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Compilation of low-energy constraints on 4-fermion operators in the SMEFT

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Abstract: We compile information from low-energy observables sensitive to flavor-conserving 4-fermion operators with two or four leptons. Our analysis includes data from e^+e^- colliders, neutrino scattering on electron or nucleon targets, atomic parity violation, parity-violating electron scattering, and the decay of pions, neutrons, nuclei and tau leptons. We recast these data as tree-level constraints on 4-fermion operators in the Standard Model Effective Field Theory (SMEFT) where the SM Lagrangian is extended by dimension-6 operators. We allow all independent dimension-6 operators to be simultaneously present with an arbitrary flavor structure. The results are presented as a multi-dimensional likelihood function in the space of dimension-6 Wilson coefficients, which retains information about the correlations. In this form, the results can be readily used to place limits on masses and couplings in a large class of new physics theories.

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JHEP08(2017)123

Contents

1 Introduction

The ongoing exploration of the high-energy frontier at the LHC strongly suggests that the only fundamental degrees of freedom at the weak scale are the Standard Model (SM) ones. Moreover, their perturbative interactions are well described by the most general renormalizable SM Lagrangian invariant under the SU(3) × SU(2) × U(1) local symmetry. A large number of precision measurements has been performed in order to test the SM predictions. The motivation is that some unknown heavy particles may affect the coupling strength or induce new effective interactions between the SM particles.

One framework designed to describe such effects in a systematic fashion goes under the name of the SM Effective Field Theory (SMEFT). In this approach, the SM particle content and symmetry structure is retained, but the usual renormalizability requirement is abandoned such that interaction terms with canonical dimensions $D > 4$ are allowed in the Lagrangian. These higher-dimensional operators encode, in a model-independent way, the effects of new particles with masses above the weak scale. One can then analyze experimental searches once and for all within this framework. The output of such analysis, namely numerical values for the Wilson coefficients of higher-dimensional operators, can then be applied to any new physics model covered by the SMEFT. Significant progress has been recently achieved concerning the automation of this EFT matching [1–6]. The efficient SMEFT program should be compared with model-dependent studies where non-trivial hadronic effects, PDFs, radiative corrections, experimental errors, cuts, etc., have to be taken into account for each model.

Assuming lepton number conservation, leading SMEFT contributions are expected to originate from dimension-6 operators [7, 8]. There is a vigorous program to characterize the effects of the dimension-6 operators on precision observables and derive constraints on their Wilson coefficients in the SMEFT Lagrangian [9–49]. Most of these analyses assume that the dimension-6 operators respect some flavor symmetry in order to reduce the number of independent parameters. On the other hand, refs. [32, 43] allowed for a completely general set of dimension-6 operators, demonstrating that the more general approach is feasible.

This paper further pursues the approach of refs. [32, 43], providing new constraints on the SMEFT where all independent dimension-6 operators may be simultaneously present with an arbitrary flavor structure. We compile information from a plethora of low-energy flavor-conserving experiments sensitive to electroweak gauge boson interactions with fermions and to 4-fermion operators with 2 leptons and 2 quarks (LLQQ) and 4 leptons (LLLL). There are two main novelties compared to the existing literature. First, precision constraints on the LLQQ operators have not been attempted previously in the flavor-generic situation. Therefore our results are relevant to a larger class of UV completions where new physics couples with a different strength to the SM generations. Note that, in particular, all models addressing the recent B-meson anomalies (see e.g. [50–54]) must necessarily involve exotic particles with flavor non-universal couplings to quarks and leptons. Our analysis provides model-independent constraints that have to be satisfied by all such constructions. Second, we include in our analysis the low-energy flavor observables (nuclear, baryon and meson decays) recently summarized in ref. [55]. At the parton level these pro-

JHEP08(2017)123

cesses are mediated by the quark transitions $d(s) \to u\ell\bar{\nu}_\ell$, hence they can probe the LLQQ operators. We will show that for certain operators the sensitivity of these observables is excellent, such that new stringent constraints can be obtained. Moreover, the low-energy flavor observables offer a sensitive probe of the W boson couplings to right-handed quarks.

Our analysis is performed at the leading order in the SMEFT. We ignore the effects of dimension-6 operators suppressed by a loop factor, except for the renormalization group running within a small subset of the LLQQ operators. Moreover all dimension-8 and higher operators are neglected, and only the linear contributions of the dimension-6 Wilson coefficients are taken into account. The corollary is that the likelihood we obtain for the SMEFT parameters is Gaussian. All in all, we provide simultaneous constraints on 61 linear combinations of the dimension-6 Wilson coefficients. In this paper we quote the central values, the 68% confidence level (CL) intervals, while the correlation matrix is provided in the attached `Mathematica` notebook [56]. That file also contains the full likelihood function in an electronic form, so that it can be more easily integrated into other analyses.

The outline of the paper is the following. section 2 introduces the theoretical framework and the necessary notation. section 3 presents the experimental input of our analysis. section 4 contains the results of our fit, in the general case and in some interesting limits. Finally section 5 discusses the interplay with LHC searches, and section 6 contains our conclusions.

2 Formalism and notation

2.1 SMEFT with dimension-6 operators

Our framework is that of the baryon- and lepton-number conserving SMEFT [7, 8]. The Lagrangian is organized as an expansion in $1/\Lambda^2$, where Λ is interpreted as the mass scale of new particles in the UV completion of the effective theory. We truncate the expansion at $\mathcal{O}(\Lambda^{-2})$, which corresponds to retaining operators up to the canonical dimension D=6 and neglecting operators with $D \geq 8$. The Lagrangian takes the form

$$\mathcal{L} = \mathcal{L}_{\rm SM} + \sum_i \frac{c_i}{v^2} O_i^{D=6}, \tag{2.1}$$

where $\mathcal{L}_{\rm SM}$ is the SM Lagrangian, $v = (\sqrt{2}G_F)^{-1/2} \simeq 246\,\mathrm{GeV}$, each $O_i^{D=6}$ is a gauge-invariant operator of dimension D=6, and c_i are the corresponding Wilson coefficients that are $\mathcal{O}(\Lambda^{-2})$. $O_i^{D=6}$ span the complete space of dimension-6 operators, see refs. [57, 58] for examples of such sets.

In order to connect the SMEFT to observables it is convenient to rewrite eq. (2.1) using the mass eigenstates after electroweak symmetry breaking. Then the effects of dimension-6 operators show up as corrections to the SM couplings between fermion, gauge and Higgs fields, or as new interaction terms not present in the SM Lagrangian. The discussion and notation below follows closely that in section II.2.1 of ref. [59]. We define the mass eigenstates such that all kinetic and mass terms are diagonal and canonically normalized. We also redefine couplings such that, at tree level, the relation between the usual SM input observables G_F, α, m_Z and the Lagrangian parameters g_L, g_Y, v is the same as in the

SM. See ref. [59] for complete definition of conventions and the complete list of interaction terms with up to 4 fields. In the following we only highlight the parts of the mass eigenstate Lagrangian directly relevant for the analysis in this paper.

One important effect from the point of view of precision measurements is the shift of the interaction strength of the weak bosons. We parametrize the interactions between the electroweak gauge bosons and fermions as

$$\begin{aligned}
\mathcal{L} \supset\, & eA^\mu \sum_{f=u,d,e} Q_f(\bar{f}_I\bar{\sigma}_\mu f_I + f_I^c\sigma_\mu \bar{f}_I^c) \\
& + \frac{g_L}{\sqrt{2}}\left[W^{\mu+}\bar{\nu}_I\bar{\sigma}_\mu(\delta_{IJ} + [\delta g_L^{We}]_{IJ})e_J + W^{\mu+}\bar{u}_I\bar{\sigma}_\mu\left(V_{IJ} + \left[\delta g_L^{Wq}\right]_{IJ}\right)d_J + \text{h.c.}\right] \\
& + \frac{g_L}{\sqrt{2}}\left[W^{\mu+}u_I^c\sigma_\mu\left[\delta g_R^{Wq}\right]_{IJ}\bar{d}_J^c + \text{h.c.}\right] \\
& + \sqrt{g_L^2+g_Y^2}\,Z^\mu \sum_{f=u,d,e,\nu} \bar{f}_I\bar{\sigma}_\mu\left((T_3^f - s_\theta^2 Q_f)\delta_{IJ} + \left[\delta g_L^{Zf}\right]_{IJ}\right) f_J \\
& + \sqrt{g_L^2+g_Y^2}\,Z^\mu \sum_{f=u,d,e} f_I^c\sigma_\mu\left(-s_\theta^2 Q_f\delta_{IJ} + \left[\delta g_R^{Zf}\right]_{IJ}\right)\bar{f}_J^c. \qquad (2.2)
\end{aligned}$$

Here, g_L, g_Y are the gauge couplings of the $\mathrm{SU}(2)_L\times\mathrm{U}(1)_Y$ local symmetry, the electric coupling is $e = g_L g_Y/\sqrt{g_L^2+g_Y^2}$, the sine of the weak mixing angle is $s_\theta = g_Y/\sqrt{g_L^2+g_Y^2}$, and $I,J = 1,2,3$ are the generation indices. For the fermions we use the 2-component spinor formalism and we follow the conventions of ref. [60], unless otherwise noted.[1] The SM fermions f_J, f_J^c are in the basis where the mass terms are diagonal, and then the CKM matrix V appears in the quark doublets as $q_I = (u_I, V_{IJ}d_J)$. The effects of dimension-6 operators are parameterized by the vertex corrections δg that in general can be flavor-violating. For flavor-diagonal interactions we will employ the shorter notation $[\delta g_{L/R}^{Vf}]_{JJ} \equiv \delta g_{L/R}^{Vf_J}$.

The vertex corrections can be expressed as linear combinations of the Wilson coefficients c_i in eq. (2.1), see appendix A for the map to the Warsaw basis. We find more transparent to recast the results of precision experiments as constraints on δg's. This is completely equivalent, provided one takes into account that not all δg's in eq. (2.2) are independent.[2] Indeed, the mapping between the vertex corrections and the Wilson coefficients implies the relations $[\delta g_L^{Z\nu}]_{IJ} - [\delta g_L^{Ze}]_{IJ} = [\delta g_L^{We}]_{IJ}$, and $[\delta g_L^{Wq}]_{IJ} = [\delta g_L^{Zu}]_{IK}V_{KJ} - V_{IK}[\delta g_L^{Zd}]_{KJ}$.

In this paper we focus on flavor-conserving observables that target flavor-diagonal Wilson coefficients. We will express the experimental constraints using the following set of independent flavor-diagonal vertex corrections:

$$\delta g_L^{Ze_I},\ \delta g_R^{Ze_I},\ \delta g_L^{We_I},\ \delta g_L^{Zu_I},\ \delta g_R^{Zu_I},\ \delta g_L^{Zd_I},\ \delta g_R^{Zd_I},\ \delta g_R^{Wq_I}. \qquad (2.3)$$

[1] Compared to [60], we use a different normalization of the antisymmetric product of the σ matrices: $\sigma_{\mu\nu} = \frac{i}{2}(\sigma_\mu\bar{\sigma}_\nu - \sigma_\nu\bar{\sigma}_\mu)$, $\bar{\sigma}_{\mu\nu} = \frac{i}{2}(\bar{\sigma}_\mu\sigma_\nu - \bar{\sigma}_\nu\sigma_\mu)$.

[2] More generally, it is often convenient to parametrize the space of dimension-6 operators using δg's and other independent parameters in the mass eigenstate Lagrangian that are in a 1-to-1 linear relation with the set of Wilson coefficients c_i [24]. One example of such parametrization goes under the name of the Higgs basis and is defined in ref. [59].

