Leptonic signatures of color-sextet scalars II: Exploiting unique large-$E_T^{\text{miss}}$ signals at the LHC

Linda M. Carpenter* and Katherine Schwind†

Department of Physics, The Ohio State University
191 W. Woodruff Ave., Columbus, OH 43210, U.S.A.

Taylor Murphy‡

Laboratoire de Physique Théorique et Hautes Énergies (LPTHE), UMR 7589
Sorbonne Université & CNRS
4 place Jussieu, 75252 Paris Cedex 05, France

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The diverse and distinct collider phenomenology of color-sextet scalars motivates thorough investigation of their effective couplings to the Standard Model at the LHC. Some of the more unique sextet signals involve not only jets but also leptons. In previous work, we proposed an LHC search for color-sextet scalars in a channel with jets and a hard opposite-sign lepton pair, which results from a dimension-six coupling. In this sequel we study the counterpart processes with neutrinos, which produce jets in association with missing transverse energy ($E_T^{\text{miss}}$) in addition to possible leptons. We consider multiple search channels, including both single and pair sextet production, all characterized by significant missing energy and some featuring distinctive kinematic features. Our multifaceted study consists of three reinterpreted existing searches and a joint-likelihood analysis designed by us to maximize HL-LHC sensitivity to single sextet production. We show that our dedicated strategy in the jets + lepton + $E_T^{\text{miss}}$ channel can supersede today’s limits from reinterpreted searches, and we make sensitivity projections for the HL-LHC. Altogether, our analysis can exclude sextet scalars lighter than 4.4 TeV or probe effective cutoffs as high as 16.8 TeV.

I. INTRODUCTION

As Run 3 of the Large Hadron Collider (LHC) gets underway, the high-energy physics community has an opportunity to assess the strategies it has employed in the ongoing search for physics beyond the Standard Model (bSM), with an eye toward maximizing the LHC’s utility in the next decade and beyond. This assessment must be conducted in view of not only the improving performance and increasingly large dataset of the LHC, but also the lack of definitive evidence for bSM physics produced by LHC Runs 1 and 2. The latter reality has already been driving a shift during the past few years from the traditional focus on elaborate constructions featuring e.g. supersymmetry to more broadly applicable frameworks, exemplified by simplified models and effective field theories (EFTs). Indeed a broader approach seems appropriate in the absence of any specific experimental hints, and given the simple fact that much of the enormous bSM parameter space still has yet to be investigated. Models featuring new states in exotic representations of the SM color gauge group(s), $SU(3)_c$, are particularly ripe for investigation at hadron colliders.

Color sextets, which transform in the six-dimensional irreducible representation (6) of $SU(3)_c$, were first studied a few decades ago [1, 2] and have garnered some renewed interest in the twenty-first century. Many works have focused on sextet diquarks [3], scalars so named for their renormalizable couplings to like-sign pairs of right-chiral (“handed”) SM quarks of the form

\[
\mathcal{L} \propto K_{ij}^{\lambda} \Phi^\dagger \sigma_{\mu\nu} q_i R^c q_j R + \text{H.c.},
\]

where $K_{ij}^{\lambda}$ ($s \in \{1, \ldots, 6\}$) are $(i \leftrightarrow j)$-symmetric coefficients in color space, and where $q, q' \in \{u, d\}$ pursuant to the weak hypercharge of the sextet diquark $\Phi$. Some of the LHC phenomenology of these particles has been studied both in a standalone context [4, 5] and embedded as messengers between the Standard Model and fermionic dark matter [6]. While the number of gauge-invariant renormalizable sextet-SM interactions is small (though even these are rich and worth further attention), there are a large number of effective interactions at relatively low mass dimensions that have just begun to be studied [7]. Effective operators that integrate out some ultraviolet degrees of freedom but retain sextets at the weak or TeV scale themselves constitute but one subclass of a huge collection of operators with light exotics that offer alternative paths to discovery of bSM physics at the LHC [8].

In this spirit, we recently undertook an effort to identify and analyze the LHC signatures of exotic effective operators containing color-sextet scalars, beginning with a proposal to search for sextets produced in association with electrons or muons [9]. There we targeted events featuring at least two hard jets and an opposite-sign lepton pair, which are generated by a dimension-six coupling between a color-sextet scalar, the SM gluon, and a right-handed quark-lepton pair. Excellent sensitivity can be achieved at the LHC by binning the invariant mass of the scalar’s decay products: in particular, a search along the lines of what we proposed in that work can probe sextet masses up to $m_\Phi = 5$ TeV, or effective cutoff scales $\Lambda$ of $O(10)$ TeV, with the planned $\mathcal{L} = 3$ ab$^{-1}$ run of the High-Luminosity (HL-)LHC.
The present work is a natural extension of that study: it targets that operator’s left-handed counterpart, which couples a color-sextet scalar to the SM quark/lepton weak doublets. This operator, in addition to producing sextet-lepton-(quark-gluon) couplings, permits the scalar to interact with a SM neutrino. These couplings open multiple interesting production and decay channels that will produce events with significant missing energy and distinct kinematics at the LHC. Since the sextet can be both pair produced and singly produced in association with a charged lepton or neutrino — and then decay to either as well — we consider all possibilities in order to achieve the best possible sensitivity for the full planned run of the HL-LHC.

Our multifaceted analysis is centered on a joint-likelihood analysis in a single-production channel we call “mostly visible”, which includes one charged lepton from either production (i.e., recoiling off of the sextet scalar) or decay. We design a search strategy for this channel and compute a joint likelihood based on a binned distribution of $m_T(j_1j_2, E_T^{\text{miss}})$, the transverse mass of the two leading jets and the missing transverse momentum, which can partially reconstruct the decayed sextet scalar if its decay involves a neutrino but also displays a low peak around the sextet mass when the neutrino is the recoiling particle. We compare our custom search to the most up-to-date existing analyses that exhibit some sensitivity to our signal processes.

This process is represented by the diagrams shown in Figure 1. Initial states provide the bulk of the production at proton colliders. To be more specific, at the $\sqrt{s} = 13$ TeV LHC, the production rate exceeds 1 pb for $m_\Phi \lesssim 750$ GeV [3, 5], thus not much smaller than the rate of color-octet scalar pair production [10]. The leading-order (LO) cross sections for this process are displayed in Figure 3, in Section III, following the requisite details about signal simulation.

### II. MODEL DISCUSSION

We begin by introducing the effective theory of color-sextet scalars to be analyzed in this work. We consider a minimal model of sextet interactions with Standard Model fermions. If the scalar has weak hypercharge $Y = 1/3$, it can couple both to quark pairs of mixed type (one up, one down) and to color-charged particles accompanied by a lepton. But these interactions cannot be simultaneously realized without inducing lepton number ($L$) non-conservation. We suppose that $L$ is conserved and that the sextet scalar has lepton number $L = -1$, so that its leading interactions with the Standard Model, apart from gauge interactions, are given by

$$
\mathcal{L} \supset \frac{1}{\Lambda^2} J^{\pm i a} \Phi_\nu G_{\mu
u} \chi_a \varepsilon_{\mu
u} \eta_{L X} + \lambda_{R_1} \eta_{R_2} \Phi \sigma_{\mu
u} \eta_{R X} + H.c.
$$

with spinor and SU(2)$_L$ indices implicit. The superscript $c$ denotes charge conjugation and subscripts $R$ denote right handedness (chirality). The tensor $\eta_{\mu\nu} = (i/2) [\gamma^\mu, \gamma^\nu]$ performs a chirality flip. The couplings $\lambda_{R_1}$ form matrices in quark and lepton generation space, with $I$ or $X = 3$ labeling the heavy generation(s). Finally, the coefficients $J$ are the generalized Clebsch-Gordan coefficients [3] required to construct gauge-invariant contractions of the direct-product representation $3 \otimes 6 \otimes 8$ in SU(3) [7].

