TeV-scale leptoquark searches at the LHC and their E_6SSM Interpretation

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Abstract

We perform a model-independent search for leptoquarks (LQs) at the Large Hadron Collider through their pair-production and subsequent decay into $t\bar{t}\tau\tau$ intermediate states. We show that, assuming full luminosity of the Run 2, a fully hadronic signal emerging from this intermediate state can surpass in sensitivity the established searches relying on leptons in the final state. Our conclusion is supported by a thorough Monte-Carlo analysis, and we advocate the deployment of our proposed search channel in the proper experimental setting of the Run 3. Furthermore, in order to highlight the full scope of this approach for constraining LQ theories, we interpret our results in the context of the string-inspired Exceptional Supersymmetric Standard Model, which naturally predicts the S_1–type scalar LQ states that we analyse here.

1 Introduction

Leptoquarks are predicted by various extensions of the Standard Model (SM), such as Technicolour [1, 2], models of quark-lepton unification [3–7], and grand unification theories (GUTs) based on $SU(5)$, $SO(10)$, E_6 [8–13], and etc. They are rather unique amongst the many new particles predicted in new physics frameworks beyond the SM for being colour-triplet bosons that carry both lepton and baryon numbers. Their other quantum numbers (i.e., spin, (fractional) electric charge, and weak isospin) can be different in different theories.

Naturally, they were extensively searched for at the *ep* colliders. In fact, at HERA, DESY, an excess was found in 1997 [14] which would have been compatible with the so-called first-generation LQs (i.e., those coupling to first-generation leptons and quarks). However, further searches at HERA (and also at Tevatron, where most of the partonic processes were initiated by quarks inside the (anti-)proton) failed to confirm the anomaly. As a result, mass limits near 300 GeV [15, 16] were set on first-generation LQs. Even weaker limits were placed on second-generation LQs. These (relatively) weak limits were owing to the fact that the partonic energy accessible at HERA and Tevatron was limited to the sub-TeV range. However, as colour triplets, LQs also interact with gluons. At the Large Hadron Collider (LHC), since the amount of gluons inside the proton can well exceed that of (both valence and sea) quarks, the availability of much higher centre-of-mass energy may enable the pair-production of third-generation LQs.

The ATLAS and CMS experiments have both, therefore, analysed the datasets from the LHC Runs 1 and 2 to put exclusion limits on the mass of the third-generation LQs, based on their $t\bar{t}\tau\tau$, $t\bar{b}\tau\nu$, and $t\bar{b}\mu\nu$ decay channels, in the TeV range [17–20]. In this paper, for LQs having couplings solely with the third generation fermions, we revisit the $t\bar{t}\tau\tau$ search channel for LQs of the S_1 type, i.e., LQs with charge $-1/3$ and spin zero. However, unlike the existing searches for this channel, which have relied on leptons in the final state, we investigate the scope of the fully hadronic final state. Following a sophisticated detector-level analysis tensioning the signal to the most relevant backgrounds, we show that such a complementary channel can offer sensitivities to this specific type of LQ comparable, if not better, than what has so far been achieved in their probes exploiting leptons. We then interpret the results of our analysis in the Exceptional Supersymmetric Standard Model (E_6SSM) [21, 22] (for a recent review, see Ref. [23]), as a theoretically well-motivated prototypical framework that naturally accommodates the S_1–type LQs.

The plan of the paper is as follows. In the next section, we will describe our phenomenological setup in a model-independent way, which will then be subjected to a Monte-Carlo analysis in the following section. In Sec. 4 we will discuss the E_6SSM and its possible scenarios that are amenable to experimental investigation by ATLAS and CMS, before concluding in Sec. 5.

2 Third-generation Scalar LQs

We consider a simple extension of the SM by adding one scalar S_1–type LQ, which is charged as $(3, 1, -1/3)$ under its gauge group, $SU(3)_C \times SU(2)_L \times U(1)_Y$. We limit our discussion to LQs with the same quantum numbers as the coloured triplet of scalars, which, along with the Higgs doublet, constitutes the fundamental representation of $SU(5)$ [24, 25]. LQs with such quantum numbers have been proposed as a solution to the a_μ and B-flavour anomalies (see, e.g., [26–28]).

We assume that our LQ, denoted by D henceforth, couples only to the third-generation fermions, with the corresponding Lagrangian being

$$\mathcal{L} = \lambda \bar{Q}_L D \bar{\ell}_L + \lambda' \bar{t}_R D \bar{\tau}_R \,, \qquad (1)$$

