

Fresh look at the LHC limits on vector leptoquarks

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Vector leptoquarks (vLQs) are popular candidates for searching for physics beyond the Standard Model (SM). In this paper, we present updated exclusion limits on various vLQ species, accounting for all the relevant production mechanisms at the LHC. In particular, we highlight the critical role of indirect production and its interference with the SM Drell-Yan process. This interference can be constructive or destructive, depending on the specific quantum numbers of the vLQ, and significantly impacts the sensitivity of current searches. Furthermore, we demonstrate that including QCD + QED mixed pair production channels leads to a noticeable shift in model-independent mass limits. Additionally, we examine the validity of the full theory with vLQs and corresponding effective operators in the high mass regime. Overall, our analysis yields a substantial improvement in the exclusion limits on vLQs compared to the existing results in the literature.

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I. INTRODUCTION

Since their emergence within the unified theories [1–7], leptoquarks (LQs) have been some of the most widely studied hypothetical particles by particle physicists. They are color-triplet, hypercharged bosons, some of which can also carry a weak charge. LQs appear naturally around the TeV scale in many top-down new physics (NP) theories (see Ref. [8] for a review on LQ phenomenology and relevant references therein). LQs' unique ability to interact simultaneously with a lepton and a quark makes the LQ search program a prominent one at the LHC. The main reason behind their recent popularity is their potential role in resolving various persistent flavor and low-energy anomalies (the literature on this topic is vast; for instance, for the muon $g - 2$ anomaly, Ref. [9] reviews leptoquark-based explanations, along with other NP scenarios).

Investigation into their collider phenomenology began even before the start of the LHC data taking [10–13], and received a boost from the observed low-energy anomalies,

especially those in B -meson decays. Various aspects of LQ phenomenology, such as model building, derivation of exclusion limits, novel signatures and search strategies, etc., have been extensively studied in the literature (a nonexhaustive list of references are Refs. [14–52]). The direct searches for their productions at the LHC have put very stringent constraints on LQ parameters, such as their masses, branching ratios (BRs), and couplings (see the summary plots from the ATLAS [53] and CMS [54] collaborations for the current limits). Pair production (PP) of LQs is used in setting model-independent limits on the masses and BRs, as it is governed by QCD interactions and is practically insensitive to the LQ-quark-lepton couplings. A model dependence enters through the t -channel lepton exchange diagrams present in the PP. Initial LQ searches at the LHC focused on the PP channels induced by strong interactions to obtain model-independent exclusion limits. On the other hand, single productions (SPs), indirect productions (IPs, t -channel LQ exchange processes), and indirect interferences [IIs, interference of the IPs with the Standard Model (SM) background] are all sensitive to the LQ couplings, and hence, are used to set limits on their couplings. A similar method can also be applied in the context of other beyond the SM particles, as shown in Refs. [55,56].

References [15,45] showed that it is possible to recast the cross section upper limits in the dilepton-dijet final states from the direct LQ PP searches and obtain competitive limits on the LQ-lepton-quark couplings. This is due to the coupling-dependent SP and IP processes, which also lead to

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dilepton-dijet events after jet radiation. Many of these events also pass the selection cuts used in the PP searches. Moreover, IP processes interfere heavily (constructively or destructively, depending on the LQ species) with the SM Drell-Yan ($qq \rightarrow \ell\ell$) background. Recently, Ref. [45] has demonstrated that even though only a small fraction of the interference events pass the cuts optimized to select the PP events, the interference contribution is important for setting the exclusion limits on the unknown LQ couplings. The IP and II contributions can also alter the high- p_T dilepton tail distributions of the dilepton resonance searches (e.g., searches for a heavy neutral gauge boson Z'). Hence, one can use, e.g., a χ^2 method to set strong limits on LQ couplings using the IP and II LQ contributions for masses even beyond the reach of the resonance searches; this method has been used in Refs. [23,24,50] to constrain LQ couplings.

Previously, in Ref. [45], we obtained the exclusion limits on the masses and couplings of all scalar LQs (sLQs) by recasting the dimuon (and related) data. There, we also counted the PP and SP contributions along with the dominant IP and II contributions to obtain precise limits. Here, we estimate similar limits on all vector LQs (vLQs) listed in Ref. [8]. As we found in the case of sLQs, a QCD + QED mixed contribution to the PP can significantly improve the model-independent mass limits on vLQs, especially for those with a high electric charge. If an LQ is heavy, the initial partons may not have enough energy to produce it resonantly; however, the indirect processes (i.e, IP and II) might still be within the reach of the LHC. In such cases, it is possible to integrate the LQ in the t -channel out, and obtain an effective field theory (EFT) description of the $qq \rightarrow \ell\ell$ processes in terms of six-dimensional operators. In the high mass limit, the actual process with LQs and its EFT description should give close results. We compare these two

descriptions and show the validity of the LQ-EFT in the high mass/scale.

The rest of the paper is organized as follows: we discuss various production channels of vLQs in Sec. II. In Sec. III, we describe the recasting of experimental data. The resulting exclusion limits are presented in Sec. IV. Finally, we summarize our findings in Sec. V.

II. PRODUCTION OF VECTOR LEPTOQUARKS AT THE LHC

Five vLQ species can directly produce the $\ell\ell jj$ final state: U_1 , $\tilde{U}_1$, V_2 , $\tilde{V}_2$, and U_3 . Another vLQ, $\tilde{U}_1$, that couples exclusively to right-handed neutrinos, cannot directly produce the $\ell\ell jj$ final state (though decays of right-handed neutrinos can lead to dilepton plus multijet final states). Following the notations of Ref. [8], we summarize the interactions of vLQs with quarks and leptons in Table I. We consider two specific alignments of the vLQ mass basis. In the *down-aligned* scenario, the vLQ mass eigenstates are aligned with the mass basis of the down-type quarks, while in the *up-aligned* scenario, they are aligned with the up-type quarks. We do not consider the most general scenario where vLQ mass eigenstates are aligned neither with the up nor with the down quark sectors. Although the results for up-aligned and down-aligned scenarios can show minor differences in some cases, for illustration, we use the down-aligned scenario in our analysis. We assume $U_{\text{PMNS}} = 1$ throughout our analysis, since the neutrino oscillation length is irrelevant to our collider study, and all neutrino flavors collectively contribute to the missing energy.

There is also a renormalizable coupling of vLQ with gluon allowed by the gauge symmetries [10,11], as follows:

$$\mathcal{L} \supset -ig_s(1 - \kappa)\chi_\mu^\dagger T^a \chi_\nu G^{a,\mu\nu}, \quad (1)$$

TABLE I. Interactions for all vLQs in both the up- and down-aligned scenarios (these notations are followed from Ref. [8]). In the later part of the paper, we switch to a more convenient notations as explained in the Appendix (the naming convention of LQs and their coupling notations are introduced in Ref. [50]). Diquark interactions are excluded as they are not pertinent to our analysis. The SM gauge group representations of the vLQs are mentioned in parentheses.

