Effective field theory for vector-like leptons
and its collider signals

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Abstract

We demonstrate that vector-like leptons in models with several high scales; \textit{e.g.} in composite Higgs models, are likely produced in a relatively large $\sqrt{s}$ region of the phase space. Likewise, they can easily decay into final states not containing Standard Model gauge bosons. This contrasts with the topology in which these new particles are being searched for at the LHC. Adopting an effective field theory approach, we show that searches for excited leptons must be used instead to test this scenario. We derive bounds on all the relevant interactions of dimension six; the most constrained ones being of about 0.05 TeV$^{-2}$. We build new observables to improve current analyses and study the impact on all single-field UV completions of the Standard Model extended with a vector-like lepton that can be captured by the effective field theory at tree level, in the current and in the high-luminosity phase of the LHC.

1 Introduction

Leptons beyond those of the Standard Model (SM), if they exist, have masses well above the electroweak (EW) scale, or else they would conflict with EW and Higgs precision data. Therefore, they cannot get their masses from the Higgs mechanism. Instead, any such new lepton $E$ must be vector-like with respect to the SM gauge symmetry $SU(3)_c \times SU(2)_L \times U(1)_Y$; namely the left-handed (LH) and right-handed (RH) components transform in the same representation, which allows an explicit mass term $M_{EE}$.

Direct searches for vector-like leptons (VLLs) are being performed at the LHC \cite{1-7}, with null results so far. This observation does not necessarily imply that there are no VLLs below the TeV scale. It can rather be that, contrary to what all the aforementioned experimental analyses presume, the actual VLLs (i) are mostly single produced, (ii) populate mainly the phase space of relatively large $\sqrt{s}$ and (iii) do not decay to SM gauge bosons.

This is indeed the case in several theoretical frameworks; most importantly in some composite Higgs models (CHM). The latter involve a new strong sector that confines around the scale $f_* \sim \text{TeV}$. While vector resonances are expected to have masses of order $\Lambda \sim g_* f_*$, with $g_* \gg 1$ being the coupling between composite resonances, fermionic resonances should rather lie at a scale closer to $f_*$, generating the hierarchy $m_E \ll \Lambda$. One reason is that EW precision data (EWPD) and flavour constraints are much stronger for vector than for fermionic resonances \cite{8}. One additional reason is that the Higgs mass in CHMs is much more sensitive to $m_E$ than to $m_V$, particularly in those in which the SM leptons interact sizeably with the strong sector. Such models, in turn, are motivated by the flavour anomalies \cite{9-13}.
It is therefore very likely that the actual phenomenology of VLLs at the LHC must be described by an effective field theory (EFT)\footnote{For a recent study of the impact of higher-dimensional operators on the phenomenology of vector-like quarks, see Ref. [14].} including not only the dimension-four Yukawa interaction $\sim y_L^2 HE$ (with $L_e$ and $H$ being the LH lepton doublet and the Higgs boson, respectively), but also dimension-six interactions suppressed by $1/\Lambda^2$. We note however that $y$ modifies the $Z$ coupling to the SM leptons and it is therefore very constrained by EWPD; $y \lesssim 0.1$\footnote{For a recent study of the impact of higher-dimensional operators on the phenomenology of vector-like quarks, see Ref. [14].}. Hence, the single production $pp \rightarrow E\ell$ mode populating the relatively large $\sqrt{s}$ phase space dominates, because the production cross section grows as $\sigma \sim s/\Lambda^4$. Likewise, the non-resonant decay channel $E \rightarrow \ell q\bar{q}$ can dominate over $E \rightarrow Z/h\ell$ (or $W\nu$). Naive dimensional analysis tells us that this happens provided $y \lesssim 0.1(m_E/\Lambda)^2$.

A thorough inspection of the experimental literature reveals that the search of Ref. [16], originally conceived for excited leptons, might be used to test this scenario. There are however severe limitations to translate the bounds obtained in that paper to our framework. To start with, only one dimension-six operator is considered in that experimental analysis. Second, it only considers the decay $E \rightarrow \ell q\bar{q}$, with $q$ a light quark, neither a $b$ nor a top. And third, the bounds obtained in that search cannot be translated to UV models with cut-off below 10 TeV. In order to overcome these weaknesses, we recast the experimental analysis in full detail and apply it to the entire EFT, for the different decay channels of $E$ in a wide range of masses, while keeping strict control of the EFT validity.

The article is organised as follows. In section 2, we introduce the EFT for the SM extended with $E$ (ESMEFT), and discuss its effects on single $E$ production and the subsequent decay. In section 3, we recast the most up-to-date search for excited leptons and analyse the impact of the different effective operators involving $E$ on its production and decay. We derive master formulae that can be used to automatically predict the number of events expected in any of the signal regions of the experimental search for arbitrary combinations of operators (all of which produce $E$ at very different regions of the phase space). We discuss the validity of the EFT and derive global bounds on the Wilson coefficients of the EFT accordingly. In section 4, we discuss modifications of the current analysis that improve the sensitivity to the ESMEFT, at current and future luminosities. In section 5, we discuss different UV completions of the ESMEFT, particularly all those extending the SM+$E$ renormalizable Lagrangian with just one field, and apply our analyses to constrain their parameter spaces. We conclude in section 6. We dedicate appendix A to a discussion of the technical details on the perturbative unitarity limits that we use when studying the validity of the EFT.

2 Theoretical setup

We extend the SM with an $SU(2)_L$ singlet VLL $E = E_R + E_L$ with hypercharge $Y = -1$. The leading (renormalizable) Lagrangian reads

$$L = \bar{E}(i\gamma^\mu - M_E)E - (y_L^2 HE + \text{h.c.}) \ .$$

At dimension six, the following contact interactions contribute to $pp \rightarrow E\ell$:

$$L = f_{ue} (\bar{u}_R \gamma^\mu u_R) (\bar{e}_R \gamma^\mu e_R) + f_{de} (\bar{d}_R \gamma^\mu d_R) (\bar{e}_R \gamma^\mu e_R) + f_{qe} (\bar{q}_L \gamma^\mu q_L) (\bar{e}_R \gamma^\mu e_R) + f_{qdl} (\bar{q}_L \gamma^\mu d_L) (\bar{e}_L \gamma^\mu e_L) + f_{qul} (\bar{q}_L \gamma^\mu u_L) \epsilon (\bar{e}_L \gamma^\mu e_L) + \text{h.c.} \ ,$$

where $f_i \equiv c_i/\Lambda^2$. As usual, $e_R$ denotes the SM lepton singlet; and $u_R$ and $d_R$ and $q_L$ represent the SM singlet quarks and the LH doublet, respectively. We also define $\epsilon = i\sigma_2$, with $\sigma_2$ being the second Pauli matrix.