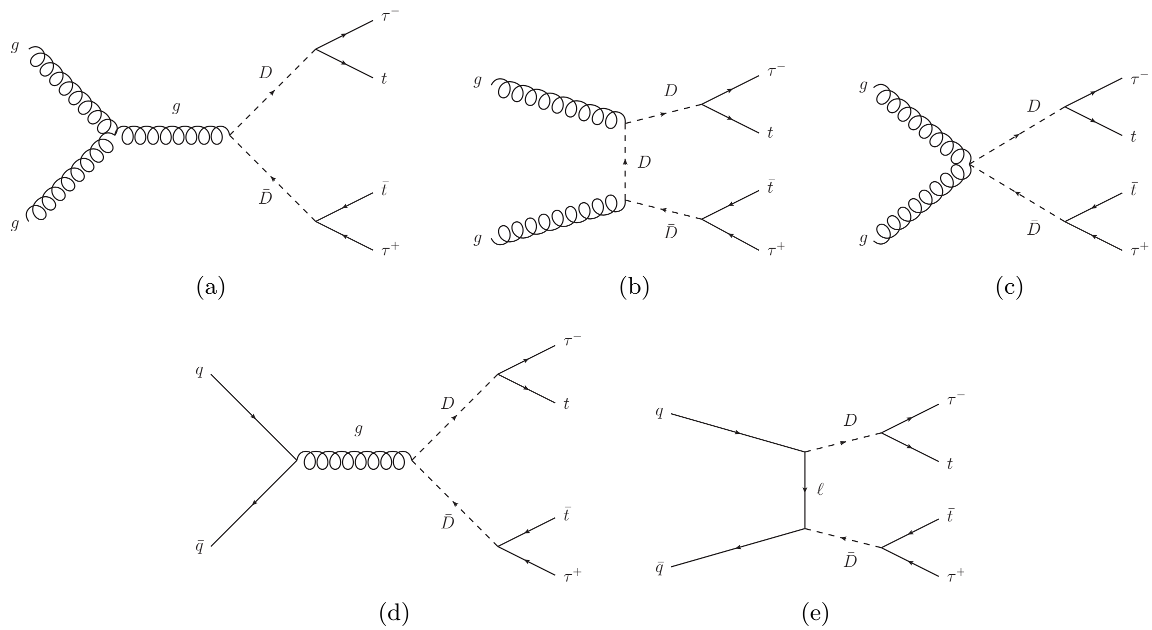

Figure 1: Pair-production of scalar LQs at the LHC.

where $Q_L = (t, b)_L^T$, and $\ell_L = (\nu_\tau, \tau)_L^T$. This implies that the left-handed t quark couples to the left-handed τ lepton, while the right-handed t couples to the right-handed τ. The mass of the t quark at the one-loop level can thus cause the τ chirality to flip. Based on the above interaction terms, the D can decay into $t\tau$ and $b\nu_\tau$ pairs. However, we assume $\lambda' \gg \lambda$, so that the decay to $t\tau$ is significantly enhanced. This LQ can be produced in pairs at the LHC through the processes shown in Fig. 1.

Given that the D couples only to the third-generation fermions, the diagram (e) in Figure 1 is negligible (unless the corresponding λ and λ' couplings are extremely large, which is never the case here). The cross section for the D pair production is therefore almost completely determined by QCD, and can be considered to be model-independent.[1] We can write this in the form

$$\sigma(pp \to t\bar{t}\tau\tau) = \sigma(pp \to D\overline{D}) \times \mathrm{BR}(D \to t\tau^-) \times \mathrm{BR}(\overline{D} \to \bar{t}\tau^+)\,, \tag{2}$$

where BR stands for branching ratio. In our setup, we have $\mathrm{BR}(D \to t\tau^-) = \mathrm{BR}(\overline{D} \to \bar{t}\tau^+) \simeq 1$, and the leading order (LO) partonic cross sections $\sigma(gg \to D\overline{D})$ and $\sigma(q\bar{q} \to D\overline{D})$ are given by [29]

$$\sigma(gg \to D\overline{D}) = \frac{\alpha_s^2 \pi}{96\hat{s}} \left[\beta(41 - 31\beta^2) + (18\beta^2 - \beta^4 - 17)\log\frac{1+\beta}{1-\beta}\right], \quad \text{and} \tag{3}$$

$$\sigma(q\bar{q} \to D\overline{D}) = \frac{2\alpha_s^2 \pi}{27\hat{s}}\beta^3, \tag{4}$$

where $\beta = (1 - 4M_D^2/\hat{s})^{1/2}$, and $\sqrt{\hat{s}}$ is the partonic center of mass energy.

In Tab. 1 we show the LO and next-to-LO (NLO) pair-production cross sections, computed with MadGraph5 v3.2.0 [30] and Prospino v2.1 [31], using the NNPDF 3.1 parton distribution functions (PDFs) at LO [32]. Note that our cross section values are somewhat smaller than those in [19], mainly due to the different PDF sets used, but they agree well with those in Ref. [33], which also uses NNPDF 3.1.

[1]The contribution from the model-dependent diagram (e) can, in fact, be distinguished from the others as one expects two b–tagged jets close to the beam direction.

m_D (GeV)	$\sigma_{\rm LO}$ (fb)	$\sigma_{\rm NLO}$ (fb)
1000	3.22	5.73
1100	1.57	2.86
1200	0.79	1.50
1300	0.42	0.79
1400	0.22	0.44
1500	0.12	0.25
1600	0.070	0.146

Table 1: The LO and NLO production cross sections for the $pp \to D\overline{D}$ process for a range of the D mass at the $\sqrt{s} = 13$ TeV LHC. The uncertainty related to the variation of the renormalisation and factorisation scales is $\pm 20\%$.

Decay Mode	Mass Limit [GeV]	Experiment
$t\bar{t}\tau\tau$	900	CMS [17]
	1400	ATLAS [19]
$tb\tau\nu$	950	CMS [18]
	1250	ATLAS [20]
	1220	ATLAS [19]

Table 2: 95% CL lower mass limits by the ATLAS and CMS Collaborations from pair-produced third-generation scalar LQs.

Our simplified extension of the SM with a S_1–type LQ can be considered an effective low-energy limit of some GUT framework. In the simplest GUTs, however, our D would have non-zero quark-lepton as well as quark-quark couplings, which can lead to rapid proton decay. To ensure the model's phenomenological viability, all the couplings of the D to quark pairs can be set to zero, implying the conservation of baryon and lepton numbers.

Furthermore, GUT models have a large separation of scales, and to stabilise the ensuing hierarchy, supersymmetry is often preferred. If we supersymmetrise our model, we need to introduce two chiral superfields, one of which has λ–type and the other λ'–type superpotential couplings. The 'exotic' particle spectrum then consists of two scalars and one fermion. The scalar states may have non-zero mixing, and one of the resulting physical states decays dominantly to $t\tau$, while the other has nearly equal BRs in the $t\tau$ and $b\nu$ decay channels. The decay of their superpartner fermion LQ leads to third-generation fermions and missing transverse momentum. Such a low-energy effective model can be embedded in a supersymmetric (SUSY) GUT framework such as the E_6SSM, which is free from the severe constraints from proton stability noted above.

