-scale heavy particles and, hence, the currently relatively poor discovery potential for such new heavy particles. In particular, the detection of NP scales $\Lambda > 5$ TeV and/or heavy new particles with masses of several TeV, will require in general higher energy colliders with higher luminosities. For example, for a LQ mass of $M_{\phi_3} \sim 4$ TeV, the opposite-charge $\phi_3\phi_3^*$ pair-production cross section (via $gg, q\bar{q} \to \phi_3\phi_3^*$) at the 13 TeV LHC is $\sigma(pp \to \phi_3\phi_3^*) \sim 10^{-6}$ fb. The new same-charge $\phi_3\phi_3$ signal discussed above is also too small at the 13 TeV LHC for $M_{\phi_3} \sim 4$ TeV; $\sigma(pp \to \phi_3\phi_3) \sim 10^{-4}$ fb, if the NP scale is $\Lambda \sim 10$ TeV. Therefore, heavy LQ with masses of several TeV are not accessible at the 13 TeV LHC with or without the new EFT interactions from the higher dimensional effective operators.

A better sensitivity to multi-TeV LQ and, in particular, to the LQ EFT dynamics presented in this work, can be obtained at future higher energy hadron colliders such as the HE-LHC and the FCC-hh mentioned above. In Figs. 6–8 we plot the same-charge LQ pair-production cross sections $pp \to \phi_3\phi_3$ for both the down-type and up-type LQ (i.e., the underlying hard-processes being $dd \to \phi_3\phi_3$ and $uu \to \phi_3\phi_3$, respectively), as well as the opposite-charge LQ pair-production (QCD) cross section $pp \to \phi_3\phi_3^*$ (via $gg, q\bar{q} \to \phi_3\phi_3^*$), for a NP scale of $\Lambda = 5$, 10 and 15 TeV. Here also, for consistency with the EFT framework, all cross sections are calculated with an invariant mass cut on the LQ pair $M_{\phi_3\phi_3} < \Lambda$, i.e., $M_{\phi_3\phi_3} < 5$, 10, 15 TeV for $\Lambda = 5$, 10, 15 TeV, respectively. We note that the cross sections in Figs. 6–8 for a third generation LQ are insensitive to the Yukawa couplings in Eq. (2), so the results for first and second generation LQ are expected to be comparable.

We see that the production rate of positively charged up-type LQ pair (in the EFT framework) can reach $\sigma(pp \to \phi_3\phi_3) \sim \mathcal{O}(1)$ fb at the 100 TeV FCC-hh, for a rather heavy LQ with $M_{\phi_3} \sim 7$ TeV and a NP scale of $\Lambda \sim 15$ TeV, whereas the corresponding opposite-charged $\phi_3\phi_3^*$ signal (i.e., for $M_{\phi_3} \sim 7$ TeV) is expected to be about 2 orders of magnitudes smaller. A 27 TeV HE-LHC is also sensitive to a several TeV LQ and a NP scale of $\mathcal{O}(10)$ TeV, e.g., expecting an $\mathcal{O}(1)$ fb cross section for pair production of positively charged up-type LQ pair when $M_{\phi_3} \sim 4$ TeV and a NP scale of $\Lambda \sim 10$ TeV.

VII. SUMMARY

We have explored the phenomenology of the EFT expansion of a low-energy TeV-scale framework, where the "light" degrees of freedom contain the SM fields and a down-type scalar LQ $\phi(3,1,-\frac{1}{3})$ or an up-type LQ $\phi(3,1,\frac{2}{3})$.

We found that there are only two dimension five operators that can be assigned to the down-type LQ $\phi(3,1,-\frac{1}{3})$ and four dimension five operators for the up-type LQ $\phi(3,1,\frac{2}{3})$; all these dimension five operators violate lepton number by two units. We have also identified the distinct underlying heavy physics that can generate these operators at tree-level.

We have shown that these dimension five operators can generate sub-eV Majorana neutrino masses at 1-loop and 2-loops, where, in the 2-loops case, the effective NP scale can be as low as $[\Lambda/\mathrm{TeV}]/f \sim 5$, where f is the corresponding Wilson coefficient derived from the underlying heavy theory. We also found that the dimension five operator involving the down-type LQ, $\bar{d}d^c\phi^2$, which is relevant to current collider phenomenology, may mediate neutrinoless double beta decay.