Remarkably, non four-fermion interactions lead to processes suppressed by loop or Yukawa factors or do not grow with energy; and they can therefore be neglected. (Evidently, although our research has been triggered by previous studies of CHMs, this EFT describes any new physics scenario involving such VLL, irrespectively of whether any other new physics is much heavier or not; it is hence more generic than the usual approach to the phenomenology of VLLs.)

The relation between interaction eigenstates $e$, $E$ and the mass eigenstates $e^-, E^-$ read

$$e_R = \cos \theta_R E^-_R + \sin \theta_R E^+_R \ ,$$
$$E_R = -\sin \theta_R e^-_R + \cos \theta_R E^+_R \ ,$$

$$e_L = \cos \theta_L E^-_L + \sin \theta_L E^+_L \ ,$$
$$E_L = -\sin \theta_L e^-_L + \cos \theta_L E^+_L \ ,$$

where $\theta_R \equiv \theta_L \equiv \frac{1}{2} \sin^{-1} \frac{m_E}{\Lambda}$. \hfill (3) \hfill (4)
for the right and left chiral fields, respectively, where
\[ \sin \theta_L \to \frac{yv}{\sqrt{2}m_E}, \quad \sin \theta_R \to 0, \] in the limit \( y \to 0, m_e \to 0 \). The relation between \( M_E \) and the physical mass \( m_E \) reads
\[ M_E = \sqrt{m_E^2 - \frac{y^2v^2}{2}}, \tag{6} \]
again, in the same limit. In what follows we shall denote \( \cos \theta_L \) and \( \sin \theta_L \) by \( c_L \) and \( s_L \), respectively.

The mixing between the SM charged leptons and \( E \) modifies the coupling of the \( Z \) boson to the left current:
\[ \frac{e}{s_Wc_W} g_L e_L \gamma^\mu e_L Z_\mu = \frac{e}{s_Wc_W} (g_L^{SM} + \delta g_L) e_L \gamma^\mu e_L Z_\mu, \tag{7} \]
where \( g_L^{SM} \) is the corresponding coupling in the SM, \( \delta g_L = (yv/\sqrt{2}m_E)^2/2 \), and \( s_W \) and \( c_W \) are the sine
and cosine of the Weinberg angle, respectively. EWPD provide the following constraint on the mixing between the SM fermions and the new heavy VLL at the 95% CL [13]:
\[ |s_L| = \left| \frac{yv}{\sqrt{2}m_E} \right| < 0.021 \ (0.030), \tag{8} \]
for \( E \) mixing with electrons (muons). Taking for reference \( m_E = 0.5 \) TeV, the bound on \( y \) then reads
\[ |y| < 0.06 \ (0.09). \tag{9} \]
The regime \( y \ll 1 \) is therefore justified. The usual regime in which effective operators are ignored corresponds to \( \Lambda \to \infty \). In both cases, the single production cross section triggered by \( u\bar{u} \), to leading order in \( \theta_W \) and having neglected \( m_Z \ll \sqrt{s} \), reads
\[ \frac{d\sigma}{d\tilde{p}} = \left( \frac{1 - m_E^2}{s} \right) \left\{ - \frac{\pi^2}{3} \frac{\alpha^2}{s_W^2} \left( (s+t)(m_E^2 - s) + \frac{1}{3} \frac{1}{\Lambda^4} \left[ s (s - m_E^2) \left( \frac{c_{qul}^2}{4} + c_{ue}^2 \right) + st \left( 2c_{ue}^2 - \frac{1}{2} c_{qul}c_{tul} \right) \right] \right) \right\} \tag{10} \]
Likewise, for the counterpart driven by \( \tilde{d} \tilde{d} \) annihilation, we have the following result for the differential cross section:
\[ \frac{d\sigma}{d\tilde{p}} = \left( \frac{1 - m_E^2}{s} \right) \left\{ - \frac{\pi^2}{3} \frac{\alpha^2}{s_W^2} \left( (s+t)(m_E^2 - s) + \frac{1}{3} \frac{1}{\Lambda^4} \left[ s (s - m_E^2) \left( \frac{c_{qul}^2}{4} + c_{de}^2 \right) + st \left( c_{qe}^2 + 2stc_{de} \right) \right] \right) \right\} \tag{11} \]
Integration over \( \theta \) can be performed by noticing that \( t = m_E^2 - 2p_i \left( \sqrt{m_E^2 + p_f^2} - p_f \cos \theta \right) \), with \( p_i = \sqrt{s}/2 \) and \( p_f = (s - m_E^2) / (2\sqrt{s}) \).

In Fig. 1 we present the total single production cross section for fixed values of the Wilson coefficient \( f_{qe} \) and assuming the maximum experimentally allowed value for \( y \). For comparison, the red line shows the cross section for \( \Lambda \to \infty \), in which the only contribution comes from the \( Z \) exchange. This \( s \)-channel contribution suppressed by the gauge boson propagator scales as \( \sigma \sim 1/s \). On the contrary, in the EFT, \( \sigma \sim s/\Lambda^4 \) and therefore the cross section grows with the energy.
Together with the $y$ suppression, this effect makes the effective interactions dominate the cross section in the large $\sqrt{s}$ region, even for $f_{qe}$ as small as $\sim 0.01 \text{ TeV}^{-2}$.

The Yukawa coupling $y$ in Eq. (1) triggers also the two-body decay of $E$ into SM gauge bosons, $E \to Z/h\ell$ and $E \to W\nu$. In the limit of $y \to 0$ and $m_Z \ll m_E$, the decay width reads

$$\Gamma = \frac{y^2}{16\pi^2 c_W^2} \left( \frac{v}{m_Z} \right)^2 m_E . \quad (12)$$

Concerning the three-body decay of $E$, let us first note that, if its interactions are flavour universal, then it couples equally to all quarks and leptons, and therefore $E$ decays mostly into $\ell q q$, with $q$ being either a light, a bottom or a top quark; because there are three (colour) copies of each quark. Likewise, if similarly to the Higgs boson, $E$ couples hierarchically to all fermions according to their masses, then its decays to three leptons is again sub-dominant. In light of this observation, we will neglect the mode $E \to \ell \ell \ell$ hereafter.