3 The fully hadronic $t\bar{t}\tau\tau$ final state at the LHC Run 2

The S_1–type LQ under consideration here is subject to stringent limits from the ATLAS and CMS collaborations, most recent ones of which are given in Table 2. These limits correspond to the BR($t\tau$) at the 95% confidence level (CL), and are obtained either directly when $t\tau$ is the exclusive decay channel of the LQ, or translated from the measurements in the $b\nu$ decay channel, when it is additionally open, as $1 - \text{BR}(b\nu)$. For the $t\tau$ decays of the pair-produced scalar LQs, these searches have so far employed final states with e and μ leptons as well as hadrons from the τ decays [19,34]. For collider analyses of some other possible LQ decay channels, see [35, 36] and references therein.

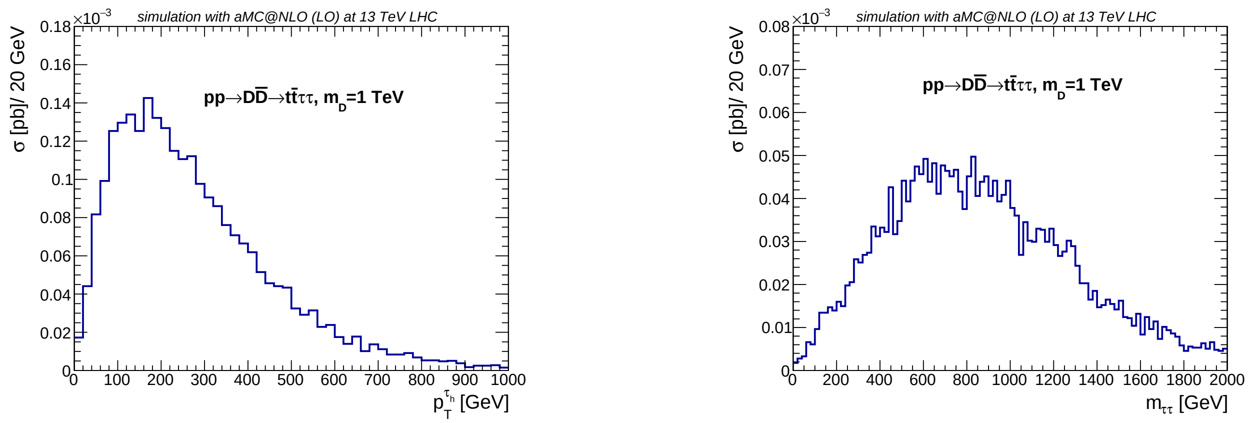

Figure 2: p_T of the τ_h (left) and the di–tau effective mass (right) for the $D\overline{D}$ signal.

In this study, we propose to complement the ATLAS and CMS probes of the $t\bar{t}\tau\tau$ intermediate state with the fully hadronic topology emerging when the decays of both the τ leptons and both the t quarks result in only jets in the final state. The analysis should select events with two τ_h (i.e., τ decaying hadronically) and at least six jets, with two of them being b–tagged. The fully hadronic final state has two advantages compared with the semi-leptonic final state. First, it allows one to fully reconstruct the D mass. Second, it increases the number of the signal events. We will show that the level of the background can be reduced to that of the expected signal using τ– and b–tagging techniques, and by exploiting the large mass of the τ_h pair.

To investigate the scope of the fully hadronic final state, we performed a signal-to-background analysis for the pair-production of LQs with a mass of 1 TeV each. The signal and the dominant $t\bar{t}jj$ background cross sections were calculated with Madgraph_aMC@NLO at the LO, with selections $\Delta R(jj) > 0.4$, $p_T^j > 20\,\mathrm{GeV}$, and $|\eta^j| < 2.4$ for the jet transverse momentum and pseudorapidity, respectively, at the parton level. τ leptons in the signal sample were forced to decay hadronically, $\tau^\pm \to \pi^\pm\pi^0\nu_\tau$. We used the efficiencies and fake rates for τ– and b–tagging, and the $\tau\tau$ mass resolution for the Z and Higgs bosons from the CMS publications (cited in the following), implying

- $p^\tau_{1(2)} = 0.7$ (0.5) efficiency of τ–tagging for real τ, with $p_T^{\tau_\mathrm{h}} > 100\,\mathrm{GeV}$ ($< 100\,\mathrm{GeV}$), and $p_{\tau-\mathrm{fake}} = 3 \times 10^{-3}$ mistagging rate for the jets (using the DeepTau algorithm [37]);
- $p_b = 0.8$ efficiency of b–tagging for real b–jets, and $p_{b-\mathrm{fake}} = 10^{-2}$ mistagging rate for non–b–jets (using the DeepJet algorithm [38]);
- 15% di–τ mass resolution for the Z and Higgs bosons, i.e., 14 GeV and 19 GeV, respectively (using the SVFit algorithm [39]).

Fig. 2 shows the distribution of $p_T^{\tau_\mathrm{h}}$ (left), and the $\tau\tau$ effective mass, $m_{\tau\tau}$, (right) for the signal. For the $t\bar{t}jj$ events, we show in the left panel of Fig. 3 the distribution of the p_T of the parton jets, i.e., the two partons in addition to $t\bar{t}$, and in the right panel the di–jet effective mass, m_{jj}, reconstructed from these parton jets.

We exploit the difference between the $m_{\tau\tau}$ and m_{jj} distributions to suppress the dominant $t\bar{t}jj$ background while preserving a high efficiency for the signal. The event selection strategy we propose is the following. Events must contain eight jets with two of them b–tagged. Six jets in the event, including the two b–tagged jets, must be associated with the two t quarks. The remaining