LQ model	Down-aligned interactions	Up-aligned interactions
$U_1(\mathbf{3}, 1, 2/3)$	$(Vx_1^{LL})_{ij}\bar{u}_L^i \nu_L^j U_{1,\mu} + x_{1ij}^{LL}\bar{d}_L^i \gamma^\mu e_L^j U_{1,\mu} + x_{1ij}^{RR}\bar{d}_R^i \gamma^\mu e_R^j U_{1,\mu}$	$x_{1ij}^{LL}\bar{u}_L^i \nu_L^j U_{1,\mu} + (V^\dagger x_1^{LL})_{ij}\bar{d}_L^i \gamma_\mu e_L^j U_{1,\mu} + x_{1ij}^{RR}\bar{d}_R^i \gamma_\mu e_R^j U_{1,\mu}$
$\tilde{U}_1(\mathbf{3}, 1, 5/3)$		$\tilde{x}_{1ij}^{RR}\bar{u}_R^i \gamma^\mu e_R^j \tilde{U}_1$
$V_2(\mathbf{\bar{3}}, 2, 5/6)$	$-x_{2ij}^{RL}(\bar{d}_R^i \gamma^\mu \nu_L^j V_{2,\mu}^{1/3} - \bar{d}_R^i \gamma^\mu e_L^j V_{2,\mu}^{4/3}) + (V^* x_2^{LR})_{ij}\bar{u}_L^i \gamma^\mu e_R^j V_{2,\mu}^{1/3} - x_{2ij}^{LR}\bar{d}_L^i \gamma^\mu e_R^j V_{2,\mu}^{4/3}$	$-x_{2ij}^{RL}(\bar{d}_R^i \gamma^\mu \nu_L^j V_{2,\mu}^{1/3} - \bar{d}_R^i \gamma^\mu e_L^j V_{2,\mu}^{4/3}) + x_{2ij}^{LR}\bar{u}_L^i \gamma^\mu e_R^j V_{2,\mu}^{1/3} - (V^T x_2^{LR})_{ij}\bar{d}_L^i \gamma^\mu e_R^j V_{2,\mu}^{4/3}$
$\tilde{V}_2(\mathbf{\bar{3}}, 2, -1/6)$	$-\tilde{x}_{2ij}^{RL}(\bar{u}_R^i \gamma^\mu e_L^j \tilde{V}_2^{1/3} - \bar{u}_R^i \gamma^\mu \nu_L^j \tilde{V}_2^{-2/3})$	
$U_3(\mathbf{3}, 3, 2/3)$	$-x_{3ij}^{LL}\bar{d}_L^i \gamma^\mu e_L^j U_{3,\mu}^{2/3} + \sqrt{2}x_{3ij}^{LL}\bar{d}_L^i \gamma^\mu \nu_L^j U_{3,\mu}^{-1/3} + (Vx_3^{LL})_{ij}\bar{u}_L^i \gamma^\mu \nu_L^j U_{3,\mu}^{2/3} + \sqrt{2}(Vx_3^{LL})_{ij}\bar{u}_L^i \gamma^\mu e_L^j U_{3,\mu}^{5/3}$	$x_{3ij}^{LL}\bar{u}_L^i \gamma^\mu \nu_L^j U_{3,\mu}^{2/3} + \sqrt{2}x_{3ij}^{LL}\bar{u}_L^i \gamma^\mu e_L^j U_{3,\mu}^{5/3} + \sqrt{2}(V^\dagger x_3^{LL})_{ij}\bar{d}_L^i \gamma^\mu \nu_L^j U_{3,\mu}^{-1/3} - (V^\dagger x_3^{LL})_{ij}\bar{d}_L^i \gamma^\mu e_L^j U_{3,\mu}^{2/3}$

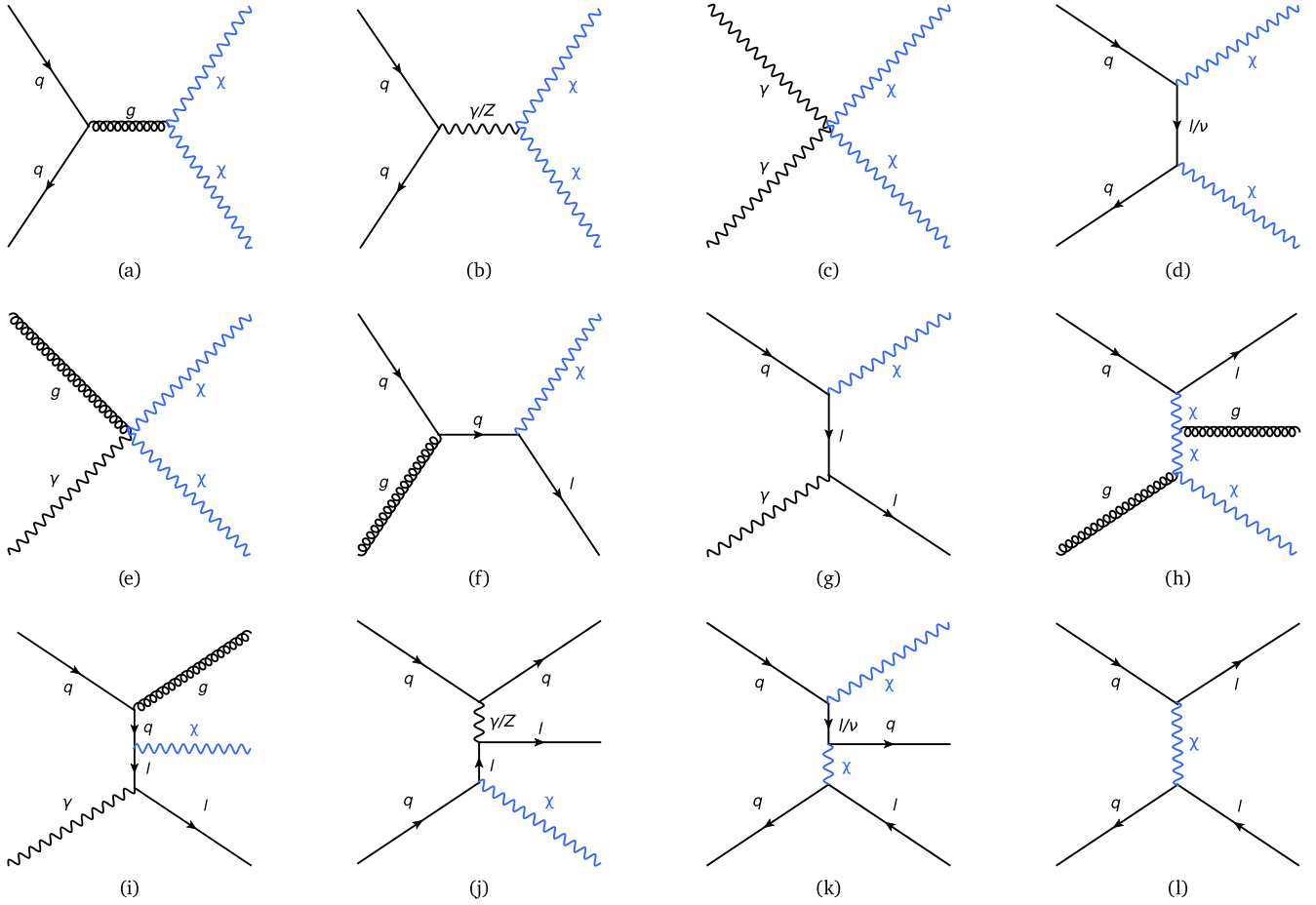


FIG. 1. Representative Feynman diagrams for pair (a)–(e), single (f)–(k) and indirect productions (l) of vLQs. A generic vLQ is denoted by χ .

where χ_μ denotes a generic vLQ and $G^{a,\mu\nu}$ denotes the field strength tensor of the gluon. The coupling κ can influence the production at the LHC.

At the LHC, vLQs produced via resonant (PP and SP) and nonresonant (IP and II) processes can contribute to the $\ell\ell jj$ final state. The interference can be constructive or destructive, depending on the type of vLQ involved. To understand how all production mechanisms contribute to the $\ell\ell jj$ final state, we systematically discuss all the production processes of vLQs (similar to the discussion presented in Ref. [58] for sLQs).

Pair production: pair-produced vLQs (same or different species) can decay to the $\ell\ell jj$ final state. The leading tree-level contributions are as follows:

- (1) $\mathcal{O}(\alpha_s^2 \alpha_\ell^0 \alpha_\chi^0)$: these model-independent, gg or qq fusion-initiated processes are purely QCD-driven. (Here, $\alpha_x = x^2/4\pi$ where x denotes a vLQ-quark-lepton coupling.) Figure 1(a) shows a representative diagram.
- (2) $\mathcal{O}(\alpha_s^2 \alpha_\ell^0 \alpha_\chi^0)$: these are qq [Fig. 1(b)] or $\gamma\gamma$ [Fig. 1(c)] initiated processes. Their contributions depend on the electric charge of the vLQ involved.
- (3) $\mathcal{O}(\alpha_s^0 \alpha_\ell^0 \alpha_\chi^2)$: in these processes, two on-shell vLQs are produced via the exchange of a lepton

in the t -channel [see Fig. 1(d)]. The vLQ-quark-lepton couplings control their contributions. The corresponding cross sections exhibit significant sensitivity to the initial-state quarks involved, owing to variations in their parton distribution functions (PDFs). The asymmetric vLQ PP, as introduced in Ref. [59], can also contribute at this order.

- (4) $\mathcal{O}(\alpha_s^1 \alpha_\ell^1 \alpha_\chi^0)$: the $g\gamma$ fusion processes are QCD + QED mixed processes that produce a pair of vLQs [Fig. 1(e)]. The interaction originates from the vLQ kinetic terms and is independent of NP couplings. These contributions become significant when the vLQ has a high electric charge.
- (5) $\mathcal{O}(\alpha_s^1 \alpha_\ell^0 \alpha_\chi^1)$: these contributions arise from the interference between diagrams of $\mathcal{O}(\alpha_s^2 \alpha_\ell^0 \alpha_\chi^0)$ and $\mathcal{O}(\alpha_s^0 \alpha_\ell^0 \alpha_\chi^2)$. This interference is destructive for all vLQ species.
- (6) $\mathcal{O}(\alpha_s^0 \alpha_\ell^1 \alpha_\chi^1)$: contributions at this order stem from the interference between diagrams of $\mathcal{O}(\alpha_s^0 \alpha_\ell^0 \alpha_\chi^2)$ and $\mathcal{O}(\alpha_s^1 \alpha_\ell^1 \alpha_\chi^0)$. Their impact on the total PP cross section is negligible.

