

Mutually elusive: Vectorlike antileptons and leptoquarks

Innes Bigaran^{1,2}, Bogdan A. Dobrescu², and Alessandro Russo³

¹*Department of Physics and Astronomy, Northwestern University,
2145 Sheridan Road, Evanston, Illinois 60208, USA*

²*Particle Theory Department, Fermilab, P.O. Box 500, Batavia, Illinois 60510, USA*

³*ETH Zürich, CH-8093 Zürich, Schafmattstrasse 16, Zurich, Switzerland*



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We study the properties of vectorlike fermions that have the same gauge charges as the Standard Model lepton doublets, but opposite lepton number. These antileptons undergo decays mediated by heavier scalar leptoquarks, while the symmetries of this renormalizable model protect the vectorlike fermions and the leptoquarks from standard decays probed so far at colliders. We derive upper limits on the new Yukawa couplings imposed by flavor-changing processes, including $B \rightarrow K \nu \bar{\nu}$ and $B_s - \bar{B}_s$ mixing, and show that they are compatible with prompt antilepton decays at the LHC for wide parameter ranges. If the new particles couple predominantly to second-generation quarks, then their collider probes involve multiple jets and two taus or neutrinos, and are hampered by large backgrounds. If couplings to third-generation quarks are large, then the collider signals involve top quarks, and can be probed more efficiently at the LHC. Even in that case, both the vectorlike fermion doublet and the leptoquarks remain more elusive than in models with standard decays.

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I. INTRODUCTION

Chirality is an essential characteristic of all Standard Model (SM) fermions. The left- and right-handed components of the SM fermions transform differently under the electroweak gauge group, and thus the quarks and leptons acquire masses by coupling to the Higgs field. Vectorlike fermions are distinct from chiral fermions in that the transformation properties under the SM gauge group are identical for both chiralities. A Dirac mass for a vectorlike fermion does not violate any SM symmetry. Therefore, the size of their masses is not tied to the Higgs doublet. If elementary vectorlike fermions exist, they represent a new form of matter, and dedicated searches for these are a major goal for the LHC experiments.

A vectorlike fermion that is a color singlet and carries lepton number $L = +1$ is referred to as a vectorlike lepton. It is typically assumed that the vectorlike leptons have the same charges under the $SU(3)_c \otimes SU(2)_W \otimes U(1)_Y$ gauge group as a SM lepton, namely $(1, 2, -1/2)$ for a weak doublet, or $(1, 1, -1)$ for a weak singlet. In that case, a Yukawa coupling to the SM Higgs doublet and a SM lepton field is allowed by the gauge symmetry, and induces

a mass mixing between the vectorlike lepton and the SM lepton [1–3]. As a result the vectorlike lepton decays into a SM lepton and a W , Z , or Higgs boson. Searches for electroweak production of vectorlike leptons that decay under these hypotheses have been performed by the CMS collaboration [4,5], and a particularly stringent lower mass limit, of 1.05 TeV, has been set on the vectorlike lepton doublet.

In this article, we study the properties of a vectorlike fermion ψ that transforms as $(1, 2, -1/2)$ under the SM gauge group, but which carries lepton number $L = -1$. As this is the opposite of the lepton number carried by the SM leptons, we generically refer to ψ as a vectorlike antilepton (VAL). We assume that ψ has a Yukawa interaction to a heavier leptoquark scalar (ξ) and a SM quark, and that all renormalizable interactions of ψ preserve lepton number. Consequently, the coupling of ψ to any SM lepton and the Higgs doublet is forbidden. Thus, the current searches for vectorlike leptons do not apply to ψ , as there are no standard decays (i.e., into a SM lepton and a SM boson). Instead, electroweak production of a ψ pair is followed, in our model, by three-body decays mediated by the leptoquark.

The reverse is also true; instead of decaying into a SM lepton and a quark (for a review, see [6]), the leptoquark predominantly decays into a VAL and a SM quark, leading to a cascade decay. As a result, the VAL and the scalar leptoquark mutually elude present searches for vectorlike leptons or leptoquarks at the LHC, and motivate new dedicated searches.

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Certain elements of the renormalizable model proposed here have been encountered in other contexts. Three-body decays of vectorlike quarks via a heavy scalar leptoquark have been analyzed in [7]. Three-body decays of a vectorlike lepton through a leptoquark of spin-1 have been studied in [8–11], leading to $\tau b \bar{b}$ final states searched for at CMS [12]; we will show that this is different than the final states obtained in the case of the VAL. Nonstandard decays of a vectorlike lepton, into a SM lepton and a light pseudoscalar, are discussed in [13] (see also [14]); in that model, though, the standard decays remain important.

On a more general note, vectorlike fermions appear in various frameworks beyond the SM, including bound states in models of quark and lepton compositeness [15], or higher multiplets of grand unified groups [16], so if they exist, vectorlike fermions would likely provide a window towards deeper principles. Thus, it is important to analyze a wide range of possible properties of vectorlike fermions, and to search for a variety of signals in experiments. The VALs studied here are as motivated as the vectorlike leptons that have mass mixing with the SM leptons, and thus are worth exploring in more detail.

In Sec. II we present the interactions of the VAL and of the leptoquark, analyze constraints from flavor processes, and compute the decay length of ψ . In Sec. III we study the LHC phenomenology of these particles, focusing on signals involving quarks of either second or third generation. We summarize our results in Sec. IV, where we also list the most important LHC searches that can test this model.

II. VECTORLIKE ANTILEPTON DOUBLET

Let us consider an extension of the SM that includes a weak-doublet vectorlike fermion $\psi = (\psi^0, \psi^-)^T$ that carries lepton number $L = -1$ and transforms under the $SU(3)_c \otimes SU(2)_W \otimes U(1)_Y$ gauge group as $(1, 2, -1/2)$. Since the SM lepton doublets ℓ_L^j ($j = 1, 2, 3$ is the generation index) have the same gauge charges as ψ , but opposite lepton number (by convention the electron has lepton number $+1$), it is appropriate to refer to ψ as an antilepton. A Dirac mass term for the ψ doublet $(-m_\psi \bar{\psi}_L \psi_R + \text{H.c.})$ is allowed both by the SM gauge symmetry (since the left- and right-handed components have the same charges) and by lepton number conservation.

We assume that the interactions of ψ preserve lepton number. Thus, any mixing between ψ and the SM leptons can be safely neglected, because it is forbidden by lepton number conservation. Note that the LEP searches for new charged particles imply $m_\psi \gtrsim 100$ GeV [17,18]. Even if the SM neutrinos acquire Majorana masses (which is not necessary, as the neutrinos may have Dirac masses), the lepton number breaking effects are suppressed by powers of $m_\nu/m_\psi \lesssim 10^{-12}$.

Although both the left- and right-handed components of ψ have $L = -1$, one could use instead their charge-conjugate

fields (ψ_L^c and ψ_R^c) which have $L = +1$. Thus, ψ^c may be viewed as a vectorlike lepton of wrong-sign hypercharge (compared to the SM lepton doublets), so the mixing between ψ^c and the SM leptons is still forbidden.

In the absence of additional fields or of higher-dimensional operators, the lightest particle of the ψ doublet remains stable. We introduce, however, a leptoquark scalar ξ that transforms as $(3, 2, +1/6)$ under $SU(3)_c \otimes SU(2)_W \otimes U(1)_Y$ that carries lepton number $L = -1$. Gauge and Lorentz invariances allow the following Yukawa interaction of ξ with ψ and with the SM up-type quark singlets (labeled u_{j_R} in the gauge eigenstate basis):

$$\lambda_{j_\psi} \xi^\dagger \bar{u}_{j_R}^c \psi_R + \text{H.c.} \quad (1)$$

Here λ_{j_ψ} are three dimensionless couplings. If the leptoquark mass satisfies $M_\xi > M_\psi$, then the above Yukawa interactions still do not allow the decay of the lightest particle of the ψ doublet. This is due to a global $U(1)_\psi$ symmetry under which both ψ and ξ carry charge $+1$, while the SM fields are singlets.