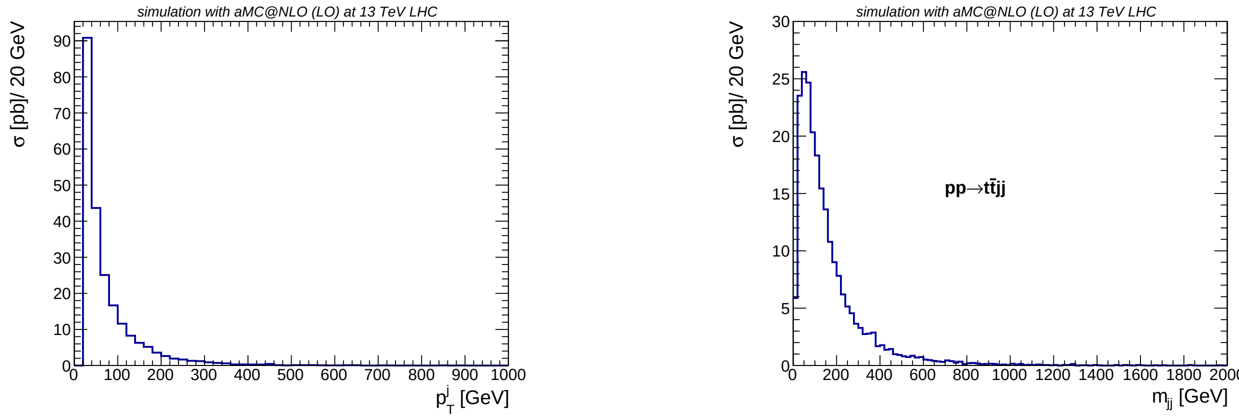

Figure 3: p_T of the parton jet (left) and the di–jet effective mass (right) for the $t\bar{t}jj$ background.

two jets have to be τ–tagged, and their effective mass should be greater than 200 GeV. In our estimate of the efficiency of the di–jet mass selection for the $t\bar{t}jj$ background we assume correct association of the jets with the t quarks for 100% of the events. It was shown, however, in [40] that the fraction of the $t\bar{t}jj$ events with correct j–to–t association is 60%. It could thus potentially make the suppression factor of the di–jet mass selection used in our estimates weaker.

For the estimation of the efficiency of the di–jet mass selection we used the generator-level $\tau_{\rm h}$ for the signal, and parton jets with $p_T > 40$ GeV and $|\eta| < 2.4$ for the $t\bar{t}jj$ background. This corresponds to the CMS di–τ high-level trigger selections used for the Run 2 [41]. Table 3 shows the cut-flow for the signal and background cross sections before and after the selections, along with the selection efficiencies. The selections considered here are the double–b– and double–τ–tagging, and the constraint on the di–τ mass.

For the signal cross section we have used two values. The cross section corresponding to the upper limit obtained in the ATLAS analysis [19] at m_D=1 TeV, and (in parenthesis) the cross section of the $pp \to D\overline{D}$ process in the E_6SSM at the 13 TeV LHC. The latter includes the NLO corrections calculated with Prospino [31]. For the background, we have used the CMS measurements of the inclusive $t\bar{t}jj$ and $t\bar{t}b\bar{b}$ cross sections [40,42], the inclusive cross sections for the $t\bar{t}Z(\to \tau\tau)$ [43] and $t\bar{t}W(\to jj)$ [44] processes, and the $Z(\to \tau\tau) + 6j$ [45] and $W(\to jj) + 6j$ [46] cross sections, as functions of the jet multiplicity. The cross section for the Higgs boson production, in association with two t quarks, followed by its $\tau\tau$ decay was taken from Ref. [47].

For two background processes, $t\bar{t}W$ and $W + 6j$, we did not evaluate the efficiency of the di–τ mass selection, and therefore give the upper limit on the corresponding cross sections after the selections. Note that the di–τ mass selection cuts eliminates the $t\bar{t}Z$ and $Z + 6j$ backgrounds due to more than 5σ difference between the di–jet threshold and m_Z. It likewise eliminates the $t\bar{t}H$ background, since the jj threshold is also about 4σ above the m_H.

As one can see from Table 3, the selection cuts used eliminate all the backgrounds except the dominant $t\bar{t}jj$. Further suppression of this background can be done by exploiting the possibility of the full reconstruction of the D mass as the effective mass of the $t\tau$ pair. The association of the two τ's and the two t quarks to the $D\overline{D}$ pair can be done as follows. Out of the two $\tau_{\rm h}$ objects, we take the one which gives the $t\tau$ mass, $\mathrm{m}_1^{\mathrm{t}\tau}$, closest to the D mass (1 TeV here). At this step, the t quark out of the two is taken randomly. The other $\tau_{\rm h}$ and the other t quark then give the mass of the second LQ, $m_2^{t\tau}$. Fig. 4 shows the distribution of the $m_1^{t\tau}$ (left) and $m_2^{t\tau}$ (right) for the $t\bar{t}jj$ events after the selections $p_T^j > 40$ GeV, $|\eta^j| < 2.4$, and $m_{jj} > 200$ GeV.

Process	σ [fb]	bb–tagging	$\tau\tau$–tagging	$m_{\tau\tau} > 200\,\text{GeV}$	σ^{sel}[fb]
Signal	0.5 (5.73)	p_b^2	$(p_1^\tau)^2$	0.97	0.15 (1.74)
$t\bar{t}jj$	275×10^3	p_b^2	$(p_{\tau-\text{fake}})^2$	0.45	0.71
$t\bar{t}Z$	$950 \times \text{BR}(Z \to \tau\tau)$	p_b^2	$(p_2^\tau)^2$	$< 5.7 \times 10^{-7}$	$< 2.6 \times 10^{-6}$
$t\bar{t}W$	$770 \times \text{BR}(W \to qq')$	p_b^2	$(p_{\tau-\text{fake}})^2$		$< 3.1 \times 10^{-3}$
$t\bar{t}H$	32	p_b^2	$(p_2^\tau)^2$	6.3×10^{-5}	3.2×10^{-4}
$Z + 6j$	50	$\text{C}_6^2 p_{b-\text{fake}}^2 (1 - p_{b-\text{fake}})^4$	$(p_2^\tau)^2$	$< 5.7 \times 10^{-7}$	$< 1.0 \times 10^{-8}$
$W + 6j$	$600 \times \text{R}_{\mu\nu}^{qq'}$	$\text{C}_8^2 p_{b-\text{fake}}^2 (1 - p_{b-\text{fake}})^6$	$(p_{\tau-\text{fake}})^2$		$< 1.0 \times 10^{-4}$

Table 3: The signal and background cross sections before and after the selections, and the selection efficiencies. The C_n^k are the binomial coefficients $\binom{n}{k}$. See text for details.