The total PP cross section is expressed by summing up all the contributions of different orders, as follows:

$$\begin{aligned}\sigma_{\text{PP}}(M_{\ell_q}, x) = & \sigma_{\text{PP}}^{200}(M_{\ell_q}) + \sigma_{\text{PP}}^{110}(M_{\ell_q}) + \sigma_{\text{PP}}^{020}(M_{\ell_q}) \\ & + x^2 \overline{\sigma_{\text{PP}}^{101}(M_{\ell_q})} + x^2 \overline{\sigma_{\text{PP}}^{011}(M_{\ell_q})} \\ & + x^4 \overline{\sigma_{\text{PP}}^{002}(M_{\ell_q})},\end{aligned}\quad (2)$$

where σ_{PP}^{ijk} represents the contribution of order $\alpha_s^i \alpha_e^j \alpha_x^k$ to the PP process. The line over σ_{PP}^{ijk} indicates that those contributions are evaluated for coupling $x = 1$, i.e., $\overline{\sigma_{\text{PP}}^{ijk}(M_{\ell_q})} = \sigma_{\text{PP}}^{ijk}(M_{\ell_q}, x = 1)$.

Single production: the resonant productions of a single vLQ either with a lepton [i.e., the two-body single production (2BSP)] or a lepton and a jet [i.e., the three-body single production (3BSP)] can result in the $\ell\ell jj$ final state.

2BSP: There are two leading-order processes of this type:

- (1) $\mathcal{O}(\alpha_s^1 \alpha_e^0 \alpha_x^1)$, initiated by the qg fusion [Fig. 1(f)].
- (2) $\mathcal{O}(\alpha_s^0 \alpha_e^1 \alpha_x^1)$, initiated by the $q\gamma$ fusion [Fig. 1(g)].

3BSP: two types of diagrams can lead to 3BSP. In one type, a hard separable jet—either initial state radiation (ISR) or final state radiation (FSR)—is emitted from a 2BSP diagram, thus qualifying as a 3BSP process:

- (1) $\mathcal{O}(\alpha_s^{1+(1)} \alpha_e^0 \alpha_x^1)$: this contribution originates from the emission of a hard jet, arising from the $\mathcal{O}(\alpha_s^1 \alpha_e^0 \alpha_x^1)$ term of the 2BSP process [Fig. 1(h)]. The jet can be emitted by any colored particle present in the diagram.
- (2) $\mathcal{O}(\alpha_s^{(1)} \alpha_e^1 \alpha_x^1)$: the contribution at this order arises from the emission of a hard jet from the $\mathcal{O}(\alpha_s^0 \alpha_e^1 \alpha_x^1)$ diagrams of the 2BSP process [Fig. 1(i)].
- (3) $\mathcal{O}(\alpha_s^{(1)} \alpha_e^{1+(1)} \alpha_x^1)$: this contribution arises when an initial-state quark undergoes a splitting process, producing a hard quark and an electroweak boson. The boson then interacts with another quark, resulting in the production of a vLQ and a lepton [Fig. 1(j)].

Additionally, there are 3BSP processes that cannot be mapped to an ISR or FSR off a 2BSP process (see Ref. [15] for details):

- (4) $\mathcal{O}(\alpha_s^0 \alpha_e^0 \alpha_x^3)$: contributions at this order are governed solely by the NP coupling. These processes originate from the qg initial state [Fig. 1(k)].
- (5) Another contribution in the 3BSP processes arises when a vLQ is produced along with a lepton-jet pair from an off-shell vLQ, i.e., $pp \rightarrow \ell_q \ell_q^* \rightarrow \ell_q \ell_j$. This process occurs at one higher order in α_x than the six PP subcategories and has a different kinematics than that of the PP process.

To obtain the total SP contribution to the $\ell\ell jj$ final state, we follow the method described in Ref. [15]. Radiations from the Born-level 2BSP processes can resemble the 3BSP processes, potentially leading to overcounting if both processes are included separately. The total SP cross section

can be written as in the notation of Ref. [15] as follows:

$$\begin{aligned}\sigma_{\text{SP}} = & \underbrace{\left(\sigma_{2\text{BB}}^{\text{LO}}(\ell_q \ell) + \overbrace{\sigma_{(\ell_q \ell)}^{\text{virtual}} + \sigma_{2\text{BR}}^{\text{soft+collinear}}(\ell_q \ell j)}^{\text{Divergent terms}} + \sigma_{3\text{BNS}}^{\text{soft}}(\ell_q \ell j) \right)}_{\text{Negligible contribution}} \\ & + \underbrace{\left(\sigma_{2\text{BR}}^{\text{hard}}(\ell_q \ell j) + \sigma_{3\text{BNS}}^{\text{hard}}(\ell_q \ell j) \right)}_{\text{Main contribution}(\approx 3\text{BSP})} + \dots.\end{aligned}\quad (3)$$

The suffixes denote different contributions: 2BB represents two-body Born-level processes (i.e., 2BSP), 2BR indicates 2BSP with a radiation jet, and 3BNS denotes the NP contribution in 3BSP. The contribution to the $\ell\ell jj$ final state from SPs mainly arises from the 3BSP processes. Therefore, we can write the total SP cross section as follows:

$$\begin{aligned}\sigma_{\text{SP}}(M_{\ell_q}, x) = & x^2 \{ \overline{\sigma_{\text{SP}}^{021}(M_{\ell_q})} + \overline{\sigma_{\text{SP}}^{201}(M_{\ell_q})} + \overline{\sigma_{\text{SP}}^{111}(M_{\ell_q})} \} \\ & + x^6 \overline{\sigma_{\text{SP}}^{003}(M_{\ell_q})}.\end{aligned}\quad (4)$$

Indirect production: in this production channel, two leptons are generated via the t -channel exchange of a vLQ. Notably, the main contribution to the $\ell\ell jj$ final state originates from a single order:

- (1) $\mathcal{O}(\alpha_s^{(2)} \alpha_e^0 \alpha_x^2)$: this process contributes to the $\ell\ell jj$ final state when two hard jets emit from the diagram shown in Fig. 1(l).

Indirect interference: this contribution arises from the interference between the IP processes and the SM $\ell\ell$ processes (via a Z/γ boson). Note that the kinematics of the interference contribution is different from that of the IP process alone.

- (1) $\mathcal{O}(\alpha_s^{(2)} \alpha_e^1 \alpha_x^1)$: this contribution plays a decisive role in setting the exclusion limits, primarily due to its large interference with the SM. The nature of the interference—whether destructive or constructive—depends on the vLQ species involved.

The total nonresonant (NR) vLQ production cross section from both the IP and II can be expressed as

$$\sigma_{\text{NR}}(M_{\ell_q}, x) = x^2 \overline{\sigma_{\text{II}}^{011}(M_{\ell_q})} + x^4 \overline{\sigma_{\text{IP}}^{002}(M_{\ell_q})}.\quad (5)$$

In Fig. 2, we show the cross sections of various production modes of U_1 at different orders as functions of its mass. The U_1 interferes destructively with the SM background for both the LL and RR -type couplings. This is shown in Fig. 3.