JHEP08(2017)123

Chirality conserving $(I,J=1,2,3)$	Chirality violating $(I,J=1,2,3)$
$[O_{\ell q}]_{IIJJ}=(\bar{\ell}_I\bar{\sigma}_\mu\ell_I)(\bar{q}_J\bar{\sigma}^\mu q_J)$	$[O_{\ell equ}]_{IIJJ}=(\bar{\ell}^j_I\bar{e}^c_I)\epsilon_{jk}(\bar{q}^k_J\bar{u}^c_J)$
$[O^{(3)}_{\ell q}]_{IIJJ}=(\bar{\ell}_I\bar{\sigma}_\mu\sigma^i\ell_I)(\bar{q}_J\bar{\sigma}^\mu\sigma^i q_J)$	$[O^{(3)}_{\ell equ}]_{IIJJ}=(\bar{\ell}^j_I\bar{\sigma}_{\mu\nu}\bar{e}^c_I)\epsilon_{jk}(\bar{q}^k_J\bar{\sigma}_{\mu\nu}\bar{u}^c_J)$
$[O_{\ell u}]_{IIJJ}=(\bar{\ell}_I\bar{\sigma}_\mu\ell_I)(u^c_J\sigma^\mu\bar{u}^c_J)$	$[O_{\ell edq}]_{IIJJ}=(\bar{\ell}^j_I\bar{e}^c_I)(d^c_J q^j_J)$
$[O_{\ell d}]_{IIJJ}=(\bar{\ell}_I\bar{\sigma}_\mu\ell_I)(d^c_J\sigma^\mu\bar{d}^c_J)$	
$[O_{eq}]_{IIJJ}=(e^c_I\sigma_\mu\bar{e}^c_I)(\bar{q}_J\bar{\sigma}^\mu q_J)$	
$[O_{eu}]_{IIJJ}=(e^c_I\sigma_\mu\bar{e}^c_I)(u^c_J\sigma^\mu\bar{u}^c_J)$	
$[O_{ed}]_{IIJJ}=(e^c_I\sigma_\mu\bar{e}^c_I)(d^c_J\sigma^\mu\bar{d}^c_J)$	

Table 1. Flavor-conserving *2-lepton-2-quark* operators in the SMEFT Lagrangian of eq. (2.1).

One flavor $(I=1,2,3)$	Two flavors $(I<J=1,2,3)$
$[O_{\ell\ell}]_{IIII}=\frac{1}{2}(\bar{\ell}_I\bar{\sigma}_\mu\ell_I)(\bar{\ell}_I\bar{\sigma}^\mu\ell_I)$	$[O_{\ell\ell}]_{IIJJ}=(\bar{\ell}_I\bar{\sigma}_\mu\ell_I)(\bar{\ell}_J\bar{\sigma}^\mu\ell_J)$
	$[O_{\ell\ell}]_{IJJI}=(\bar{\ell}_I\bar{\sigma}_\mu\ell_J)(\bar{\ell}_J\bar{\sigma}^\mu\ell_I)$
$[O_{\ell e}]_{IIII}=(\bar{\ell}_I\bar{\sigma}_\mu\ell_I)(e^c_I\sigma^\mu\bar{e}^c_I)$	$[O_{\ell e}]_{IIJJ}=(\bar{\ell}_I\bar{\sigma}_\mu\ell_I)(e^c_J\sigma^\mu\bar{e}^c_J)$
	$[O_{\ell e}]_{JJII}=(\bar{\ell}_J\bar{\sigma}_\mu\ell_J)(e^c_I\sigma^\mu\bar{e}^c_I)$
	$[O_{\ell e}]_{IJJI}=(\bar{\ell}_I\bar{\sigma}_\mu\ell_J)(e^c_J\sigma^\mu\bar{e}^c_I)$
$[O_{ee}]_{IIII}=\frac{1}{2}(e^c_I\sigma_\mu\bar{e}^c_I)(e^c_I\sigma^\mu\bar{e}^c_I)$	$[O_{ee}]_{IIJJ}=(e^c_I\sigma_\mu\bar{e}^c_I)(e^c_J\sigma^\mu\bar{e}^c_J)$

Table 2. Flavor-conserving *4-lepton operators* in the SMEFT Lagrangian of eq. (2.1).

The vertex corrections correspond to 24 linear combinations of dimension-6 Wilson coefficients, 3 of which are complex (those entering δg_R^{Wq}). We consider only CP-conserving observables, thus the imaginary part enters at the quadratic level and is neglected. To simplify the notation we will omit Re in front of complex Wilson coefficients.

In this paper we will also discuss constraints on flavor-diagonal 4-fermion operators in the SMEFT Lagrangian of eq. (2.1). We work with the same set of 4-fermion operators as in ref. [57] and employ a similar notation.[3] The main focus is on the flavor-conserving 2-lepton-2-quark dimension-6 operators (LLQQ) summarized in table 1, and defined in the flavor basis where the up-quark Yukawa matrices are diagonal. Overall, there are $10\times3\times3=90$ such operators, of which 27 (the chirality-violating ones) are complex. In the latter case the corresponding Wilson coefficient is complex, and the Hermitian conjugate operator is included in eq. (2.1). For the sake of combining our results with those of ref. [43], we also list in table 2 the 27 flavor-conserving 4-lepton operators (LLLL), 3 of which are complex ($[O_{\ell e}]_{IJJI}$).

[3] One difference is that for operators with the $SU(2)_L$ singlet contraction of fermionic currents we omit the superscript $^{(1)}$. We also rename $\mathcal{Q}_{qe}\to\mathcal{Q}_{eq}$ so that the first (last) two flavor indices of all LLQQ operators correspond to the leptons (quarks).

JHEP08(2017)123

All in all, our analysis eyes 147 linear combinations of dimension-6 operators displayed in eq. (2.3), table 1, and table 2. The observables discussed in this paper will not depend on all of them, and thus we will be able to constrain only a limited number of the combinations. In particular the operators involving the 3rd generation fermions are currently, with a few exceptions, poorly constrained by experiment. Nevertheless, the constraints we derive are robust, in the sense that they do not involve any strong assumptions about the unconstrained operators, other than the validity of the SMEFT description at the weak scale. We assume that our results are not invalidated by $\mathcal{O}\left(\frac{1}{16\pi^2\Lambda^2}\right)$ corrections, which arise at one loop in the SMEFT and inevitably introduce dependence of our observables on other D=6 Wilson coefficients. We will also treat V as the unit matrix when it multiplies dimension-6 Wilson coefficients. This ignores all contributions to observables where the Wilson coefficients are multiplied by an off-diagonal CKM element.[4]

Last, we will also particularize our results to more restrictive scenarios, such as the so-called flavor-universal SMEFT, where dimension-6 operators respect the $\mathrm{U}(3)^5$ global flavor symmetry acting in the generation space on the SM fermion fields q, ℓ, u^c, d^c, e^c.

2.2 Weak interactions below the weak scale

Precision experiments with a characteristic momentum transfer $Q \ll m_Z$ can be conveniently described using the low-energy effective theory where the SM W and Z bosons are integrated out. In this framework, weak interactions between quark and leptons are mediated by a set of 4-fermion operators. Within the SM, these operators effectively appear due to the exchange of W and Z bosons at tree level or in loops, and their coefficients can be calculated by the standard matching procedure. Once the SM is extended by dimension-6 operators, these coefficients may be modified, either due to modified propagators and couplings of W and Z, or due to the presence of contact 4-fermion operators in the SMEFT Lagrangian.

Below we define the low-energy operators that are relevant for the precision measurements we include in our analysis. We follow the PDG notation [61] (section 10), and we present the matching between the coefficients of the low-energy operators and the parameters of the SMEFT.

2.2.1 Charged-current (CC) interactions: $qq'\ell\nu$

The low-energy CC interactions of leptons with the 1st generation quarks are described by the effective 4-fermion operators:

$$\begin{aligned}\mathcal{L}_{\rm eff} \supset -\frac{2\tilde{V}_{ud}}{v^2}\Bigg[&\left(1+\bar{\epsilon}_L^{de_J}\right)(\bar{e}_J\bar{\sigma}_\mu\nu_J)(\bar{u}\bar{\sigma}^\mu d)+\epsilon_R^{de}(\bar{e}_J\bar{\sigma}_\mu\nu_J)(u^c\sigma^\mu\bar{d}^c) \\ &+\frac{\epsilon_S^{de_J}+\epsilon_P^{de_J}}{2}(e_J^c\nu_J)(u^c d)+\frac{\epsilon_S^{de_J}-\epsilon_P^{de_J}}{2}(e_J^c\nu_J)(\bar{u}\bar{d}^c)+\epsilon_T^{de_J}(e_J^c\sigma_{\mu\nu}\nu_J)(u^c\sigma_{\mu\nu}d)+\text{h.c.}\Bigg].\end{aligned} \tag{2.4}$$

[4] Such an approach is not completely satisfactory, since the Cabibbo angle is not small enough to always justify neglecting it. However, including the new physics contributions suppressed by the Cabibbo angle would require extending our analysis to include flavor-violating observables, which we leave for future publications. On the other hand, one naively expects the neglected operators to be severely constrained by other observables where the CKM suppression is not present, which would justify our approximation.

To make contact with low-energy flavor observables, we defined the rescaled CKM matrix element $\tilde{V}_{ud}$ [55]. It is distinct from the *actual* V_{ud}, i.e., the 11 element of the unitary matrix V that appears in the Lagrangian after rotating quarks to the mass eigenstate basis. The two are related by $V_{ud} = \tilde{V}_{ud}(1+\delta V_{ud})$ where δV_{ud} is chosen such as to impose the relation $\bar{\epsilon}_L^{de} = -\epsilon_R^{de}$ in eq. (2.4).[5]

Let us note that in general $\tilde{V}_{ud}$ is also different from the phenomenological value obtained within the SM, which we will denote by $V_{ud}^{\rm PDG}$. Currently this value comes from superallowed nuclear beta decays [62] that depend on the vector couplings via the combination $\bar{\epsilon}_L^{de} + \epsilon_R^{de}$. By setting $\bar{\epsilon}_L^{de} = -\epsilon_R^{de}$, this nonstandard effect has been conveniently absorbed into the definition of $\tilde{V}_{ud}$. However, the relevant transitions also depend, each in a different way, on the scalar coefficient ϵ_S^{de}. Thus $\tilde{V}_{ud}$ and $V_{ud}^{\rm PDG}$ only coincide if ϵ_S^{de} vanishes, whereas in general it is not possible to redefine away all new physics contributions through $\tilde{V}_{ud}$. For this reason we treat $\tilde{V}_{ud}$ as a free parameter that is fit together with the EFT Wilson coefficients [55]. In principle the difference between $\tilde{V}_{ud}$ and $V_{ud}^{\rm PDG}$ must be taken into account every time the latter is used to calculate any given SM prediction. In practice, this effect will be negligible in most cases, given the strong constraints on ϵ_S^{de} from the same nuclear decay data, cf. eq. (3.17).