We have then focused on collider phenomenology of both the down and up-type scalar LQ in the EFT framework. In particular, motivated by the current anomalies in B-decays, we have suggested an approximate Z_3 generation symmetry and studied the signals of third generation down-type and up-type LQs (ϕ_3) at the LHC. We found that the dimension five operators may give rise to striking asymmetric, same-charge dilepton final states in the reactions $pp \to \phi_3\phi_3$ for both the down and up-type scalar LQs, that have low background.

For example, for the third generation down-type LQ with a mass $M_{\phi_3} \sim 1$ TeV and a NP scale $\Lambda \sim 5$ TeV, the resulting same-sign lepton signature is $pp \to \phi_3\phi_3 \to \tau^-\tau^- + 4\cdot j + 2\cdot j_b$ ($j =$ light jet and $j_b =$ b-jet), which is expected to yield about 500 such $\tau^-\tau^-$ events at the 13 TeV LHC with a luminosity of 300 fb^{-1}. For the third generation up-type LQ, we expect about 6000 events of same-sign positively charged $\tau^+\tau^+$ from the process $pp \to \phi_3\phi_3 \to \tau^+\tau^+ + 2\cdot j_b$, if $\Lambda \sim 5$ TeV. Moreover, for similar parameters, the same-charge up-type $\phi_3\phi_3$ pair production process can also generate events with pairs of same-charge top quarks $pp \to tt + \not{E}_T$ (when each LQ decays via $\phi_3 \to t\nu$), leading to about 50 same-sign dilepton events $pp \to \ell^+\ell^+ + 2\cdot j_b + \not{E}_T$ (when each top-quark decays leptonically via $t \to W^+b \to \ell^+\nu_\ell b$), for any of the three charged leptons, $\ell = e, \mu, \tau$.

We have also defined a double lepton-charge asymmetry that may be useful for detection and disentangling these same-sign lepton signals.

Finally, since the LQ production cross sections sharply drop with the LQ mass at the 13 TeV LHC, due to its limited phase-space for producing multi-TeV heavy particles, we have also calculated the projected same-charge LQ pair production cross sections, $\sigma(pp \to \phi_3\phi_3)$, at 27 and 100 TeV hadron colliders; the future planned HE-LHC and FCC-hh, respectively. As expected, we find that these future higher energy hadron colliders can extend the sensitivity to the LQ EFT dynamics up to masses of $M_\phi \gtrsim 5$ TeV and a NP scale of $\Lambda \sim 15$ TeV.

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APPENDIX: DIMENSION SIX OPERATORS FOR THE DOWN-TYPE SCALAR LQ $\phi(3,1,-\frac{1}{3})$

There are several classes of dimension six operators which correspond to the generic form of Eq. (5), which will be listed here.

The only ϕ^6 operator is

$$O^{(6)}_{\phi^6} = (\phi^*\phi)^3. \tag{A1}$$

There are no operators of the form $\phi^5 H$ due to gauge invariance and out of the $\phi^4 H^b D^{2-b}$ type operators there are only two nonredundant gauge invariant operators corresponding to the $b = 0$ and $b = 2$ cases,

$$O^{(6)}_{\phi^4 H^2} = (H^\dagger H)(\phi^*\phi)^2, \qquad O^{(6)}_{\phi^4 D^2} = |\phi|^2|D\phi|^2. \tag{A2}$$

Out of the operators that contain a ϕ^3 factor, the ones of the form $\phi^3 H^b D^{3-b}$ are absent since they violate either gauge (b odd) or Lorentz (b even) invariance. On the other hand, in the class $\phi^3\psi^2$ operators there are four gauge invariant combinations which can be constructed, all of the form $|\phi|^2\phi\bar{\psi}_L\psi_R$,

$$\begin{aligned} O^{(6)}_{\phi^3\psi^2} \in |\phi|^2(\bar{q}l^c\phi), \qquad & |\phi|^2(\bar{u}e^c\phi), \\ |\phi|^2(\overline{q^c}q\phi), \qquad & |\phi|^2(\overline{u^c}d\phi), \end{aligned} \tag{A3}$$

where the last two $\phi^3\psi^2$ operators above violate both baryon and lepton number. The operators that contain a ϕ^2