In our results, we do not include any higher-order QCD correction factors to the vLQ production cross section. Since these are vector particles, unlike sLQs, their contributions can depend on the new gauge interactions beyond the tree level. One can calculate

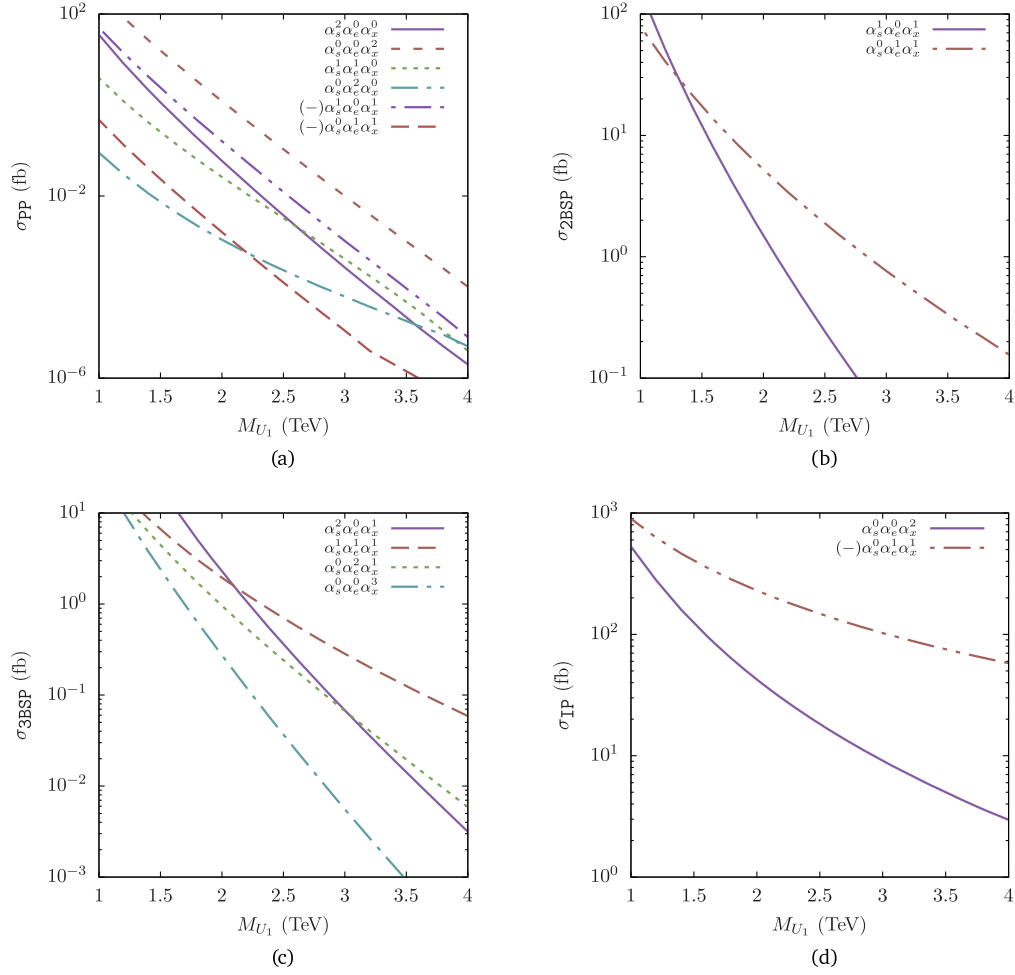


FIG. 2. PP, SP, IP, and II contributions for U_1 (subject to a $M_{\mu^+\mu^-} \geq 120$ GeV cut) as functions of M_{U_1} for $x_{10,12}^{LL} = 1$ at different orders of α_s , α_e , and α_x . The constructive/destructive interference contributions are indicated with $(\pm)$ signs in front.

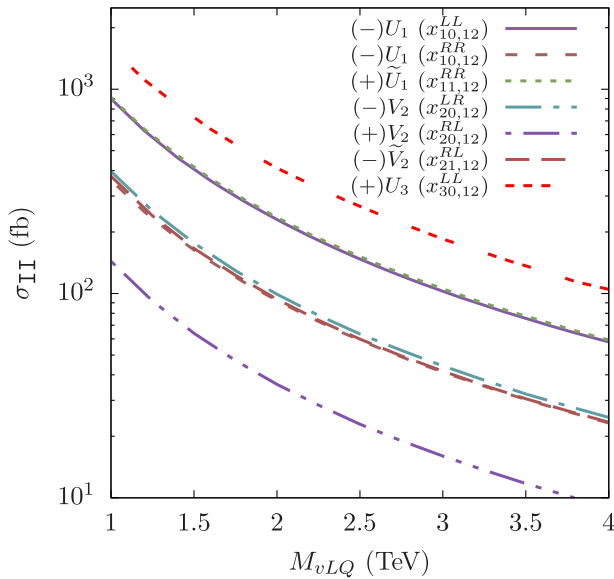


FIG. 3. Indirect interference contributions for different vLQs.

the higher-order corrections to vLQ production cross sections in a given UV complete framework [60–65] (see Refs. [66,67] for next-to-leading-order (NLO) corrections incorporating lepton density functions inside the proton). Similarly, NLO corrections in the context of sLQs can be found in Refs. [68–71]. Since our results are obtained using the leading-order (LO) cross sections, they are subject to large uncertainties. Therefore, the exclusion limits obtained below should be interpreted with care.

III. RECASTING THE LHC SEARCHES

A. Direct-search data

Both the ATLAS and CMS collaborations have searched for LQs in the PP channel. Here, we use the observed upper limit on the cross section in the $\mu\mu jj$ channel from the ATLAS measurement [57] to constrain the vLQ parameter space. Alternatively, the corresponding analysis by the CMS Collaboration [72] can also be used; we have verified that both yield similar bounds.

TABLE II. Cross sections (σ), the numbers of events ($\mathcal{N}$) surviving the $\mu\mu jj$ -selection cuts [57] for $\mathcal{L}_{\text{exp}} = 139 \text{ fb}^{-1}$ of integrated luminosity, and the cut efficiencies (ϵ) for three benchmark masses of U_1 . As in Fig. 2, we have set $x_{10,12}^{LL} = 1$ to obtain the numbers. We see that the importance of the indirect modes increases with the mass of the LQ.

Mass (U_1) (TeV)	Pair production			Single production			Indirect production			Indirect interference		
	σ_{PP} (fb)	ϵ_{PP}	$\mathcal{N}_{\text{PP}}$	σ_{SP} (fb)	ϵ_{SP}	$\mathcal{N}_{\text{SP}}$	σ_{IP} (fb)	ϵ_{IP}	$\mathcal{N}_{\text{IP}}$	σ_{II} (fb)	ϵ_{II}	$\mathcal{N}_{\text{II}}$
1.5	14.72	0.35	179.00	30.4	0.15	317.0	129.4	0.029	521.6	-406.2	0.01	-564.6
2.5	1.015×10^{-1}	0.53	1.87	1.3	0.21	19.0	18.86	0.028	73.4	-147.7	0.01	-205.3
4.0	1.045×10^{-4}	0.62	2.2×10^{-3}	6.7×10^{-2}	0.27	1.25	3.1	0.025	11.0	-57.8	0.01	-80.3

Our procedure to recast the direct bounds is as follows: we first pass all kinematically distinct production channels through the analysis cuts (tuned and validated to mimic those used in Ref. [57]) to determine the selection efficiencies. With those, we estimate the total number of the $\mu\mu jj$ events from vLQ by adding the contributions from these channels. We get the exclusion limits by equating the total vLQ yield (as a function of model parameters) to the observed number of events reported by the experiment, using the following relation:

$$\frac{\mathcal{N}_{\text{obs}}(M_{\ell_q})}{\mathcal{L}_{\text{exp}}} = \sigma_{\text{obs}}(M_{\ell_q}) \times \epsilon_{\text{exp}}(M_{\ell_q})$$

$$= \sum_i \sigma_i(M_{\ell_q}, \vec{x}) \times \beta_i(\vec{x}) \times \epsilon_i(M_{\ell_q}). \quad (6)$$

The sum is performed over various contributions: PP, SP, IP, and II. The symbols ϵ denotes selection cut efficiencies, β denotes BRs, and $\mathcal{L}_{\text{exp}}$ denotes the integrated luminosity. We use $\mathcal{N}_{\text{obs}}$ for $M_{\ell_q} = 2 \text{ TeV}$ for a conservative estimation, as choosing a lower mass point leads to weaker limits on vLQ parameters. Table II shows the number of events coming from different processes.

The vLQ interaction Lagrangians were implemented in FeynRules [73] in Ref. [50]; we use the Universal FeynRules Output model files [74] for event generation. We generate all Monte Carlo events at the LO using MadGraph5 [75]. We use the NNPDF2.3QED PDFs [76] with the dynamical renormalization and factorization scales. These events are then passed through PYTHIA8 [77] for parton showering and hadronization, and subsequently through Delphes [78] for detector simulation. Jets are clustered by the anti- k_T algorithm [79], as implemented in FastJet [80].

B. Dilepton-search data

As previously noted, high- p_T dilepton resonance searches can also constrain the vLQ parameter space. We use the latest CMS dimuon data [81] to derive bounds on vLQ couplings as functions of their masses. Both resonant and nonresonant vLQ production yield high- p_T leptons, affecting the tail of the dimuon spectrum. Adopting the χ^2 analysis method from Ref. [23], we derive 2σ exclusion limits by fitting the dimuon invariant mass distribution for fixed M_{ℓ_q} values. The CMS analysis

[81] does not impose any condition on the number of hard jets accompanying the lepton pair. Hence, all vLQ production modes—which produce dileptons—can contribute to the signal. We incorporate all such contributions within our statistical framework to derive exclusion limits. Since nonresonant production (IP and II) cross sections decrease at a slower rate with mass than the resonant channels (PP and SP), the exclusion limits in the high mass region are essentially determined by the IP and II contributions.