Renormalizable gauge-invariant terms that explicitly break $U(1)_\psi$ can be written as

$$\lambda_{j_\tau} \xi \bar{d}_{j_R} \ell_L^3 + \text{H.c.} \quad (2)$$

We assume that the above $U(1)_\psi$ breaking is suppressed, so the Yukawa couplings satisfy $\lambda_{j_\tau} \ll \lambda_{j_\psi}$. The Yukawa interaction (2) involves only the third-generation lepton doublet because we impose that τ -lepton number (L_τ) is conserved in this model, assigning $L_\tau = -1$ to both ψ and ξ . Focusing on the interaction with $\ell_L^3 = (\nu_{\tau_L}, \tau_L)$ is motivated in part by the larger τ Yukawa coupling to the Higgs doublet in the SM, which suggests that new physics might also have larger couplings to τ than to μ or e . A consequence of this coupling is that we will focus on collider signatures containing taus, which have larger backgrounds than those with final states containing muons or electrons.

Assigning baryon number $B = 1/3$ to ξ , and $B = 0$ to ψ , ensures that B is preserved by all the interactions introduced here. We point out that scalars that transform as $(3, 2, +1/6)$ under the SM gauge group have been studied in various contexts, such as squarks (labeled $\tilde{q}$) in supersymmetric extensions of the SM, or leptoquark doublets (sometimes labeled $\tilde{R}_2$ [6] or Φ_2 [19]). We use here the notation ξ to emphasize that the scalar participates in the peculiar interaction (1).

Besides the gauge interactions of ψ and ξ , and the Yukawa interactions discussed above, the model includes a single renormalizable interaction, namely $\lambda_H \xi^\dagger \xi H^\dagger H$. This Higgs portal coupling for the leptoquark changes both the Higgs production via gluon-fusion and the Higgs decay widths to $\gamma\gamma$ and γZ (see [6] for a review). However, the ensuing constraints on the λ_H coupling remain relatively mild for now [20].

A. Constraints from flavor processes

In the mass eigenstate basis for the up-type quarks, the Yukawa interactions (1) take the form

$$\xi^\dagger (\lambda_t \bar{t}_R^c + \lambda_c \bar{c}_R^c + \lambda_u \bar{u}_R^c) \psi_R. \quad (3)$$

Here, and in what follows, it is implicit that all the interaction terms are accompanied by their Hermitian conjugates in the Lagrangian. Note that there is a linear relation between the Yukawa couplings introduced in (1) and (3) through the 3×3 matrix U_R that relates the physical and gauge eigenstate bases for the right-handed up-type quarks. As U_R is not determined by any existing measurements, we focus on the physical couplings λ_t , λ_c , λ_u instead of λ_{j_ψ} .

The terms involving the up and charm quarks give rise to flavor-changing neutral current processes involving D mesons. In particular, there are contributions to $D^0 - \bar{D}^0$ meson mixing from box diagrams with ψ and ξ are running in the loop, such as the one shown in Fig. 1. The $\Delta C = 2$ operator generated by the box diagrams is given by

$$f(m_\psi/M_\xi) \frac{\lambda_c^2 \lambda_u^{*2}}{16\pi^2 M_\xi^2} (\bar{u}_R^\alpha c_L^\beta) (\bar{u}_R^\beta c_L^\alpha), \quad (4)$$

where α, β are color indices, and $f(m_\psi/M_\xi)$ is a function with values roughly of order one. A discussion of the above operator, labeled Q_3 , can be found in [21,22]. The comparison of the measured $D^0 - \bar{D}^0$ mixing with the SM prediction (which suffers from large QCD uncertainties related to the charm quark) imposes that the coefficient of the $(\bar{c}_L u_R)^2$ operator has a real part smaller than $6 \times 10^{-8} \text{ TeV}^{-2}$ and an imaginary part smaller by an extra factor of about 5 [23]. Based on that constraint, we estimate

$$\{\text{Re}(\lambda_c \lambda_u^*), \text{Im}(\lambda_c \lambda_u^*)\} \lesssim \{O(3), O(1)\} \frac{10^{-3} M_\xi}{1 \text{ TeV}}. \quad (5)$$

In what follows we assume $|\lambda_u| \ll |\lambda_c|$, so we neglect any effects due to λ_u . Thus, we do not need a more precise estimate of the above constraint.

Searches for rare top-quark decays can in principle constrain the product $|\lambda_t \lambda_c|$. For example, the $t \rightarrow cg$ and $t \rightarrow c\gamma$ decays would be induced by ψ and ξ running in a

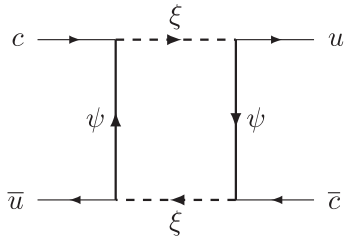


FIG. 1. Box diagram contributing to $D^0 - \bar{D}^0$ meson mixing, with the ψ VAL and ξ scalar running in the loop.

loop. This constraint is weak given the current sensitivity of LHC experiments, because the new t branching fractions are suppressed by both the loop factor and the large ξ mass.

The Yukawa interactions of ξ to the quarks of charge $-1/3$, shown in (2), can be written in the mass eigenstate basis for the down-type quarks as

$$\xi(\lambda_b \bar{b}_R + \lambda_s \bar{s}_R + \lambda_d \bar{d}_R) \ell_L^3. \quad (6)$$

Any of the three pairs of down-type quark flavors in these interactions is constrained by flavor-changing neutral processes. For example, there are contributions to $K^0 - \bar{K}^0$ meson mixing from box diagrams (one of them shown in Fig. 2) with $\xi = (\xi_u, \xi_d)^\top$ and the SM lepton doublet ℓ_L^3 running in the loop. This process imposes constraints on the real and imaginary parts of $\lambda_s \lambda_d^*$ similar to (5), but stronger [24],

$$\{|\text{Re}(\lambda_s \lambda_d^*)|, |\text{Im}(\lambda_s \lambda_d^*)|\} \lesssim \{3 \times 10^{-2}, 10^{-4}\} \frac{M_\xi}{1 \text{ TeV}}. \quad (7)$$

Another set of flavor constraints arises from the tree-level $s \rightarrow d \nu_\tau \bar{\nu}_\tau$ transition mediated by the leptoquark of electric charge $-1/3$ (i.e., ξ_d), as shown in Fig. 3. This contributes to the $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ decay, and imposes limits [24] on the real and imaginary parts of $\lambda_s \lambda_d^*$ that have a different M_ξ dependence than the limits from $K^0 - \bar{K}^0$ mixing. Conservatively the $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ limits can be written as

$$|\lambda_s \lambda_d| \lesssim 4 \times 10^{-4} \left(\frac{M_\xi}{1 \text{ TeV}} \right)^2. \quad (8)$$

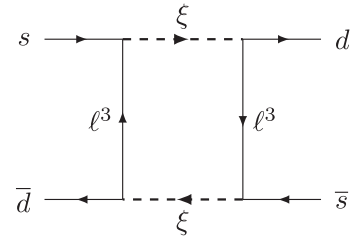


FIG. 2. Box diagram contributing to $K^0 - \bar{K}^0$ meson mixing, with the ξ scalar leptoquark and the SM third-generation lepton doublet ℓ^3 running in the loop.

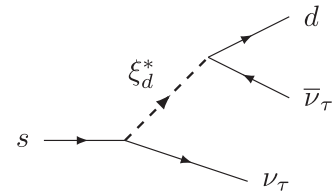


FIG. 3. Tree-level $s \rightarrow d \nu_\tau \bar{\nu}_\tau$ transition, which induces $K \rightarrow \pi \nu \bar{\nu}$ decays, mediated by an off-shell leptoquark of electric charge $-1/3$. The width of this transition depends on the Yukawa couplings of the ξ doublet to the s_R and d_R quarks.

We assume that $|\lambda_d| \ll |\lambda_s|$, so that the constraints from kaon measurements are satisfied without imposing an upper limit on $|\lambda_s|$. Combining this assumption with the earlier one made in the discussion of the bounds from $D^0 - \bar{D}^0$, we will neglect the couplings of ξ to first-generation quarks.

Even when the leptoquark does not couple to first-generation quarks, there are nontrivial constraints on the $\lambda_b \lambda_s^*$ product of couplings. These arise from comparing the measurements of the $B_s - \bar{B}_s$ meson mixing with the computation [25] of the box diagrams with ξ and ℓ^3 in the loop (replace the s and $\bar{d}$ fields by b and $\bar{s}$ in Fig. 2), which gives

$$|\lambda_b \lambda_s| \lesssim 3.6 \times 10^{-2} \frac{M_\xi}{1 \text{ TeV}}. \quad (9)$$

Here we used the fit result from [26] for the B_s mass difference: $\Delta m(B_s)/\Delta m(B_s)_{\text{SM}} = 1.14 \pm 0.08$.