factor can be divided into two categories: the ones proportional to gauge invariant factor $|\phi|^2$ and the ones that contain ϕ^2 or $(\phi^*)^2$. The former case is straight forward, since it includes all operators involving an SU(3) singlet $\phi^\dagger\phi$ of the form,

$$O^{(6)}_{\phi^2 SM^4} \in |\phi|^2 \mathcal{O}^4_{\rm SM}, \tag{A4}$$

where $\mathcal{O}^4_{\rm SM}$ includes all the dimension 4 renormalizable terms of the SM Lagrangian. In addition, there are operators involving the SU(3) octet $\phi^\dagger\phi$ states of the form,

$$(\phi^\dagger \lambda^a D_\mu \phi)(\bar{q}\lambda^a \gamma^\mu q), \qquad (\phi^\dagger \lambda^a \phi) B^{\mu\nu} G^a_{\mu\nu}, \tag{A5}$$

where λ^a are the SU(3) Gell-Mann matrices and $B^{\mu\nu}$ is the SM SU(1) field strength. The latter case (i.e., operators which contain a ϕ^2 factor) is more elaborate, but it can be shown that there are only two nonredundant gauge invariant operators of this class, both in the form $\phi^2\psi^2 D$, where ψ^2 is composed out of one quark and one lepton,

$$O^{(6)}_{\phi^2\psi^2 D} \in \epsilon^{abc}\phi_a (D_\mu \phi)_b \bar{\ell}\gamma^\mu q_c, \qquad \epsilon^{abc}\phi_a (D_\mu \phi)_b \bar{e}\gamma^\mu d_c, \tag{A6}$$

where here a, b, c are color indices. Finally, the dimension six operator which contain only one LQ field have to be of the form $\phi\psi^2 H^b D^{2-b}$, where $0 \le b \le 2$ and ψ^2 is either a quark-lepton or quark-quark pair. For the $b = 2 (b = 1)$ case we find six(five) gauge invariant operators:

$$O^{(6)}_{\phi\psi^2 H^2} \in |H|^2 \phi^\dagger \bar{q} q^c, \qquad |H|^2 \phi \bar{q} l^c, \qquad |H|^2 \phi \bar{d^c} u, \qquad |H|^2 \phi \bar{u} e^c, \qquad \phi^\dagger (\bar{q} H)(H^\dagger q^c), \qquad \phi (\bar{q} H)(H^\dagger l^c), \tag{A7}$$

$$O^{(6)}_{\phi\psi^2 HD} \in (\bar{q} H)\gamma^\mu u^c D_\mu \phi^\dagger, \qquad (\bar{q}\tilde{H})\gamma^\mu d^c D_\mu \phi^\dagger, \qquad (\bar{q}\tilde{H})\gamma^\mu e^c D_\mu \phi, \qquad (\bar{l} H)\gamma^\mu u^c D_\mu \phi, \qquad (\bar{l}\tilde{H})\gamma^\mu d^c D_\mu \phi, \tag{A8}$$

where we have omitted the color indices and the antisymmetric tensor ϵ_{abc} in the above operators containing $3 \otimes 3 \otimes 3$ and $\bar{3} \otimes \bar{3} \otimes \bar{3}$ states.

The case of $b = 0$, i.e., operators of the type $\phi\psi^2 D^2$, contain four possible combinations of ψ^2 fields of the form,

$$O^{(6)}_{\phi\psi^2 D^2} \in D^2 \times \bar{q} q^c \phi^*, \qquad D^2 \times \bar{q} l^c \phi, \qquad D^2 \times \bar{d^c} u \phi, \qquad D^2 \times \bar{u} e^c \phi, \tag{A9}$$

where the notation above indicates that the two derivatives are to act on any of the fields; note though that $D_\mu D^\mu$ acting on a field gives a redundant operator, but $[D_\mu, D_\nu]$ does not. Thus, for example, $D^2 \times \bar{q} l^c \phi^\dagger$ corresponds to

$$D^2 \times \bar{q} l^c \phi \to (\bar{q} D_\mu l^c) D^\mu \phi, \qquad (\bar{q}\sigma_{\mu\nu} l^c) B^{\mu\nu}\phi, \qquad (\bar{q}\sigma_{\mu\nu}\sigma^I l^c) W^{\mu\nu}_I \phi, \qquad (\bar{q}\sigma_{\mu\nu}\lambda^A l^c) G^{\mu\nu}_A \phi. \tag{A10}$$

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