At tree level, the low-energy parameters are related to the SMEFT parameters as

$$\begin{aligned}
\delta V_{ud} &= -\delta g_L^{Wq_1} - \delta g_R^{Wq_1} + \delta g_L^{W\mu} - \frac{1}{2}[c_{\ell\ell}]_{1221} + [c_{lq}^{(3)}]_{1111}, \\
\epsilon_R^{de} = -\bar{\epsilon}_L^{de} &= \delta g_R^{Wq_1}, \\
\bar{\epsilon}_L^{d\mu} &= -\delta g_R^{Wq_1} + \delta g_L^{W\mu} - \delta g_L^{We} + [c_{lq}^{(3)}]_{1111} - [c_{lq}^{(3)}]_{2211}, \\
\epsilon_S^{de_J} &= -\frac{1}{2}\left([c_{lequ}]^*_{JJ11} + [c_{ledq}]^*_{JJ11}\right), \\
\epsilon_P^{de_J} &= -\frac{1}{2}\left([c_{lequ}]^*_{JJ11} - [c_{ledq}]^*_{JJ11}\right), \\
\epsilon_T^{de_J} &= -\frac{1}{2}[c_{lequ}^{(3)}]^*_{JJ11}\,,
\end{aligned} \tag{2.5}$$

As indicated earlier, at $\mathcal{O}(\Lambda^{-2})$ we treat the CKM matrix as the unit matrix. In this limit, the effective parameters in eq. (2.4) depend only on flavor-diagonal vertex corrections and 4-fermion operators. See appendix B for more general expressions where non-diagonal elements of V are retained. Note also that the rescaled CKM matrix is no longer unitary. In particular we have $|\tilde{V}_{ud}|^2 + |V_{us}|^2 \approx 1 + \Delta_{\rm CKM}$, where

$$\Delta_{\rm CKM} = -2\delta V_{ud} = 2\delta g_L^{Wq_1} + 2\delta g_R^{Wq_1} - 2\delta g_L^{W\mu} + [c_{\ell\ell}]_{1221} - 2[c_{lq}^{(3)}]_{1111}. \tag{2.6}$$

Although the extraction of the V_{us} element is also affected by dimension-6 operators, their contribution to this unitarity test is suppressed by V_{us} and therefore it can be neglected in our approximation ($V \approx 1$ at order Λ^{-2}). See eq. (B.5) for the complete expression.

[5]The bar in the $\bar{\epsilon}_L^{de_J}$ coefficient reminds the reader that this coefficient is not the usual $\epsilon_L^{de_J}$ (see e.g. ref. [55]) where the shift of new physics effects into $\tilde{V}_{ud}$ is not carried out. These two are trivially related by $V_{ud}\,(1+\epsilon_L^{de_J}) = \tilde{V}_{ud}\,(1+\bar{\epsilon}_L^{de_J})$.

JHEP08(2017)123

2.2.2 Neutral-current (NC) neutrino interactions: $qq\nu\nu$

The low-energy NC neutrino interactions with light quarks are described by the effective 4-fermion operators:

$$\mathcal{L}_{\rm eff} \supset -\frac{2}{v^2}(\bar{\nu}_J\bar{\sigma}^\mu\nu_J)\left(g_{LL}^{\nu_J q}\bar{q}\bar{\sigma}_\mu q + g_{LR}^{\nu_J q} q^c\sigma_\mu\bar{q}^c\right). \tag{2.7}$$

At tree level, the low-energy parameters are related to the SMEFT parameters as

$$\begin{aligned}
g_{LL}^{\nu_J u} &= \frac{1}{2} - \frac{2s_\theta^2}{3} + \delta g_L^{Zu} + \left(1 - \frac{4s_\theta^2}{3}\right)\delta g_L^{Z\nu_J} - \frac{1}{2}([c_{lq}]_{JJ11} + [c_{lq}^{(3)}]_{JJ11}),\\
g_{LR}^{\nu_J u} &= -\frac{2s_\theta^2}{3} + \delta g_R^{Zu} - \frac{4s_\theta^2}{3}\delta g_L^{Z\nu_J} - \frac{1}{2}[c_{lu}]_{JJ11},\\
g_{LL}^{\nu_J d} &= -\frac{1}{2} + \frac{s_\theta^2}{3} + \delta g_L^{Zd} - \left(1 - \frac{2s_\theta^2}{3}\right)\delta g_L^{Z\nu_J} - \frac{1}{2}([c_{lq}]_{JJ11} - [c_{lq}^{(3)}]_{JJ11}),\\
g_{LR}^{\nu_J d} &= \frac{s_\theta^2}{3} + \delta g_R^{Zd} + \frac{2s_\theta^2}{3}\delta g_L^{Z\nu_J} - \frac{1}{2}[c_{ld}]_{JJ11}.
\end{aligned} \tag{2.8}$$

The experiments probing these couplings usually normalize the NC cross section using its CC counterpart. Thus, it is convenient to define the following combinations of effective couplings:

$$(g_{L/R}^{\nu_J})^2 \equiv \frac{(g_{LL/LR}^{\nu_J u})^2 + (g_{LL/LR}^{\nu_J d})^2}{\left(1 + \bar{\epsilon}_L^{de_J}\right)^2}, \qquad \theta_{L/R}^{\nu_J} \equiv \arctan\left(\frac{g_{LL/LR}^{\nu_J u}}{g_{LL/LR}^{\nu_J d}}\right), \tag{2.9}$$

where we took into account that SMEFT dimension-6 operators modify in general both NC and CC processes. Let us notice that additional (linear) effects in the normalizing CC process due to ϵ_R^{de} and $\epsilon_{S,P,T}^{de_J}$ can be neglected because they are suppressed by the ratio $m_u m_d/E^2$ and m_{e_J}/E respectively. The effect due to the possible difference between $\tilde{V}_{ud}$ and $V_{ud}^{\rm PDG}$ can also be safely neglected here, given the limited precision of the neutrino scattering experiments included in our fit. Last, the same holds for the δV_{ud} contribution that appears if the unitarity of the CKM matrix is used in the SM determination.

2.2.3 Neutral-current charged-lepton interactions: $qq\ell\ell$

We parametrize[6] the 4-fermion operators with 2 charged leptons and 2 light quarks as

$$\begin{aligned}
\mathcal{L} \supset\ & \frac{1}{2v^2}\left[g_{AV}^{e_J q}(\bar{e}_J\gamma_\mu\gamma_5 e_J)(\bar{q}\gamma_\mu q) + g_{VA}^{e_J q}(\bar{e}_J\gamma_\mu e_J)(\bar{q}\gamma_\mu\gamma_5 q)\right]\\
&+\frac{1}{2v^2}\left[g_{VV}^{e_J q}(\bar{e}_J\gamma_\mu e_J)(\bar{q}\gamma_\mu q) + g_{AA}^{e_J q}(\bar{e}_J\gamma_\mu\gamma_5 e_J)(\bar{q}\gamma_\mu\gamma_5 q)\right],
\end{aligned} \tag{2.10}$$

[6] For the parity-violating electron couplings, another frequently used notation is $g_{AV}^{eq} \equiv C_{1q}$, $g_{VA}^{eq} \equiv C_{2q}$.

JHEP08(2017)123

where we momentarily switch to the Dirac notation with $\gamma_5\psi_L = -\psi_L$, $\gamma_5\psi_R = +\psi_R$. At tree level, the parameters $g^{e_i q}_{XY}$ are related to the SMEFT parameters as

$$\begin{aligned}
g^{e_J u}_{AV} &= -\frac{1}{2} + \frac{4}{3}s^2_\theta - \left(\delta g^{Zu}_L + \delta g^{Zu}_R\right) + \frac{3-8s^2_\theta}{3}\left(\delta g^{Ze_J}_L - \delta g^{Ze_J}_R\right) \\
&\quad + \frac{1}{2}\left[c^{(3)}_{lq} - c_{lq} - c_{lu} + c_{eq} + c_{eu}\right]_{JJ11}, \\
g^{e_J d}_{AV} &= \frac{1}{2} - \frac{2}{3}s^2_\theta - \left(\delta g^{Zd}_L + \delta g^{Zd}_R\right) - \frac{3-4s^2_\theta}{3}\left(\delta g^{Ze_J}_L - \delta g^{Ze_J}_R\right) \\
&\quad + \frac{1}{2}\left[-c^{(3)}_{lq} - c_{lq} - c_{ld} + c_{eq} + c_{ed}\right]_{JJ11}, \\
g^{e_J u}_{VA} &= -\frac{1}{2} + 2s^2_\theta - \left(1-4s^2_\theta\right)\left(\delta g^{Zu}_L - \delta g^{Zu}_R\right) + \left(\delta g^{Ze_J}_L + \delta g^{Ze_J}_R\right) \\
&\quad + \frac{1}{2}\left[c^{(3)}_{lq} - c_{lq} + c_{lu} - c_{eq} + c_{eu}\right]_{JJ11}, \\
g^{e_J d}_{VA} &= \frac{1}{2} - 2s^2_\theta - \left(1-4s^2_\theta\right)\left(\delta g^{Zd}_L - \delta g^{Zd}_R\right) - \left(\delta g^{Ze_J}_L + \delta g^{Ze_J}_R\right) \\
&\quad + \frac{1}{2}\left[-c^{(3)}_{lq} - c_{lq} + c_{ld} - c_{eq} + c_{ed}\right]_{JJ11}, \\
g^{e_J u}_{AA} &= \frac{1}{2} + \delta g^{Zu}_L - \delta g^{Zu}_R - \delta g^{Ze_J}_L + \delta g^{Ze_J}_R + \frac{1}{2}\left[-c^{(3)}_{lq} + c_{lq} - c_{lu} - c_{eq} + c_{eu}\right]_{JJ11}, \\
g^{e_J d}_{AA} &= -\frac{1}{2} + \delta g^{Zd}_L - \delta g^{Zd}_R + \delta g^{Ze_J}_L - \delta g^{Ze_J}_R + \frac{1}{2}\left[c^{(3)}_{lq} + c_{lq} - c_{ld} - c_{eq} + c_{ed}\right]_{JJ11}.
\end{aligned} \tag{2.11}$$

We do not display the expressions for $g^{e_i q}_{VV}$ here because they will not be needed in the following.

2.2.4 Four-lepton interactions: $\ell\ell\ell\ell$ and $\ell\ell\nu\nu$

Although the main focus of this work are the LLQQ operators, in this section we provide a few expressions concerning 4-lepton operators that will be needed in our subsequent phenomenological analysis. First, we parametrize the ν-e interaction in the effective theory below the weak scale as:

$$\mathcal{L} \supset -\frac{1}{v^2}(\bar\nu_J\bar\sigma_\mu\nu_J)\left[\left(g^{\nu_J e_I}_{LV} + g^{\nu_J e_I}_{LA}\right)(\bar e_I\bar\sigma_\mu e_I) + \left(g^{\nu_J e_I}_{LV} - g^{\nu_J e_I}_{LA}\right)(e^c_I\sigma_\mu\bar e^c_I)\right]. \tag{2.12}$$

Matching to the SMEFT one finds

$$\begin{aligned}
g^{\nu_J e_I}_{LV} &= \delta_{IJ} - \frac{1}{2} + 2s^2_\theta + \delta_{IJ}\left(2\delta g^{We_I}_L - \delta g^{We}_L - \delta g^{W\mu}_L + \frac{1}{2}[c_{\ell\ell}]_{1221}\right) \\
&\quad - \left(1-4s^2_\theta\right)\delta g^{Z\nu_J}_L + \delta g^{Ze_I}_L + \delta g^{Ze_I}_R - \frac{1}{2}\left(x_{IJ} + [c_{\ell e}]_{JJII}\right), \\
g^{\nu_J e_I}_{LA} &= \delta_{IJ} - \frac{1}{2} + \delta_{IJ}\left(2\delta g^{We_I}_L - \delta g^{We}_L - \delta g^{W\mu}_L + \frac{1}{2}[c_{\ell\ell}]_{1221}\right) \\
&\quad - \delta g^{Z\nu_J}_L + \delta g^{Ze_I}_L - \delta g^{Ze_I}_R - \frac{1}{2}\left(x_{IJ} - [c_{\ell e}]_{JJII}\right),
\end{aligned} \tag{2.13}$$

where $x_{IJ} = [c_{\ell\ell}]_{IIJJ}$ if $I \le J$ or $x_{IJ} = [c_{\ell\ell}]_{JJII}$ otherwise.