Other constraints on $\lambda_b \lambda_s^*$ are set by the measurement of the $B^+ \rightarrow K^+ \nu \bar{\nu}$ branching fraction (B_{K^+}) [27], and by the existing limit on the $B^0 \rightarrow K^{*0} \nu \bar{\nu}$ branching fraction ($B_{K^{*0}}$) [28], which get contributions from the tree-level $b \rightarrow s \nu \tau \bar{\nu}$ transition [25,29] (as in Fig. 3 with appropriate quark replacements). The recent measurement $B_{K^+} = (23 \pm 7) \times 10^{-6}$ performed by Belle II [27] is 2.7σ higher than the SM prediction. Thus, one can derive both an upper limit and a lower limit on $|\lambda_b \lambda_s|$ at the 95% CL. However, the lower limit needs further experimental confirmation of potential new physics contributions to $B^+ \rightarrow K^+ \nu \bar{\nu}$, and thus we will not consider it as a constraint on this model.¹

Using input parameters from [30], we find that the upper limit imposed by the $B_{K^{*0}}$ limit is stronger than the one from the B_{K^+} measurement, and is given by

$$|\lambda_b \lambda_s| \lesssim 3.8 \times 10^{-2} \left(\frac{M_\xi}{1 \text{ TeV}} \right)^2. \quad (10)$$

Due to the different dependence on M_ξ , this constraint is weaker than (9) for $M_\xi > 950$ GeV. Given that we assume small $U(1)_\psi$ breaking, Yukawa couplings of ξ to b or s quarks are expected to satisfy $|\lambda_b|, |\lambda_s| \lesssim O(0.2)$, so the constraints (9) and (10) are rather mild.

B. Decays of the antileptons

For $m_\psi \ll M_\xi$ we can integrate out the leptoquark from (3) and (6) to obtain the following 4-fermion

¹Reconciling the measurement for B_{K^+} and the experimental limit on $B_{K^{*0}}$ may eventually require a new physics coupling to right-handed b and s quarks, generating a right-handed vector interaction, see e.g., [30]. We will not explore here the range of parameters in our model capable of reconciling these experimental results.

interactions that violate ψ number but preserve lepton number,

$$\frac{1}{M_\xi^2} [(\lambda_t \bar{\ell}_R^c + \lambda_c \bar{\ell}_R^c) \psi_R] [(\lambda_b \bar{b}_R + \lambda_s \bar{s}_R) \ell_L^3]. \quad (11)$$

The Lorentz indices here are contracted inside the square brackets, while the $SU(3) \times SU(2)_W$ indices are contracted between the two square brackets. These operators induce three-body decays of the VALs. There are four final states that involve a charm antiquark, with partial widths given at leading order by

$$\begin{aligned} \Gamma(\psi^- \rightarrow \bar{c} s \bar{\nu}_\tau) &\simeq \Gamma(\psi^0 \rightarrow \bar{c} s \tau^+) \simeq |\lambda_c \lambda_s|^2 \Gamma_0, \\ \Gamma(\psi^- \rightarrow \bar{c} b \bar{\nu}_\tau) &\simeq \Gamma(\psi^0 \rightarrow \bar{c} b \tau^+) \simeq |\lambda_c \lambda_b|^2 \Gamma_0, \end{aligned} \quad (12)$$

where small corrections due to the mass splitting between the two components of ψ are neglected, and we defined

$$\Gamma_0 = \frac{m_\psi^5}{2048 \pi^3 M_\xi^4}. \quad (13)$$

For $m_\psi \gtrsim m_t + m_\tau \approx 175$ GeV, decays that involve a top antiquark are kinematically allowed,

$$\Gamma(\psi^- \rightarrow \bar{t} s \bar{\nu}_\tau) \simeq \Gamma(\psi^0 \rightarrow \bar{t} s \tau^+) \simeq |\lambda_t \lambda_s|^2 \Gamma_0 F_1, \quad (14)$$

with the suppression due to the top mass taken into account through a function of fermion masses,

$$F_1 = 1 - 8 \frac{m_t^2}{m_\psi^2} + O(m_t^4/m_\psi^4, m_\tau^2/m_\psi^2). \quad (15)$$

For $m_\psi \gtrsim m_t + m_b + m_\tau \approx 180$ GeV, there are also decays that involve both a top antiquark and a b quark,

$$\Gamma(\psi^- \rightarrow \bar{t} b \bar{\nu}_\tau) \simeq \Gamma(\psi^0 \rightarrow \bar{t} b \tau^+) \simeq |\lambda_t \lambda_b|^2 \Gamma_0 F_2, \quad (16)$$

where

$$F_2 \approx F_1 + O(m_b^2/m_\psi^2). \quad (17)$$

If $|\lambda_c| \ll |\lambda_t|$, the decay modes of ψ that involve its Yukawa coupling to the top quark may dominate even for $m_\psi < m_t$. In that case, the top quark is off-shell, and the ψ VALs have four-body decays; $\psi^- \rightarrow W^- \bar{b} s \bar{\nu}_\tau$ and $\psi^0 \rightarrow W^- \bar{b} b \tau^+$.

Let us now consider the limit $|\lambda_t| \ll |\lambda_c|$, which implies that the total widths of ψ^- and ψ^0 are dominated by the partial widths (12). The ensuing decay length in the rest frame for the ψ VALs is

$$c\tau_\psi = \frac{1.25 \times 10^{-7} \text{ cm}}{|\lambda_c|^2 (|\lambda_s|^2 + |\lambda_b|^2)} \left(\frac{M_\xi}{1 \text{ TeV}} \right)^4 \left(\frac{100 \text{ GeV}}{m_\psi} \right)^5. \quad (18)$$

Even after taking into account the typical boost of the ψ pair produced at the LHC, the ψ VALs decay promptly into

visible final states provided $c\tau_\psi \lesssim 10^{-2}$ cm. We focus on this scenario in what follows.

A small mass splitting between the two components of the ψ doublet is induced at one loop. The contributions from the Z boson or photon running in the loop have been computed [31], and give $\delta m_\psi = m_{\psi^-} - m_{\psi^0} > 0$. This mass splitting is only logarithmically dependent on m_ψ , and continuously grows from 250 MeV to 350 MeV when m_ψ grows from 100 GeV to 1 TeV. This mass splitting allows a charged-current decay, $\psi^- \rightarrow \psi^0 \pi^-$, with a width

$$\Gamma(\psi^\pm \rightarrow \psi^0 \pi^\pm) \simeq \frac{G_F^2}{\pi} f_\pi^2 (\delta m_\psi)^3 \left(1 - \frac{m_\pi^2}{(\delta m_\psi)^2}\right)^{1/2}, \quad (19)$$

where $f_\pi \approx 130$ MeV is the pion decay constant, and G_F is the Fermi constant. The decay length in the rest frame associated with $\psi^- \rightarrow \psi^0 \pi^-$ is above 0.6 cm, and thus much longer than the one given in (18). Consequently, the branching fraction for $\psi^- \rightarrow \psi^0 \pi^-$ is small enough so we can neglect this process.

An implication of the expression for the decay length in (18) is that the couplings λ_s and λ_b , although expected to be small due to the $U(1)_\psi$ symmetry, cannot both be arbitrarily small if the ψ VALs decay promptly in the detector. In Fig. 4, the unshaded region corresponds to prompt ψ decays ($c\tau_\psi < 10^{-2}$ cm) and is allowed by the B -meson constraints (9) and (10). The mass ratio is fixed there at $M_\xi/m_\psi = 5$, which is near the lower end that allows² us to integrate out ξ and use the 4-fermion operators (11). For larger $M_\xi/m_\psi \equiv r_m$, all the lines in Fig. 4 move up by the following factors: $\sqrt{r_m/5}$ for the upper solid line, $r_m/5$ for the dashed line, and $(r_m/5)^2$ for the lower solid and dotted lines.