These 144 coefficients are displayed in Figure 3, in Section III, following the requisite details about signal simulation. This process is represented by the diagrams shown in Figure 1. The $gg$ initial states provide the bulk of the production at proton colliders. To be more specific, at the $\sqrt{s} = 13$ TeV LHC, the production rate exceeds 1 pb for $m_\Phi \lesssim 750$ GeV [3, 5], thus not much smaller than the rate of color-octet scalar pair production [10]. The leading-order (LO) cross sections for this process are displayed in Figure 3, in Section III, following the requisite details about signal simulation.

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1 The dot $\cdot$ denotes an antisymmetric contraction of the SU(2)$_L$ doublets, so that the weak singlet is explicitly written as $\overline{Q}_A e^{AB} L_{LB}$ with $A, B \in \{1, 2\}$ and $e$ the two-dimensional totally antisymmetric symbol.
B. Single scalar production

While sextet pair production is unavoidable, we are principally interested in the single-production processes available in this specific realization. This interest stems from the competitive cross sections (in some parameter space) and distinct kinematic features of resonant sextet production, which—as we should later—makes single production the superior experimental target, increasingly so for heavier sextets. The kinematic features of resonant sextet production, which—as above—permit competitive cross sections (in some parameter space) and distinct partial widths to each charged-lepton channel can be expressed as

\[ \Gamma(\Phi \to \ell \bar{\nu}_{\ell} g) = \frac{1}{(2\pi)^3} \frac{1}{8m_{\Phi}^2} \left[ \left( \frac{\lambda_{L}^{I \chi}}{\Lambda^2} \right)^2 + \left( \frac{\lambda_{R}^{I \chi}}{\Lambda^2} \right)^2 \right] \times \int_{m_{\ell}^2}^{(m_{\Phi}-m_{\nu})^2} ds_{13} \mathcal{F}(s_{13}, m_{\Phi}, m_{u}, m_{\ell}), \]  

where \( s_{13} = (p_1 + p_3)^2 \) is the invariant squared mass of the lepton-gluon subsystem (an arbitrary choice), and where

\[ \mathcal{F}(s_{13}, m_{\Phi}, m_{u}, m_{\ell}) \approx \frac{1}{s_{13}^3} \left( s_{13} - m_{\Phi}^2 \right)^3 \left( m_{u}^2 - m_{\nu}^2 - s_{13} \right) \times \left[ m_{u}^4 - 2m_{u}^2 (\bar{s} + m_{\Phi}^2) + (\bar{s} - m_{\Phi}^2)^2 \right]^{1/2} \]  

is proportional to the squared transition amplitude integrated over \( s_{23} = (p_2 + p_3)^2 \). The integral in (9) does not admit an analytical result, except in the massless-fermion limit with \( \mathcal{F}(s_{13}) \approx s_{13}(m_{\Phi}^2 - s_{13})^2 \) in which

\[ \Gamma(\Phi \to \ell \bar{\nu}_{\ell} g) \approx \frac{2}{3} \frac{1}{(8\pi)^3} \left[ \left( \frac{\lambda_{L}^{I \chi}}{\Lambda^2} \right)^2 + \left( \frac{\lambda_{R}^{I \chi}}{\Lambda^2} \right)^2 \right] m_{\Phi}^5, \]  

but can be evaluated numerically in the general case. Of course, if any element of \( \lambda_{L} \) is non-vanishing, then \( \Phi \) can also decay to a gluon, a down-type quark, and a neutrino. The partial widths for these channels, \( \Gamma(\Phi \to \nu_{X} d_{1} g) \), are given by (9) and (11) with the replacements

\[ m_{u} \to m_{d}, \quad m_{\ell} \to m_{\nu}, \quad 0, \]  

and

\[ \frac{m_{\Phi}^4}{\lambda_{L}^2} \left( \frac{\lambda_{R}^{I \chi}}{\Lambda^2} \right)^2 \to \left( \frac{\lambda_{L}^{I \chi}}{\Lambda^2} \right)^2 \]  

are of equal size, the total charged-lepton and neutrino branching ratios are each nearly 50%. We restrict ourselves to parameter space in which the sextet scalar decays promptly, which is the case for all but the lightest sextets with very high (\( O(10) \) TeV) EFT cutoffs.

D. Computer implementation and benchmark scenarios

For analytic and numerical investigation, we implement this model in \textsc{FeynRules} version 2.3.43 [13, 14], a package for
An investigation of the opposite scenario, with only elements needed flavors $I, X \in \{1, 2\}$ and $\lambda_R^{IX} = 0$. 

\begin{equation}
\lambda_L^{IX} = 1 \quad \text{for} \quad I, X \in \{1, 2\}
\end{equation}

An investigation of the opposite scenario, with only elements needed flavors $I, X \in \{1, 2\}$ and $\lambda_R^{IX} = 0$. 

Before we move on, it is appropriate to discuss the range of validity of the effective operator (2) and the resulting allowed values of the cutoff $\Lambda$. It is widely understood that effective operators can only accurately describe physical processes taking place at energy scales lower than the scale of the degrees of freedom integrated out of some ultraviolet theory in order to produce the infrared model. This physical constraint implies some relationship between the scale $\Lambda$, the energy scale of a physical process of interest, and perhaps the mass(es) of the infrared modes. In previous work, we computed the perturbative unitarity bound [22–24] on the right-chiral part of the operator (2) and used a simulation-driven approach to avoid parameter space violating the limit. A similar analysis can be performed for the left-chiral part, which gives a lower bound proportional to

\begin{equation}
\Lambda_{\min}^2 \propto \lambda_L^{IX}(\hat{s} - m_{\Phi}^2)
\end{equation}

with $\hat{s}$ the partonic center-of-mass energy of the sextet scalar production process. We use the full result, which can be found by comparison with [9], to define the edge of the valid EFT parameter space in the $(m_\Phi, \Lambda)$ plane for the purpose of limit calculation in Section VI.

\section{III. SIGNAL SIMULATION AND STATISTICAL ANALYSIS}

In this section we lay the groundwork for the analyses in Sections IV and V by detailing the simulation of the signal processes and describing the statistical analysis procedures. The intent is to relegate most technical details here in order to leave Sections IV–VI free to focus on phenomenology.

\subsection{A. Signal samples}

Our multifaceted analysis requires a number of signal samples with a variety of final states in a large parameter space. Most of the important samples are for resonant sextet scalar production. In this category, we simulate

\begin{itemize}
  \item $pp \rightarrow \ell^\pm \Phi$, $\Phi \rightarrow \bar{v}d\bar{g} \ (\ell \in \{e, \mu\})$
  \item $pp \rightarrow \nu \Phi$, $\Phi \rightarrow \ell^\pm \bar{u}d\bar{g}$, and
  \item $pp \rightarrow \nu \Phi$, $\Phi \rightarrow \bar{v}d\bar{g}$
\end{itemize}

and their conjugate processes.\footnote{Recall that the unusual presence of two anti-fermions in the final states arises from the color conjugation of the quark in (2).} We categorize the first two processes as “mostly visible” and the latter as “semi-invisible”, as a shorthand for the lepton/neutrino content. All processes feature visible hard jets. We prepare both a set of samples inclusive with respect to (light) lepton flavor and separate sets including only $e^\pm$ or $\mu^\pm$, since the latter are required to pass selections for two recast searches discussed in Section IV. We reiterate that the couplings $\lambda_L^{IX}$ are set to unity only for the needed flavors $I, X$ in each case. Heavy quarks and leptons are not considered in this study. These single-production
processes are simulated in MG5\_AMC version 3.3.1 at leading order for an LHC center-of-mass energy of $\sqrt{s} = 13$ TeV. The scattering amplitudes are convolved with the NNPDF2.3 LO set of parton distribution functions [25] with renormalization and factorization scales $\mu_R, \mu_F$ fixed equal to $m_\Phi$. Showering and hadronization is performed by Pythia 8 version 8.244 [26]. For computation of the yields, the signals are normalized to the cross sections reported by MG5\_AMC after appropriate color-factor correction [9]. As demonstrated in Section II, the normalizations are proportional to $\Lambda^{-4}$; rescaling is done wherever required in order to explore the $\Lambda$ dimension of the $(m_\Phi, \Lambda)$ EFT parameter space. Some example cross sections for production in association with a single lepton or neutrino flavor are displayed in Figure 3. These single-production cross sections are computed assuming first-generation SM fermions only and an EFT cutoff $\Lambda = 3$ TeV. This choice of cutoff renders these results self-consistent within the EFT framework, following the discussion in Section II. In this scenario, the cross sections are of $O(10-100)$ fb at the TeV scale.