Last, we parameterize the parity-violating self-interaction of electrons in the effective theory below the weak scale as

$$\mathcal{L} \supset \frac{1}{2v^2} g_{AV}^{ee} \left[-(\bar{e}\bar{\sigma}_\mu e)(\bar{e}\bar{\sigma}_\mu e) + (e^c \sigma_\mu \bar{e}^c)(e^c \sigma_\mu \bar{e}^c) \right] , \tag{2.14}$$

with the following SMEFT expression

$$g_{AV}^{ee} = \frac{1}{2} - 2s_\theta^2 - 2\left(1 - 2s_\theta^2\right)\delta g_L^{Ze} - 4s_\theta^2 \delta g_R^{Ze} - \frac{1}{2}[c_{\ell\ell}]_{1111} + \frac{1}{2}[c_{ee}]_{1111} \,. \tag{2.15}$$

2.3 Renormalization and scale running of the Wilson coefficients

In general the Wilson coefficients display renormalization-scale dependence that is to be canceled in the observables by the opposite dependence in the quantum corrections to the matrix elements. Let us first discuss the QCD running, which can have a numerically significant impact due to the magnitude of the strong coupling constant α_s. This effect is further enhanced by the large separation of scales of the experiments discussed in this work (from low-energy precision measurements to LHC collisions). Among the coefficients involved in our analysis, only the three chirality-violating ones, $c_{lequ}, c_{ledq}, c_{lequ}^{(3)}$ (i.e. $\epsilon_{S,P,T}^{d\ell}$ in the low-energy EFT), present a non-zero 1-loop QCD anomalous dimension, namely [63]

$$\frac{d\,\vec{x}(\mu)}{d\log\mu} = \frac{\alpha_s(\mu)}{2\pi} \begin{pmatrix} -4 & 0 & 0 \\ 0 & -4 & 0 \\ 0 & 0 & 4/3 \end{pmatrix} \vec{x}(\mu), \tag{2.16}$$

where $\vec{x}$ refers to the SMEFT coefficients $\vec{c} = (c_{ledq},\, c_{lequ},\, c_{lequ}^{(3)})$ if the scale μ is above the weak scale or to the low-energy EFT coefficients $\vec{\epsilon} = (\epsilon_S^{d\ell},\, \epsilon_P^{d\ell},\, \epsilon_T^{d\ell})$ below it. We find that higher-loop QCD corrections to the running are numerically significant, and we include them in our analysis.[7]

On the other hand we neglect in this work the electromagnetic/weak running of the SMEFT Wilson coefficients, which is expected to have a much smaller numerical importance simply due to the smallness of the corresponding coupling constants. There is however one exception to this, namely the chirality-violating operators discussed above, for two reasons: (i) contrary to the QCD running, the QED/weak running involves mixing between these operators; (ii) pion decay makes possible to set bounds of order 10^{-7} on the pseudoscalar coupling $\epsilon_P^{d\ell}(\mu_{\rm low})$, which can give important bounds on scalar and tensor via mixing despite the smallness of α_{em}. In order to take into account this effect, eq. (2.16) has to be replaced by

$$\frac{d\,\vec{x}(\mu)}{d\log\mu} = \left(\frac{\alpha_{em}(\mu)}{2\pi}\gamma_x + \frac{\alpha_s(\mu)}{2\pi}\gamma_s \right) \vec{x}(\mu)\,, \tag{2.17}$$

[7] We use the 3-loop QCD anomalous dimension [64], and we include the threshold corrections at m_b and m_t extracted from refs. [65] and [66] for scalar and tensor operators respectively. See ref. [67] for further details.

JHEP08(2017)123

where we will use the 1-loop QED (electroweak) anomalous dimension, $\gamma_x = \gamma_{em(w)}$, to evolve the coefficients $\vec{\epsilon}$ ($\vec{c}$) below (above) the weak scale [67–70]:

$$\gamma_{\rm em} = \begin{pmatrix} \frac{2}{3} & 0 & 4 \\ 0 & \frac{2}{3} & 4 \\ \frac{1}{24} & \frac{1}{24} & -\frac{20}{9} \end{pmatrix}, \qquad \gamma_{\rm w} = \begin{pmatrix} -\frac{4}{3c_\theta^2} & 0 & 0 \\ 0 & -\frac{11}{6c_\theta^2} & \frac{15}{c_\theta^2} + \frac{9}{s_\theta^2} \\ 0 & \frac{5}{16c_\theta^2} + \frac{3}{16s_\theta^2} & \frac{1}{9c_\theta^2} - \frac{3}{2s_\theta^2} \end{pmatrix}, \tag{2.18}$$

where we neglect terms suppressed by Yukawa couplings [70, 71]. Integrating numerically the coupled differential renormalization group equations we find

$$\begin{pmatrix} \epsilon_S^{d\ell} \\ \epsilon_P^{d\ell} \\ \epsilon_T^{d\ell} \end{pmatrix}_{(\mu = m_Z)} = \begin{pmatrix} 0.58 & 1.42\times10^{-6} & 0.017 \\ 1.42\times10^{-6} & 0.58 & 0.017 \\ 1.53\times10^{-4} & 1.53\times10^{-4} & 1.21 \end{pmatrix} \begin{pmatrix} \epsilon_S^{d\ell} \\ \epsilon_P^{d\ell} \\ \epsilon_T^{d\ell} \end{pmatrix}_{(\mu = 2\,\mathrm{GeV})}, \tag{2.19}$$

$$\begin{pmatrix} c_{ledq} \\ c_{lequ} \\ c_{lequ}^{(3)} \end{pmatrix}_{(\mu = 1\,\mathrm{TeV})} = \begin{pmatrix} 0.84 & 0 & 0 \\ 0 & 0.84 & 0.16 \\ 0 & 3.3\times10^{-3} & 1.04 \end{pmatrix} \begin{pmatrix} c_{ledq} \\ c_{lequ} \\ c_{lequ}^{(3)} \end{pmatrix}_{(\mu = m_Z)}. \tag{2.20}$$

These results use the QCD beta function and anomalous dimensions up to 3 loops, and we included the bottom and top quark thresholds effects, see ref. [67] for details. The diagonal entries would change by $\sim 12\%$ if just 1-loop QCD running were included, while two-loop results differ by only $\sim 1.5\%$. In our subsequent analysis we will use the numerical results in eq. (2.19) and eq. (2.20).

3 Low-energy experiments

3.1 Neutrino scattering

Neutrino scattering experiments measure the ratio of neutral- and charged-current neutrino or anti-neutrino scattering cross sections on nuclei:

$$R_{\nu_i} = \frac{\sigma(\nu_i N \to \nu X)}{\sigma(\nu_i N \to \ell_i^- X)}, \qquad R_{\bar\nu_i} = \frac{\sigma(\bar\nu_i N \to \bar\nu X)}{\sigma(\bar\nu_i N \to \ell_i^+ X)}. \tag{3.1}$$

At leading order and for isoscalar nucleus targets (equal number of protons and neutrons) one has the so-called Llewellyn-Smith relations [72]:

$$R_{\nu_i} = (g_L^{\nu_i})^2 + r(g_R^{\nu_i})^2, \qquad R_{\bar\nu_i} = (g_L^{\nu_i})^2 + r^{-1}(g_R^{\nu_i})^2, \tag{3.2}$$

where r is the ratio of ν to $\bar\nu$ charged-current cross sections on N that can be measured separately, and the effective couplings $g_{L/R}^{\nu_i}$ are defined in eq. (2.9). In some experiments the beam is a mixture of neutrinos and anti-neutrinos, and the following ratio is measured

$$R_{\nu_i\bar\nu_i} = \frac{\sigma(\nu_i N \to \nu X) + \sigma(\bar\nu_i N \to \bar\nu X)}{\sigma(\nu_i N \to \ell_i^- X) + \sigma(\bar\nu_i N \to \ell_i^+ X)} = (g_L^{\nu_i})^2 + (g_R^{\nu_i})^2. \tag{3.3}$$

Experiment	Observable	Experimental value	SM value	Ref.
CHARM ($r = 0.456$)	R_{ν_μ}	0.3093 ± 0.0031	0.3156	[74]
	$R_{\bar{\nu}_\mu}$	0.390 ± 0.014	0.370	[74]
CDHS ($r = 0.393$)	R_{ν_μ}	0.3072 ± 0.0033	0.3091	[75]
	$R_{\bar{\nu}_\mu}$	0.382 ± 0.016	0.380	[75]
CCFR	κ	0.5820 ± 0.0041	0.5830	[76]

Table 3. The results of muon-neutrino scattering experiments most relevant for constraining dimension-6 operators in the SMEFT. The SM values of R_{ν_μ} and κ include subleading corrections [77], whereas those of $R_{\bar{\nu}_\mu}$ are the tree-level values, which should be sufficient taking into account the larger experimental errors.

ν_e data. The CHARM experiment [73] made a measurement of electron-neutrino scattering cross sections:

$$R_{\nu_e \bar{\nu}_e} = 0.406^{+0.145}_{-0.135}, \tag{3.4}$$

where the uncertainties quoted here and everywhere else in this work are 1-sigma (68%C.L.) errors. To avoid dealing with asymmetric errors we approximate it as $R_{\nu_e \bar{\nu}_e} = 0.41 \pm 0.14$, and we estimate the SM expectation as $R^{\rm SM}_{\nu_e \bar{\nu}_e} = 0.33$. To our knowledge, this weakly constraining measurement is currently the best probe of the electron-neutrino neutral-current interactions.

ν_μ data. For the muon-neutrino scattering the experimental data are much more abundant and precise. We summarize the relevant results in table 3. The observable κ measured in CCFR probes the following combinations of couplings [76]:

$$\kappa = 1.7897 (g_L^{\nu_\mu})^2 + 1.1479 (g_R^{\nu_\mu})^2 - \frac{0.0916 \left[(g_{LL}^{\nu_\mu u})^2 - (g_{LL}^{\nu_\mu d})^2 \right] + 0.0782 \left[(g_{LR}^{\nu_\mu u})^2 - (g_{LR}^{\nu_\mu d})^2 \right]}{(1 + \bar{\epsilon}_L^{d\mu})^2} . \tag{3.5}$$

The additional small dependence on the difference of the up and down effective couplings appears when one takes into account that the target (in this case iron) is not exactly isoscalar. For the reasons explained in ref. [61], in our fits we do not take into account the results of the NuTeV experiment.

The observables in table 3 constrain 3 independent combinations of the SMEFT coefficients. Rather then combining these results ourselves, we use the PDG combination [61] that also uses additional experimental input [78] from neutrino induced coherent neutral pion production from nuclei [79, 80] and elastic neutrino-proton scattering [81, 82]. Although their precision is quite limited, their inclusion allows one to constrain the 4 muon-neutrino effective couplings to quarks [77]. The results of the latest PDG fit are [61]:

$$\begin{aligned} (g_L^{\nu_\mu})^2 &= 0.3005 \pm 0.0028, & (g_R^{\nu_\mu})^2 &= 0.0329 \pm 0.0030, \\ \theta_L^{\nu_\mu} &= 2.50 \pm 0.035, & \theta_R^{\nu_\mu} &= 4.56^{+0.42}_{-0.27}. \end{aligned} \tag{3.6}$$

The correlations are quoted to be small in ref. [61] and in the following we neglect them. We symmetrize the uncertainty on θ_R taking the larger of the errors, so as to avoid dealing with

JHEP08(2017)123

asymmetric errors. The corresponding SM predictions are given in table 4. To evaluate their dimension-6 EFT corrections in eq. (2.8) we will use $s_\theta^2 = 0.23865$, which is the central value in the $\overline{MS}$ scheme at low energies [61]. We neglect the error of the SM predictions when it is much smaller than the experimental uncertainties; otherwise we combine it in quadrature.