The Yukawa coupling of the b quark to the ξ leptoquark is fixed in Fig. 4 at $|\lambda_b| = |\lambda_s|$. For a different λ_b , the dashed and upper solid lines scale as $|\lambda_s/\lambda_b|^{1/2}$, while the dotted and lower solid lines scale as $\sqrt{2}/(1 + |\lambda_b/\lambda_s|^2)^{1/2}$. If instead of the $|\lambda_t/\lambda_c| \ll 1$ limit assumed in Fig. 4, we consider the opposite limit of dominant λ_t , then the region of prompt decays is reduced. This is because these decay widths decrease due to the additional phase-space suppression for final states with a top quark.

For $m_\psi < m_t$, the top quark can decay through an off-shell ξ into three-body final states, such as $t \rightarrow \bar{\psi}^0 b \tau^+$ and $t \rightarrow \psi^+ b \bar{\nu}_\tau$. The branching fractions of these processes are proportional to $|\lambda_t \lambda_b|^2$, and are highly phase-space suppressed. Nevertheless, it is worth searching for peculiar cascade decays of the top quark, for example $t \rightarrow \bar{\psi}^0 b \tau^+ \rightarrow (c \bar{b} \tau^-) b \tau^+$.

²For a smaller mass ratio, higher-order terms in $(M_\xi/m_\psi)^{-1}$ would need to be included. The overall impact would be a shorter decay length of ψ , and also stronger constraints from flavor processes, so the allowed region of prompt decays in Fig. 4 would shift towards smaller $|\lambda_s|$ values.

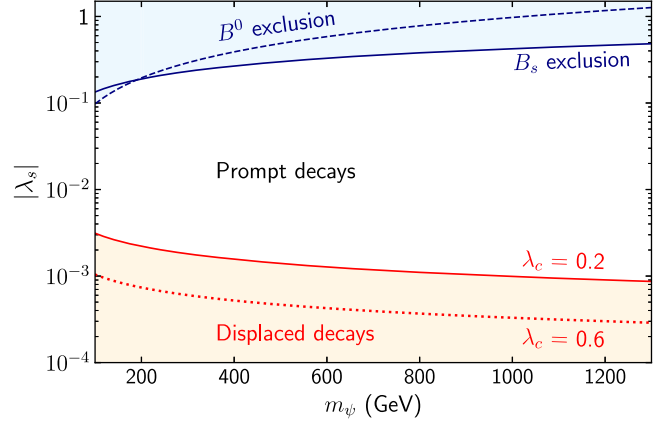


FIG. 4. Constraints on $|\lambda_s|$ as a function of the ψ mass, where λ_s is the Yukawa coupling of the s quark to the ξ leptoquark. Here we fix the b quark coupling $\lambda_b = \lambda_s$. The unshaded region is allowed and corresponds to prompt ψ decays, for $M_\xi = 5m_\psi$. The upper shaded (blue) region is ruled out by ξ contributions to $B_s - \bar{B}_s$ mixing (above the solid blue line) and $B^0 \rightarrow K^{*0} \nu \bar{\nu}$ (above the dashed blue line). The lower shaded (orange) region corresponds to ψ displaced vertices (rest-frame decay length $c\tau_\psi < 10^{-2}$ cm) for a c -quark coupling $\lambda_c = 0.2$ and $|\lambda_t/\lambda_c| \ll 1$. The edge of that region moves to the dotted red line when $\lambda_c = 0.6$.

III. COLLIDER PHENOMENOLOGY

The global $U(1)_\psi$ symmetry of (1) implies that direct production of ψ VALs at colliders can proceed only in pairs, through electroweak processes. Although the explicit $U(1)_\psi$ breaking allows in principle single leptoquark production, followed by $\xi \rightarrow \psi c$ or $\xi \rightarrow \psi t$, the small couplings in (6) and the need for a τ in the initial state effectively suppress this channel. Effects of a ξ_u exchanged in the t -channel [32] lead to $s\bar{s} \rightarrow \tau^+ \tau^-$ and $b\bar{b} \rightarrow \tau^+ \tau^-$ partonic processes (the latter is further PDF suppressed), but these are suppressed³ by $|\lambda_s|^4$ or $|\lambda_b|^4$.

It is also possible to pair produce the leptoquarks, with one of them decaying to a ψ and the other one decaying through the $U(1)_\psi$ breaking couplings (6). This would lead to a final state with a single ψ , but we will neglect it in what follows as the associated branching fractions are likely to be very small.

In this section, we will study pair production of the new particles, followed by decays via couplings involving SM quarks of either the second or third generation. We will not address here the case where the couplings to the second and third generation quarks are comparable, which would imply sizable branching fractions for final states involving t and s quarks, or b and c quarks.

³Recall that $|\lambda_s|, |\lambda_b| \ll 1$ because it is natural to impose that the global $U(1)_\psi$ symmetry is approximately preserved. If that assumption is relaxed, then traditional leptoquark searches [6,33] may dominate.

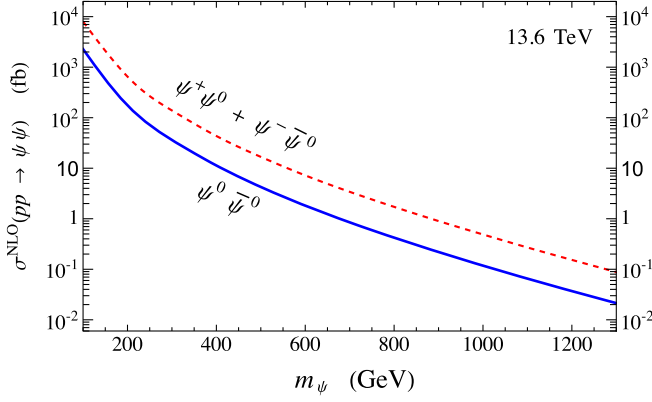


FIG. 5. NLO cross sections for pair production of a weak-doublet fermion through s -channel off-shell electroweak bosons at the 13.6 TeV LHC. The sum of the cross sections for $pp \rightarrow W^{+*} \rightarrow \psi^+ \psi^0$ and $pp \rightarrow W^{-*} \rightarrow \psi^- \bar{\psi}^0$ is given by the dashed red line. The cross section for $pp \rightarrow Z^* \rightarrow \psi^0 \bar{\psi}^0$ is given by the solid blue line, and is almost equal to the cross section for $\psi^+ \psi^-$ production (see Fig. 6).

A. Pair production of fermion doublet at the LHC

The ψ^0 and ψ^- fermions, which carry lepton number $L = -1$, have antiparticles $\bar{\psi}^0$ and ψ^+ fermions that carry lepton number $L = +1$. There are four main mechanisms for ψ production at hadron colliders,

$$\begin{aligned} u\bar{d} &\rightarrow W^{+*} \rightarrow \psi^+ \psi^0, \\ d\bar{u} &\rightarrow W^{-*} \rightarrow \psi^- \bar{\psi}^0, \\ q\bar{q} &\rightarrow Z^* \rightarrow \psi^0 \bar{\psi}^0, \\ q\bar{q} &\rightarrow \gamma^*, \quad Z^* \rightarrow \psi^+ \psi^-. \end{aligned} \quad (20)$$

Each of the above s -channel electroweak boson is off-shell, and only the dominant partonic processes are shown here.

To compute the cross sections for production of each of the four fermion pairs in proton-proton collisions, we generated model files using FeynRules [34] and run simulations using MadGraph_aMC@NLO [35]. The PDF set used here is NNPDF4.0 NLO [36]. The renormalization and factorization scales were set dynamically by MadGraph_aMC@NLO.

In Fig. 5 we show the ψ pair production cross sections, at the next-to-leading order (NLO) in α_s , as a function of the ψ mass, for pp collisions at $\sqrt{s} = 13.6$ TeV, which is the center of mass energy of the current Run 3 at the LHC. The sum of the cross sections for the $pp \rightarrow W^{+*} \rightarrow \psi^+ \psi^0$ and $pp \rightarrow W^{-*} \rightarrow \psi^- \bar{\psi}^0$ processes⁴ is shown as the dashed red line, while the cross section for $pp \rightarrow Z^* \rightarrow \psi^0 \bar{\psi}^0$ is shown as the solid blue line.

⁴For brevity, processes such as $pp \rightarrow W^{-*} + X \rightarrow \psi^- \psi^0 + X$, where X represents hadronic activity not related to the partonic processes (20), are denoted here without the X .