Figure 3 also depicts the leading-order cross sections of sextet scalar pair production, which we mentioned in Section II and include in our analysis in the interest of completeness. We simulate pair production followed by mostly visible and semi-invisible decays in view of the analyses to be reinterpreted in our model framework (viz. Sections IV and V). The events are generated at LO in MG5\_AMC and showered with Pythia 8, and the signal normalizations are computed with renormalization and factorization scales set again to $\mu_R = \mu_F = m_\Phi$. Note that pair production, which proceeds through gauge interactions, has an inclusive rate independent of the EFT parameters in (2). The exclusive cross sections depend on $\lambda_L^{X\Lambda^{-2}}$, but as noted in Section II our simple benchmark values imply branching fractions $\beta(\Phi \rightarrow \ell^+\ell^- g) \approx \beta(\Phi \rightarrow \ell^+\ell^- g) \approx 1/2$ for a given generation.

B. Limit setting and significance calculation

On a technical level, this study contains two kinds of analysis. There are three cases in which we identify an existing search executed by one of the LHC collaborations that may be sensitive to our color-sextet scalar signals containing, at minimum, jets and $E_T^{miss}$. We reinterpret the null results of these searches within our model framework in order to impose limits on $m_\Phi$ and $\Lambda$ where appropriate, which are valid today — based on the integrated luminosity used for each analysis — and then projected into the future via luminosity rescaling. These reinterpretations are performed using MadAnalysis 5 (MA5) version 1.9.20 [27], which in its reconstruction mode (~R) can run any of the validated recasted LHC analyses available on the MA5 Public Analysis Database (PAD) [28-30].

When called in reconstruction mode, MA5 accepts as input a HepMC event record produced by Pythia 8. Once one or more analyses on the PAD are chosen for execution by the user, MA5 simulates the ATLAS or CMS detector (whichever is needed for the analysis at hand) either by calling Delphes 3 version 3.4.3 [31, 32] or, for more recent analyses, by using its inbuilt Simplified Fast Detector Simulation (SFS) module [33]. In either case, the simulated detector response is parametrized by an analysis-specific card tuned during the development of the recast in order to optimize the agreement between recast and official results. Meanwhile, object reconstruction is performed by FastJet version 3.4.0 [34]. The reconstructed event sample is then analyzed by the SampleAnalyzer core of MA5, which imposes the analysis selection criteria and records the sample efficiency under each cut. With this done, MA5 finally computes the upper limit(s) at 95% confidence level (C.L.) [35] on an arbitrary new-physics cross section with the efficiency of the provided sample. When an analysis comprises multiple signal regions, a limit is computed for each. MA5 computes the expected limit, based on the expected background yields provided by the experimental collaboration, and the observed limit based on the true event yield. If the user supplies a signal cross section at the import step, MA5 further computes the significance $1 - CL_s$ of the sample, which is computed at 95% C.L. if $1 - CL_s \geq 0.95$. Our final note is on luminosity rescaling for extrapolation to HL-LHC: MA5 provides a number of options to estimate the expected limits at 95% C.L. from a present-day analysis hypothetically run in the future on a larger dataset. Any new integrated luminosity can be chosen by the user for extrapolation. The signal and (central) background yields are rescaled linearly with the luminosity. The user can choose how the background uncertainties are rescaled; we use the default option, whereby the systematic uncertainties are rescaled linearly, the statistical uncertainties are rescaled; we use the default option, whereby the systematic uncertainties are rescaled linearly, the statistical uncertainties are rescaled; we use the default option, whereby the systematic uncertainties are rescaled linearly, the statistical uncertainties are rescaled; and these components are added in quadrature [36].

The reinterpretation of existing analyses very helpfully paints a picture of the real-life experimental status of our theory. As alluded to in multiple places, however, the crux of this work is a dedicated search we tailor to the unique features of our leptonic sextet processes. We use a different workflow and our own statistical tools to evaluate the sensitivity of our custom analysis. In particular, we use the MadAnalysis 5 SFS module
to perform detector simulation and object reconstruction for all of our signal and background samples. We use an SFS card optimized for the ATLAS detector. We then feed the (relevant) reconstructed event files to MA5 once again, this time imposing our own selection criteria and recording the efficiencies as usual.

At this stage, our workflow diverges from what is outlined above for reinterpretation because we carry out a joint-likelihood analysis and use measures of the significance for exclusion and discovery that are not natively available in MadAnalysis 5. To be specific, we compute the significance $Z$ of a signal $s$ (whose uncertainty we neglect) with strength modifier $\mu$, given an expected background $\langle b \rangle$ with non-negligible uncertainty $\sigma_b$, according to the asymptotic (large-$m$) formula $Z_m^\mu = (q_m^\mu)^{1/2}$ [37], where [38]

$$q_m^\mu = -2 \ln \frac{L(m \mid \mu, \hat{b})}{L(m \mid \hat{\mu}, b)} \hat{\mu} \leq \mu,$$ \hspace{1cm} (15)

is a frequentist test statistic for one-sided limits based on the ratio of profile likelihoods [39]

$$L(m \mid \mu, b) = \prod_{i=1}^{N_{\text{bin}}} \left( \frac{m_i^{s(b)}}{m_i !} e^{-(\mu s_i + b_i)} \right) \times \frac{1}{\sqrt{2\pi} \sigma_{b,i}} \exp \left\{ -\frac{1}{2} \left( \frac{b_i - (\langle b_i \rangle)^2}{\sigma_{b,i}^2} \right) \right\}. \hspace{1cm} (16)$$

This expression gives the joint likelihood for $N_{\text{bin}}$ bins, which is relevant for our custom analysis with cuts binned in an observable that displays peaks at $m_{\Phi}$ for our signal samples. In this construction, the total event count $m_i$ in each bin $i$ of some observable is assumed to follow a Poisson distribution, while each background $b_i$ is taken as a nuisance parameter constrained (in principle) by a Gaussian-distributed measurement in some control region [40]. The bin yields are assumed to be uncorrelated. The quantity $\hat{b} = \hat{b}(\mu)$ in (15) is the conditional maximum-likelihood (ML) estimator of the background $b$ for a given $\mu$, while the pair $(\hat{\mu}, \hat{b})$ are unconditional ML estimators of the likelihood $L$. It is also possible to define the likelihood for an individual bin by rewriting (16) as

$$L(m \mid \mu, b) = \prod_{i=1}^{N_{\text{bin}}} L(m_i \mid \mu, b_i). \hspace{1cm} (17)$$

Within this framework, we compute the median significance for exclusion and discovery of our signal using the Asimov datasets according to

$$Z_{\text{excl}} \equiv Z_{\mu=1}^{m(z(b))}$$ and $$Z_{\text{disc}} \equiv Z_{\mu=0}^{m(x+b)(b)},$$ \hspace{1cm} (18)

taking $Z_{\text{excl}} = 1.68$ and $Z_{\text{disc}} = 5$ as our exclusion and discovery thresholds. In this calculation, the background uncertainties $\sigma_b$ follow from the theoretical cross section uncertainties $\delta\sigma_{\text{theo}}$, which are listed in Section V. We compute the test statistic $q_m^\mu$ using both the joint likelihood (16) and the individual likelihoods for each bin in (17). The significance(s) (18) derived in the first case give the most optimistic estimates of collider sensitivity to our bSM signal, whereas the significance(s) from the most sensitive individual bin give the most conservative projections. We provide both results in Sections V and VI in order to give an idea of the range of limits one can expect from our dedicated analysis.