This implies that the operators relevant for analysing the decay of $E$ are also precisely those in Eq. (2). The differential decay widths for $E \to \ell u \bar{d}$ reads

$$\frac{d\Gamma'}{dE_1dE_2} = \frac{3}{128\pi^3 m_E} \left[ 2E_1 m_E (m_E^2 - 2E_1 m_E) (q_{lu}^2 + 4q_{qe}^2) + 2f_{luq} f_{qul} (m_E^2 - 2E_1 m_E) (m_E^2 - 2E_3 m_E) + 2E_3 f_{qul}^2 m_E (m_E^2 - 2E_3 m_E) + 8E_2 f_{ue}^2 m_E (m_E^2 - 2E_2 m_E) \right]. \quad (13)$$

Analogously, for $E \to \ell d \bar{d}$ we have

$$\frac{d\Gamma'}{dE_1dE_2} = \frac{3}{128\pi^3 m_E} \left[ 8E_1 f_{q_e}^2 m_E (m_E^2 - 2E_1 m_E) + 8E_2 f_{ue}^2 m_E (m_E^2 - 2E_2 m_E) + 2E_3 f_{q_d}^2 m_E (m_E^2 - 2E_3 m_E) \right], \quad (14)$$

where $E_1$, $E_2$ and $E_3$ are the energies of $u(d)$, $\pi(\bar{d})$ and $\ell$, respectively. Upon integrating over the whole phase space, we arrive at

$$\Gamma' = \frac{m_E^5}{2048\pi^3} \left[ f_{luq}^2 + f_{luq} f_{qul} + f_{qul}^2 + f_{qdl}^2 + 4 (2f_{q_e}^2 + f_{ue}^2 + f_{de}^2) \right]. \quad (15)$$
Assuming $O(1)$ couplings and all quarks, the comparison between $\Gamma'$ and $\Gamma$ reveals that the three-body decay dominates for $y \lesssim 0.2 (m_E/\Lambda)^2$. Namely, $y \lesssim 0.008 (0.02)$ for $\Lambda/m_E \sim 5 (3)$. This value of $y$ is very close to the EWPD bound, it is therefore very likely that $E$ decays predominantly via EFT operators. Hereafter we study the regime $y \to 0$, and focus only on the case $\ell = \mu$. Departures from this assumption are discussed in section 6.

3 Collider signatures

In the regime $y \to 0$, the single production of $E$ and its subsequent decays proceed as depicted in Fig. 2. The experimental analysis of Ref. [16] is optimised for the light quark channel, shown in the left panel (the one with $q\bar{q}$).

In general terms, it requires first two isolated leptons with $p_T > 35$ GeV (25 GeV) and $|\eta| < 1.44$ or $1.56 < |\eta| < 2.50$ ($|\eta| < 2.4$) for electrons (muons). (Isolation is defined by the requirement that the sum of the $p_T$ of all tracks within $\Delta R = 0.3$ of a lepton is smaller than 5 GeV.) Likewise, it requires at least two anti-$k_T$ ($R = 0.4$) jets with $p_T > 50$ GeV. The leading lepton is also required to have $p_T > 230$ GeV (53 GeV) for electrons (muons). Finally, the invariant mass of the two leptons must be above 500 GeV.

The discriminating variable is the invariant mass of the two leptons and the two leading jets, $m_{\ell\ell jj}$. It is split into five energy bins: [0.5–1.5] TeV, [1.5–2.5] TeV, [2.5–3.5] TeV, [3.5–4.5] TeV, [4.5–10] TeV.

In order to determine limits on $f$, we use the energy bin [1.5–2.5] TeV, so that our EFT can be used to describe a wide range of UV models. (If we use all bins, models with $\Lambda < 10$ TeV can not be studied using the EFT approach.) Within this energy region, even $f$ of order $O(1)$ TeV$^{-2}$ are allowed by perturbative unitarity constraints; see appendix A ($m_{\ell\ell jj}$ can be used as a proxy for the partonic centre-of-mass energy $\sqrt{s}$).

Following Eqs. (10) and (11), the cross section, and therefore the number of events in each of these bins, can be written as

$$N = \frac{1}{\Lambda^2} \left[ T_{1u}^d \left( \frac{c_{qul}^2}{4} + c_{ue}^2 \right) + T_{2u}^d \left( \frac{c_{luq}^2}{4} + c_{ue}^2 + c_{qe}^2 \right) + T_{3u}^d \left( 2c_{ue} - \frac{1}{2} c_{qul}^2 \right) + T_{1d}^d \left( \frac{c_{qld}^2}{4} + c_{de}^2 \right) + T_{2d}^d \left( c_{de}^2 + c_{qe}^2 \right) + 2T_{3d}^d \right],$$

where the coefficients $T_{i}^d$, $q = u, d$, $i = 1, 2, 3$, are bin as well as mass dependent and must be obtained from simulation. To this aim we have generated signal events using MadGraph v5 [17] and Pythia v8 [18] for the three cases: $E \to \ell q\bar{q}$, $\ell b\bar{b}$, $\ell t\bar{t}$. To extract $T_{1u}^d$, $T_{2u}^d$ and $T_{3u}^d$, we turn on $c_{qul}$, $c_{luq}$ and $c_{ue}$, respectively. Furthermore, we set $c_{de} \neq 0$ to realise the decay of $E$ to the down-type quarks, whereas the semi-leptonic decay of $E$ to a pair of tops is triggered by the operator responsible for the production of $E$. All other operator coefficients are set to zero. To obtain $T_{1d}^d$, $T_{2d}^d$ and $T_{3d}^d$, we turn on $c_{qld}$, $c_{qe}$ and $c_{de}$, respectively. The same operators trigger the decay of $E$ to the down-type quarks. To allow for the decay of $E$ to a pair of tops, we switch on $c_{ue}$, except for the second case, when $c_{qe} \neq 0$ already ensures such a decay.