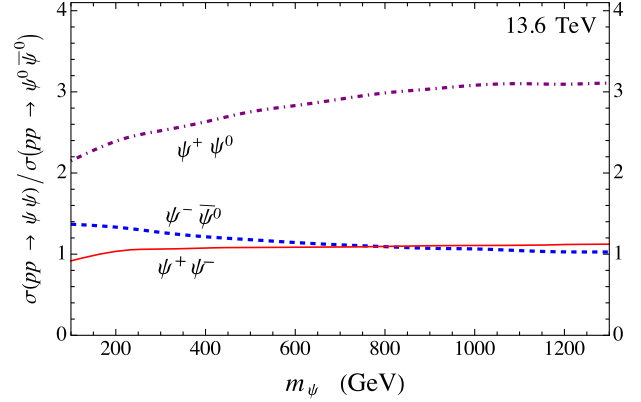


FIG. 6. Ratios of cross sections for ψ pair production ($\psi^+ \psi^0$, $\psi^- \bar{\psi}^0$, $\psi^+ \psi^-$) to the $pp \rightarrow Z^* \rightarrow \psi^0 \bar{\psi}^0$ cross section at a center-of-mass energy of 13.6 TeV.

The cross sections for the $pp \rightarrow \gamma^*$, $Z^* \rightarrow \psi^+ \psi^-$ and $pp \rightarrow Z^* \rightarrow \psi^0 \bar{\psi}^0$ processes differ by less than 10% for m_ψ in the 0.1–1 TeV range. Furthermore, the cross sections for $pp \rightarrow W^{-*} \rightarrow \psi^- \bar{\psi}^0$ and $pp \rightarrow Z^* \rightarrow \psi^0 \bar{\psi}^0$ differ by less than 15% for $m_\psi > 0.6$ TeV. These differences are accidentally small due to the interplay between the PDFs and the electroweak couplings of the quarks and ψ . In Fig. 6 we show the following ratios of cross sections: $\sigma(pp \rightarrow \gamma^*, Z^* \rightarrow \psi^+ \psi^-) / \sigma(pp \rightarrow \psi^0 \bar{\psi}^0)$ as a solid red line, $\sigma(pp \rightarrow W^{-*} \rightarrow \psi^- \bar{\psi}^0) / \sigma(pp \rightarrow \psi^0 \bar{\psi}^0)$ as a dashed blue line, and $\sigma(pp \rightarrow W^{+*} \rightarrow \psi^+ \psi^0) / \sigma(pp \rightarrow \psi^0 \bar{\psi}^0)$ as a purple dash-dotted line.

Note that the ratio of the cross sections for the processes mediated by W^+ and W^- grows from about 1.7 to 2.9 when m_ψ is increased from 0.1 to 1 TeV because the ratio of the u and d PDFs is approximately 2 at low parton-momentum fraction (x) where the valence quarks of the protons are counted.

The effective couplings (11) allow four decay modes for each of the ψ fermions. To reduce the number of final states, we now adopt a flavor ansatz for the coupling of the VALs and leptoquarks to quarks: ξ and ψ couple predominantly to one generation of SM quarks at a time.

B. Couplings to second-generation quarks

In this section, we assume that the couplings to second-generation quarks dominate, i.e., $|\lambda_b|^2 \ll |\lambda_s|^2$ and $|\lambda_t|^2 \ll |\lambda_c|^2$, and study the collider signals of the ψ and ξ particles.

1. Signals of the ψ antileptons with taus and jets

The main decay modes of the VALs in the case of dominant couplings to second-generation quarks are

$$\begin{aligned} \psi^0 &\rightarrow \bar{c}s\tau^+, & \bar{\psi}^0 &\rightarrow c\bar{s}\tau^-, \\ \psi^- &\rightarrow \bar{c}s\bar{\nu}_\tau, & \psi^+ &\rightarrow c\bar{s}\nu_\tau. \end{aligned} \quad (21)$$

The process of the largest cross section involving ψ VALs at proton-proton colliders is then

$$pp \rightarrow W^{\pm*} \rightarrow \psi^{\pm} \psi^{(\pm)*} \rightarrow (\tau^{\pm} jj)(\bar{\nu} jj), \quad (22)$$

as can be seen in the first diagram of Fig. 7. The signal is thus $\tau + 4j + \cancel{E}$, where j stands for any hadronic jet that is not b -tagged. Two of the four jets can in principle be tagged as c -jets, but that would reduce the signal rate because currently the efficiency of c -tagging is relatively low. The missing transverse energy ($\cancel{E}$) due to the neutrino is roughly $m_{\psi}/3$, so it provides limited means of rejecting the background.

An irreducible background to the $\tau + 4j + \cancel{E}$ signal is SM $W + 4j$ production followed by $W \rightarrow \tau\nu$, which has an inclusive cross section⁵ times branching fraction of the order of 70 pb [37] when the transverse momentum of each jet is above 30 GeV. Since this is at least an order of magnitude larger than the signal even for m_{ψ} near 100 GeV, and there are additional challenging backgrounds (such as pure QCD in the case of hadronic tau decays), these charged-current channels can be probed only if the experimental collaborations would perform advanced, dedicated searches.

Another channel is the neutral-current process

$$pp \rightarrow Z^* \rightarrow \psi^0 \bar{\psi}^0 \rightarrow (\tau^+ jj)(\tau^- jj), \quad (23)$$

which is shown in the second diagram of Fig. 7. This $\tau^+ \tau^- + 4j$ signal has a smaller cross section (see Fig. 5), but the backgrounds are also significantly smaller. Most notable backgrounds include SM production of $Z + 4j$ followed by $Z \rightarrow \tau^+ \tau^-$, or $W^+ W^- + 4j$ followed by each W boson decaying to $\tau\nu$. The small cross sections of these backgrounds, and the additional opportunities at rejecting the background, make $\tau^+ \tau^- + 4j$ the most sensitive channel.

A CMS search for a $\tau^+ \tau^- + 4b$ signal [12], arising from pair production of a vectorlike lepton doublet that decays through an off-shell vector leptoquark [8–11], has set a lower mass limit of 500 GeV. That CMS search, based on Run 2 data, relied on at least three b -tags to reduce the backgrounds. Given that our $\tau^+ \tau^- + 4j$ signal does not include b -jets, the highest m_{ψ} that a dedicated search with Run 2 data can explore is substantially below 500 GeV.

The remaining ψ pair production channel,

$$pp \rightarrow \gamma^*, \quad Z^* \rightarrow \psi^+ \psi^- \rightarrow (\nu jj)(\bar{\nu} jj), \quad (24)$$

leads to a difficult $4j + \cancel{E}$ signal. Searches at the LHC in similar channels are sensitive to cross sections [39] larger

⁵The cross sections quoted in [37] are smaller because they are computed in the presence of additional cuts. We thank Stefan H  che for providing the $W + 4j$ cross sections computed with sherpa [38] in the presence of only the jet $p_T > 30$ GeV and $\eta < 5$ cuts.

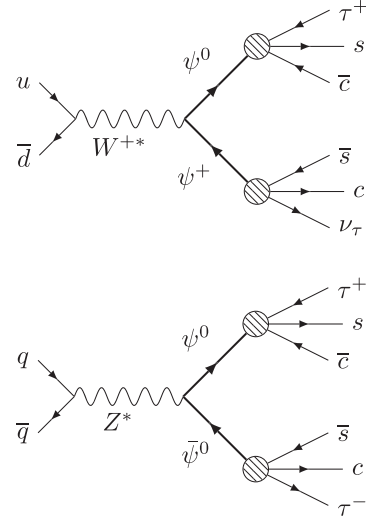


FIG. 7. Representative diagrams for the production and decay of the VALs. The blobs represent the effective interactions mediated by a heavy ξ leptoquark. The charged-current process leads to a $\tau^+ + 4j + \cancel{E}$ signal, while the process mediated by an off-shell Z leads to a $\tau^+ \tau^- + 4j$ signal.

by an order of magnitude compared to the one for $\psi^+ \psi^-$ given in Figs. 5 and 6. Nevertheless, future dedicated searches for the VALs may combine this channel with $\tau^+ \tau^- + 4j$ and $\tau + 4j + \cancel{E}$ to increase the sensitivity.