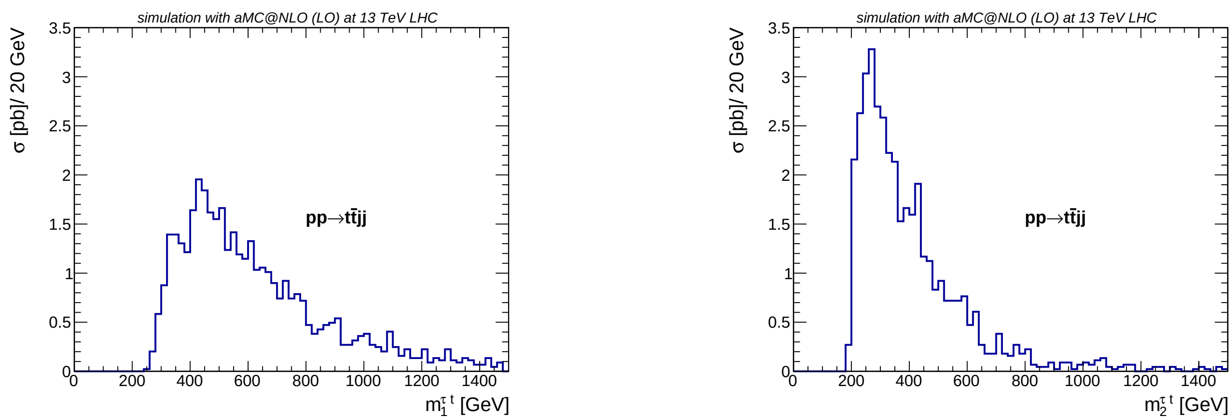

Figure 4: Distribution of the $m_1^{t\tau}$ (left) and $m_2^{t\tau}$ (right) for the $t\bar{t}jj$ events after the selections $p_T^j > 40\,\text{GeV}$, $|\eta^j| < 2.4$, and $m_{jj} > 200\,\text{GeV}$.

One can see that selecting a large $m_2^{t\tau}$ can sufficiently suppress the $t\bar{t}jj$ background. The detector mass resolution of the t quark in a fully hadronic decay mode is $14\,\text{GeV}$ (8% of the m_t) at the CMS [48]. We assume conservatively a detector resolution of 10% for the $t\tau$ mass, so that $\sigma_D^{\text{exp}} = 100\,\text{GeV}$ for $m_D = 1\,\text{TeV}$, and the natural width of D to be negligible compared to it. The efficiency of the selection $m_2^{t\tau} > m_D - 2\sigma_D^{\text{exp}} = 800\,\text{GeV}$ is 0.06 for the $t\bar{t}jj$ background, and 0.95 for the signal. It leads to cross sections of 0.14 (1.65) fb for the signal and 0.04 fb for the $t\bar{t}jj$ background after all the selections imposed so far. It corresponds to 20 (231.0) expected signal events and 6 background events for an integrated luminosity of 140 fb^{-1}, resulting in a signal significance, $2(\sqrt{S+B} - \sqrt{B})$ [49], of 5.2 (> 5). Moreover, the expected upper limit on the number of the signal events is 6.2 corresponding to the cross section of 0.16 fb, which is better than the 0.5 fb limit at $m_{S_1} = 1\,\text{TeV}$ from the ATLAS analysis [19].

Two factors related to the detector reconstruction effects can potentially weaken the performance of our analysed channel. First, as mentioned already, the association of the jets and b–jets to the t quarks is not expected to be 100% correct. And second, the jet energy smearing can reduce the impact of the $m_{jj} > 200\,\text{GeV}$ selection cut. Such effects can be estimated with a full detector simulation and reconstruction. We, however, do not expect our reported results to change drastically, and emphasise that the fully hadronic $t\bar{t}\tau\tau$ final state can nicely complement the current Run 2 analyses employing final states with leptons.

4 Interpretation in the E_6SSM

As an application of our analysis to a top-down framework, we look at how it constrains the parameter space of the E_6SSM, which is one of the best-motivated E_6-inspired SUSY models. The simple extension of the SM discussed in section 2 may be embedded in this model, with the sparticle mass scale lying in the few TeV range.

In the E_6SSM, the E_6 group is broken down to the SM gauge group plus an extra $U(1)_N$ symmetry around the GUT scale M_X, where

$$U(1)_N = \frac{1}{4}U(1)_\chi + \frac{\sqrt{15}}{4}U(1)_\psi \,. \tag{5}$$

The $U(1)_\psi$ and $U(1)_\chi$ symmetries are associated with the subgroups $E_6 \supset SO(10) \times U(1)_\psi \supset SU(5) \times U(1)_\chi \times U(1)_\psi$. It is assumed that in the E_6SSM the matter parity $Z_2^M = (-1)^{3(B-L)}$ is also preserved.

In this SUSY model the anomalies are automatically cancelled if the low-energy particle spectrum includes three complete 27-dimensional representations of E_6 (27_i with $i = 1, 2, 3$). Each 27-plet contains one generation of ordinary matter, a SM singlet field Φ_i carrying non–zero $U(1)_N$ charge, up- and down-type Higgs doublets H_i^u and H_i^d, as well as exotic scalar quarks (squarks) D_i and $\overline{D}_i$ with electric charge $\pm 1/3$. In order to avoid rapid proton decay via these exotic states, one must impose some additional symmetry. The simplest options are a Z_2^L symmetry, which implies that all supermultiplets except the lepton ones are even, or a Z_2^B symmetry, under which exotic quark and lepton supermultiples are odd, whereas the others remain even [21, 22]. In the first case the D_i and $\overline{D}_i$ are diquarks (Model I). In the second scenario the exotic squarks carry baryon ($B_D = 1/3$ and $B_{\overline{D}} = -1/3$) and lepton ($L_D = 1$ and $L_{\overline{D}} = -1$) numbers simultaneously so that they are LQs (Model II).