IV. A “SEMI-INVISIBLE” SEXTET: JETS + $E_T^{\text{miss}}$

The first channel we consider excludes charged leptons and thus includes only hard jets and missing transverse energy. This search channel can be sensitive to signal processes of the type displayed in Figure 2c, which we have been referring to as semi-invisible. Since the jets + $E_T^{\text{miss}}$ signal provides a classic search channel for many popular bSM constructions, including dark matter models and supersymmetric models, there exist limits based on the full Run 2 dataset. We therefore begin by exploring the extent to which existing searches constrain our model. In particular, the MadAnalysis 5 PAD includes a Run 2 search performed by the ATLAS Collaboration for new phenomena in final states with jets and significant missing transverse energy [41]. This analysis was announced as ATLAS-CONF-2019-040 and subsequently relabeled ATLAS-SUSY-2018-22. The data correspond to an integrated luminosity of $L = 139$ fb$^{-1}$. ATLAS interpreted the null results of this search, ATLAS-CONF-2019-040, for typical scenarios involving squark and gluino pair production; we reinterpret the analysis within our model framework following the workflow described in Section III.

This ATLAS analysis requires at least two jets, the harder of which must have at least $p_T(j_1) > 200$ GeV, and at least $E_T^{\text{miss}} > 300$ GeV, with adequate separation between the two or three hardest jets and the missing transverse momentum. The analysis (which consists of a multi-bin subanalysis and a boosted decision tree subanalysis, only the first of which is recast) imposes more stringent requirements on jet multiplicity in some bins. There are furthermore selections on the $E_T^{\text{miss}}$ significance and the effective mass $m_{\text{eff}}$, defined as the scalar sum of $E_T^{\text{miss}}$ and all jets with $p_T(j) > 50$ GeV. These selections powerfully suppress the Standard Model backgrounds, some of which are very large, most notably Z boson production with hard jets followed by the invisible decay $Z \to \nu\bar{\nu}$. Since some bins in the multi-bin analysis allow more than three hard jets, the search is sensitive to both single and pair production of our color-sextet scalar. We therefore apply this recast analysis to the semi-invisible single-production samples and the pair samples.
meanwhile, shows the limits on pair production. The observed fraction of sextets to neutrino modes is set to

direct comparison to the analyses in Section V, the branching factor(s) of the scalar to neutrino modes, are independent of the EFT cutoff $\Lambda$ — provided that it is low enough for the resonance to decay promptly — because the production is due to gauge interactions. We find it interesting that for pair production, which can be directly compared to the supersymmetric scenario considered by ATLAS, the efficiencies of the sextets samples are not too dissimilar to those of gluino pair production followed by the effective three-body decay $\tilde{g} \rightarrow q\bar{q}'\nu$ via a heavy off-shell squark. The limits are of course not coincident because its pair-production rate is lower than that of the color-octet gluino due to color and spin.

While the limits from the semi-invisible channel are non-trivial, comfortably disfavoring sextet scalars lighter than $\sim 1$ TeV, they turn out to be the weakest bounds calculated in this work. We furthermore find that a custom cut-and-count analysis in this channel cannot significantly improve upon existing Run 2 limits on single sextet production, such that extending the constraints in the channel requires a larger dataset. We briefly demonstrate this difficulty using the selections listed in Table I, which are aimed — but unsuccessful — at exceeding the sensitivity of ATLAS-CONF-2019-040 with a non-equivalent set of cuts. We furthermore show in Figure 5 some signal and background distributions of observables one might naturally consider by ATLAS, the efficiencies of the sextets samples are not too dissimilar to those of gluino pair production followed by the effective three-body decay $\tilde{g} \rightarrow q\bar{q}'\nu$ via a heavy off-shell squark. The limits are of course not coincident because its pair-production rate is lower than that of the color-octet gluino due to color and spin.

simulated with both sextets decaying to neutrinos. The results are displayed in Figure 4.

Figure 4a shows the upper observed and expected limits at 95% C.L. on the single-production cross section for a range of sextet masses at the TeV scale. These limits are compared to the cross section in a benchmark scenario with sextet couplings to first- and second-generation SM fermions and an EFT cutoff of $\Lambda = 4$ TeV. As discussed in Section II, and for direct comparison to the analyses in Section V, the branching fraction of sextets to neutrino modes is set to $\beta = 1/2$. In this benchmark we obtain an observed lower limit on the sextet mass of $m_\Phi \approx 1424$ GeV, marginally stronger than the expected limit of 1124 GeV due to underfluctuations in the data. Figure 4b, meanwhile, shows the limits on pair production. The observed limit at 95% C.L., which is very close to the expected limit, is at $m_\Phi \approx 1262$ GeV. It is important to note that the pair-production limits, while sensitive to the branching fraction(s) of the scalar to neutrino modes, are independent of the EFT cutoff $\Lambda$ — provided that it is low enough for the resonance to decay promptly — because the production is due to gauge interactions. We find it interesting that for pair production, which can be directly compared to the supersymmetric scenario considered by ATLAS, the efficiencies of the sextets samples are not too dissimilar to those of gluino pair production followed by the effective three-body decay $\tilde{g} \rightarrow q\bar{q}'\nu$ via a heavy off-shell squark. The limits are of course not coincident because its pair-production rate is lower than that of the color-octet gluino due to color and spin.

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| Selection criterion | Selection ranges |
|---------------------|------------------|
| Jets, anti-$k_t$ | $R = 0.4$ |
| $N_{jet} \geq 2$ |
| $p_T(j) > 100$ GeV, $|p_T(j)| < 2.5$ |
| Charged leptons | veto $e, \mu$ with $p_T > 15$ GeV |
| $b$-tagged jets | veto $b$-jets with $p_T > 15$ GeV |
| Leading jet momenta | $p_T(j_1) > 250$ GeV |
| $p_T(j_2) > 250$ GeV |
| Missing energy, $E_T^{miss}$ | Require $E_T^{miss} > 750$ GeV |

TABLE I: Selection criteria of attempted search for color-sextet scalar produced in association with a neutrino and decaying to $dg\nu$. |
instance, shows that the missing energy significance

$$S(E_{T}^{\text{miss}}) = \frac{E_{T}^{\text{miss}}}{\sqrt{H_{T}}} \text{ with } H_{T} = \sum_{i=1}^{N_{\text{jet}}} p_{T}(j_{i}) \quad (19)$$

is actually more sharply peaked for $Z$+jets than for a $m_{\Phi} = 2$ TeV signal, though the signal distribution has a slightly fatter tail. The culprit is the unusual contribution to the missing transverse momentum from both the recoiling neutrino and the invisible decay product of the sextet; see once more Figure 2c. We therefore find no well motivated selection to impose in addition to, or instead of, the quite stringent requirements of ATLAS-CONF-2019-040. As a result, the limits from the search defined in Table I fail to improve upon the existing bounds: at $L = 139$ fb$^{-1}$, for $m_{\Phi} = 1$ TeV, we can rule out a maximum EFT cutoff of $\Lambda \approx 4$ TeV, and for $m_{\Phi} = 2$ TeV we are sensitive to no more than $\Lambda \approx 3$ TeV. These limits are quite close to the expected limits from ATLAS-CONF-2019-040 (viz. Section VI), but no better. At any rate, as we just noted and will demonstrate in the next section, the limits from the semi-invisible channel are far superseded by existing and custom searches in the mostly visible channel, so we set aside the all-neutrino processes until Section VI.

V. A “MOSTLY VISIBLE” SEXTET: JETS + $\ell + E_{T}^{\text{miss}}$

The next channel we consider is characterized by hard jets and a single lepton (electron or muon) accompanied by a sizeable amount of missing transverse energy. This channel is populated by signal events of the type displayed in Figures 2a and 2b, but also by signals in popular bSM scenarios involving dark matter and pair-produced leptoquarks. In this section, we first design a custom search for color-sextet scalars singly produced in this mode. We identify the Standard Model processes with significant overlap for the mostly-visible final state and detail their simulation, define selection criteria to exploit the kinematic distinctiveness of our signals, explore a few signal and background distributions, and compare their yields after the common selection. We then recast two existing analyses sensitive to our mostly visible sextet in preparation for a comparison of all bounds in Section VI.