We will not attempt to predict the range of m_{ψ} that the ATLAS and CMS Collaborations can explore, but we can safely conclude that the mass reach in the ψ pair production channels is much smaller than in the case of a vectorlike lepton doublet that decays into two SM particles via mass mixing with the tau (the current lower mass limit in that case is 1.05 TeV [5]).

The cross section for pair production of VALs at a muon collider or an $e^+ e^-$ collider is large enough [40,41], and the backgrounds are much smaller than at hadron colliders. Therefore, a lepton collider with $\sqrt{s} \geq 2m_{\psi}$ would allow a thorough investigation of the VAL properties.

2. Signals of the ξ leptoquarks with taus and jets

Given that the breaking of the global $U(1)_{\psi}$ symmetry is small, and that the couplings to second-generation quarks are assumed to be dominant in this section, the largest interaction of the leptoquark doublet ξ is provided by the $\lambda_c \xi^{\dagger} \bar{c}_R^c \psi_R$ term in (3). This implies that the dominant decay modes for the ξ_u and ξ_d leptoquarks are

$$\xi_u \rightarrow c\psi^0, \quad \xi_d \rightarrow c\psi^-, \quad (25)$$

for $M_{\xi} > m_{\psi}$. Each ψ fermion produced in a ξ decay subsequently undergoes a three-body decay into SM fermions, mediated by an off-shell ξ . As a result, leptoquark

pair production is followed by the following cascade decays:

$$\begin{aligned} pp &\rightarrow \xi_u \xi_u^\dagger \rightarrow (c\psi^0)(\bar{c}\bar{\psi}^0) \rightarrow (\tau^+ jjj)(\tau^- jjj), \\ pp &\rightarrow \xi_d \xi_d^\dagger \rightarrow (c\psi^-)(\bar{c}\bar{\psi}^+) \rightarrow (\bar{\nu}_\tau jjj)(\nu_\tau jjj). \end{aligned} \quad (26)$$

Thus, pair production of the leptoquark of electric charge $+2/3$ (ξ_u) leads to a $\tau^-\tau^+ + 6j$ signal with a topology shown in Fig. 8.

Existing searches for leptoquarks coupling to a tau and a quark assume a $(\tau b)(\tau b)$ or $(\tau t)(\tau t)$ signal, and set lower mass limits of 1.19 TeV [42] and 1.43 TeV [5,43], respectively. Compared to the existing $(\tau b)(\tau b)$ search, our $\tau^-\tau^+ + 6j$ signal does not include b -jets and has much softer taus and jets due to the larger number of final-state particles. Consequently, the backgrounds are much larger, so we expect that the mass reach of even a dedicated search for ξ_u with the Run 2 data would be below 1 TeV.

Pair production of the leptoquark of electric charge $-1/3$ (ξ_d) leads to a $6j + E$ signal. Current leptoquark searches in the jets + E channel [44] assume a $(\nu j)(\nu j)$ signal, and set a lower limit mass limit of 1.14 TeV. As our signal is plagued by much larger backgrounds, due to larger jet transverse momenta and more missing transverse energy, we again expect that M_ξ can be substantially below 1 TeV. Thus, the ξ leptoquarks are clearly more elusive than the leptoquarks searched so far in LHC experiments.

C. Couplings to third-generation quarks

Let us now consider the case where the couplings to third-generation quarks dominate, i.e., the Yukawa couplings in Eqs. (3) and (6) satisfy $|\lambda_s|^2 \ll |\lambda_b|^2$ and $|\lambda_c|^2 \ll |\lambda_t|^2$. As this scenario leads to signals involving t and b quarks, the ψ and ξ particles are less elusive than in the case of couplings to second-generation quarks discussed in Sec. III B.

1. Signals of the ψ antileptons with top quarks

If the VALs are sufficiently heavy, $m_\psi > m_t + m_b + m_\tau \approx 180$ GeV, then their main decay modes are

$$\psi^0 \rightarrow \bar{t}b\tau^+, \quad \psi^- \rightarrow \bar{t}b\bar{\nu}_\tau. \quad (27)$$

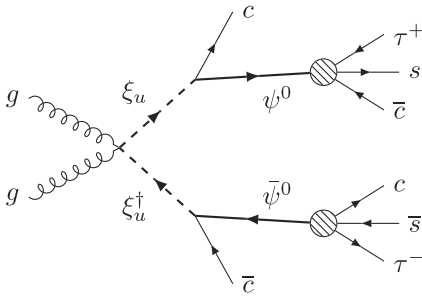


FIG. 8. Representative diagram for pair production of the ξ_u leptoquark followed by cascade decays, leading to a $\tau^-\tau^+ + 6j$ signal at hadron colliders.

In this case the main signal at proton-proton colliders is $\tau^\pm W^+ W^- + 4b + E$, and arises from the diagram with an off-shell W^+ shown in Fig. 9 and the analogous diagram with an off-shell W^- .

Tagging three or four of the b -jets substantially reduces several backgrounds (e.g., $t\bar{t}$ plus a W or a Z , and $WZb\bar{b}$) without significantly depleting the signal. For a b -tagging efficiency ϵ_b in the 0.8–0.9 range [45,46], the combined efficiency of at least three b -tags, $\epsilon_{3b} = \epsilon_b^3(4 - 3\epsilon_b)$, is in the 82%–95% range.

Nevertheless, irreducible backgrounds with four b -jets exist in the SM, including $t\bar{t}h^0$ or $WW + 4b$. The presence in the final state of two muons or electrons of same charge, due to leptonic decays of the τ or W bosons, suppresses these backgrounds, but at the cost of also reducing the signal. The combined branching fractions for the $\tau^\pm W^+ W^- + 4b + E$ final state to include two same-sign charged leptons ($\ell^\pm \ell^\pm$ where ℓ stands for an electron or a muon), or three charged leptons ($\ell^\pm \ell^+ \ell^-$), are

$$\begin{aligned} B_{\ell^\pm \ell^\pm} &= 2B_{\tau\ell}B(W \rightarrow \ell\bar{\nu})(2 + B_{\tau\ell}) \approx 8.3\%, \\ B_{3\ell} &= 2B_{\tau\ell}B(W \rightarrow \ell\bar{\nu})^2(2 + B_{\tau\ell})^2 \approx 2.0\%, \end{aligned} \quad (28)$$

where $B_{\tau\ell}$ is the average of the τ branching fractions into $\mu\nu\bar{\nu}$ and $e\nu\bar{\nu}$, and lepton universality is taken as a good approximation in the W branching fractions [47].

Even these $\ell^\pm \ell^\pm + (\geq 3)b$ and $3\ell + (\geq 3)b$ signals have SM backgrounds, especially from $t\bar{t}t\bar{t}$ and $t\bar{t}t$ productions. There are also important backgrounds due to leptons produced in meson decays or in photon conversions, and due to the misidentification of the lepton charge or of the b -tags.

Despite all these complicated backgrounds, we could have a rough estimate for the Run 3 reach in ψ pair production by comparing to the $pp \rightarrow t\bar{t}t\bar{t}$ process, which has recently been observed by the ATLAS and CMS Collaborations [48,49]. With 138–139 fb $^{-1}$ of data at $\sqrt{s} = 13$ TeV, the measured $t\bar{t}t\bar{t}$ cross section (σ_{4t}) has a central value of about 18 fb for CMS and 24 fb for ATLAS, and are within 1σ of each other and consistent with the SM

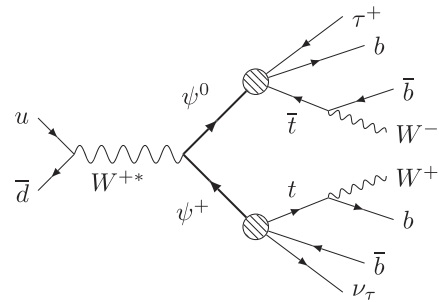


FIG. 9. Pair production of VALs through an off-shell W^+ , followed by decays induced by couplings to third-generation quarks. This process leads to a $\tau^+ W^+ W^- + 4b + E$ signal.

prediction [50]. The combined leptonic branching fractions for ψ signals, given in (28), are smaller than the analogous ones for $t\bar{t}\bar{t}$ ($B_{\ell^\pm\ell^\pm} \approx 11\%$, $B_{3\ell} \approx 4.3\%$). Furthermore, the backgrounds are larger at $\sqrt{s}=13.6$ TeV compared to 13 TeV, so even with a larger integrated luminosity of 300 fb^{-1} the cross section sensitivity to ψ pair production cannot be smaller than roughly 20 fb (*i.e.*, about the measured σ_{4t} in Run 2). This corresponds to a mass reach $m_\psi \lesssim 500$ GeV, as can be seen in Fig. 5.