Recently, the CMS Collaboration has performed a search for nonresonant production of LQs in the dilepton channel, taking into account the interference with the SM dilepton processes [82]. Their analysis covers both sLQs and vLQs, and they presented exclusion limits in the mass-coupling plane. Our exclusion limits derived from the high- p_T dimuon searches agree well with their results.

IV. EXCLUSION LIMITS

In the limit $x \rightarrow 0$, vLQ mass exclusion limits are solely determined by the SM gauge couplings, and the PP becomes the only production channel at the LHC. In Table III, we summarize such model-independent exclusion limits on the masses of various vLQ species. The QCD column shows the limits obtained from the $\mathcal{O}(\alpha_s^2 \alpha_e^0 \alpha_x^0)$ contribution to the PP. The variation in limits across different vLQs arises from differences in their BRs.

TABLE III. Comparison of model-independent mass exclusion limits (in GeV) on various vLQ models with and without QED contributions, for $\kappa = 0$ and $\kappa = 1$. We assume vLQ decays through only one small coupling (shown in parentheses). The limits remain the same for second-generation quarks.

Model	$\kappa = 1$		$\kappa = 0$	
	QCD	QCD + QED	QCD	QCD + QED
$U_1(x_{10,12}^{LL})$	1746	1806	2038	2051
$U_1(x_{10,12}^{RR})$	1989	2056	2279	2295
$\tilde{U}_1(x_{11,12}^{RR})$	1989	2319	2279	2410
$V_2(x_{20,12}^{LR})$	2091	2275	2383	2451
$V_2(x_{20,12}^{RL})$	1989	2217	2279	2360
$\tilde{V}_2(x_{21,12}^{RL})$	1989	2010	2279	2283
$U_3(x_{30,12}^{LL})$	2027	2337	2310	2436

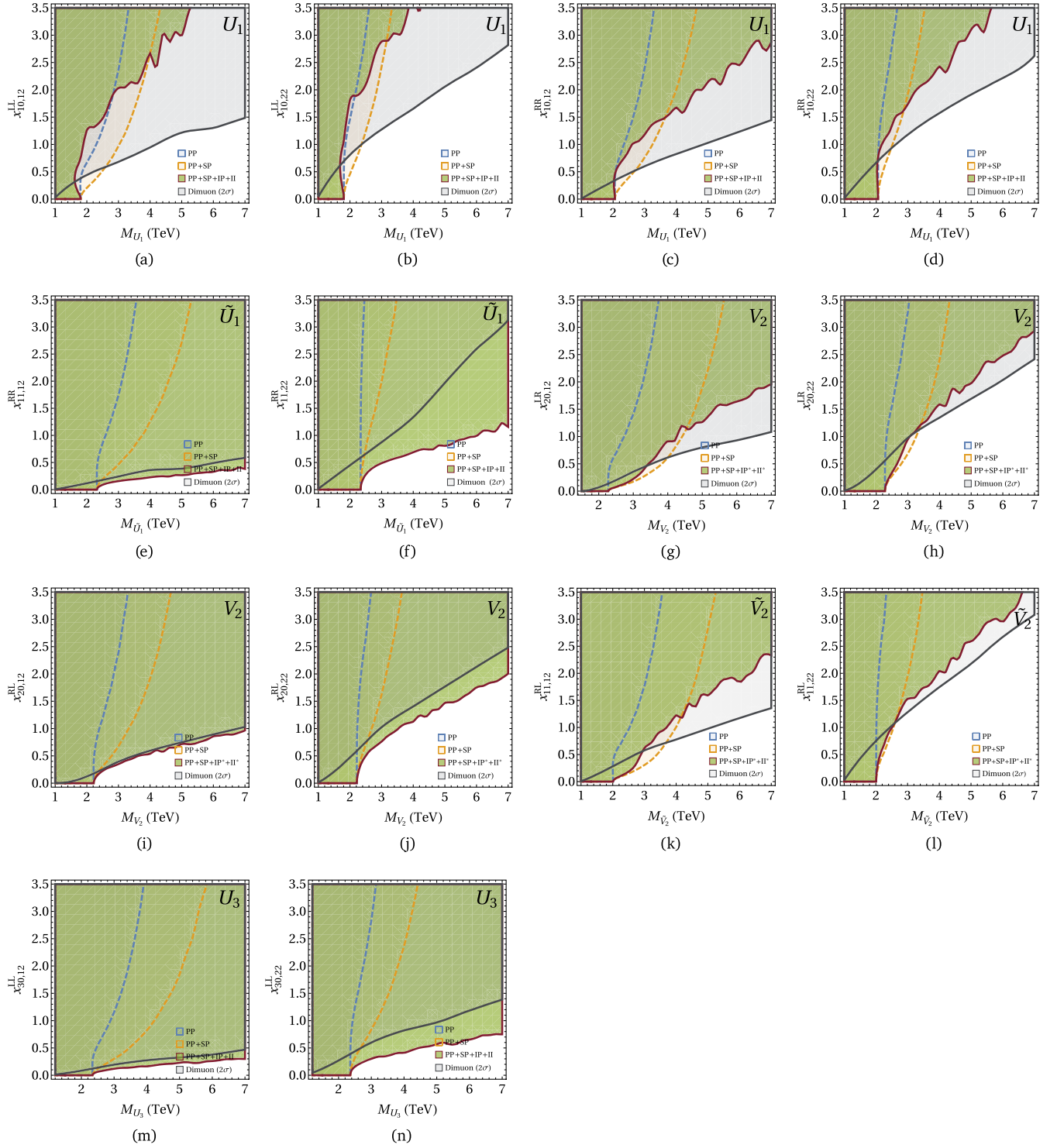


FIG. 4. Exclusion limits on all vLQs (a–d) for U_1 , (e,f) for $\tilde{U}_1$, (g–j) for V_2 , (k,l) for $\tilde{V}_2$, and (m,n) for U_3 when only one coupling is nonzero (see Appendix for notations). The contributions from the PP, PP + SP, and PP + SP + IP + II processes are separately marked. In the case of V_2 and $\tilde{V}_2$, we obtained the IP and II contributions in an EFT framework. These results are marked as IP* and II*.

So far, in the literature, the QCD + QED mixed contributions to the vLQ mass exclusion limits were usually ignored. These additional contributions at $\mathcal{O}(\alpha_s^1 \alpha_e^1 \alpha_x^0)$ enhance the PP cross section, improving the bounds. We present results for

two representative κ values [Eq. (1)]—the QCD + QED limits are stronger than the pure QCD limits for vLQs with larger electric charges. However, due to large uncertainties in the photon PDF in NNPDF, the limits can vary

significantly. For example, for $\tilde{U}_1$, they range from 2016 GeV to 2990 GeV.¹

A. The singlets (U_1 and $\tilde{U}_1$)

The $\mu\mu jj$ final state can emerge from U_1 interactions mediated by the following couplings: $x_{10,12}^{LL}, x_{10,22}^{LL}, x_{10,12}^{RR}$, and $x_{10,22}^{RR}$. Similarly, the couplings, $x_{11,12}^{RR}$ and $x_{11,22}^{RR}$ for $\tilde{U}_1$ can also lead to the same final state. In Figs. 4(a)–4(f), we present the (95% confidence limit) exclusion limits on these couplings as functions of mass. These limits are obtained by recasting the LQ search results presented in Ref. [57] in the $\mu\mu jj$ channels. We make a comparison of how the single coupling limits change if we consider only the PP, PP + SP, or all contributions, i.e., PP + SP + IP + II, together. We also present 2σ exclusion limits on couplings using the high- p_T dimuon (invariant mass distribution) tail on the same plot.

Due to the destructive nature of the U_1 interference contribution, the dimuon-search limits are more stringent than the direct-search limits. A similar effect is also observed for the S_1 in Ref. [45]. However, for $\tilde{U}_1$, the direct-search limits are stronger, owing to the constructive interference, similar to the case of $\tilde{S}_1$. For fixed values of M_{ℓ_q} and x , the IP and II cross sections are larger for those cases where first-generation quarks participate in the initial states, due to the larger PDFs compared to the second generation. As a result, vLQs coupling to first-generation quarks face tighter constraints. In the high-mass regime, where the limits are dominated by the IP and II contributions—which depend on the coupling strength x rather than on BRs—left- and right-handed couplings yield comparable exclusions. Some differences arise from the II terms, as the Z boson couples differently to the left- and right-handed fermions. The exclusion limits in the two-coupling plane for some fixed values of M_{ℓ_q} are presented in Figs. 5(a)–5(g).