A. Single-lepton background processes

There are many Standard Model processes that may be expected to populate this search channel, some with quite large cross sections. The background processes we consider in our analysis are listed in Table II. This table lists the most accurate available cross sections at a center-of-mass energy of $\sqrt{s} = 13$ TeV, along with estimates of the theoretical uncertainties, given in most cases by adding the scale and parton distribution function (PDF) uncertainties in quadrature. These cross sections are exclusive whenever a decay is specified in the table. The largest backgrounds a priori are leptonically decaying $W$ boson production in association with hard jets, $\bar{t}t$ production, and single-top production with a light quark ($qt$).
Weak boson pair production leading to final states with exactly one lepton is a potentially important background that survives the $b$-jet veto that decimates the top quark processes. On the other hand, diboson events with jets and missing energy can be discarded. The same applies to $Z$+jets since we simultaneously demand exactly one lepton and large $E_T^{\text{miss}}$. For completeness, we provide some simulation details for each sample used in our analysis.

We simulate the $W + j$ets process in MG5$_\text{aMC}$ version 3.3.1 at LO with up to two additional partons included in the hard-scattering matrix element. The scattering amplitude is convolved with the NNPDF2.3 LO set of parton distribution functions [25] with renormalization and factorization scales $\mu_R, \mu_F$ fixed equal to $m_W$. The LO hard-scattering events are matched to parton showers in the MLM scheme [52] at a scale of 30 GeV with the aid of PYTHIA 8 version 8.244 [26]. The signal is normalized to the cross section at next-to-next-to-leading-order (NNLO) in the strong coupling with next-to-leading order (NLO) electroweak corrections, as computed by FEWZ version 3.1 [43, 44] with the NNPDF3.0 NNLO parton distribution functions [53]. Pair production of weak bosons decaying to one lepton, which includes the processes $WW \rightarrow \ell \nu jj$ and $WZ \rightarrow \ell \nu jj$, is simulated at LO in MG5$_\text{aMC}$, including interference from virtual photons $\gamma^*$. For efficiency optimization, these are divided into single-lepton and zero-lepton samples. The signal normalizations are computed at NLO in the strong coupling using POWHEG-Box version 2 with the CT10 NLO PDF set [54–57]. The renormalization and factorization scales are set dynamically to $\mu_R = \mu_F = m_{V_jV_j}/2$.

Top pair production is simulated at LO in MG5$_\text{aMC}$. We use a top quark mass of $m_t = 172.5$ GeV for all relevant processes.

We again use the NNPDF 2.3 LO parton distribution functions but this time set $\mu_R = \mu_F = m_t$. The normalization of this sample is computed by Top++ 2.0 for our choice of $m_t$ at NNLO, including resummation of next-to-next-to-leading logarithmic (NNLL) soft-gluon terms [47], with the NNPDF3.0 NNLO PDF set. Production of a single top quark in association with a weak boson is simulated at LO in MG5$_\text{aMC}$ along similar lines as previous samples. The $t\bar{t} + Z$ cross section is computed at NLO in QCD in MG5$_\text{aMC}$ using the NNPDF3.0 NLO PDF set, with $\mu_R = \mu_F = m_t$; whereas the $t\bar{t} + W$ cross sections are computed at NLO in QCD using the parton-level integrator MCFM [58] with the MSTW2008 NLO PDF set. We note that the best uncertainty estimates we can find [48] indicate larger uncertainties for the $t\bar{t} + V$ processes than for the other backgrounds in this analysis.

Production of a single top quark in association with a light quark is simulated at LO in MG5$_\text{aMC}$ within the five-flavor scheme (bottom quark mass set to $m_b = 0$), since the parton-level process is $q b \rightarrow q' t$. The sample is normalized to the cross section computed at NNLO in QCD [49] with $\mu_R = \mu_F = m_t$ and with the five-flavor CT14 NNLO PDF set [59]. $tW$ production is similarly simulated in the five-flavor scheme, since the parton-level process is $gb \rightarrow tW^*$. Our sample is produced at LO and the normalization is computed at NLO in QCD with NNLL resummation [51] with the five-flavor NNPDF3.1 NLO PDF set [60]. Finally, the relatively small $tZ$ process is simulated at LO in QCD in MG5$_\text{aMC}$, and the cross section is computed at NLO in QCD using the same setup with $\mu_R = \mu_F = m_t$ and using the NNPDF3.0 NLO PDF set.

| Background process | $\sigma_{\text{LHC}}^{13\text{TeV}}$ [pb] | $\delta\sigma_{\text{theo}}$ [%] | Notes |
|-------------------|---------------------------------|-----------------|-------|
| $W + \text{jets}, \ W \rightarrow \ell \nu$ | 20,400.0 | 1.30 | NNLO QCD + NLO EW [42–44] |
| Diboson ($WW/WZ/ZZ$) | 55.73 | 6.00 | NLO QCD [45, 46] |
| $t\bar{t}$ | 833.9 | 4.37 | NNLO QCD + NNLL [47] |
| $t\bar{t} + Z$ | 0.7587 | 29.8 | NLO QCD from MG5$_\text{aMC}$ |
| $t\bar{t} + W$ | 0.5662 | 1.49 | NLO QCD [19, 48] |
| $t$-channel single top ($qt$) | 214.2 | 3.30 | NLO QCD from MG5$_\text{aMC}$ |
| $tW$ | 79.3 | 1.49 | NLO QCD + NNLL [51] |
| $tZ$ ($qt + \nu\bar{\nu}$ incl. non-resonant $Z$) | 0.1399 | | |

TABLE II: Leading background processes generating final states with two hard jets, one (charged) lepton, and considerable missing transverse energy. Single $W$ and $t$ denote aggregate particle + antiparticle processes while $q$ denotes light quarks. LHC cross sections for $\sqrt{s} = 13$ TeV are displayed at indicated precision along with total theoretical uncertainties and event descriptions where helpful. Cross sections are inclusive unless decay is specified.
TABLE III: Common selection criteria of proposed search for color-sixtet scalar produced in association with a lepton (neutrino) and decaying to $dg\nu\,(ug\ell)$.

| Selection criterion | Selection ranges |
|---------------------|------------------|
| Jets, anti-$k_T$ $R = 0.4$ | \(N_{\text{jet}} \geq 2\) $p_T(j) > 100\text{ GeV}, |\eta(j)| < 2.5$ |
| Charged leptons | $N_{\ell} = 1$ $p_T(\ell) > 15\text{ GeV}$ |
| $b$-tagged jets | veto $b$-jets with $p_T > 15\text{ GeV}$ |
| Leading jet momenta | $p_T(j_1) > 250\text{ GeV}$ $p_T(j_2) > 200\text{ GeV}$ |
| Lepton momentum | $p_T(\ell) > 250\text{ GeV}$ |
| Missing energy, $E_T^{\text{miss}}$ | require $E_T^{\text{miss}} > 300\text{ GeV}$ |
| Jet-$E_T^{\text{miss}}$ separation | $\Delta\phi(j_1, p_T^{\text{miss}}) \notin [2\pi/3, 4\pi/3]$ |

B. Selection criteria; distributions and yields

The common selections for our dedicated analysis are listed in Table III. The two leading jets from our signal events are expected to have roughly equal transverse momentum and should be much harder than the leading jets from Standard Model backgrounds. Our baseline selection therefore requires at least two jets with anti-$k_T$ radius parameter $R = 0.4$ to have transverse momentum greater than 100 GeV and pseudorapidity $|\eta| < 2.5$. We then require the leading and second-leading jets to have $p_T$ greater than 250 GeV and 200 GeV, respectively. Many of the processes in Table II involving W bosons also involve top quarks; since our sextet scalar does not couple to third-generation SM fermions, we impose a veto on $b$-tagged jets. The combination of the $b$-jet veto and jet $p_T$ requirements (especially the cut on the second-leading jet momentum) so powerfully suppresses the events including top quarks that, at this stage of the analysis, the background is dominated by $W +$ jets, with single-lepton diboson events still surviving but subleading. From here, the selection is therefore dedicated almost entirely to suppressing the leptonically decaying $W$ events.