The reach in m_ψ can be increased by adding the $pp \rightarrow Z^* \rightarrow \psi^0\bar{\psi}^0 \rightarrow \tau^+\tau^-W^+W^- + 4b$ process, which has a smaller production cross section (by a factor of about 4, see Figs. 5 and 6) but larger leptonic branching fractions ($B_{\ell^\pm\ell^\pm} \approx 17\%$, $B_{3\ell} \approx 10\%$). This process also leads to final states with hadronic τ decays that would further increase the mass reach. Even though the $pp \rightarrow \psi^+\psi^- \rightarrow t\bar{t}b\bar{b} + \cancel{E}$ process does not produce two same-sign leptons, it can be useful due to the missing transverse energy, which suppresses the $t\bar{t}b\bar{b}$ background.

Combining all these channels and optimizing the event selections would improve the sensitivity, so that the largest m_ψ that can be probed in Run 3 is likely in the 600–700 GeV range. Note that this is again significantly below the lower mass limit of 1.05 TeV set by Run 2 searches [5] for vectorlike lepton doublets undergoing standard decays.

2. Signals of the ξ leptoquarks with top quarks

Consider next the ξ leptoquark signals in the case of ψ decays involving third-generation quarks shown in (27). Pair production of the charge 2/3 leptoquark followed by cascade decays through the ψ VALs leads, as shown in the diagram from Fig. 10, to a $\tau^+\tau^-b\bar{b} + 4t$ final state. Thus, the signal is $\tau^+\tau^- + 6b + 4W$, and includes a large number (up to 6) of charged leptons in addition to the six b -jets. These provide various ways of eliminating the backgrounds. On the other hand, requiring many leptons implies small combined branching fractions which would reduce the signal too much. Moreover, the large multiplicity of particles in the final state implies that most of them are relatively soft. Following the discussion in Sec. III C 1, we expect that the reach in ξ pair production cross section is of the same order of magnitude as for $4t$, so around 20 fb. This corresponds to $M_\xi \lesssim 0.9$ TeV [33,51,52], which is substantially below the Run 2 lower mass limits (in the 1.2–1.4 TeV range [5,42,43,53]) for leptoquarks that have two-body decays into SM particles.

3. Signals for m_ψ in the 100–180 GeV range

Finally, if m_ψ is in the ~ 100 –180 GeV range, then the ψ decays shown in (27), which involve a top quark, are kinematically closed. In that case, ψ^0 and ψ^- may undergo four-body decays through an off-shell top quark ($\psi^0 \rightarrow W^-b\bar{b}\tau^+$ and $\psi^- \rightarrow W^-b\bar{b}\nu_\tau$). If these are the main decay modes, then the additional phase-space suppression

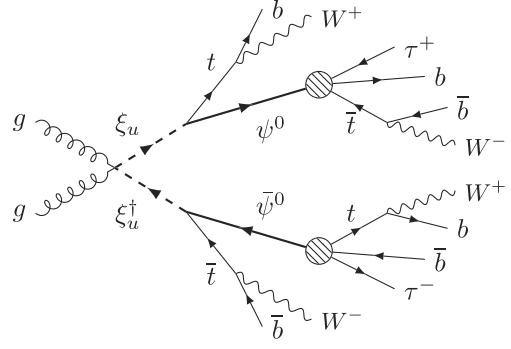


FIG. 10. Representative diagram for pair production of the ξ_u leptoquark followed by cascade decays via the ψ VALs, leading to a $\tau^-\tau^+b\bar{b} + 4t$ signal at hadron colliders.

is likely to make the decay length long enough so that a highly-ionizing charged track due to $\psi^\pm$ can be observed. Another possibility is that $|\lambda_c|$ is large enough (while satisfying $|\lambda_c|^2 \ll |\lambda_t|^2$ so that the couplings to third-generation quarks dominate), such that the three-body decays involving a charm quark,

$$\psi^0 \rightarrow \bar{c}b\tau^+, \quad \psi^- \rightarrow \bar{c}b\bar{\nu}_\tau, \quad (29)$$

have larger widths than the 4-body decays, and are prompt. We focus on this case in the remainder of this section.

Pair production of ψ then includes the following signals at the LHC:

$$\begin{aligned} pp &\rightarrow W^{\pm*} \rightarrow \psi^\pm \bar{\psi}^{(\mp)*} \rightarrow \tau^\pm b\bar{b}jj + \cancel{E}, \\ pp &\rightarrow Z^* \rightarrow \psi^0\bar{\psi}^0 \rightarrow \tau^+\tau^-b\bar{b}jj. \end{aligned} \quad (30)$$

Although the presence of b -jets avoids some of the backgrounds discussed in Sec. III B, these signals remain challenging. The above charge-current process has very large irreducible backgrounds from $t\bar{t}$ and $Wb\bar{b}jj$ productions. Even the neutral-current process may be hard to observe due to $t\bar{t}jj$ and $W + \text{jets}$ backgrounds.

It may be that the discovery modes for the VALs in this case are via the decays of the leptoquarks. Note that the latter decays involve top quarks even when ψ is below the m_t threshold. The most interesting process for m_ψ below the top threshold is

$$pp \rightarrow \xi_u \xi_u^\dagger \rightarrow (t\psi^0)(\bar{t}\bar{\psi}^0) \rightarrow (\tau^+tbj)(\tau^-\bar{t}\bar{b}j). \quad (31)$$

The signal includes four b -jets and up to four charged leptons, so if M_ξ is roughly in the 0.5–1 TeV range, then it is possible that a dedicated search in this channel would discover both ξ_u and ψ^0 in Run 3 of the LHC.

IV. CONCLUSIONS

The extension of the SM proposed here includes two new fields, which are weak doublets: a vectorlike fermion ψ

transforming as $(1, 2, -1/2)$ under the SM gauge group, and a scalar leptoquark ξ transforming as $(3, 2, +1/6)$. Gauge invariance allows a single type of renormalizable interaction involving both these fields, namely a Yukawa coupling to the SM up-type weak-singlet quarks, as displayed in (3). Likewise, there is a single type of interaction involving only ξ and SM fields: a Yukawa interaction of a down-type weak-singlet quark with a SM lepton doublet (assumed here to be the third-generation one) and ξ , of the form shown in (6). Both these types of Yukawa interactions conserve lepton number provided both ψ and ξ carry lepton number $L = -1$. Since the ψ doublet has the gauge charges of an often studied vectorlike lepton, but opposite lepton number, we use the term vectorlike antilepton, or VAL.

Interactions (3) preserve a global $U(1)_\psi$ symmetry under which ψ and ξ carry the same charge, so it is natural to expect that the interactions (6), which break that symmetry, have smaller couplings. This coupling hierarchy implies that the ξ scalars elude direct searches for leptoquark decays into two-body SM final states, while allowing the three-body decays of the VALs to be prompt for a range of parameters, as shown in Fig. 4.

Constraints from flavor-changing processes, such as $D^0 - \bar{D}^0$ mixing, $K^0 - \bar{K}^0$ mixing, and $K^+ \rightarrow \pi^+ \nu \bar{\nu}$, suggest that the Yukawa couplings of the new particles to first-generation quarks are suppressed. Even the couplings to second- and third-generation quarks are constrained, but only mildly (see Fig. 4), by flavor-changing processes including $B_s - \bar{B}_s$ mixing, $B^0 \rightarrow K^{*0} \nu \bar{\nu}$, and $B^+ \rightarrow K^+ \nu \bar{\nu}$.

The electroweak production of vectorlike fermion pairs at the LHC, computed at NLO in α_s for $\sqrt{s} = 13.6$ TeV and

displayed in Fig. 5, applies to a VAL doublet, but also to a vectorlike lepton doublet. While the current lower mass limit is already above 1 TeV for a vectorlike lepton [5], it is substantially less stringent for a VAL. If the couplings of the new particles are predominantly to the second-generation quarks, then the main LHC signals are $\tau^+ \tau^- + 4j$ and $\tau + 4j + \cancel{E}$ for VALs, and $\tau^+ \tau^- + 6j$ for leptoquarks. These signals with taus and jets are difficult to observe due to large backgrounds, but with advanced experimental techniques can be probed in the current Run 3 of the LHC.