B. The doublets (V_2 and $\tilde{V}_2$)

The V_2 can lead to the $\mu\mu jj$ final state through $x_{20,12}^{LR}, x_{20,22}^{LR}, x_{20,12}^{RL}$, and $x_{20,22}^{RL}$ couplings. With LR couplings, both $V_2^{1/3}$ and $V_2^{4/3}$ components contribute to the $\mu\mu jj$ final state with a negative II, whereas only $V_2^{4/3}$ contributes for RL couplings with a positive II. In Figs. 4(g)–4(l), we present the (95% confidence limit) exclusion limits [57] on these couplings as functions of mass. The exclusion limits in the two-coupling plane are presented in Figs. 5(h)–5(n).²

¹Currently, we are investigating the PDF dependence of this correction, which we will report elsewhere.

²In the case of V_2 and $\tilde{V}_2$, a straightforward evaluation of the interference contributions with our FeynRules models fails in MadGraph. While we are currently communicating with the authors of these packages to address the issue, here, we bypass the problem by taking an effective-operator approach by integrating out V_2 and $\tilde{V}_2$ as described later in Sec. IV F.

C. The triplet: U_3

For U_3 , the $x_{30,12}^{LL}$ and $x_{30,22}^{LL}$ couplings can lead to the $\mu\mu jj$ final state. Among the three components of U_3 , only $U_3^{2/3}$ and $U_3^{5/3}$ interact with charged leptons. The contribution to the II from the $U_3^{2/3}$ component is destructive, while that from $U_3^{5/3}$ is constructive. However, the overall effect of the II remains positive for the involvement of quarks of the first and second generations in the initial states due to $\sqrt{2}$ -times higher values of coupling for the positively interfering $U_3^{5/3}$ component (and also higher PDF for the first-generation case). For the third-generation case, the sign of the II is negative, as the positively interfering contribution is missing due to the absence of the top quark PDF. The single coupling limits as functions of mass are shown in Figs. 4(m)–4(n), whereas the two-coupling limits for some fixed masses are shown in Fig. 5(o).

D. Heavy jets: Third-generation quarks

So far, we have only considered light jets in the final states to obtain the exclusion limits. To obtain limits on the vLQ couplings with the third-generation quarks, we use the CMS $\mu\mu bb$ [83] and $\mu\mu tt$ [84] searches in the context of LQ searches. If vLQs couple to the bottom quark (this happens for U_1 , V_2 , and U_3), there is a possibility of IP and II productions (through the t -channel vLQ exchange, $bb \rightarrow \mu\mu$) leading to a muon pair in the final state. Additionally, we need two b -tagged jets for recasting the $\mu\mu bb$ results. Two b quarks in the final state can come from the gluon-splitting (four-flavor scheme). Although the probability of getting a bb pair from a gluon is suppressed compared to lighter quarks, it is not negligibly small. In the five-flavor scheme, we get a b quark directly from a proton, and the small probability is encoded in their small PDFs. We have worked in the five-flavor scheme and observed that substantial fractions of the SP, IP, and II events pass the selection cuts in the $\mu\mu bb$ search [83]. This leads to a stronger exclusion limit in the mass-coupling plane for vLQs [see Fig. 6(e)] that constructively interfere with the SM $bb \rightarrow \mu\mu$ processes. On the other hand, in the case of destructive interference, the exclusion limits become much weaker [see Figs. 6(a), 6(b), 6(d), and 6(h)].

If vLQs couple to the top quark, we do not get any contribution from the coupling-dependent IP and II processes, as the top quark PDF essentially vanishes. However, the 3BSP, like $pp \rightarrow \ell_q \mu t \rightarrow \mu t \mu t$, can produce the $\mu\mu tt$ final state. Combining PP and SP, we obtain the exclusion limits by recasting $\mu\mu tt$ results from Ref. [84] for $\tilde{U}_1$, V_2 , $\tilde{V}_2$, and U_3 as presented in Fig. 6.

E. Extrapolating the limits

The single coupling limits in Figs. 4 and 6 show that the mass exclusion limits become stronger as the value of the

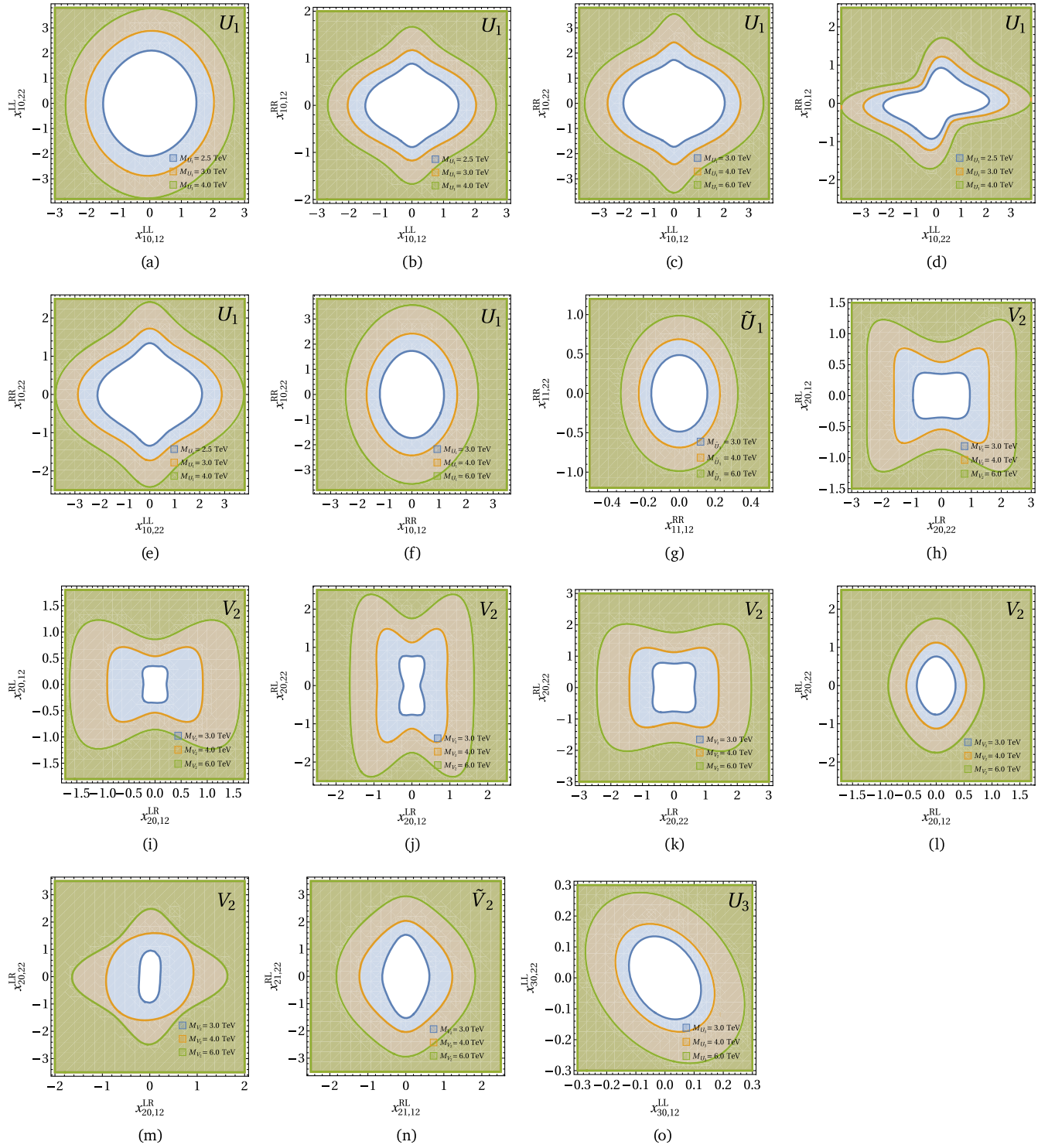


FIG. 5. Same as Fig. 4, except here, two couplings are nonzero: (a–f) for U_1 , (g) for $\tilde{U}_1$, (h–m) for V_2 , (n) for $\tilde{V}_2$, and (o) for U_3 .

couplings increases. In Table IV, we present the perturbative mass limits (when coupling approaches $x \approx \sqrt{4\pi}$) in the single coupling scenarios (see Ref. [85] for perturbative unitarity bounds on couplings).