The single lepton in the signal events should be similarly hard and provide a good discriminant from the $W +$ jets background (though it will not by itself sufficiently suppress this process). The lepton $p_T$ distributions are displayed in Figure 6 for the backgrounds surviving all cuts preceding that on $p_T(\ell)$ (of which only $W +$ jets is visible on this scale), as well as for the pair of signal events corresponding to Figures 2a and 2b for a sextet mass of $m_\Phi = 2\text{ TeV}$ and an EFT cutoff of $\Lambda = 4\text{ TeV}$. This figure shows a clear difference not only between signals and backgrounds but also between the two types of signal event. In particular, we find that the lepton is much harder when it recoils off of the sextet at the production stage than when it is a decay product of the scalar resonance. We find that a requirement of $p_T(\ell) > 250\text{ GeV}$ already suppresses over 90% of the remaining $W +$ jets events while accepting most signal events, especially for heavy sextets.

The lepton $p_T$ cut, followed by a requirement of more than 300 GeV of missing transverse energy, effectively eliminate the already decimated top quark backgrounds. The $E_T^{\text{miss}}$ cut also culls some remaining $W +$ jets and diboson events, but at this stage the analysis is still fairly weak because of the sheer size of the first background. The final common selection takes another important step toward suppressing this background by forbidding the leading jet from being antiparallel to the missing transverse momentum. Figure 7 shows that the backgrounds surviving all cuts up to and including the lepton $p_T$ requirement have distributions of azimuthal $j_1 - p_T^{\text{miss}}$ separation, $\Delta\phi(j_1, p_T^{\text{miss}})$, peaked around $\pi$, whereas the signals — particularly the lepton-first events corresponding to Figure 2a — are much flatter in this observable. We therefore find that requiring $\Delta\phi(j_1, p_T^{\text{miss}}) < 2\pi/3$, or $> 4\pi/3$, is quite effective.

Finally we consider the possibility of strengthening the analysis by finding an observable on which to make binned cuts. One simple approach is to identify an observable that is sharply peaked on at least one side of the sextet scalar mass, such that a set of non-overlapping bins can be established and a joint likelihood can be computed according to the procedure in Section III. Our situation is interesting in that each of the processes represented by Figures 2a and 2b has its own well suited observable that does not apply particularly well to the other. In particular, if the lepton is a decay product of the sextet, then the sextet is fully reconstructible using the invariant mass $m_{ij,fs}$ of the leading jets and lepton. As shown in Figure 3, however, this neutrino-first process suffers from lower cross sections than its lepton-first counterpart, whose $m_{ij,fs}$ distribution is not peaked at $m_\Phi$. It turns out to be better to preferentially target the larger lepton-first process in Figure 2a, in which the sextet scalar is partially reconstructible with the transverse mass $m_T(j_{1,2}, E_T^{\text{miss}})$ of the system composed of the
FIG. 7: Azimuthal separation between leading jet and missing transverse momentum in signal ($m_\Phi = 2$ TeV) and background events surviving the selections in Table III up to $p_T(\ell) > 250$ GeV. Distributions are normalized to unity.

FIG. 8: Transverse mass of the hard-jet system $j_1j_2$ and missing transverse momentum in signal ($m_\Phi = 2$ TeV) and background events surviving all common selections in Table III. Distributions are normalized to total yields at $L = 139$ fb$^{-1}$.

to denote the transverse mass of a pair $\{A, B\}$, keeping in mind that $A$ or $B$ could itself be a system (such as in the current discussion, where $A = j_1j_2$). Here, in analogy with the previous paragraph, $\Delta \phi$ is the azimuthal separation between transverse momentum vectors. We show in Figure 8 the distributions of $m_T(j_1j_2, E_T^{\text{miss}})$ for the signal and background processes surviving all common selections. This figure shows not only the discrepancy in the yields of the two signal processes but also the clear shelf at $m_\Phi$ in the transverse mass distribution of the lepton-first signal. In the interest of simplicity, our final selection requires $m_T(j_1j_2, E_T^{\text{miss}}) > 250$ GeV and then cuts the transverse mass in 250 GeV-wide bins, terminating at $m_T = 5.0$ TeV. This bin width is particularly well suited to relatively light sextets, whose $m_T$ peaks are quite narrow but which suffer the bulk of the remaining $W +$ jets background. We show in Table IV the background and signal yields in each bin of the $m_T$ selection (this time for an integrated luminosity of $L = 3$ ab$^{-1}$, for some variety). The bulk of the $m_\Phi = 2$ TeV lepton-first events indeed populate the bins just below $m_T = 2.0$ TeV, though there is also a gentle peak below this threshold for neutrino-first events. As discussed in Section III, we use the yields for all signals in all bins to compute the signal significance both for the single most sensitive bin and considering the full $m_T$ distribution, assuming no correlations in the interest of extreme simplicity. The results are discussed in Section VI.

C. Existing single-lepton analyses

With our dedicated analysis now established, we turn (in the interest of completeness) to two analyses capable of constraining our model using parts of the LHC Run 2 dataset. The first of these is a search by the CMS Collaboration for pair production of scalar leptoquarks coupling exclusively to first-generation Standard Model fermions [61]. This search was announced as CMS-EXO-17-009 and superseded by CERN-EP-2018-265. The relevant part of this analysis is the “$e\nu jj$” channel, in which one leptoquark decays (promptly) to a quark and an electron with branching fraction $\beta$ and the other leptoquark decays to a quark and a neutrino with branching fraction $1 - \beta$. The luminosity for this analysis is relatively low, $L = 35.9$ fb$^{-1}$, but a follow-up analysis including a neutrino channel has not
been released.\(^5\)

The \(evj\) selection requires exactly one electron, at least two jets, and \(E^\text{miss}_T > 100\) GeV. Since this analysis targets first-generation leptoquarks, a muon veto is applied. Since signal events are assumed to result from production of two degenerate particles, the transverse masses \(m_T(j_1, E^\text{miss}_T), m_T(j_2, e)\) are computed with \(\{j_1, j_2\}\) chosen from the two hardest jets \(\{j_1, j_2\}\) such that the difference in transverse masses is minimized. A number of thresholds are declared for the jet-electron transverse mass, constituting a set of inclusive (overlapping) bins. There is finally an inclusive binned selection on \(S_T\), the scalar sum of transverse momenta of the electron, the two leading jets, and the missing transverse momentum. Much like the ATLAS search in the semi-invisible channel discussed in Section IV, there is nothing in this search \(a\ priori\) that favors single or pair production of our sextet, since both modes can produce a single electron and a single neutrino. Sextet pair production is expected to produce more than two hard jets from the two three-body decays, but the jet multiplicity requirement in CMS-EXO-17-009 is bounded only from below. We therefore apply the recast of this search, available in the MA5 PAD, to our mostly visible single-production samples (exclusive to electrons) and a set of pair-production samples with one sextet decaying to an electron and the other to a neutrino. The results for single production are displayed in Figure 9.

In analogy with Section IV, this figure shows the upper observed and expected limits at 95% C.L. on the single-production cross section compared to the theoretical rate in a benchmark with EFT cutoff \(\Lambda = 4\) TeV. Here the electron and neutrino branching fractions are both set to \(\beta = 1/2\). In this bench-

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\(^5\) A light-generation leptoquark pair analysis based on \(\mathcal{L} = 139\) fb\(^{-1}\) has been released by the ATLAS Collaboration [62], but it treats only the fully visible decays LQ \(\rightarrow fj\).