The exotic states in the E_6SSM may also give rise to flavour changing processes. In particular, the scalar components of the supermultiplets H_i^u and H_i^d can interact with the SM fermions of different generations, thus contributing to the amplitude of K^0–$\overline{K}^0$ oscillations, and resulting in new channels of muon decay, such as $\mu \to e^- e^+ e^-$. The non-diagonal flavour transitions can be suppressed by imposing Z_2^H symmetry, under which all the matter supermultiplets except one SM singlet superfield ($\Phi \equiv \Phi_3$) and one pair of H_i^u and H_i^d (say $H_d \equiv H_3^d$ and $H_u \equiv H_3^u$) are odd [21,22]. The discrete Z_2^H symmetry can only be an approximate one because it forbids all operators that allow the lightest exotic quark or squark to decay.

The appropriate suppression of the flavour-changing processes can be achieved when all Z_2^H symmetry-violating couplings are less than 10^{-4}. Within this original E_6SSM, only H_u, H_d, and Φ form the Higgs sector, whereas the Higgs-like doublets H_α^d and H_α^u as well as the SM singlets Φ_α ($\alpha = 1, 2$) do not acquire vacuum expectation values (VEVs). The VEV of $\Phi = \varphi/\sqrt{2}$ breaks $U(1)_N$ gauge symmetry, generating the mass of the Z' boson,

$$m_{Z'} \approx g_1' \tilde{Q}_\Phi \, \varphi \,, \tag{6}$$

and the effective μ–term. In Eq. (6) g_1' and $\tilde{Q}_\Phi = \dfrac{5}{\sqrt{40}}$ are the $U(1)_N$ gauge coupling and the $U(1)_N$ charge of the superfield Φ, respectively. The LHC constraints require the extra $U(1)_N$ gauge boson to be heavier than 4.5 TeV [50, 51]. To satisfy this constraint φ should be larger than 12 TeV.

The conservation of the Z_2^M symmetry and R-parity in the E_6SSM ensures that the lightest SUSY particle (LSP) is stable. In this case the LSP and the next-to-LSP (NLSP) are mostly composed of the fermion components of the superfields Φ_α. Using the approach discussed in [52–54], it was shown that the masses of the LSP and NLSP ($\tilde{H}_1^0$ and $\tilde{H}_2^0$) are smaller than 60–65 GeV [55]. The couplings of these fermions to the SM particles tend to be rather small. Nevertheless if the LSP had a mass close to $M_Z/2$, it could account for some of the observed cold dark matter (DM) relic density [55]. In these scenarios, the lightest Higgs boson decays mainly into either $\tilde{H}_1^0$ or $\tilde{H}_2^0$, while all other BRs are highly suppressed. Such scenarios have been already excluded by the LHC experiments. When the LSP and NLSP are considerably lighter than M_Z, the annihilation cross section of $\tilde{H}_1^0\tilde{H}_1^0 \to$ SM particles becomes too small, resulting in a DM density which is substantially larger than its measured value.

Since 2006 several modifications of the E$_6$SSM have been explored [21, 22, 56–68]. The implications of the $U(1)_N$ extensions of the MSSM have been studied for the Z–Z' mixing [69], the neutralino sector [69–71], leptogenesis [72–75] and electroweak symmetry breaking [70, 76, 77], the renormalisation group (RG) flow of couplings [70, 78], the renormalisation of VEVs [79, 80], non-standard neutrino models [81], the DM [55, 67, 68, 82], and the signatures associated with the inert neutralino states [83, 84]. Within the E$_6$SSM the upper bound on the lightest Higgs mass near the quasi-fixed point was examined in [85, 86]. The corresponding quasi-fixed point is an intersection of the invariant and quasi-fixed lines [87, 88]. The particle spectrum in the constrained E$_6$SSM (cE$_6$SSM) and its modifications has been analysed in [89–95]. The degree of the fine tuning and the threshold corrections were explored in [96, 97] and [98], respectively.

The exotic matter in the E$_6$SSM may result in distinctive LHC signatures [21, 22, 59, 62, 91, 99–102], and can give rise to non-standard Higgs decays [55, 57, 86, 103–109]. In the E$_6$SSM with approximate Z_2^H symmetry, one can also impose an exact Z_2^S symmetry, which implies that only components of the Φ_α superfields are odd [63]. In this case the LSP and NLSP become massless and decouple. The presence of such massless fermions does not affect the Big Bang Nucleosynthesis if the Z' boson is rather heavy [63]. In this variant of the model, one of the lightest R-parity odd state is stable and may account for some of the observed DM matter relic abundance. Neglecting non-renormalisable interactions and all the Z_2^H symmetry-violating couplings, the superpotential of this model can be written as

$$\begin{aligned} W_{\mathrm{E_6SSM}} &= \lambda\Phi(H_u H_d) + \lambda_\alpha \Phi(H_\alpha^d H_\alpha^u) + \kappa_i \Phi(D_i \overline{D}_i) \\ &+ \frac{1}{2} M_i N_i^c N_i^c + h_{ij} N_i^c (H_u L_j) + W_{\mathrm{MSSM}}(\mu = 0)\,, \end{aligned} \tag{7}$$

where N_i^c and L_j are supermultiplets of the right–handed neutrinos and left–handed leptons.

In the case of Model II, the Z_2^H symmetry-violating terms, that permit the lightest exotic quark or squark to decay, can be presented in the following form

$$W_2 = g_{ijk} e_i^c u_j^c D_k + \overline{g}_{ijk}(Q_i L_j)\overline{D}_k \ + g'_{ijk} d_i^c N_j^c D_k\,. \tag{8}$$

Here e_i^c, u_k^c, d_i^c, and Q_i are supermultiplets of the right-handed charged leptons, the right-handed up-type quarks, the right-handed down-type quarks, and the left-handed quark doublets, respectively.