The Monte Carlo events are subsequently passed through a recast version of the experimental analysis that we have implemented using dedicated routines based on Fastjet v3 [19] and ROOT v6 [20,21]. We do not include detector simulation. We have validated the analysis using the dominant background given by Drell-Yan production merged up to two extra matrix element partons, finding good agreement with the numbers provided in Ref. [16] (see Fig. 7 therein).

The coefficients $T_{i}^d$ obtained in the way described above are shown in Tabs. 1, 2 and 3 for $m_E = 500$, 700 and 900 GeV, respectively. We focus on $\ell = \mu$; the (small) differences for electrons due to the different detector response are succinctly discussed in section 6.
We have assumed $m_\ell = 500$ GeV and $\mathcal{B}(E \to \mu \ell) = 1$ (top), $\mathcal{B}(E \to \mu \bar{b}) = 1$ (middle), and $\mathcal{B}(E \to \mu \bar{t}) = 1$ (bottom). We also display the SM prediction, the data and the maximal allowed signal $s_{\text{max}}$ in each bin (for muons in the final state). This latter number is computed using the CL$_{s}$ method, taking into account the uncertainty on the background displayed in the table as well as 15% uncertainty on the signal; see the text for details.

Using these tables, we have compared the predicted number of signal events in each bin as derived from Eq. (16) to that obtained directly from simulation for $\mathcal{O}(100)$ different combinations of Wilson coefficients. The latter is always contained in a band of $\pm$15% bins in $2\ell 2j$ mass [TeV]

| Bins in $2\ell 2j$ mass [TeV] | 0.5 - 1.5 | 1.5 - 2.5 | 2.5 - 3.5 | 3.5 - 4.5 | 4.5 - 10 |
|-----------------------------|-----------|-----------|-----------|-----------|---------|
| $T_1^2/\mathcal{O}^2$       |           |           |           |           |         |
| $T_2^2/\mathcal{O}^2$       |           |           |           |           |         |
| $T_3^2/\mathcal{O}^2$       |           |           |           |           |         |

Table 1: Coefficients $T_i^2$, $q = u (d)$, in TeV$^4$ and rounded to two significant figures for $pp \to \mu^+\mu^- jj$ obtained upon recasting the experimental analysis of Ref. [16] for $\sqrt{s} = 13$ TeV and total integrated luminosity $\mathcal{L} = 77.4$ fb$^{-1}$. We have assumed $m_E = 500$ GeV and $\mathcal{B}(E \to \mu \ell) = 1$ (top), $\mathcal{B}(E \to \mu \bar{b}) = 1$ (middle), and $\mathcal{B}(E \to \mu \bar{t}) = 1$ (bottom). We also display the SM prediction, the data and the maximal allowed signal $s_{\text{max}}$ in each bin (for muons in the final state). This latter number is computed using the CL$_{s}$ method, taking into account the uncertainty on the background displayed in the table as well as 15% uncertainty on the signal; see the text for details.

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| Bins in $2\ell 2j$ mass [TeV] | 0.5 - 1.5 | 1.5 - 2.5 | 2.5 - 3.5 | 3.5 - 4.5 | 4.5 - 10 |
|-----------------------------|-----------|-----------|-----------|-----------|---------|
| $T_1^2/\mathcal{O}^2$       |           |           |           |           |         |
| $T_2^2/\mathcal{O}^2$       |           |           |           |           |         |
| $T_3^2/\mathcal{O}^2$       |           |           |           |           |         |

Table 2: Coefficients $T_i^2$, $q = u (d)$, in TeV$^4$ and rounded to two significant figures for $pp \to \mu^+\mu^- jj$ obtained upon recasting the experimental analysis of Ref. [16] for $\sqrt{s} = 13$ TeV and total integrated luminosity $\mathcal{L} = 77.4$ fb$^{-1}$. We have assumed $m_E = 700$ GeV and $\mathcal{B}(E \to \mu \ell) = 1$ (top), $\mathcal{B}(E \to \mu \bar{b}) = 1$ (middle), and $\mathcal{B}(E \to \mu \bar{t}) = 1$ (bottom).
around the former\footnote{Note that in deriving Eq. (16), we have neglected the impact of the different effective operators triggering the decay of $E$ on the efficiency of the analysis. The difference in the efficiencies for selecting single produced events in two samples that differ only by the operator driving the decay of $E$ is small. Moreover, this difference tends to vanish if the two operators are linear combinations of $\{O_{luq}, O_{qe}, O_{ul} \}$ or $\{O_{de}, O_{qe} \}$. The reason is that the differential $E$ decay widths (see Eqs. (13) and (14)) driven by two operators within the same set differ only by $E_1 \leftrightarrow E_2$, while the cuts are the same for all jets. We have checked that this fact reflects well on the simulation.}. We therefore take 15% as the systematic error in our prediction for the number of signal events.

We also report in Tab. 1 the number of observed events as well as the number of expected SM events as given in Ref. [16]. Using the CL$_s$ method\footnote{Note that in deriving Eq. (16), we have neglected the impact of the different effective operators triggering the decay of $E$ on the efficiency of the analysis. The difference in the efficiencies for selecting single produced events in two samples that differ only by the operator driving the decay of $E$ is small. Moreover, this difference tends to vanish if the two operators are linear combinations of $\{O_{luq}, O_{qe}, O_{ul} \}$ or $\{O_{de}, O_{qe} \}$. The reason is that the differential $E$ decay widths (see Eqs. (13) and (14)) driven by two operators within the same set differ only by $E_1 \leftrightarrow E_2$, while the cuts are the same for all jets. We have checked that this fact reflects well on the simulation.}, including the aforementioned 15% uncertainty on the signal as well as the uncertainties on the background, we derive the maximum number of allowed signal events in each bin. These numbers are also shown in the table.

The situation is very different for the top channel. The larger number of jets in the final state, together with the relatively small top quark leptonic branching ratio, makes the corresponding $I$s even more than a factor of two smaller.

In Fig. 3 we show the limits on each of the operators of the EFT for $m_E = 500$ GeV and for $m_E = 900$ GeV. For setting bounds on $f_{luq}$ we marginalise over $f_{qul}$ (as they interfere among themselves); and vice versa. Note also that for these maximum values of $f$, the energy bin used in the analysis, $[1.5 - 2.5]$ TeV, is well within the energy regime of validity of the EFT in light of perturbative unitarity constraints; see appendix A.