The main LHC signals in the case of dominant couplings to third-generation quarks are $\tau^\pm W^+ W^- + 4b + \cancel{E}$ and $\tau^+ \tau^- W^+ W^- + 4b$ for VALs, and $\tau^+ \tau^- + 6b + 4W$ for leptoquarks. The presence of several charged leptons from τ or W decays makes these signals easier to detect than the ones with just taus and jets. Nevertheless, the large multiplicities of these signals imply that most objects in the final states are relatively soft. Thus, even in the case of third-generation couplings, the vectorlike fermions and leptoquarks make each other more elusive than in the case of standard two-body decays.

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- [1] A. Falkowski, D. M. Straub, and A. Vicente, *J. High Energy Phys.* **05** (2014) 092.
 - [2] N. Kumar and S. P. Martin, *Phys. Rev. D* **92**, 115018 (2015).
 - [3] P. N. Bhattiprolu and S. P. Martin, *Phys. Rev. D* **100**, 015033 (2019).
 - [4] A. M. Sirunyan *et al.* (CMS Collaboration), *Phys. Rev. D* **100**, 052003 (2019).
 - [5] A. Tumasyan *et al.* (CMS Collaboration), *Phys. Rev. D* **105**, 112007 (2022).
 - [6] I. Dor sner, S. Fajfer, A. Greljo, J. F. Kamenik, and N. Ko snik, *Phys. Rep.* **641**, 1 (2016).
 - [7] B. A. Dobrescu and F. Yu, *J. Phys. G* **45**, 08LT01 (2018).
 - [8] L. Di Luzio, A. Greljo, and M. Nardecchia, *Phys. Rev. D* **96**, 115011 (2017).
 - [9] M. Bordone, C. Cornella, J. Fuentes-Martin, and G. Isidori, *Phys. Lett. B* **779**, 317 (2018).
 - [10] L. Di Luzio, J. Fuentes-Martin, A. Greljo, M. Nardecchia, and S. Renner, *J. High Energy Phys.* **11** (2018) 081.
 - [11] A. Greljo and B. A. Stefanek, *Phys. Lett. B* **782**, 131 (2018).
 - [12] A. Tumasyan *et al.* (CMS Collaboration), *Phys. Lett. B* **846**, 137713 (2023).
 - [13] E. Bernreuther and B. A. Dobrescu, *J. High Energy Phys.* **07** (2023) 079.
 - [14] Q.-H. Cao, J. Guo, J. Liu, Y. Luo, and X.-P. Wang, *arXiv:2311.12934*.
 - [15] B. A. Dobrescu, *Phys. Rev. Lett.* **128**, 241804 (2022).
 - [16] J. L. Hewett and T. G. Rizzo, *Phys. Rep.* **183**, 193 (1989).
 - [17] P. Achard *et al.* (L3 Collaboration), *Phys. Lett. B* **517**, 75 (2001).
 - [18] ALEPH, DELPHI, L3, OPAL, LEP Collaborations, *Eur. Phys. J. C* **73**, 2463 (2013).
 - [19] A. Crivellin and L. Schnell, *Comput. Phys. Commun.* **271**, 108188 (2022).
 - [20] A. Crivellin, D. M uller, and F. Saturnino, *J. High Energy Phys.* **11** (2020) 094.

- [21] L. Silvestrini, *Effective Field Theory in Particle Physics and Cosmology: Lecture Notes of the Les Houches Summer School* (Oxford Academic, Oxford, 2020).
- [22] F. Gabbiani, E. Gabrielli, A. Masiero, and L. Silvestrini, *Nucl. Phys.* **B477**, 321 (1996).
- [23] G. Isidori, in *2012 European School of High-Energy Physics* (2014), pp. 69–105, <https://inspirehep.net/literature/1217584>.
- [24] R. Mandal and A. Pich, *J. High Energy Phys.* **12** (2019) 089.
- [25] D. Bečirević, S. Fajfer, and N. Košnik, *Phys. Rev. D* **92**, 014016 (2015).
- [26] M. Bona *et al.* (UTfit Collaboration), *Rend. Lincei Sci. Fis. Nat.* **34**, 37 (2023).
- [27] I. Adachi *et al.* (Belle-II Collaboration), [arXiv:2311.14647](https://arxiv.org/abs/2311.14647).
- [28] J. Grygier *et al.* (Belle Collaboration), *Phys. Rev. D* **96**, 091101 (2017); **97**, 099902(A) (2018).
- [29] T. E. Browder, N. G. Deshpande, R. Mandal, and R. Sinha, *Phys. Rev. D* **104**, 053007 (2021).
- [30] R. Bause, H. Gisbert, and G. Hiller, *Phys. Rev. D* **109**, 015006 (2024).
- [31] S. D. Thomas and J. D. Wells, *Phys. Rev. Lett.* **81**, 34 (1998).
- [32] U. Haisch, L. Schnell, and S. Schulte, *J. High Energy Phys.* **11** (2022) 106.
- [33] B. Diaz, M. Schmaltz, and Y.-M. Zhong, *J. High Energy Phys.* **10** (2017) 097.
- [34] A. Alloul, N. D. Christensen, C. Degrande, C. Duhr, and B. Fuks, *Comput. Phys. Commun.* **185**, 2250 (2014).
- [35] J. Alwall, R. Frederix, S. Frixione, V. Hirschi, F. Maltoni, O. Mattelaer, H. S. Shao, T. Stelzer, P. Torrielli, and M. Zaro, *J. High Energy Phys.* **07** (2014) 079.
- [36] R. D. Ball *et al.* (NNPDF Collaboration), *Eur. Phys. J. C* **82**, 428 (2022).
- [37] F. R. Anger, F. Febres Cordero, S. Höche, and D. Maître, *Phys. Rev. D* **97**, 096010 (2018).
- [38] E. Bothmann *et al.* (Sherpa Collaboration), *SciPost Phys.* **7**, 034 (2019).
- [39] A. M. Sirunyan *et al.* (CMS Collaboration), *J. High Energy Phys.* **10** (2019) 244.
- [40] P. N. Bhattiprolu, S. P. Martin, and A. Pierce, [arXiv:2308.08386](https://arxiv.org/abs/2308.08386).
- [41] F. Baspehlivan, B. Dagli, O. E. Delialioglu, and S. Sultansoy, [arXiv:2201.08251](https://arxiv.org/abs/2201.08251).
- [42] G. Aad *et al.* (ATLAS Collaboration), *Phys. Rev. D* **104**, 112005 (2021).
- [43] G. Aad *et al.* (ATLAS Collaboration), *J. High Energy Phys.* **06** (2021) 179.
- [44] A. M. Sirunyan *et al.* (CMS Collaboration), *Eur. Phys. J. C* **80**, 3 (2020).
- [45] A. M. Sirunyan *et al.* (CMS Collaboration), *J. Instrum.* **13**, P05011 (2018).
- [46] E. Bols, J. Kieseler, M. Verzetti, M. Stoye, and A. Stakia, *J. Instrum.* **15**, P12012 (2020).
- [47] R. L. Workman *et al.* (Particle Data Group), *Prog. Theor. Exp. Phys.* **2022**, 083C01 (2022).
- [48] G. Aad *et al.* (ATLAS Collaboration), *Eur. Phys. J. C* **80**, 1085 (2020).
- [49] A. Hayrapetyan *et al.* (CMS Collaboration), *Phys. Lett. B* **847**, 138290 (2023).
- [50] M. van Beekveld, A. Kulesza, and L. M. Valero, *Phys. Rev. Lett.* **131**, 211901 (2023).
- [51] T. Mandal, S. Mitra, and S. Seth, *Phys. Rev. D* **93**, 035018 (2016).
- [52] M. Kramer, T. Plehn, M. Spira, and P. M. Zerwas, *Phys. Rev. D* **71**, 057503 (2005).
- [53] G. Aad *et al.* (ATLAS Collaboration), *J. High Energy Phys.* **05** (2021) 093.