The VEV of the Φ induces the masses of the exotic fermions, which in the leading approximation are given by

$$\mu_{D_i} = \frac{\kappa_i}{\sqrt{2}}\,\varphi\,, \qquad\qquad \mu_{\tilde{H}_\alpha} = \frac{\lambda_\alpha}{\sqrt{2}}\,\varphi\,, \tag{9}$$

where μ_{D_i} are the exotic quark masses, while $\mu_{\tilde{H}_\alpha}$ are the masses of the states composed of the fermionic components of the H_α^d and H_α^u. These masses are determined by the values of the Yukawa couplings κ_i and λ_α.

The breakdown of the gauge symmetry in the E$_6$SSM may give rise to a substantial mixing between the scalar components of the supermultiplets D_i and $\overline{D}_i$. Since we choose the field basis such that the Yukawa couplings of the D_i and $\overline{D}_i$ to Φ are flavour diagonal, which leads to a mixing only between the exotic squarks from the same family. As a consequence, the calculation of the exotic squark masses reduces to the diagonalisation of three 2×2 matrices

$$
\begin{gathered}
M^2(i) = \begin{pmatrix} M_{11}^2(i) + \Delta_{11}(i) & \mu_{D_i} X_{D_i} + \Delta_{12}(i) \\ \mu_{D_i} X_{D_i} + \Delta_{12}(i) & M_{22}^2(i) + \Delta_{22}(i) \end{pmatrix} , \\
\text{with} \quad M_{11}^2(i) = m_{D_i}^2 + \mu_{D_i}^2 + \Delta_D \,, \quad M_{22}^2(i) = m_{\overline{D}_i}^2 + \mu_{D_i}^2 + \Delta_{\overline{D}} \,, \\
X_{D_i} = A_{\kappa_i} - \frac{\lambda}{\sqrt{2}\varphi} v_1 v_2 \,, \quad \text{and} \quad \Delta_\phi = \frac{g_1'^2}{2}\Big(\tilde{Q}_{\overline{D}} v_1^2 + \tilde{Q}_D v_2^2 + \tilde{Q}_\Phi \varphi^2\Big)\tilde{Q}_\phi \,,
\end{gathered}
\tag{10}
$$

where $i = 1, 2, 3$, v_1 and v_2 are the VEVs of the Higgs doublets H_d and H_u, $\tilde{Q}_D = -\frac{2}{\sqrt{40}}$ and $\tilde{Q}_{\overline{D}} = -\frac{3}{\sqrt{40}}$ are the $U(1)_N$ charges of the D_i and $\overline{D}_i$, while $\Delta_{lm}(i)$ $(l, m = 1, 2)$ are the contributions of loop corrections. $m_{D_i}^2$ and $m_{\overline{D}_i}^2$ in the above equation are the soft scalar masses of the D_i and $\overline{D}_i$, whereas A_{κ_i} are the trilinear scalar couplings associated with the Yukawa couplings κ_i. The parameters $m_{D_i}^2$, $m_{\overline{D}_i}^2$, and A_{κ_i} break global SUSY.

The $U(1)_N$ D–term contributions to the masses of the exotic squarks are set by $M_{Z'}^2$, i.e.,

$$
\Delta_D \approx -\frac{1}{5} M_{Z'}^2 \,, \qquad \Delta_{\overline{D}} \approx -\frac{3}{10} M_{Z'}^2 \,. \tag{11}
$$

These contributions are negative because the $U(1)_N$ charge of the Φ and the $U(1)_N$ charges of the D_i and $\overline{D}_i$ are opposite. On the other hand, the $U(1)_N$ D–term contributions to the masses of the ordinary squarks and sleptons are positive. As a result, in some parts of the E$_6$SSM parameter space the exotic squarks tend to be substantially lighter than the superpartners of the SM fermions.

The magnitude of mixing in the exotic squark sector is governed by the mixing parameters X_{D_i} as well as the masses μ_{D_i}. If μ_{D_i} are sufficiently large, the mixing effects can be so substantial that the corresponding lightest exotic squarks may be among the lightest SUSY particles in the E$_6$SSM. In the Model II, the lightest scalar LQ, which we denote by D in order to identify it with the S_1–type LQ in our low-energy SM extension discussed earlier, is a linear superposition of the scalar components of the D_1 and $\overline{D}_1$ supermultiplets. Its mass is given by

$$
\begin{aligned}
m_{S_1}^2 \;=\; & \frac{1}{2}\Bigg[M_{11}^2(1) + \Delta_{11}(1) + M_{22}^2(1) + \Delta_{22}(1) \\
& - \sqrt{\Big(M_{11}^2(1) + \Delta_{11}(1) - M_{22}^2(1) - \Delta_{22}(1)\Big)^2 + 4\Big(\mu_{D_1} X_{D_1} + \Delta_{12}(1)\Big)^2}\,\Bigg] .
\end{aligned}
\tag{12}
$$

To simplify our analysis of the D in the E$_6$SSM, we set $A_{\kappa_1} = 0$. Because of this the mixing between the scalar components of the D_1 and $\overline{D}_1$ is almost always small unless $M_{11}^2(1) \approx M_{22}^2(1)$. In this case, D is almost entirely the scalar component of either D_1 or $\overline{D}_1$, and m_D^2 is defined by the smallest diagonal element of the mass matrix in Eq. (10), i.e., either by $M_{11}^2(1)$ or $M_{22}^2(1)$. Moreover, we focus on the part of the E$_6$SSM parameter space where m_D is smaller than the mass of the lightest fermion LQ. Therefore, the decay modes of the D are entirely determined by the Yukawa interactions in Eq. (8).