We note that LLQQ (and 4-lepton) operators can also be probed via matter effects in neutrino oscillations, see e.g. [83, 84]. However, the resulting constraints are not available in the model-independent form where all 4-fermion operators can be present simultaneously. Moreover, neutrino oscillations probe linear combinations of lepton-flavor-diagonal operators and of the off-diagonal ones (which we marginalize over). For these reasons, we do not include the oscillation constraints in this paper.

3.2 Parity violation in atoms and in scattering

Atomic parity violation (APV) and parity-violating electron scattering experiments access the parity-violating effective couplings of electrons to quarks g_{AV}^{eq} and g_{VA}^{eq}. In particular, APV and elastic scattering on a target with Z protons and N neutrons probe its so-called weak charge Q_W that is given by

$$Q_W(Z, N) = -2\left((2Z + N)g_{AV}^{eu} + (Z + 2N)g_{AV}^{ed}\right), \tag{3.7}$$

up to small radiative corrections [61, 77]. The most precise determination is performed in ^{133}Cs, where $Q_W(55, 133 - 55) \approx -376g_{AV}^{eu} - 422g_{AV}^{ed}$. Taking into account recent reanalyses [85] of the measured parity-violating transitions in cesium atoms [86], the latest edition of the PDG Review [61] quotes

$$Q_W^{\mathrm{Cs}} = -72.62 \pm 0.43, \tag{3.8}$$

where the SM prediction is $Q_{W,\mathrm{SM}}^{\mathrm{Cs}} = -73.25 \pm 0.02$ [61]. Other APV measurements, e.g. with thallium atoms, probe slightly different combinations of the g_{AV}^{eq} couplings, although with larger errors.

Instead, a very different linear combination of g_{AV}^{eu} and g_{AV}^{ed} is precisely probed by measurements of the weak charge of the proton, $Q_W^{\mathrm{p}} = Q_W(1, 0)$, in scattering experiments with low-energy polarized electrons. The QWEAK experiment [87] finds

$$Q_W^{\mathrm{p}} = 0.064 \pm 0.012, \tag{3.9}$$

where the SM prediction is $Q_{W,\mathrm{SM}}^{\mathrm{p}} = 0.0708 \pm 0.0003$ [61].

In order to access the effective couplings g_{VA}^{eq} one needs to resort to deep-inelastic scattering of polarized electrons. Currently, the most precise of these is the PVDIS experiment [88] that studies electron scattering on deuterium targets. The experiment is sensitive to the following two linear combinations of effective couplings [88]:

$$\begin{aligned} A_1^{\mathrm{PVDIS}} &= 1.156 \times 10^{-4}\left(2g_{AV}^{eu} - g_{AV}^{ed} + 0.348(2g_{VA}^{eu} - g_{VA}^{ed})\right) \\ A_2^{\mathrm{PVDIS}} &= 2.022 \times 10^{-4}\left(2g_{AV}^{eu} - g_{AV}^{ed} + 0.594(2g_{VA}^{eu} - g_{VA}^{ed})\right). \end{aligned} \tag{3.10}$$

The measured values are [88]

$$A_1^{\rm PVDIS} = (-91.1 \pm 4.3) \times 10^{-6}, \qquad A_2^{\rm PVDIS} = (-160.8 \pm 7.1) \times 10^{-6}, \tag{3.11}$$

where the SM predictions are $A_{1,SM}^{\rm PVDIS} = -(87.7 \pm 0.7) \times 10^{-6}$, $A_{2,SM}^{\rm PVDIS} = -(158.9 \pm 1.0) \times 10^{-6}$ [88].

The PDG combines the results of APV, QWEAK, and PVDIS experiments into correlated constraints on 3 linear combinations of g_{VA}^{eq} and g_{AV}^{eq} [61]:

$$\begin{pmatrix} g_{AV}^{eu} + 2g_{AV}^{ed} \\ 2g_{AV}^{eu} - g_{AV}^{ed} \\ 2g_{VA}^{eu} - g_{VA}^{ed} \end{pmatrix} = \begin{pmatrix} 0.489 \pm 0.005 \\ -0.708 \pm 0.016 \\ -0.144 \pm 0.068 \end{pmatrix}, \qquad \rho = \begin{pmatrix} & -0.94 & 0.42 \\ & & -0.45 \\ & & \end{pmatrix}. \tag{3.12}$$

To disentangle g_{VA}^{eu} and g_{VA}^{ed} one needs more input from earlier (less precise) measurements of parity-violating scattering. We include two results provided by the SAMPLE collaboration [89]:

$$g_{VA}^{eu} - g_{VA}^{ed} = -0.042 \pm 0.057, \qquad g_{VA}^{eu} - g_{VA}^{ed} = -0.12 \pm 0.074, \tag{3.13}$$

from the scattering of polarized electrons on deuterons in the quasi-elastic kinematic regime at two different values of the beam energy. Combining the likelihood obtained from eq. (3.12) with the SAMPLE results we find the following constraints:

$$\begin{pmatrix} \delta g_{AV}^{eu} \\ \delta g_{AV}^{ed} \\ \delta g_{VA}^{eu} \\ \delta g_{VA}^{ed} \end{pmatrix} = \begin{pmatrix} 0.0033 \pm 0.0054 \\ -0.0047 \pm 0.0051 \\ -0.041 \pm 0.081 \\ -0.032 \pm 0.11 \end{pmatrix}, \qquad \rho = \begin{pmatrix} & -0.98 & -0.37 & -0.27 \\ & & 0.37 & 0.27 \\ & & & 0.94 \\ & & & \end{pmatrix}. \tag{3.14}$$

Here δg_{XY}^{eq} are shifts of the effective couplings away from the SM values, whose dependence on the dimension-6 Wilson coefficients can be read off from eq. (2.11).

There are also results concerning effective muon couplings to quarks. A CERN SPS experiment [90] measured a DIS asymmetry using polarized muon and anti-muon scattering on an isoscalar carbon target. The results can be recast as the measurement of the observable $b_{\rm SPS}$ defined as

$$b_{\rm SPS} = \frac{3}{5e^2v^2}\left(g_{AA}^{\mu d} - 2g_{AA}^{\mu u} + \lambda(g_{VA}^{\mu d} - 2g_{VA}^{\mu u})\right), \tag{3.15}$$

where λ is the muon beam polarization fraction. Two measurements of $b_{\rm SPS}$ at different beam energies and polarization fractions were carried out [90]:

$$\begin{aligned} b_{\rm SPS} &= -(1.47 \pm 0.42) \times 10^{-4}\,{\rm GeV}^{-2} \text{ for } \lambda = 0.81 \;\Rightarrow\; b_{\rm SPS}^{\rm SM} = -1.56 \times 10^{-4}\,{\rm GeV}^{-2}, \\ b_{\rm SPS} &= -(1.74 \pm 0.81) \times 10^{-4}\,{\rm GeV}^{-2} \text{ for } \lambda = 0.66 \;\Rightarrow\; b_{\rm SPS}^{\rm SM} = -1.57 \times 10^{-4}\,{\rm GeV}^{-2}. \end{aligned} \tag{3.16}$$

JHEP08(2017)123

3.3 Low-energy flavor

The partonic process $d_j \to u_i \ell \bar{\nu}_\ell$ underlies a plethora of (semi)leptonic hadron decays. Ref. [55] studied $d(s) \to u \ell \bar{\nu}_\ell$ transitions, such as nuclear, baryon and meson decays, within the SMEFT framework and obtained bounds for 14 combinations of effective low-energy couplings between light quarks and leptons ($\epsilon_i^{d_I e_J}$). Ignoring the CKM mixing at $\mathcal{O}(\Lambda^{-2})$, the effective couplings of strange quarks depend only on flavor-off-diagonal Wilson coefficients (see appendix B). Marginalizing over them, we obtain the likelihood for 6 combinations of effective couplings together with the $\tilde{V}_{ud}$ CKM parameter:[8]

$$\begin{pmatrix} \tilde{V}_{ud} \\ \Delta_{\rm CKM} \\ \epsilon_R^{de} \\ \epsilon_S^{de} \\ \epsilon_P^{de} \\ \epsilon_T^{de} \\ \Delta_{LP}^{d} \end{pmatrix} = \begin{pmatrix} 0.97451(38) \\ -(1.2 \pm 8.4) \cdot 10^{-4} \\ -(1.3 \pm 1.7) \cdot 10^{-2} \\ (1.4 \pm 1.3) \cdot 10^{-3} \\ (4.0 \pm 7.8) \cdot 10^{-6} \\ (1.0 \pm 8.0) \cdot 10^{-4} \\ (1.9 \pm 3.8) \cdot 10^{-2} \end{pmatrix}, \quad \rho = \begin{pmatrix} 1. & 0.88 & 0. & 0.82 & 0.01 & 0. & 0.01 \\ 0.88 & 1. & 0. & 0.73 & 0.01 & 0. & 0.01 \\ 0. & 0. & 1. & 0. & -0.87 & 0. & -0.87 \\ 0.82 & 0.73 & 0. & 1. & 0.01 & 0. & 0.01 \\ 0.01 & 0.01 & -0.87 & 0.01 & 1. & 0. & 0.9995 \\ 0. & 0. & 0. & 0. & 0. & 1. & 0. \\ 0.01 & 0.01 & -0.87 & 0.01 & 1. & 0. & 1. \end{pmatrix}, \quad (3.17)$$

in the $\overline{MS}$ scheme at $\mu = 2\,\mathrm{GeV}$. The effective couplings ϵ were defined in section 2.2.1, and $\Delta_{LP}^d \approx \bar{\epsilon}_L^{de} - \bar{\epsilon}_L^{d\mu} + 24\epsilon_P^{d\mu}$. See appendix B for the complete likelihood [55] that also involves the effective couplings of the strange quarks and allows one to constrain some off-diagonal Wilson coefficients. Using eq. (2.19) we can run these results up to the weak scale, where the matching with the SMEFT is carried out, cf. eq. (2.5) and eq. (2.6).

It is useful to recall the physics behind these bounds [55]. Roughly speaking, $\tilde{V}_{ud}$ and $\epsilon_{R,S,P,T}^{de}$ were obtained comparing the total rates of various superallowed nuclear decays and $\pi \to e\nu_e$, as well as using various differential distributions in $\pi \to e\nu\gamma$ and neutron decay. The comparison with $\Gamma(\pi \to \mu\nu_\mu)$ provides us with Δ_{LP}^d, and the combination of the obtained $\tilde{V}_{ud}$ with V_{us}, extracted from (semi)leptonic kaon decays, makes possible to extract Δ_{CKM}.

3.4 Quark pair production in e^+e^- collisions

Electron-positron colliders operating at center-of-mass energies above or below the Z mass provide complementary information about 4-fermion operators containing electrons. Unlike the low-energy experiments discussed above, they also probe flavor-conserving operators with strange, charm and bottom quarks. Typically, the experiments quote the total measured cross section for $\sigma_q \equiv \sigma(e^+e^- \to q\bar{q})$ and the asymmetry $A_{FB}^q = \frac{\sigma_q^{\rm FB}}{\sigma_q}$, where σ_q^{FB} is the difference between the cross sections with the electron going forward and backward in the center-of-mass frame. In the presence of dimension-6 operators, at $\mathcal{O}(\Lambda^{-2})$ these cross

[8]There is a small (but nonzero) correlation with the effective couplings of strange quarks that we marginalized over. This must be taken into account when going to specific scenarios. The full likelihood is available in ref. [55].