### 4 Improvements and prospects

Extending the aforementioned experimental analysis with cuts on appropriate new observables can make it more sensitive to the ESMEFT. One such observable is the invariant mass of the reconstructed $E$. Note that, because $E$ is heavy, it carries less momentum that the lepton in $pp \rightarrow E\ell$. Therefore, this lepton is typically the hardest one. This effect is strengthened by the fact that when $E$ decays it releases energy to several particles.

Thus, one can reconstruct the four-momentum of $E$ as the sum of the four-momenta of the softest lepton and the two hardest jets. The invariant mass of this

| $E \rightarrow \mu d$ | $E \rightarrow \mu b$ | $E \rightarrow \mu t$ |
|----------------------|----------------------|----------------------|
| $I_1^3/\mathcal{N}$ | $I_2^3/\mathcal{N}$ | $I_3^3/\mathcal{N}$ |
| 120 (35) 400 (170) 210 (87) 72 (25) 29 (8.1) | 140 (51) 380 (170) 190 (81) 66 (23) 23 (5.7) | 110 (100) 230 (140) 100 (50) 34 (14) 12 (2.8) |
| 43 (27) 150 (84) 68 (36) 25 (11) 9.7 (3.5) | 52 (33) 140 (78) 65 (33) 22 (8.7) 7.2 (2.3) | 64 (39) 100 (53) 38 (19) 12 (4.3) 2.7 (1.1) |
| 60 (17) 200 (86) 100 (43) 37 (13) 15 (4.0) | 68 (25) 190 (86) 94 (41) 33 (11) 11 (2.6) | 56 (52) 120 (69) 50 (24) 18 (6.8) 5.2 (1.5) |

Table 3: Coefficients $I_i^3$, $q = u$ ($d$), in TeV$^3$ and rounded to two significant figures for $pp \rightarrow \mu^+\mu^-jj$ obtained upon recasting the experimental analysis of Ref. [16] for $\sqrt{s} = 13$ TeV and total integrated luminosity $L = 77.4$ fb$^{-1}$. We have assumed $m_E = 900$ GeV and $B(E \rightarrow \mu d) = 1$ (top), $B(E \rightarrow \mu b) = 1$ (middle), and $B(E \rightarrow \mu t) = 1$ (bottom).
Figure 3: The global limits on the EFT coefficients $f$ for $m_E = 0.5$ TeV (left) and $m_E = 0.9$ TeV (right), using the second bin defined in Tabs. [4] and [5].

Figure 4: Normalized distribution of $m_E^{\text{rec}}$ right after the cut on $m_{\ell^{+}\ell^{-}} > 500$ GeV, in the signal for $\mathcal{B}(E \rightarrow \ell qq) = 1$ and in the main background.

Object, $m_E^{\text{rec}}$, peaks well around the actual $m_E$ when $\mathcal{B}(E \rightarrow lqq) \sim 1$; see Fig. [4]. (For this figure we assume $m_E = 700$ GeV. Given the low sensitivity of our previous results to $m_E$, and because it is in between the two extreme cases, $m_E = 500$ GeV and $m_E = 900$ GeV considered before, we restrict to this value hereafter.) The main background, ensuing from $Z +$ jets is also shown for comparison.

We extend the current analysis with the extra cut $650 \text{GeV} < m_E^{\text{rec}} < 750 \text{GeV}$. In good approximation, the fraction of signal events that do not only pass all previous analysis cuts but also this extra one is bin and operator independent and of about 0.6. In the background, however, this number goes down to $\sim 0.1$.

The search has to be modified in a different way if one aims to be more sensitive to the case $\mathcal{B}(E \rightarrow \ell b\bar{b}) \sim 1$ or to $\mathcal{B}(E \rightarrow \ell t\bar{t}) \sim 1$. In the bottom channel, we require the presence of exactly two $b$-tagged jets. (In our simulation, $b$-jet candidates are selected among those jets with a $B$-meson within a cone of radius $\Delta R = 0.5$; the $b$-tagging efficiency is subsequently set to 0.7.) We then reconstruct $E$ as the sum of the two leading $b$-jets and the softest lepton. The invariant mass of the reconstructed $E$ is shown in Fig. [5] for both the signal and the main background, which in this case is $t\bar{t}$ (because the $b$-tagging requirement reduces $Z +$ jets to negligible levels). In this case we require $550 \text{GeV} < m_E^{\text{rec}} < 700 \text{GeV}$. The fraction of signal events surviving the new cuts is $\sim 0.25$, while for the background we get $\sim 0.05$.

Finally, in the top channel, in addition to requiring exactly two $b$-jets, we demand the presence of at least three light jets. We subsequently reconstruct $E$ as the sum of the softest lepton, the two $b$-jets and the main three light jets. The corresponding $m_E^{\text{rec}}$ is depicted in Fig. [6] in the signal and in $t\bar{t}$. We require in this case $500 \text{GeV} < m_E^{\text{rec}} < 800 \text{GeV}$. The fraction of signal (background) events surviving the new extra cuts is $\sim 0.2 (0.05)$. These numbers reflect the smaller difference between signal and background in this case.
Figure 5: Normalized distribution of \( m^\text{rec}_E \) right after the cut on \( m_{\ell^+\ell^-} > 500 \text{ GeV} \) and after requiring exactly two \( b \)-jets, in the signal for \( \mathcal{B}(E \rightarrow \mu b \bar{b}) = 1 \) and in the main background.

Table 4: Values of \( s_{\text{max}} \) in four signal regions of the improved analyses with collected luminosity of \( \mathcal{L} = 77.4 \text{ fb}^{-1} \) (HL-LHC with \( \mathcal{L} = 3 \text{ ab}^{-1} \)).

| \( E \rightarrow \ell q\bar{q} \) | \( E \rightarrow \ell b\bar{b} \) | \( E \rightarrow \ell t\bar{t} \) |
|---|---|---|
| 0.5 − 1.5 | 14 (210) | 5 (84) | 4 (68) |
| 1.5 − 2.5 | 14 (210) | 6 (101) | 4 (68) | 4 (68) |
| 2.5 − 3.5 | 14 (210) | 6 (101) | 4 (68) | 4 (68) |
| 3.5 − 4.5 | |

Using the CL\( s \) method, assuming again a 15\% uncertainty on the signal and the same uncertainties as before for the background, and assuming the data to be well described by the SM, we obtain the values of \( s_{\text{max}} \) shown in Tab. 4. They also include the numbers for the high-luminosity phase of the LHC (HL-LHC), in which the collected luminosity will reach \( \mathcal{L} = 3 \text{ ab}^{-1} \). Using these numbers, we demonstrate that the improved analyses can strengthen the sensitivity on \( f \) by more than 50\%; see Tab. 5.