F. Effective operator approach

A vLQ is exchanged in the t -channel in the IP and II processes, producing a lepton pair in the final state. If the vLQ mass is sufficiently high ($M_{\ell_q} \gg \sqrt{s}$), one can

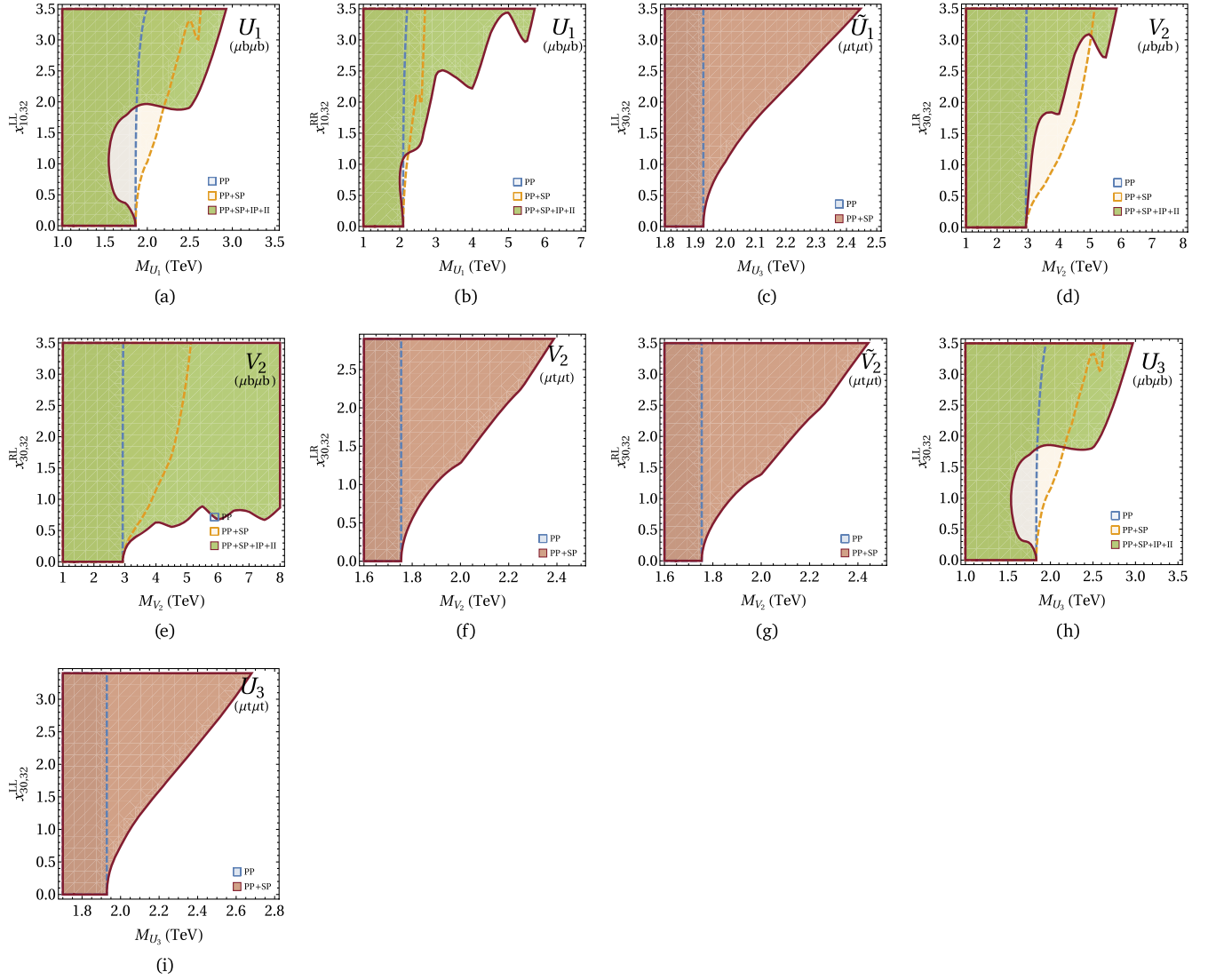


FIG. 6. Single-coupling exclusion limits when the quark involved is a third-generation quark.

TABLE IV. Extrapolated maximum mass exclusion limit (TeV) for $x \approx \sqrt{4\pi}$ for different sLQs. To obtain these, we have assumed only one coupling is nonzero at a time.

sLQ	x	Limit (TeV)	x	Limit (TeV)	Limit (TeV)	
					$(\mu b \mu b)$	$(\mu t \mu t)$
U_1	$x_{10,12}^{LL}$	5.3	$x_{10,22}^{LL}$	4.0	$x_{10,32}^{LL}$	2.8
	$x_{10,12}^{RR}$	8.1	$x_{10,22}^{RR}$	5.6	$x_{10,32}^{RR}$	6.0
$\tilde{U}_1$	$x_{11,12}^{RR}$	57.0	$x_{11,22}^{RR}$	21.5	$x_{11,32}^{RR}$	—
V_2	$x_{20,12}^{LR}$	12.0	$x_{20,22}^{LR}$	8.2	$x_{20,32}^{LR}$	6.0
	$x_{20,12}^{RL}$	27.0	$x_{20,22}^{RL}$	13.0	$x_{20,32}^{RL}$	40.0
$\tilde{V}_2$	$x_{21,12}^{RL}$	9.5	$x_{21,22}^{RL}$	6.6	$x_{21,32}^{RL}$	—
U_3	$x_{30,12}^{LL}$	85.0	$x_{30,22}^{LL}$	33.0	$x_{30,32}^{LL}$	2.8

integrate out the vLQs to obtain a description in terms of effective operators with appropriate Wilson coefficients (functions of vLQ parameters). This approach is closely related to that explored in Refs. [86,87], where similar EFT operators are studied. The effective operator approach should give us the same bounds as the full theory for heavy vLQs. We consider only the dilepton final states in the one-coupling scenarios and write down the effective operators obtained by integrating out the t -channel vLQ propagator. In Table V, we present a list of these operators and their Fierz-transformed forms for all vLQs.

Due to their small cross sections in the high mass region, the PP and SP processes hardly contribute to the exclusion limits, and the IP and II processes play the determining role. Hence, the limits from the effective operator approach agree well with the ones obtained by combining only the IP and II contributions. In Fig. 7, we see that the exclusion

TABLE V. Effective operators obtained by integrating out the t -channel vLQ (in IP processes) producing a lepton pair and their Fierz transformed versions.

vLQ	Effective operator	Fierz transformed effective operator
U_1	$-\frac{C_1}{\Lambda^2}([\bar{d}\gamma^\mu P_L \ell][\bar{\ell}\gamma_\mu P_L d] + [\bar{u}\gamma^\mu P_L \nu][\bar{\nu}\gamma_\mu P_L u])$ $-\frac{C_2}{\Lambda^2}[\bar{d}\gamma^\mu P_R \ell][\bar{\ell}\gamma_\mu P_R d]$	$-\frac{C_1}{\Lambda^2}([\bar{d}\gamma^\mu P_L d][\bar{\ell}\gamma_\mu P_L \ell] + [\bar{u}\gamma^\mu P_L u][\bar{\nu}\gamma_\mu P_L \nu])$ $-\frac{C_2}{\Lambda^2}[\bar{d}\gamma^\mu P_R d][\bar{\ell}\gamma_\mu P_R \ell]$
$\tilde{U}_1$	$-\frac{C_1}{\Lambda^2}[\bar{u}\gamma^\mu P_R \ell][\bar{\ell}\gamma_\mu P_R u]$	$-\frac{C_1}{\Lambda^2}[\bar{u}\gamma^\mu P_R u][\bar{\ell}\gamma_\mu P_R \ell]$
V_2	$-\frac{C_1}{\Lambda^2}([\bar{d}^C\gamma^\mu P_L \ell][\bar{\ell}\gamma_\mu P_L d^C] + [\bar{d}^C\gamma^\mu P_L \nu][\bar{\nu}\gamma_\mu P_L d^C])$ $-\frac{C_2}{\Lambda^2}([\bar{u}^C\gamma^\mu P_R \ell][\bar{\ell}\gamma_\mu P_R u^C] + [\bar{d}^C\gamma^\mu P_R \ell][\bar{\ell}\gamma_\mu P_R d^C])$	$\frac{C_1}{\Lambda^2}([\bar{d}\gamma^\mu P_R d][\bar{\ell}\gamma_\mu P_L \ell] + [\bar{d}\gamma^\mu P_R d][\bar{\nu}\gamma_\mu P_L \nu])$ $+\frac{C_2}{\Lambda^2}([\bar{u}\gamma^\mu P_L u][\bar{\ell}\gamma_\mu P_R \ell] + [\bar{d}\gamma^\mu P_L d][\bar{\ell}\gamma_\mu P_R \ell])$
$\tilde{V}_2$	$-\frac{C_1}{\Lambda^2}([\bar{u}^C\gamma^\mu P_L \ell][\bar{\ell}\gamma_\mu P_L u^C] + [\bar{u}^C\gamma^\mu P_L \nu][\bar{\nu}\gamma_\mu P_L u^C])$	$\frac{C_1}{\Lambda^2}([\bar{u}\gamma^\mu P_R u][\bar{\ell}\gamma_\mu P_L \ell] + [\bar{u}\gamma^\mu P_R u][\bar{\nu}\gamma_\mu P_L \nu])$
U_3	$-\frac{C_1}{\Lambda^2}([\bar{d}\gamma^\mu P_L \ell][\bar{\ell}\gamma_\mu P_L d] + [\bar{u}\gamma^\mu P_L \nu][\bar{\nu}\gamma_\mu P_L u])$ $+2[\bar{u}\gamma^\mu P_L \ell][\bar{\ell}\gamma_\mu P_L u] + 2[\bar{d}\gamma^\mu P_L \nu][\bar{\nu}\gamma_\mu P_L d]$	$-\frac{C_1}{\Lambda^2}([\bar{d}\gamma^\mu P_L d][\bar{\ell}\gamma_\mu P_L \ell] + [\bar{u}\gamma^\mu P_L u][\bar{\nu}\gamma_\mu P_L \nu])$ $+2[\bar{u}\gamma^\mu P_L u][\bar{\ell}\gamma_\mu P_L \ell] + 2[\bar{d}\gamma^\mu P_L d][\bar{\nu}\gamma_\mu P_L \nu]$