JHEP08(2017)123

sections are modified as follows

$$\delta\sigma_q = \frac{1}{8\pi s}\left[-e^2(g_L^2+g_Y^2)\frac{s}{s-m_Z^2}(\delta A_{Fq}+\delta A_{Bq})+(g_L^2+g_Y^2)^2\frac{s^2}{(s-m_Z^2)^2}(\delta B_{Fq}+\delta B_{Bq})\right]$$
$$+\frac{1}{8\pi v^2}(g_L^2+g_Y^2)\frac{s}{s-m_Z^2}\left(\hat{g}_L^{Ze}\hat{g}_L^{Zq}c_{LL}+\hat{g}_L^{Ze}\hat{g}_R^{Zq}c_{LR}+\hat{g}_R^{Ze}\hat{g}_L^{Zq}c_{RL}+\hat{g}_R^{Ze}\hat{g}_R^{Zq}c_{RR}\right)$$
$$-\frac{1}{8\pi v^2}e^2Q_q\left(c_{LL}+c_{LR}+c_{RL}+c_{RR}\right), \tag{3.18}$$
$$\delta\sigma_q^{\rm FB} = \frac{3}{32\pi s}\left[-e^2(g_L^2+g_Y^2)\frac{s}{s-m_Z^2}(\delta A_{Fq}-\delta A_{Bq})+(g_L^2+g_Y^2)^2\frac{s^2}{(s-m_Z^2)^2}(\delta B_{Fq}-\delta B_{Bq})\right]$$
$$+\frac{3}{32\pi v^2}(g_L^2+g_Y^2)\frac{s}{s-m_Z^2}\left(\hat{g}_L^{Ze}\hat{g}_L^{Zq}c_{LL}+\hat{g}_R^{Ze}\hat{g}_R^{Zq}c_{RR}-\hat{g}_L^{Ze}\hat{g}_R^{Zq}c_{LR}-\hat{g}_R^{Ze}\hat{g}_L^{Zq}c_{RL}\right)$$
$$-\frac{3}{32\pi v^2}e^2Q_q\left(c_{LL}+c_{RR}-c_{LR}-c_{RL}\right), \tag{3.19}$$

where $\sqrt{s}$ is the center-of-mass energy of the e^+e^- collision, $\hat{g}^{Zf}\equiv T_f^3-s_\theta^2Q_f$ (i.e., the SM values), and

$$\delta A_{Fq} = Q_q\left(\delta g_L^{Ze}\hat{g}_L^{Zq}+\delta g_R^{Ze}\hat{g}_R^{Zq}+\hat{g}_L^{Ze}\delta g_L^{Zq}+\hat{g}_R^{Ze}\delta g_R^{Zq}\right), \tag{3.20}$$
$$\delta A_{Bq} = Q_q\left(\delta g_L^{Ze}\hat{g}_R^{Zq}+\delta g_R^{Ze}\hat{g}_L^{Zq}+\hat{g}_L^{Ze}\delta g_R^{Zq}+\hat{g}_R^{Ze}\delta g_L^{Zq}\right),$$
$$\delta B_{Fq} = \hat{g}_L^{Ze}\left(\hat{g}_L^{Zq}\right)^2\delta g_L^{Ze}+\hat{g}_R^{Ze}\left(\hat{g}_R^{Zq}\right)^2\delta g_R^{Ze}+\left(\hat{g}_L^{Ze}\right)^2\hat{g}_L^{Zq}\delta g_L^{Zq}+\left(\hat{g}_R^{Ze}\right)^2\hat{g}_R^{Zq}\delta g_R^{Zq},$$
$$\delta B_{Bq} = \hat{g}_L^{Ze}\left(\hat{g}_R^{Zq}\right)^2\delta g_L^{Ze}+\hat{g}_R^{Ze}\left(\hat{g}_L^{Zq}\right)^2\delta g_R^{Ze}+\left(\hat{g}_R^{Ze}\right)^2\hat{g}_L^{Zq}\delta g_L^{Zq}+\left(\hat{g}_L^{Ze}\right)^2\hat{g}_R^{Zq}\delta g_R^{Zq}.$$

For the up-type quark production, $q=u_J$, the four-fermion Wilson coefficients c_{XY} in eq. (3.18) and eq. (3.19) are given by

$$c_{LL}=[c_{\ell q}]_{11JJ}-[c_{\ell q}^{(3)}]_{11JJ}, \quad c_{LR}=[c_{\ell u}]_{11JJ}, \quad c_{RL}=[c_{eq}]_{11JJ}, \quad c_{RR}=[c_{eu}]_{11JJ}, \tag{3.21}$$

while for the down-type quark production, $q=d_J$,

$$c_{LL}=[c_{\ell q}]_{11JJ}+[c_{\ell q}^{(3)}]_{11JJ}, \quad c_{LR}=[c_{\ell d}]_{11JJ}, \quad c_{RL}=[c_{eq}]_{11JJ}, \quad c_{RR}=[c_{ed}]_{11JJ}. \tag{3.22}$$

The operators $O_{\ell equ}$, $O_{\ell equ}^{(3)}$ and $O_{\ell eqd}$ do not enter at $\mathcal{O}(\Lambda^{-2})$ because they do not interfere with the SM amplitudes due to the different chirality structure.

The LEP-2 experiment studied e^+e^- collisions at energies above the Z-pole, ranging from $\sqrt{s}=130$ Gev to $\sqrt{s}=209$ GeV. Available data includes the total cross section $\sigma(q\bar{q})\equiv\sum_{q=u,d,s,c,b}\sigma_q$ [91], as well as the total cross section and forward-backward asymmetry for the charm and for the bottom quark production [92]. This amounts to 5 distinct observables, each measured at different $\sqrt{s}$. From eq. (3.18) and eq. (3.19), given the energy dependence, each of these observables should resolve 4 different combinations of dimension-6 Wilson coefficients.[9] In practice, the energy range scanned by LEP-2 is not large enough to efficiently disentangle these different combinations. Therefore, in our fit

[9]Note that two of these combinations involve only vertex corrections though.

we also include earlier, less precise measurements of heavy quark production below the Z-pole. Specifically, we include the measurements from the VENUS [93] and TOPAZ [94] collaborations of the $c\bar{c}$ and $b\bar{b}$ pair production at $\sqrt{s} = 58\,\mathrm{GeV}$ (total cross sections and FB asymmetries).

3.5 Other measurements

To increase the power of our global analysis, in this section we will combine the observables described above with those considered previously in refs. [32, 43]. At this point there are more parameters than observables, hence more experimental input is needed. The SMEFT corrections to low-energy observables typically depend on linear combinations of 4-fermion Wilson coefficients and vertex corrections δg. The latter can be independently constrained by the so-called pole observables where a single W or Z boson is on-shell. We use the set of pole observables described in ref. [32]. As advertised in that reference, all diagonal δg can be simultaneously constrained with a very good precision.[10] Moreover, we use the low-energy and e^+e^- collider observables probing 4-lepton operators. Our analysis closely resembles that in ref. [43] with the following differences:

1. Instead of combining ourselves the results of different experiments measuring the scattering of muon neutrinos on electrons, we use the PDG combination for the low-energy ν_μ-e couplings from table 10.8 of ref. [61]:
$$g_{LV}^{\nu_\mu e} = -0.040 \pm 0.015, \qquad g_{LA}^{\nu_\mu e} = -0.507 \pm 0.014, \tag{3.23}$$
with the correlation coefficient $\rho = -0.05$.

2. Instead of recasting the weak mixing angle measured in parity-violating electron scattering [95], we use the PDG result for the parity-violating effective self-coupling of electrons [61]:
$$g_{AV}^{ee} = 0.0190 \pm 0.0027. \tag{3.24}$$

3. To evaluate SMEFT corrections to e^+e^- collider observables we use the electroweak couplings at the scale m_Z (instead of 200 GeV).

4. We add the measurement of the τ polarization $\mathcal{P}_\tau$ and its FB asymmetry $A_\mathcal{P}$ in $e^+e^- \to \tau^+\tau^-$ production at $\sqrt{s} = 58\,\mathrm{GeV}$ by the VENUS collaboration [96]:
$$\mathcal{P}_\tau = 0.012 \pm 0.058, \qquad A_\mathcal{P} = 0.029 \pm 0.057. \tag{3.25}$$
The analytic expressions for $\mathcal{P}_\tau$ and $A_\mathcal{P}$ in function of the SMEFT parameters and $\sqrt{s}$ are easy to obtain but are too long to be quoted here. Instead, we give the numerical expressions at $\sqrt{s} = 58\,\mathrm{GeV}$:
$$\begin{aligned}\delta\mathcal{P}_\tau \approx{}& -0.87\delta g_L^{Ze} - 0.93\delta g_R^{Ze} + 0.17\delta g_L^{Z\tau} + 0.25\delta g_R^{Z\tau} \\ &+0.21[c_{ee}]_{1133} + 0.32[c_{le}]_{1133} - 0.34[c_{le}]_{3311} - 0.20([c_{\ell\ell}]_{1133} + [c_{\ell\ell}]_{1331}), \\ \delta A_\mathcal{P} \approx{}& 0.13\delta g_L^{Ze} + 0.19\delta g_R^{Ze} - 0.65\delta g_L^{Z\tau} - 0.70\delta g_R^{Z\tau} \\ &+0.16[c_{ee}]_{1133} - 0.25[c_{le}]_{1133} + 0.24[c_{le}]_{3311} - 0.15([c_{\ell\ell}]_{1133} + [c_{\ell\ell}]_{1331}).\end{aligned} \tag{3.26}$$

[10] The observables in ref. [32] do not constrain δg_R^{Zt}, which is however not needed in our analysis.

5. We include the constraints from the trident production $\nu_\mu \gamma^* \to \nu_\mu \mu^+ \mu^-$ [97–99]. Dimension-6 operators modify the trident cross section as

$$\frac{\sigma_{\text{trident}}}{\sigma_{\text{trident, SM}}} \approx 1 + 2\frac{g_{LV}^{\nu_\mu\mu,\text{SM}}\delta g_{LV}^{\nu_\mu\mu} + g_{LA}^{\nu_\mu\mu,\text{SM}}\delta g_{LA}^{\nu_\mu\mu}}{(g_{LV}^{\nu_\mu\mu,\text{SM}})^2 + (g_{LA}^{\nu_\mu\mu,\text{SM}})^2}. \tag{3.27}$$

The first 3 modifications lead to negligible numerical differences compared to the fit in ref. [43]. The 4th one allows us to break the degeneracy between $[c_{\ell e}]_{1133}$ and $[c_{\ell e}]_{3311}$ and improve constraints on other 4-lepton operators involving τ. The last modification leads to a constraint on one linear combination of 4-muon dimension-6 operators.

4 Global fit

4.1 Scope

The main goal of this paper is to provide model-independent constraints on flavor-diagonal 2-lepton-2-quark operators summarized in table 1. Among the chirality-conserving ones, most of the observables considered in this paper probe the operators involving the 1st generation leptons. There are 21 such operators and for an easy reference we list here their Wilson coefficients:

$$[c_{\ell q}]_{11JJ},\ [c_{\ell q}^{(3)}]_{11JJ},\ [c_{\ell u}]_{11JJ},\ [c_{\ell d}]_{11JJ},\ [c_{eq}]_{11JJ},\ [c_{eu}]_{11JJ},\ [c_{ed}]_{11JJ}\,, \qquad J = 1, 2, 3. \tag{4.1}$$

Scattering of muons and muon neutrinos on nucleons gives us access to chirality-conserving operators involving 2nd generation leptons and 1st generation quarks. There are 7 such operators:

$$[c_{\ell q}]_{2211},\ [c_{\ell q}^{(3)}]_{2211},\ [c_{\ell u}]_{2211},\ [c_{\ell d}]_{2211},\ [c_{eq}]_{2211},\ [c_{eu}]_{2211},\ [c_{ed}]_{2211}. \tag{4.2}$$

Finally, the likelihood in eq. (3.17) summarizing the constraints from low-energy flavor observables gives us also access to chirality-violating operators involving 1st and 2nd generation leptons and 1st generation quarks. There are 6 such operators:

$$[c_{lequ}]_{JJ11}\,,\ [c_{ledq}]_{JJ11}\,,\ [c_{lequ}^{(3)}]_{JJ11}\,, \qquad J = 1, 2\,, \tag{4.3}$$

which should be understood as evaluated at the renormalization scale $\mu = m_Z$ unless otherwise stated.