5 Applications

The single-field extensions of the SM+E that contribute to the EFT at tree level are summarised in Tab. 6. The names of the new scalars follow Ref. 23, and those of the new vectors Ref. 24. In general, more than one EFT operator is generated. Assuming \( m_E = 700 \text{ GeV} \), we can use Eq. (16) together with Tab. 2 and the values of \( s_{\text{max}} \) reported in Tabs. 1 and 4 to derive bounds on the space of couplings for a fixed mass of the heavy mediator (set to 5 TeV), taking all operators into account.

Assuming for simplicity that all couplings not involving \( E \) are equal, we show these results for the scalar mediators in Fig. 7. We also show for comparison the bounds from low-energy data and dijet searches 23. Interestingly, e.g., in the case of \( \omega_1 \), we see that for sufficiently large values of \( y_{\ell q} \), the bound on \( y_{\ell u} = y^u_{\ell u} \) from our study is about 6 times more stringent than that from other data, and it can be improved by a factor of two at the HL-LHC. Despite not being explicitly shown, results for vector boson extensions of the SM+E are similar.

In good approximation, our results can also be easily extended to four-fermion operators involving only second and third generation quarks. For example, due to the PDF suppression, the cross section for single \( E \) production initiated by bottom quarks is about two orders of magnitude smaller than that initiated by down quarks. Therefore, it is expected that values of the Wilson coefficients \( f \) ten times larger can be probed at the LHC. Note that the EFT is still valid in this case if we still restrict to \( m_{\ell ij} < 2.5 \text{ TeV} \); see appendix A.

This observation can be used to explore the sensitivity
Relevant fermionic current

\[
J_E = \mu \frac{i\sigma_2}{2} u_L \gamma^a T^a \overline{u}_L + y^u_i \mu \frac{i\sigma_2}{2} \overline{u}_L \gamma^a T^a u_R
\]

\[
J = y^E \overline{E} \sigma_{\mu\nu} \gamma^\mu E + y^d \overline{d} \sigma_{\mu\nu} \gamma^\mu d + y^e \overline{e} \sigma_{\mu\nu} \gamma^\mu e
\]

\[
J = y^\omega \overline{\omega} \sigma_{\mu\nu} \gamma^\mu \omega + y^\nu \overline{\nu} \sigma_{\mu\nu} \gamma^\mu \nu + y^\gamma \overline{\gamma} \sigma_{\mu\nu} \gamma^\mu \gamma
\]

\[
J = y^{\Pi_2} \overline{\Pi_2} \sigma_{\mu\nu} \gamma^\mu \Pi_2 + y^{\Pi_3} \overline{\Pi_3} \sigma_{\mu\nu} \gamma^\mu \Pi_3
\]

\[
J = y^{\Pi_5} \overline{\Pi_5} \sigma_{\mu\nu} \gamma^\mu \Pi_5 + y^{\Pi_6} \overline{\Pi_6} \sigma_{\mu\nu} \gamma^\mu \Pi_6
\]

Table 5: Bounds on the Wilson coefficients rounded to two significant figures, in TeV$^{-2}$, in the current and improved (future) analyses. We have assumed $m_E = 700$ GeV and used the energy bin [1.5 – 2.5] TeV.

\[
\begin{array}{c|c|c|c}
\hline
B(E \rightarrow \mu q \overline{q}) & B(E \rightarrow \mu b \overline{b}) & B(E \rightarrow \mu t \overline{t}) \\
\hline
f_{ue} & 0.060, 0.037 & 0.060, 0.038 & 0.076, 0.050 \\
\hline
f_{de} & 0.079, 0.049 & 0.081, 0.051 & 0.100, 0.072 \\
\hline
f_{qe} & 0.048, 0.030 & 0.049, 0.031 & 0.060, 0.042 \\
\hline
f_{qdl} & 0.110, 0.066 & 0.110, 0.067 & 0.120, 0.085 \\
\hline
f_{qul} & 0.130, 0.082 & 0.130, 0.083 & 0.160, 0.110 \\
\hline
f_{f_{luq}} & 0.220, 0.140 & 0.220, 0.140 & 0.240, 0.170 \\
\hline
\end{array}
\]

Table 6: The relevant Lagrangian for a scalar $\sigma$ is $L = \partial_\mu \sigma \partial^\mu \sigma - m_\sigma^2 \sigma^2 - (\sigma^2 J_{\sigma} + \text{h.c.})$. For a vector $V$ instead we have $L = -\partial_\mu V^\mu \partial^\mu V^\nu + m_V^2 V^\mu V^\nu - (V^\mu J^\nu + \text{h.c.})$. For each row in the top (bottom) part of the table, $m = m_\sigma$ ($m_V$).

to other models. For concreteness, following Ref. [13], let us consider the SM+E extension with a full singlet vector boson $V$ with mass $m_V$ and couplings

\[
L = V^\mu \left[ g_{Vq_\mu q_\nu} \overline{q}_L q_\nu + g_{Vq} \overline{q}_L \epsilon_{\gamma\mu\nu} + g_{VE} \overline{E} \gamma_{\mu\nu} + \text{h.c.} \right] + \ldots
\]

The ellipsis encode terms not relevant for us, such as light lepton couplings to $V$, etc.

We fix $g_{Vq_\mu q_\nu} \sim 0.05 m_V^2 / \text{TeV}^2$. In the original reference this value is motivated by the flavour anomalies [25,31]. We keep the strong coupling $g_{VE}\epsilon$ free; while in the original reference it is fixed to 2.5. (The phenomenology studied there is not very sensitive to the value of this coupling.)