We examine the BRs of the decays of the D assuming that Z_2^H is mainly broken by the operators involving quarks and leptons of the third generation. In other words, we ignore all the

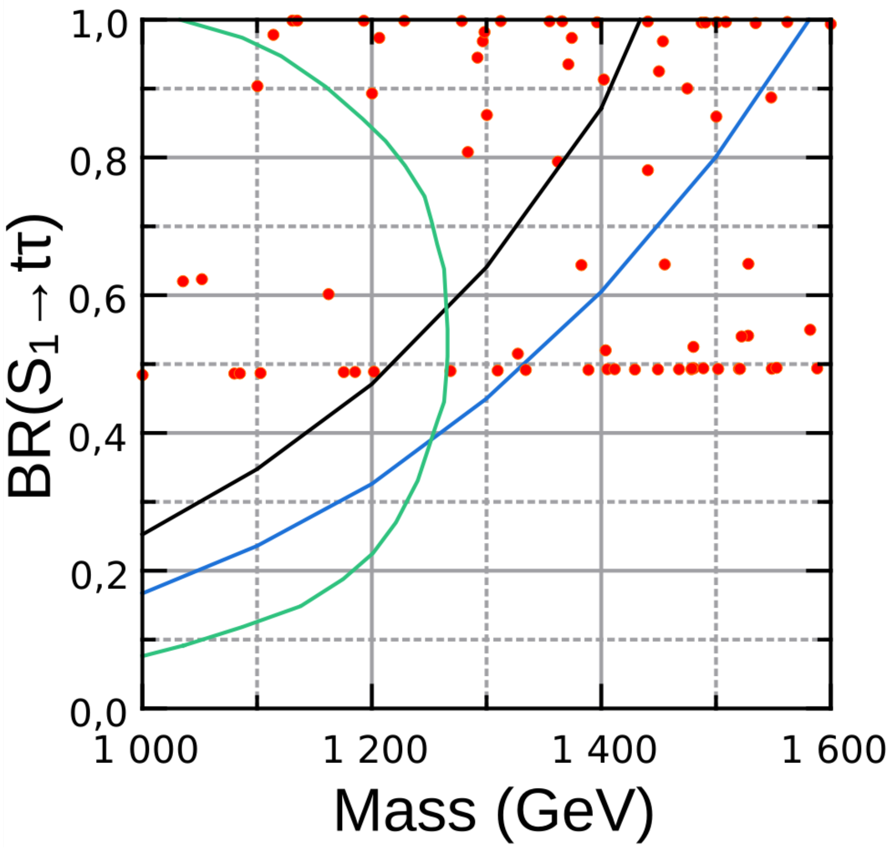

Figure 5: The exclusion bound (blue) in the $\{m_{S_1}, \mathrm{BR}(D \to t\tau)\}$ plane. The area on the left hand side of the curves is excluded. The black and green lines are from the ATLAS analyses [19] and [20], respectively. The red points correspond to some particular configurations of the E_6SSM parameter space, explained in the text.

Yukawa couplings g_{ij1} and $\overline{g}_{ij1}$, except g_{331} and $\overline{g}_{331}$, so that the D decays only into $t\tau$ and $b\nu$, and as a result $\mathrm{BR}(D \to b\nu) = 1 - \mathrm{BR}(D \to t\tau)$. To compute these BRs we incorporated the model in the Mathematica package SARAH v4.14.4 [110–116], which evaluates the interaction vertices, and writes down model files for the FORTRAN code SPheno v4.0.5 [117, 118] for performing phenomenological studies.

In Fig. 5, we show the expected exclusion contour from our signal-to-background analysis in the $\{m_{S_1}, \mathrm{BR}(D \to t\tau)\}$ plane, together with those from two recent ATLAS results. Also plotted are a number of E_6SSM parameter space points, obtained for different values of $m^2_{D_1}$, $m^2_{\overline{D}_1}$, and κ_1. When the D is predominantly the scalar component of the $\overline{D}_1$, it couples to the doublets and thus has $\mathrm{BR}(D \to t\tau)$ close to 50%. When it is mostly the D_1 it couples to singlets and decays solely to $t\tau$ if the right-handed neutrino is heavier than the D. Hence a lot of the points have the $\mathrm{BR}(D \to t\tau)$ close to either 50% or 100%. In the intermediate cases $M^2_{11}(1) \approx M^2_{22}(1)$ and there can be sizeable mixing between the scalar components of the D_1 and $\overline{D}_1$.

We see in the figure that, when the D is $\overline{D}_1$–like, our proposed analysis of the fully hadronic final state resulting from its pair-production can exclude its mass up to about 1340 GeV. The exclusion bound is even stronger, close to 1580 GeV, if the D is predominantly the D_1 instead.

5 Summary and Conclusions

LQs, being coloured, can be produced efficiently at the LHC, and are thus among the potential signatures of BSM physics. We have studied the possibility of searching third-generation LQs decaying to $t\tau$ in the fully hadronic channel. This analysis would complement the existing searches relying on the semi-leptonic channel, and has the advantage of larger statistics, as well as the possibility of reconstructing the mass of the LQ.

In the fully hadronic channel, we expect larger backgrounds from QCD processes. But we have shown that, with suitable event selection, the noise from the $t\bar{t}jj$ background gets under control, while the other backgrounds are negligible. We have found that this channel has a large estimated sensitivity for the Run 2 data, leading to an expected exclusion limit of up to 1580 GeV (assuming $\mathrm{BR}(S_1 \to t\tau) = 1$) on the LQ mass, thereby besting the ATLAS result based on the semi-leptonic channel by about 150 GeV.

As an example of a model with such a signature, we have discussed a variant of the E_6SSM in which the fundamental representation contains scalar (and fermion) LQs. The D-term contributions to the scalar LQ masses are negative, unlike those to other sfermions, and they can hence be among the lightest BSM states of this SUSY scenario. If their couplings to the fermions of the first two generations are vanishing, one of the two lightest scalar LQs can decay dominantly to $t\tau$ and the other to $t\tau$ and $b\nu$, with nearly equal BR in each of these channels. We have then demonstrated how the fully hadronic channel is expected to improve the sensitivity of the LHC searches to the pair-production of the lightest LQ in this model.

Acknowledgments

The work of SK is partially supported by STDF under grant number 37272. SM is supported in part through the NExT Institute and the STFC Consolidated Grant ST/L000296/1. HW is supported by the Carl Trygger Foundation under grant No. CTS18:164. MA acknowledges the receipt of the grant from the Abdus Salam International Centre for Theoretical Physics, Trieste, Italy.

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