limits in the effective operator approach and the full theory are essentially the same in the high mass region, validating these two approaches in the high mass limits. From Fig. 7, we see that both approaches give the same result for $M_{\ell_q} \approx \Lambda \gtrsim 3$ TeV for vLQs coupling to the first-generation quarks. For vLQs coupling to higher-generation quarks, we noticed that the point of matching shifts to the higher-mass side.

As mentioned before, we present the limits in Fig. 4 using the effective operator approach for V_2 and $\tilde{V}_2$. This would give us a reasonable estimation of the exclusion limits for those vLQs, especially when they couple to the first and second generation quarks, for $M_{\ell_q} \gtrsim 3-4$ TeV. (Note that the PP and SP events for those vLQs are generated with MadGraph without any issue, and the cross sections are cross-checked with the other vLQs having similar interactions.)

V. SUMMARY AND DISCUSSIONS

In this paper, we derived exclusion limits on vLQs using the latest LHC searches for LQs in the $\mu\mu jj$

channel. PP searches typically provide largely model-independent bounds on vLQ masses and BRs. In contrast, constraints on the vLQ-quark-lepton couplings also arise from the high- p_T tail of the dilepton invariant mass spectrum (or the monolepton plus missing transverse energy, $\ell + \cancel{E}_T$, distribution). This study extended our earlier work [45], where we had advocated for a systematic combination of events from various production modes—PP, SP, IP, and II—for sLQs leading to the same (or experimentally indistinguishable) final states. To accurately determine the number of signal events at a given parameter point (specified by mass and couplings), our combined-signal approach offers a consistent method for estimating limits from processes contributing to the $\mu\mu jj$ channel or the $\mu\mu$ invariant mass spectrum.

For illustrations, we presented exclusion limits on all vLQs in the mass-coupling and two-coupling planes (for fixed masses). Additionally, we used the LHC $\mu b\mu b$ and $\mu t\mu t$ data to constrain the vLQ couplings to third-generation quarks. Our analysis can be extended to include any

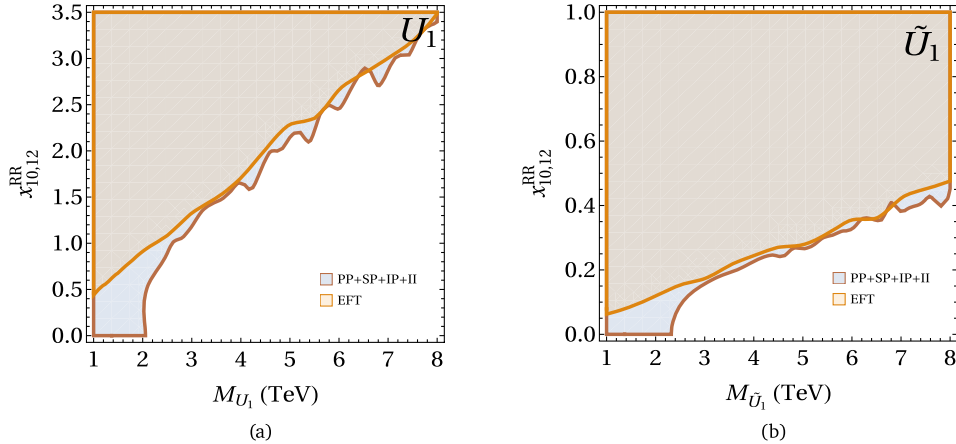


FIG. 7. Exclusion limits from the effective operator approach and the full theory in the high mass limit.

number of couplings simultaneously for all vLQs listed in Table I. (This will be incorporated in a forthcoming version of our `TooLQit` package [50].)

Typically, model-independent mass limits on vLQs are more stringent than those on sLQs, owing to the larger PP cross sections at the same mass. We derived such model-independent mass bounds using the PP on all vLQ species, and these results were summarised in Table III for two benchmark choices of κ . Furthermore, we included the impact of mixed QCD + QED contributions to the PP in our mass exclusion limits. We found that these contributions are important for vLQs with large electromagnetic charges. We have used the photon PDF from the NNPDF sets to estimate the QCD + QED effects. However, the photon PDF in the NNPDF sets has very large uncertainties. A detailed study of the impact of the associated PDF and scale uncertainties on the exclusion limits, including comparisons with photon PDFs from alternative PDF sets, is beyond the scope of the current paper (see Refs. [88–92] for precise determination of photon PDF). This will be investigated in future work.

Beyond the photon PDF, there are other sources of uncertainty; for example, scale and PDF variations can also significantly affect our LO results. For instance, for a 2.5 TeV $\tilde{U}_1$ with $x_{11,12}^{RR} = 1$, the IP process exhibits a scale uncertainty of approximately $\pm 2\%$ and a PDF uncertainty of $\pm 5\%$. In contrast, the PP process for the same benchmark shows significantly larger uncertainties: about $\pm 40\%$ from scale variation and $\pm 30\%$ from PDFs. Similarly, the 3BSP process has $\pm 35\%$ scale and $\pm 12\%$ PDF uncertainties. Hence, the limits we obtain should be considered as indicative. Since our results are obtained using the LO cross sections, they are subject to significant uncertainties. Consequently, the exclusion limits presented in this work should be interpreted with caution.

A vLQ is exchanged in the t -channel in the IP and II processes. In the high-mass limit, the vLQ can be integrated out, yielding corresponding four-fermion (two quarks-two leptons) effective operators with appropriate Wilson coefficients. In the heavy mass regime, the full theory containing the vLQ and the EFT approach should produce consistent results, thereby validating the EFT framework. In Table V, we listed all relevant effective operators—along with their Fierz-transformed versions—that arise from

integrating out the t -channel vLQs in the dilepton channel. The EFT approach is particularly powerful when the resonant production of heavy particles lies just beyond the direct reach of the collider. Our analysis can straightforwardly be extended to include scalar and vector LQs in the dilepton (same or different flavor) or $\ell + E_T$ channels. By recasting data from these channels, one can set bounds on the coefficients of the effective operators, which in turn constrain the underlying LQ parameters.

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DATA AVAILABILITY

The data that support the findings of this article are openly available [93].

APPENDIX: NOTATIONS FOR THE COUPLINGS

In our analysis, we adopted a compact and transparent notation (available at GitHub [93]) for the vLQ interactions. The generic vLQ-quark-lepton coupling is written as

$$x_{pq,ij}^{AB}. \quad (\text{A1})$$

Here, the index p specifies the $SU(2)$ representation of the sLQ, while $q = 1, 2$ distinguishes leptoquarks carrying a tilde ($\sim$) or a bar ($\bar{}$), respectively, and $q = 0$ otherwise. The indices i and j label the quark and lepton generations. The superscripts A and B indicate the chirality of the quark and lepton fields.

As an illustration, $x_{11,22}^{RR}$ corresponds to the interaction of $\tilde{U}_1$ with a right-handed second-generation quark and a right-handed second-generation lepton. Likewise, $x_{21,12}^{RL}$ denotes the coupling of the $SU(2)$ -doublet leptoquark $\tilde{V}_2$ to a right-handed first-generation quark and a left-handed second-generation lepton.

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