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| Process | Lower bound on \(m_T(j_1, j_2, E^\text{miss}_T)\) [TeV] | (recall that \(m_T\) bins are non-overlapping) |
|---------|----------------|------------------|
| \(W + \text{jets}, W \rightarrow ℓν\) | 58.0 | 0.38 |
| Diboson (1f) | 0.77 | 0.38 |
| \(pp \rightarrow ℓΦ \rightarrow ℓνjj\) | 10.6 | 0.38 |
| \(pp \rightarrow νΦ \rightarrow ℓνjj\) | 3.04 | 0.38 |

### TABLE IV: Binned yields for (surviving) background processes and the pair of \(m_Φ = 2\) TeV signals following the full selection described in Section V. These estimates are computed for \(\sqrt{s} = 13\) TeV and \(\mathcal{L} = 3\) ab\(^{-1}\). The EFT cutoff is again taken to be \(\Lambda = 4\) TeV.

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**FIG. 9:** Observed and expected limits at 95% C.L. from CMS-EXO-17-009, a Run 2 search for first-generation scalar leptoquarks in final states with jets and either an opposite-sign electron pair or a single electron and missing transverse energy. Only the latter channel applies to this figure. In a scenario with EFT cutoff \(\Lambda = 4\) TeV, the observed (expected) lower bound on \(m_Φ\) from this search is 2527 (2580) GeV.

**FIG. 10:** Observed and expected limits at 95% C.L. from CMS-EXO-17-015, a Run 2 search for dark matter in events with a muon-philic leptoquark and missing transverse energy. In a scenario with EFT cutoff \(\Lambda = 4\) TeV, the observed (expected) lower bound on \(m_Φ\) from this search is 1776 (1736) GeV.
mark we obtain an observed lower limit on the sextet mass of $m_Φ ≈ 2527$ GeV, negligibly different from the expected limit of 2580 GeV. Meanwhile, under the same branching fraction assumptions, we obtain an expected lower limit of $m_Φ ≈ 1417$ GeV and an observed limit of $m_Φ ≈ 1421$ GeV on pair-produced sextets. We display these limits later in Section VI; here we simply note that they outstrip the pair-production limits from ATLAS-CONF-2019-040 shown in Figure 4b and are indeed the strongest we can obtain for pair production in any final state with missing energy.

Finally we mention CMS-EXO-17-015 (later CERN-EP-2018-78), a search by the CMS Collaboration for dark matter in events with a leptoquark and missing transverse momentum. The data for this analysis correspond to an integrated luminosity of $L = 77.4$ fb$^{-1}$. The basic requirements are an isolated muon with $p_T > 60$ GeV and $E_T^{miss} > 100$ GeV. The leading jet is required to be hard, with $p_T > 100$ GeV, and $b$-tagged jets are vetoed for reasons similar to ours in the dedicated analysis. Electrons and $τ$ hadron candidates are also vetoed, as are additional muons if they form an opposite-sign muon pair with the leading muon whose invariant mass falls in the window $m_{μμ} = m_Z ± 10$ GeV. Other than isolation requirements, the final selection requires the transverse mass of the muon and leading jet to exceed 500 GeV. We apply the MadAnalysis 5 recast of this search to single- and pair-production samples including exactly one muon. The results for single production are displayed in Figure 10; the pair-production limits are weaker than those from CMS-EXO-17-009 and so we omit them from the discussion.

This figure can be interpreted much like its predecessors: here the theory cross section is in a benchmark with (again) an EFT cutoff of $Λ = 4$ TeV and with muon and neutrino branching fractions each set to $β = 1/2$. In this benchmark we obtain an observed lower limit on the sextet mass of $m_Φ ≈ 1776$ GeV and an expected limit of 1736 GeV, both of which are weaker than the limits imposed by the CMS search for first-generation leptoquarks in the $eνjj$ channel. This discrepancy is despite comparable acceptances of our signal events and the doubling of the dataset for the muon analysis, and is ultimately due to the multi-bin approach of the earlier leptoquark search, which (especially for bins targeting heavier leptoquarks) better suppresses the $W +$ jets and $t\bar{t}$ backgrounds.

### VI. LHC Sensitivity Projections

We finally map the results of Sections IV and V onto the $(m_Φ, Λ)$ plane and make sensitivity projections for the planned 3 ab$^{-1}$ run of the HL-LHC. This map is displayed in Figure 11. The custom selection strategy in the mostly visible channel is detailed in Section V. Our projections are made using the workflows detailed in Section III based on our samples simulated for center-of-mass energy of $\sqrt{s} = 13$ TeV. These sensitivities are therefore conservative in view of upgrades to $\sqrt{s} = 13.6$ TeV or even 14 TeV.

All of the space in Figure 11 not shaded in gray (due to perturbative unitarity violation; viz. Section II) is self-consistent EFT parameter space with a promptly decaying color-sextet scalar. The regions shaded in color are excluded by observed limits at 95% C.L. from one of the searches reinterpreted in the preceding sections. The corresponding expected limits are marked by dashed contours. Our extrapolations to an integrated luminosity of $L = 3$ ab$^{-1}$ are furthermore marked by thicker solid contours. These limits correspond directly to the benchmarks considered in Sections IV and V: this means that these are the strongest limits than can be derived from each search, since for instance diminishing the sextet branching fraction to electrons would weaken the limits from the CMS search for first-generation leptoquarks. We therefore proceed with the caveat that these limits do not apply simultaneously to any single benchmark scenario.

Figure 11 demonstrates, as discussed earlier, that the CMS search for first-generation leptoquarks (blue) imposes most of the strongest limits on single sextet production in the mostly visible channel. These limits already extend as far as $m_Φ ≈ 3.0$ TeV or $Λ ≈ 6.2$ TeV, which is even more impressive given the relatively low luminosity of this search. This analysis naturally has the most to gain from additional statistics, such that the projected reach at HL-LHC is as high as $m_Φ ≈ 3.6$ TeV or $Λ ≈ 8.75$ TeV. For single production, the CMS search for dark matter and a muon is competitive with the LQ analysis for light sextets but rapidly loses sensitivity. The ATLAS jets + $E_T^{miss}$ search is generally the least sensitive for the reasons discussed in Section IV. In the interest of completeness, we include in Figure 11 the exclusion contour at $L = 139$ fb$^{-1}$ for our unsuccessful custom jets + $E_T^{miss}$ analysis detailed in Section IV. As mentioned there, we find similar exclusions to the ATLAS-CONF-2019-040 expected limits, but nowhere do we improve upon their results. Moving on, also shown in Figure 11 are the $A$-independent limits on pair-produced scalars from the CMS leptoquark analysis. As discussed in Section IV, just as for single production, the pair-production limits from this search are stronger than those imposed by the other two searches, so we suppress the other sextet pair limits in the interest of visual clarity.

The results of our dedicated analysis in the jets + single lepton + $E_T^{miss}$ channel are drawn in orange. As in Section V, we start with the full Run 2 luminosity of $L = 139$ fb$^{-1}$ and then extrapolate to $L = 3$ ab$^{-1}$. There is a pair of contours for each projection: in each case, the thinner curve traces the values of $(m_Φ, Λ)$ for which the exclusion significance of the most sensitive individual $m_T(j_1j_2, E_T^{miss})$ bin attains $Z_{excl} = 1.68$ (“best”), and the thicker contour (“joint”) does the same for the significance computed using the joint likelihood as described in Section III. Displaying both contours demonstrates both the power of our custom selections and the additional sensitivity given by considering the full likelihood. We caution once more that these are the most optimistic estimates, since the sensitivity will wane for a real analysis with a full statistical model that includes the correlations between the non-overlapping $m_T$ bins. But we see, at minimum (without the full likelihood), that our analysis is roughly as sensitive with the current LHC dataset as the CMS leptoquark analysis will be after the HL-LHC shuts down. Our analysis moreover outstrips all existing analyses at the HL-LHC, probing singly produced sextets as heavy as $m_Φ ≈ 4.4$ TeV or cutoffs as high as $Λ ≈ 16.8$ TeV.
FIG. 11: Map of EFT parameter plane \((m_{\Phi}, \Lambda)\) showing most recent available Run 2 LHC limits and \(\mathcal{L} = 3\, \text{ab}^{-1}\) HL-LHC projected limits, compared to exclusion sensitivity \(Z_{\text{excl}}\) of our dedicated mostly visible analysis for \(\mathcal{L} = 139\, \text{fb}^{-1}\) and \(3\, \text{ab}^{-1}\). “Best” limits are derived from most sensitive \(m_{T}(j_1j_2, E_{\text{miss}}^T)\) bin, while “joint” limits use the full \(m_{T}\) distribution. Also shown, for reference, is the exclusion sensitivity from the unbinned custom semi-invisible search discussed in Section IV.