Upon integrating $V$ out, the only ESMEFT operator (relevant for single production) generated is $O_{q_e}$, with

\[
f_{q_e} = -\frac{g_{Vq_\mu q_\nu} g_{VE} \epsilon}{m_V^2} \sim -0.05 g_{VE} \epsilon \text{ TeV}^{-2}.
\]

To compare the complementarity between our current analysis and that of Ref. [13], let us assume that $E$ decays equally into SM gauge bosons and via the four-fermion operators. Thus, the region that can be probed
Figure 7: Constraints on the couplings of the scalar UV completions of the ESMEFT derived under the assumption that all couplings to the SM fields in a given model are equal. We have used Eq. (16) along with the values of $I$s from Tab. 3 and those of $s_{\text{max}}$ from Tabs. 2 and 4, assuming the light quark decay channel of $E$ and the second bin. “Current”, “improved” (“future”) refer to the developed LHC (HL-LHC) analyses described in the text. The light blue regions are excluded from EWPD or dijet searches at the LHC [23].

at the HL-LHC following Ref. [13] (see right panel of Fig. 5 therein) is depicted in blue in Fig. 8. The area below the line $m_V = 2m_E$, in which an on-shell produced $V$ decays into $EE$, is not accessible within that analysis. Due to the resonant nature of that search, the region above $m_V = 2.5$ TeV remains open.

On the other hand, within our current analysis in the bin $[1.5 - 2.5]$ TeV, we can probe values of $f_{qe}$ of order $0.3$ TeV$^{-2}$ at the HL-LHC, which corresponds to $g_{VE}\ell \sim 6$. (This value is significantly smaller if $E$ decays only via four-fermions; in which case the analysis of Ref. [13] is not sensitive to the model.) Notably, this constraint is $m_V$-independent, provided $m_V > 2.5$ TeV so that the EFT approach is valid. The corresponding bound is shown in red in Fig. 8.

Thus, we can conclude that for sufficiently large $g_{VE}\ell$, our analysis together with that in Ref. [13] can completely probe the corresponding explanation of the flavour anomalies.

3We are making the conservative assumption that within our current analysis we are equally sensitive to values of $m_E$ above our higher benchmark of $m_E = 900$ GeV. In light of the experimental results in Ref. [16], it is expected that the sensitivity to heavier $E$ could be even better.

Figure 8: Reach of the LHC to the model described in Eq. (17) using the resonant analysis of Ref. [13] (blue) versus the reach using the EFT analysis described in this article (red); see text for details.

6 Conclusions

Using an effective field theory (EFT) approach, we have demonstrated that, differently to what current searches for vector-like leptons (VLLs) $E$ assume, $E$ can be likely produced at high values of $\sqrt{s}$ via four-fermion interac-
tions at the LHC. They also decay as \( E \rightarrow \ell q \bar{q} \) with no intermediate Standard Model (SM) gauge bosons.

We have shown that there are other (few) experimental analyses, most importantly searches for excited leptons \[16\], that are very sensitive to our hypothesis. They are however limited in scope, because they focus only on the case \( q = u, d, c, s \) leaving bottom and top quarks aside, as well as a single four-fermion operator. Moreover, the statistical analysis in Ref. \[16\] does not apply to models with further particles below 10 TeV. Likewise, interpreting their bounds on ESMEFT operators involving only sea quarks breaks the EFT validity. (These objections apply also to other previous similar analyses \[2,32\].)

Thus, we have worked out the most generic base of EFT contact interactions involving \( E \) to dimension six. Upon recasting the experimental analysis of Ref. \[16\], we have obtained global bounds on all the EFT directions, for light, bottom and top quarks separately. To this aim, we have restricted to events with \( E \sim 0.015 \text{ TeV}^{-2} \). This translates to \( f \sim 0.015 \text{ TeV}^{-2} \).

In the electron channel, our bounds on the Wilson coefficients are only slightly altered. Taking, for example, \( m_E = 500 \text{ GeV} \) and using the energy bin \([1.5 - 2.5] \text{ TeV} \), the bounds on \( f_{qe} \) for muons and for electrons read, respectively, 0.046 (0.067) \( \text{ TeV}^{-2} \) and 0.048 (0.070) \( \text{ TeV}^{-2} \), for \( E \) decaying into light and bottom jets (tops).

We have also modified the current analysis with cuts on new observables (most importantly the number of \( b \)-tagged jets and the reconstructed mass of \( E \)); improving the aforementioned bounds by a factor of \( \sim 1.6 \) (\( \sim 1.4 \)) for light and bottom quarks (tops).

Finally, we have applied our findings to concrete UV completions of the SMEFT extended with \( E \). In particular, we have classified all possible single field extensions of the SM+\( E \) that induce the four-fermion interactions of interest at tree level. The limits on the couplings of these fields to purely SM currents can overcome those from low-energy data and dijet searches at the LHC by almost an order of magnitude with the improved analysis.

Altogether, our work motivates different searches for VLLs, that might be implemented by small modifications of current searches for excited leptons.

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**A Perturbative unitarity bounds**

The aim of this appendix is to discuss the validity of the EFT approach. To this end, we first sum up the perturbative unitarity condition and apply it to the tree-level EFT amplitudes \( qq \rightarrow E\ell \). More specifically, we derive constraints on the maximum partonic centre-of-mass energy \( \hat{s} \) at which the EFT is applicable as a function of the Wilson coefficient \( f \) of each operator.

The unitarity of the \( S \) matrix, \( SS^T = 1 \), together with the requirement of perturbativity imply that the partial waves \( T_j \) in the following partial wave decomposition of inelastic scattering amplitudes

\[
M = 16\pi \sum_{\ell=0}^{\infty} (2J+1) T^{(J)}_{\lambda_1',\lambda_2';\lambda_1,\lambda_2}(\hat{s}) d^{(J)}_{\mu,\nu}(\theta) ,
\]

should fulfill the condition

\[
|T^{(J)}_{\lambda_1',\lambda_2';\lambda_1,\lambda_2}(\hat{s})| \leq \frac{1}{2} \tag{20}
\]

for each \( J \in \{J_{\text{min}},J_{\text{min}}+1,\ldots\} \). In this expression, \( \lambda_{1,2} \) and \( \lambda_{1,2}' \) are the helicities of the initial and final particles, respectively; \( \mu = \lambda_1 - \lambda_2 \), \( \nu = \lambda_1' - \lambda_2' \); \( d^{(J)} \) are
the Wigner matrices in the limit of azimuthal scattering angle $\phi \to 0$, and $J_{\min} = \max(\{|\lambda_1 - \lambda_2|, |\lambda'_1 - \lambda'_2|\})$. For more details on the partial wave unitarity condition, see e.g. Ref. [33].