We will use the observables summarized in section 3 to constrain as many as possible of the 34 Wilson coefficients in eqs. (4.1)–(4.3). We will also present simultaneous constraints on these parameter, together with the vertex corrections and 4-lepton Wilson coefficients.

4.2 Flat directions

Not all linear combinations of the parameters eqs. (4.1)–(4.3) can be constrained by the observables we consider. Before venturing into a global fit, we need to count the independent constraints and determine the flat directions in the parameter space. In table 4 we have the following probes of LLQQ operators:

JHEP08(2017)123

Class	Observable	Exp. value	Ref. & Comments	SM value
$\nu_e\nu_e qq$	$R_{\nu_e\bar{\nu}_e}$	0.41(14)	CHARM [73]	0.33
$\nu_\mu\nu_\mu qq$	$(g_L^{\nu_\mu})^2$	0.3005(28)	PDG [61], $\rho \approx 1$	0.3034
	$(g_R^{\nu_\mu})^2$	0.0329(30)		0.0302
	$\theta_L^{\nu_\mu}$	2.500(35)		2.4631
	$\theta_R^{\nu_\mu}$	$4.56^{+0.42}_{-0.27}$		5.1765
PV low-E $eeqq$	$g_{AV}^{eu} + 2g_{AV}^{ed}$	0.489(5)	PDG [61], $\rho \neq 1$	0.4951
	$2g_{AV}^{eu} - g_{AV}^{ed}$	$-0.708(16)$		-0.7192
	$2g_{VA}^{eu} - g_{VA}^{ed}$	$-0.144(68)$		-0.0949
	$g_{VA}^{eu} - g_{VA}^{ed}$	$-0.042(57)$ $-0.120(74)$	SAMPLE [89]	-0.0627
PV low-E $\mu\mu qq$	$b_{\rm SPS}(\lambda = 0.81)$	$-1.47(42) \cdot 10^{-4}$	SPS [90]	$-1.56 \cdot 10^{-4}$
	$b_{\rm SPS}(\lambda = 0.66)$	$-1.74(81) \cdot 10^{-4}$		$-1.57 \cdot 10^{-4}$
$d(s) \to u\ell\nu$	$\epsilon_i^{d_j\ell}$	eq. (3.17)	Ref. [55]	0
$e^+e^- \to q\bar{q}$	$\sigma(q\bar{q})$	$f(\sqrt{s})$	LEPEWWG [91], $\rho \neq 1$	$f(\sqrt{s})$
	σ_c, σ_b		LEPEWWG [100], VENUS [93], TOPAZ [94]	
	A_{FB}^{cc}, A_{FB}^{bb}			
$\nu_\mu\nu_\mu ee$	$g_{LV}^{\nu_\mu e}$	$-0.040(15)$	PDG [61], $\rho \neq 1$	-0.0396
	$g_{LA}^{\nu_\mu e}$	$-0.507(14)$		-0.5064
$e^-e^- \to e^-e^-$	g_{AV}^{ee}	0.0190(27)	PDG [61]	0.0225
$\nu_\mu\gamma^* \to \nu_\mu\mu^+\mu^-$	$\frac{\sigma}{\sigma_{\rm SM}}$	1.58(57) 0.82(28)	CHARM [97] CCFR [98]	1
$\tau \to \ell\nu\nu$	$G_{\tau e}^2/G_F^2$	1.0029(46)	PDG [61], $\rho \approx 1$	1
	$G_{\tau\mu}^2/G_F^2$	0.981(18)		1
$e^+e^- \to \ell^+\ell^-$	$\frac{d\sigma(ee)}{d\cos\theta}$	$f(\sqrt{s})$	LEPEWWG [91], $\rho \approx 1$	$f(\sqrt{s})$
	$\sigma_\mu, \sigma_\tau, \mathcal{P}_\tau$		LEPEWWG [100], VENUS [96]	
	A_{FB}^μ, A_{FB}^τ			

Table 4. Summary of experimental input (sensitive to LLQQ and LLLL contact interactions) used in our fit. The correlations that are taken into account in our fit are specified. Each observable in $e^+e^- \to f\bar{f}$ is measured at various c.o.m. energies, which we denote in the table by $f(\sqrt{s})$. The specific numerical values can be found in the corresponding original references. We also use the set of pole observables described in ref. [32] in order to independently constrain the vertex corrections δg.

- 1 combination of the parameters in eq. (4.1) is constrained (poorly) via the only $\nu_e \nu_e qq$ measurement ($R_{\nu_e \bar{\nu}_e}$);
- 4 combinations in eq. (4.2) are constrained via $\nu_\mu \nu_\mu qq$ measurements;
- 4 new combinations in eq. (4.1) are constrained via PV low-energy $eeqq$ measurements ($g^{eq}_{VA/AV}$);
- 1 different combination in eq. (4.2) is constrained (poorly) via PV low-energy $\mu\mu qq$ measurements ($b_{\rm SPS}$), which also probe a second combination already constrained by $\nu_\mu \nu_\mu qq$ data;
- 5 additional combinations in eqs. (4.1)–(4.3) are constrained by low-energy flavor observables ($d(s) \to u \ell \nu_\ell$ transitions);[11]
- 10 additional combinations in eq. (4.1) are probed by $e^+ e^- \to q\bar{q}$ data, through the measurement of the total hadronic cross section and heavy flavor (b and c) fractions and asymmetries.

All together we have 25 constraints on 34 parameters, which leaves 9 flat directions. These can be characterized quite concisely:

$$
\begin{aligned}
&(\mathbf{F1}): && [c_{\ell u}]_{1133}, \\
&(\mathbf{F2}): && [c_{eu}]_{1133}, \\
&(\mathbf{F3}): && [c^{(3)}_{\ell q}]_{1133} = -[c_{\ell q}]_{1133}, \\
&(\mathbf{F4}): && [c^{(3)}_{\ell q}]_{1122} = [c_{\ell q}]_{1122}, \quad [c_{\ell d}]_{1122} = \left(-5 + \frac{3 g_L^2}{g_Y^2}\right)[c_{\ell q}]_{1122}, \\
& && [c_{ed}]_{1122} = \left(3 - \frac{3 g_L^2}{g_Y^2}\right)[c_{\ell q}]_{1122}, \\
&(\mathbf{F5}): && [c_{\ell q}]_{1111} = -[c_{\ell u}]_{1111} = -[c_{\ell d}]_{1111} = -[c_{eq}]_{1111} = [c_{eu}]_{1111} = [c_{ed}]_{1111}, \\
&(\mathbf{F6}): && [c_{eq}]_{2211} = -[c_{ed}]_{2211}, \\
&(\mathbf{F7}): && [c_{eq}]_{2211} = 2[c_{eu}]_{2211}, \\
&(\mathbf{F8}, \mathbf{F9}): && 0.86[c_{ledq}]_{2211} - 0.86[c_{lequ}]_{2211} + 0.012[c^{(3)}_{ledq}]_{2211} = 0.
\end{aligned} \tag{4.4}
$$

The flat directions F1, F2, F3 arise because low-energy precision measurements do not probe the top quark couplings, which may be amended one day by e^+e^- collider operating above the $t\bar{t}$ threshold. F4 is due to the insufficient information about the strange quark couplings, and it could be lifted by off Z-pole measurements of the strange asymmetry. F5 is the consequence of the fact that the parity conserving operator $(\bar{e}\gamma_\mu\gamma_5 e)\sum_q(\bar{q}\gamma_\mu\gamma_5 q)$ and the axial neutrino-quark interaction $(\bar{\nu}_L\gamma_\mu\nu_L)\sum_q(\bar{q}\gamma_\mu\gamma_5 q)$ are unconstrained by low-energy

[11] The likelihood in eq. (3.17) also independently constrains $\delta g_R^{Wq_1}$.

JHEP08(2017)123

measurements and by e^+e^- colliders. F6 and F7 are due to little data on muon scattering on nucleons. Finally, F8 and F9 appear because, with our approximations, the low-energy flavor observables probe only one combination of light quark couplings to muons (through $\pi \to \mu\nu$). The low-energy constraint on $\epsilon_P^{d\mu} = ([c_{ledq}]_{2211} - [c_{lequ}]_{2211})/2$ at $\mu = 2\,\mathrm{GeV}$ (via Δ^d_{LP} in eq. (3.17)), after taking into account the running up to m_Z, becomes a constraint on the linear combination in the last line of eq. (4.4).

In order to isolate the flat directions we define

$$
\begin{aligned}
[\hat{c}_{eq}]_{1111} &= [c_{eq}]_{1111} + [c_{\ell q}]_{1111}, \\
[\hat{c}_{\ell u}]_{1111} &= [c_{\ell u}]_{1111} + [c_{\ell q}]_{1111} - [\hat{c}_{eq}]_{1111}, \\
[\hat{c}_{\ell d}]_{1111} &= [c_{\ell d}]_{1111} + [c_{\ell q}]_{1111} - [\hat{c}_{eq}]_{1111}, \\
[\hat{c}_{eu}]_{1111} &= [c_{eu}]_{1111} - [c_{\ell q}]_{1111}, \\
[\hat{c}_{ed}]_{1111} &= [c_{ed}]_{1111} - [c_{\ell q}]_{1111}, \\
[\hat{c}^{(3)}_{\ell q}]_{1122} &= [c^{(3)}_{\ell q}]_{1122} - [c_{\ell q}]_{1122}, \\
[\hat{c}_{\ell d}]_{1122} &= [c_{\ell d}]_{1122} + \left(5 - \frac{3g_L^2}{g_Y^2}\right)[c_{\ell q}]_{1122} - [\hat{c}_{eq}]_{1111}, \\
[\hat{c}_{ed}]_{1122} &= [c_{ed}]_{1122} - \left(3 - \frac{3g_L^2}{g_Y^2}\right)[c_{\ell q}]_{1122} - [\hat{c}_{eq}]_{1111}, \\
[\hat{c}^{(3)}_{\ell q}]_{1133} &= [c^{(3)}_{\ell q}]_{1133} + [c_{\ell q}]_{1133}, \\
[\hat{c}_{eq}]_{2211} &= [c_{eq}]_{2211} + [c_{ed}]_{2211} - 2[c_{eu}]_{2211}, \\
\epsilon_P^{d\mu}(2\text{ GeV}) &= 0.86[c_{ledq}]_{2211} - 0.86[c_{lequ}]_{2211} + 0.012[c^{(3)}_{ledq}]_{2211}, \\
[\hat{c}_{\ell\ell}]_{2222} &= [c_{\ell\ell}]_{2222} + \frac{2g_Y^2}{g_L^2 + 3g_Y^2}[c_{\ell e}]_{2222}.
\end{aligned}
\tag{4.5}
$$

The last variable projects out the flat direction among 4-muon operators in the trident observable. Using these variables, the global likelihood depends on the Wilson coefficients on the right-hand sides of eqs. (4.5) only via the $\hat{c}$ and $\epsilon_P^{d\mu}(2\text{ GeV})$ combinations.[12] Moreover, the dependence on $[\hat{c}_{eq}]_{1111}$ appears only thanks to the loose $R_{\nu_e\bar{\nu}_e}$ constraint, and thus we know beforehand that there is no sensitivity to $[\hat{c}_{eq}]_{1111} \lesssim 1$.