VII. CONCLUSIONS

In this work we have compared an array of analyses targeting hard jets and significant missing transverse energy as probes of color-sexet scalar production in association with leptons and/or neutrinos at the LHC. Such processes are minimally allowed by dimension-six operators with a unique \(\text{SU}(3)_c\) color structure that has only begun to be explored recently. We have reinterpreted three existing searches within an effective framework to constrain both single and pair production of our sextet scalar. The existing limits are already at the multi-TeV scale for both the sextet mass \(m_{\Phi}\) and the cutoff \(\Lambda\) that is related to the scale of ultraviolet physics. Nevertheless, we have demonstrated that the reach of the LHC can be improved by way of a dedicated joint-likelihood analysis, in a channel with one visible lepton, with binned selections of \(m_{T}(j_1j_2, E_{\text{miss}}^T)\), the transverse mass of the two leading jets and missing transverse momentum. Our custom search can ultimately rule out sextet scalars lighter than \(m_{\Phi} \approx 4.4\, \text{TeV}\) or probe cutoffs as high as \(\Lambda \approx 16.8\, \text{TeV}\) (14\, TeV in space not already excluded for pair production).

It is worthwhile to contrast the limits obtained in this work with the experimental status of color-sexet scalars assumed to couple to the Standard Model through the conventional diquark portal exemplified by (1). Sextets of this class can be directly probed at the LHC with searches for dijet resonances (for single production) and for at least four jets (for pair production), some of which could be \(b\)-tagged to target sextets coupling to third-generation quarks. While the limits obtained by recasting such searches are highly sensitive to the couplings (hence branching fractions) of the sextet scalar, direct Run 2 searches have pushed both light-flavor [4, 5] and heavy-flavor sextet diquarks [6] to the TeV scale, and in fact above 1.5 TeV for resonances decaying with unit branching fraction to quarks subsequently hadronizing to form flavorless jets. By comparison with Figure 11, we see that these limits are roughly competitive with, but marginally stronger than, the recast limits derived from CMS-EXO-17-009, the search for pair-produced leptoquarks — which we noted are independent of the EFT cutoff \(\Lambda\). Meanwhile, the diquark couplings of sextet scalars can be probed indirectly by searches for flavor-changing neutral currents in mixings of neutral kaons and \(D\) and \(B\) mesons, to which sextets can contribute sizably at tree and loop level. The precise bounds depend on whether the sextet couples to up-type, down-type, or mixed quarks; but generally these constraints are very stringent, heavily disfavor democratic couplings to first- and second-generation quarks, and can restrict some dimensionless diquark couplings to be as small as \(O(10^{-6})\) [63].

While we have mentioned these limits to provide context, we caution the reader against viewing these diquark constraints and the results of our work as part of a single coherent picture, and particularly against relating lower bounds on the sextet diquark mass to lower bounds on our EFT cutoff \(\Lambda\). These two schemes cannot be combined naively because (a) as mentioned...
in the Introduction, the diquark operators (1) cannot coexist with the effective operators (2) without at minimum lepton number violation, and (b) if we suppose that multiple two-body and three-body production/decay channels are available, then it is not clear without a dedicated study what constraints apply to parameter space in which such channels are of comparable size. In the interest of caution, rather, we view the diquark portal and the effective framework as non-overlapping schemes to be analyzed separately. From this point of view, the interesting lesson of this work and its prequel [9] is therefore that (in appropriately simple benchmark scenarios) the LHC can be used to impose strong limits on resonant sextet scalar production with non-standard semi-invisible topologies that rival the constraints obtained from conventional multijet searches.

There are a number of other interesting possible extensions of this work. One direction involves applying our EFT results to ultraviolet-complete constructions. It is straightforward to imagine loop-induced couplings of color-charged fields to leptons and neutrinos within existing popular bSM frameworks (featuring, for example, leptoquarks) that might be good candidates for searches like ours that target semi-visible three-body decays. Moreover, UV-complete theories may produce non-democratic sextet branching fractions and more distinct kinematics that could help us further refine the strategy detailed here. An entirely different avenue of study lies in further investigating the jets + \( E_T^{\text{miss}} \) channel, for which crafting a dedicated search with superior sensitivity proved impossible using standard observables. Our sextet model provides an interesting example of a bSM process in the jets + \( E_T^{\text{miss}} \) channel with an asymmetric topology (by which we mean the four-body final state consists of three decay products and a recoiling particle instead of two pairs of two decay products). In principle, this event topology should produce signatures distinct from those of e.g. squark pair production. But in this work we were not yet able to surpass the standard jets + \( E_T^{\text{miss}} \) searches designed to target sparticle pair production. We are therefore working on new analyses aimed at exploiting the asymmetric event topology. Sextets are an interesting test case for these signals, and we could also explore other novel operators with color sex-

tets. In this work we chose to study \( SU(2)_L \) (weak)-singlet sextets. Other sextet models could contain fields with higher \( SU(2)_L \) representations and would present even more rich phenomenology.

To take a broader view, new SU(3)_c charged states, with their wide range of new signatures, remain ripe for exploration. New operator catalogs would include interesting diboson production modes and decays [8, 65, 68] and multiple heavy-flavor events [69]. Following the spirit of this work, we plan to explore other novel EFTs involving exotic color-charged states. Further work with color octets, mentioned above, is one obvious direction. One avenue for further study could be to explore asymmetric single production of exotic color octets in association with other states, as we have done here with sextets.

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**STATEMENT OF HUMAN AUTHENTICITY**

The authors affirm that all text, images, and conceptual creation put forward in this work are entirely human generated, with no input from generative artificial intelligence (AI). The authors do not consent to the use of this work in the training of generative AI models.

[1] R. S. Chivukula, M. Golden, and E. H. Simmons, Phys. Lett. B 257, 403 (1991).
[2] A. Çelikel, M. Kantar, and S. Sultansoy, Phys. Lett. B 443, 359 (1998).
[3] T. Han, I. Lewis, and T. McElmurry, J. High Energy Phys. 01, 123, arXiv:0909.2666 [hep-ph].
[4] C.-R. Chen, W. Klemm, V. Rentala, and K. Wang, Phys. Rev. D 79, 054002 (2009).
[5] T. Han, I. Lewis, and Z. Liu, J. High Energy Phys. 12, 085, arXiv:1010.4309 [hep-ph].
[6] L. M. Carpentet, T. Murphy, and T. M. P. Tait, arXiv:2205.06824 [hep-ph] (2022).
[7] L. M. Carpenter, T. Murphy, and T. M. P. Tait, Phys. Rev. D 105, 035014 (2022), arXiv:2110.11359 [hep-ph].
[8] L. M. Carpenter, T. Murphy, and M. J. Smylie, J. High Energy Phys. 08, 050, arXiv:2302.01344 [hep-ph].
[9] L. M. Carpenter, T. Murphy, and K. Schwid, Phys. Rev. D 106, 115006 (2022), arXiv:2209.04456 [hep-ph].
[10] L. M. Carpenter, T. Murphy, and M. J. Smylie, J. High Energy Phys. 11, 024, arXiv:2006.15217 [hep-ph].
[11] N. Cabibbo, Phys. Rev. Lett. 10, 531 (1963).
[12] M. Kobayashi and T. Maskawa, Prog. Theor. Phys. 49, 652 (1973).