Since the EFT amplitudes grow with the energy $\hat{s}$, so do the partial waves. We define the distinguished energy scale $\sqrt{\hat{s}^U}$ as the one that saturates the condition in Eq. [20]:

$$|T^{(J)}_{\lambda_1 \lambda_2 ; \lambda'_1 \lambda'_2}(\hat{s}^U)| = \frac{1}{2}.$$  \hspace{0.5cm} (21)

Importantly, $\hat{s}^U$ is a function of $f$. Given that for energies above $\hat{s}^U$ the EFT amplitudes are ill-defined, $\hat{s}^U$ defines the upper bound on $\hat{s}$ for which the EFT approach is valid.

Typically, the first partial wave yields the strongest unitarity bounds on $\hat{s}$. Correspondingly, we derive the bounds using $T^{(J_{\min})}$. Since in our study the global bounds on $f$ are expressed in terms of each Wilson coefficient $f$ separately, we compute the unitarity bounds using one operator at a time. The $J$-th partial wave projections are computed using the orthogonality of the Wigner functions:

$$T^{(J)}_{\mu, \nu} = \frac{1}{32\pi} \int_{-1}^{1} d\cos\theta \, d^{(J)}_{\mu, \nu}(\theta) \mathcal{M}.$$  \hspace{0.5cm} (22)

More specifically, for each operator we consider all helicity $q\bar{q} \to E\ell$ amplitudes, where $q = u, d$, that are non-vanishing in the relativistic limit. For each such helicity combination we project the amplitude $\mathcal{M}$ onto the $J_{\min}$ partial wave and derive the corresponding bound on $\hat{s}$. Finally, we identify $\hat{s}^U$ as the lowest amongst all such bounds.

Unitarity bounds for different values of $f$ are presented in Tab. 7. For example, for $f = 1$ TeV$^{-2}$ the bounds are in the range $\sqrt{\hat{s}^U} \in [5 - 7]$ TeV, depending on the operator involved.

For completeness, let us comment on the values of $c$ in $f = c/\Lambda^2$ setting $\Lambda = \sqrt{\hat{s}^U}$, for different values of $f$ and for each effective operator. Independently of the value of $f$, we obtain that $\sqrt{c} = 6.1$ for $f = f_{ue}, f_{de}, f_{qe}$; $\sqrt{c} = 5$ for $f = f_{qdit}, f_{quti}, f_{luq}$, and $\sqrt{c} = 7.1$ for $f = f_{luq}$.

Note that, for a fixed $f$, $\Lambda = \sqrt{\hat{s}^U}$ can be (roughly) identified with the upper bound on the scale $\Lambda$. (The new physics scale in a UV completion should not be significantly separated from $\sqrt{\hat{s}^U}$ because it is responsible for unitarization of the complete amplitudes.) Therefore for a given $f$, the value of $c$, assuming $\Lambda = \sqrt{\hat{s}^U}$, is an approximate upper bound on the corresponding UV coupling.

Interestingly, the aforementioned values of $c$ are (i) independent of the value of $f$ and (ii) in the range between 1 and $4\pi$, hence indeed close to the perturbative regime (as required by the perturbative unitarity condition).

Given this, a discussion on the EFT consistency of the analyses in sections 3 and 4 is in order. We note that the larger the value of $f$, the stronger the unitarity bounds. Thus, in particular, for the largest $f$ within the limits, the $\sqrt{\hat{s}^U}$ should not be lower than the chosen cut-off on the (proxy) variable $m_{\ell\ell\ell}$j. Otherwise one uses events outside the validity of the EFT amplitudes while setting limits on the effective coefficients $f$; turning them to be not suitable for EFT interpretation.

In Tab. 8 we present the unitarity bounds as function of $f$ for the values relevant for the 2.5 TeV cut-off case. Comparing the table with Fig. 3 one can see that all limits on $f$ correspond to unitarity bounds that are not lower than the 2.5 TeV cut-off. Hence the limits are EFT interpretable. More explicitly, unitarity bounds $\sqrt{\hat{s}^U}$ that e.g. correspond to $f$ from the first column in Tab. 5 read 25, 22, 27, 16, 13 and 16 TeV for $f_{ue}, f_{de}, f_{qe}, f_{qdit}, f_{quti}$ and $f_{luq}$, respectively.

| $f$ [TeV$^{-2}$] | $f_{ue}$ | $f_{de}$ | $f_{qe}$ | $f_{qdit}$ | $f_{quti}$ | $f_{luq}$ |
|-----------------|---------|---------|---------|------------|------------|---------|
| 10              | 1.9     | 1.9     | 1.9     | 1.6        | 1.6        | 2.2     |
| 1               | 6.1     | 6.1     | 6.1     | 5.0        | 5.0        | 7.1     |
| 0.1             | 19      | 19      | 19      | 16         | 16         | 22      |
| 0.01            | 61      | 61      | 61      | 50         | 50         | 71      |

Table 7: Solutions for $\sqrt{\hat{s}^U}$ (in TeV) from tree-level partial wave unitarity in the presence of a single operator at a time, for different values of the Wilson coefficients $f$. The examined processes are $u\bar{u} \to E\ell, d\bar{d} \to E\ell$.

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| $f$ [TeV$^{-2}$] | $f_{ue}$ | $f_{de}$ | $f_{qe}$ | $f_{qut}$ | $f_{qull}$ |
|-----------------|---------|---------|---------|---------|---------|
| 0.3             | 11      | 11      | 11      | 9.2     | 9.2     | 13      |
| 0.25            | 12      | 12      | 12      | 10      | 10      | 14      |
| 0.2             | 14      | 14      | 14      | 11      | 11      | 16      |
| 0.15            | 16      | 16      | 16      | 13      | 13      | 18      |
| 0.1             | 19      | 19      | 19      | 16      | 16      | 22      |
| 0.075           | 22      | 22      | 22      | 18      | 18      | 26      |
| 0.06            | 25      | 25      | 25      | 20      | 20      | 29      |
| 0.05            | 27      | 27      | 27      | 22      | 22      | 32      |

Table 8: Same as Tab. 7 but for different values of the Wilson coefficients $f$. [11]

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