4.3 Reconnaissance

We start by presenting the constraints in the case when only one of the LLQQ operators is present at a time, and all vertex corrections and 4-lepton operators vanish. We stress that this is just a warm-up exercise and not our main result. Indeed, one-by-one constraints are basis dependent and could be different if another basis of dimension-6 operators was used.

[12]Let us stress that the LLQQ coefficients in the r.h.s. of $\epsilon_P^{d\mu}(2\text{ GeV})$ in eq. (4.5) are defined at $\mu = m_Z$.

JHEP08(2017)123

$(ee)(qq)$

	$[c_{\ell q}^{(3)}]_{1111}$	$[c_{\ell q}]_{1111}$	$[c_{\ell u}]_{1111}$	$[c_{\ell d}]_{1111}$	$[c_{eq}]_{1111}$	$[c_{eu}]_{1111}$	$[c_{ed}]_{1111}$
CHARM	-80 ± 180	700 ± 1800	370 ± 880	-700 ± 1800	x	x	x
APV	27 ± 19	$\mathbf{1.6 \pm 1.1}$	$\mathbf{3.4 \pm 2.3}$	$\mathbf{3.0 \pm 2.0}$	$\mathbf{-1.6 \pm 1.1}$	$\mathbf{-3.4 \pm 2.3}$	$\mathbf{-3.0 \pm 2.0}$
QWEAK	7.0 ± 12	-2.3 ± 4.0	-3.5 ± 6.0	-7 ± 12	2.3 ± 4.0	3.5 ± 6.0	7 ± 12
PVDIS	-8 ± 12	24 ± 35	38 ± 48	-77 ± 96	-77 ± 96	-12 ± 17	24 ± 35
SAMPLE	-8 ± 45	x	-17 ± 90	17 ± 90	x	-17 ± 90	17 ± 90
$d_j \to u \ell \nu$	$\mathbf{0.38 \pm 0.28}$	x	x	x	x	x	x
LEP-2	3.5 ± 2.2	-42 ± 28	-21 ± 14	42 ± 28	-18 ± 11	-9.0 ± 5.7	18 ± 11

$(\mu\mu)(qq)$

	$[c_{\ell q}^{(3)}]_{2211}$	$[c_{\ell q}]_{2211}$	$[c_{\ell u}]_{2211}$	$[c_{\ell d}]_{2211}$	$[c_{eq}]_{2211}$	$[c_{eu}]_{2211}$	$[c_{ed}]_{2211}$
PDG ν_μ	20 ± 15	$\mathbf{4 \pm 21}$	$\mathbf{18 \pm 19}$	$\mathbf{-20 \pm 37}$	x	x	x
SPS	0 ± 1000	0 ± 3000	0 ± 1500	0 ± 3000	$\mathbf{40 \pm 390}$	$\mathbf{-20 \pm 190}$	$\mathbf{40 \pm 390}$
$d_j \to u \ell \nu$	$\mathbf{-0.4 \pm 1.2}$	x	x	x	x	x	x

Table 5. 68% C.L. constraints (in units of 10^{-3}) on chirality-conserving $(ee)(qq)$ and $(\mu\mu)(qq)$ operators from different precision experiments. The bounds are derived assuming that only one operator is present at a time. See table 4 and main text for further details about the different experiments. The best constraint in each case is highlighted in blue, while 'x' signals that the operator is not probed at tree level by that experiment or combination.

Only the global likelihood encoding the correlated constraints on all Wilson coefficients in a given basis has a model-independent meaning. The main purpose of this exercise is to compare the sensitivity of various experiments to a few particular directions in the space of Wilson coefficients.

The one-by-one constraints on chirality-conserving LLQQ operators involving 1st generation quarks are shown in table 5. One can see that atomic parity violation is the most sensitive probe for most of the operators with electrons and the first generation quarks. The exception is $[O_{\ell q}^{(3)}]_{1111}$, which contributes to charged-current transitions and can be probed in $d \to u e \bar{\nu}_e$ decays.[13] We stress however that the less sensitive experiments will be absolutely crucial to probe more independent directions in the space of dimension-6 operators. For the operators involving the 2nd generation lepton doublet the muon-neutrino scattering is a fairly sensitive probe. Again, $[O_{\ell q}^{(3)}]_{2211}$ is very precisely probed by the low-energy flavor observables because it affects the charged current. The sensitivity of low-energy experiments to the operators involving the right-handed muons is very poor. However, this is not a pressing problem, given these directions are very well probed by the LHC [22], as will be discussed in section 5. The $(ee)(qq)$ bounds shown in table 5 are in

[13]The single-operator bounds from $d(s) \to u \ell \bar{\nu}_\ell$ data shown in this section are obtained using the likelihood of eq. (3.17), which was marginalized over strange-quark couplings. Using instead the full likelihood [55] given in appendix B slightly stronger constraints (and central values closer to zero) are obtained.

JHEP08(2017)123

excellent agreement with the 1-by-1 results of ref. [22], whereas our $(\mu\mu)(qq)$ bounds are more stringent due to the inclusion of additional experimental input.

The LEP-2 constraints on operators involving 2nd generation or bottom quarks are similar as those shown in table 5. We also give 1-by-1 constraints on the chirality-violating LLQQ operators from the low-energy flavor observables:

$$\begin{pmatrix} [c_{\ell equ}]_{1111} \\ [c_{\ell edq}]_{1111} \\ [c^{(3)}_{\ell equ}]_{1111} \end{pmatrix} = \begin{pmatrix} -(0.8 \pm 2.9) \cdot 10^{-7} \\ (0.8 \pm 2.9) \cdot 10^{-7} \\ (0.5 \pm 2.0) \cdot 10^{-5} \end{pmatrix},$$

$$\begin{pmatrix} [c_{\ell equ}]_{2211} \\ [c_{\ell edq}]_{2211} \\ [c^{(3)}_{\ell equ}]_{2211} \end{pmatrix} = \begin{pmatrix} (1.7 \pm 5.8) \cdot 10^{-5} \\ -(1.7 \pm 5.8) \cdot 10^{-5} \\ -(1.2 \pm 4.1) \cdot 10^{-3} \end{pmatrix}. \tag{4.6}$$

This exceptional sensitivity arises because these operators violate the approximate symmetries of the SM, leading potentially to a large enhancement of several decays of low-mass hadrons.[14] In particular, new physics generating the pseudo-scalar $(ee)(qq)$ operator is probed up to $\Lambda/g_* \sim 100\,\text{TeV}$. Let us note that they dominate the $c^{(3)}_{\ell equ}$ bounds shown above, despite the fact that they probe them only via 1-loop QED mixing [67, 101]. For consistency with the rest of this work, these individual limits are obtained using $V = 1$ at order Λ^{-2}. Working instead with the full non-diagonal CKM matrix the limits are slightly modified, but more importantly one can set strong 1-by-1 limits in a long list of other (offdiagonal) operators.

Finally, for the sake of completeness we show the 1-by-1 bound on the W coupling to right-handed 1st-generation quarks

$$\delta g_R^{Wq_1} = -(3.9 \pm 2.9) \cdot 10^{-4}, \tag{4.7}$$

which is completely dominated by its contribution to the CKM-unitarity test of eq. (2.6).

4.4 All out

We now combine all the experimental observables summarized in table 4 along with the pole observables discussed in ref. [32], which represent a total of 264 experimental input. These provide simultaneous constraints on 61 combinations of Wilson coefficients of dimension-6 operators in the SMEFT Lagrangian (21 vertex corrections δg, 25 LLQQ and 15 LLLL operators) and on the $\tilde{V}_{ud}$ SM parameter. Marginalizing over $\tilde{V}_{ud}$ we find the

[14] More specifically they violate the approximate flavor symmetry of the SM $\mathrm{U}(1)_\ell \times \mathrm{U}(1)_e$ that suppresses the decay $\pi \to \ell\nu_\ell$ by a factor m_ℓ^2/Λ_{QCD}^2. Thus, their bounds benefit from this large Λ_{QCD}/m_ℓ chiral enhancement. This does not apply however to the tensor operator $c^{(3)}_{\ell equ}$, whose tree-level contribution to this specific decay is zero by simple Lorentz invariance considerations.

following constraints:

$$
\begin{pmatrix}
\delta g_L^{We} \\ \delta g_L^{W\mu} \\ \delta g_L^{W\tau} \\ \delta g_L^{Ze} \\ \delta g_L^{Z\mu} \\ \delta g_L^{Z\tau} \\ \delta g_R^{Ze} \\ \delta g_R^{Z\mu} \\ \delta g_R^{Z\tau} \\ \delta g_L^{Zu} \\ \delta g_L^{Zc} \\ \delta g_L^{Zt} \\ \delta g_R^{Zu} \\ \delta g_R^{Zc} \\ \delta g_L^{Zd} \\ \delta g_L^{Zs} \\ \delta g_L^{Zb} \\ \delta g_R^{Zd} \\ \delta g_R^{Zs} \\ \delta g_R^{Zb} \\ \delta g_R^{Wq_1} \\ [c_{\ell\ell}]_{1111} \\ [c_{\ell e}]_{1111} \\ [c_{ee}]_{1111} \\ [c_{\ell\ell}]_{1221} \\ [c_{\ell\ell}]_{1122} \\ [c_{\ell e}]_{1122} \\ [c_{\ell e}]_{2211} \\ [c_{ee}]_{1122} \\ [c_{\ell\ell}]_{1331} \\ [c_{\ell\ell}]_{1133} \\ [c_{\ell e}]_{1133} \\ [c_{\ell e}]_{3311} \\ [c_{ee}]_{1133} \\ [\hat{c}_{\ell\ell}]_{2222} \\ [c_{\ell\ell}]_{2332}
\end{pmatrix}
=
\begin{pmatrix}
-1.00 \pm 0.64 \\ -1.36 \pm 0.59 \\ 1.95 \pm 0.79 \\ -0.023 \pm 0.028 \\ 0.01 \pm 0.12 \\ 0.018 \pm 0.059 \\ -0.033 \pm 0.027 \\ 0.00 \pm 0.14 \\ 0.042 \pm 0.062 \\ -0.8 \pm 3.1 \\ -0.15 \pm 0.36 \\ -0.3 \pm 3.8 \\ 1.4 \pm 5.1 \\ -0.35 \pm 0.53 \\ -0.9 \pm 4.4 \\ 0.9 \pm 2.8 \\ 0.33 \pm 0.17 \\ 3 \pm 16 \\ 3.4 \pm 4.9 \\ 2.30 \pm 0.88 \\ -1.3 \pm 1.7 \\ 1.01 \pm 0.38 \\ -0.22 \pm 0.22 \\ 0.20 \pm 0.38 \\ -4.8 \pm 1.6 \\ 1.5 \pm 2.1 \\ 1.5 \pm 2.2 \\ -1.4 \pm 2.2 \\ 3.4 \pm 2.6 \\ 1.5 \pm 1.3 \\ 0 \pm 11 \\ -2.3 \pm 7.2 \\ 1.7 \pm 7.2 \\ -1 \pm 12 \\ -2 \pm 21 \\ 3.0 \pm 2.3
\end{pmatrix}
\times 10^{-2},
\qquad
\begin{pmatrix}
[c^{(3)}_{\ell q}]_{1111} \\ [\hat{c}_{eq}]_{1111} \\ [\hat{c}_{\ell u}]_{1111} \\ [\hat{c}_{\ell d}]_{1111} \\ [\hat{c}_{eu}]_{1111} \\ [\hat{c}_{ed}]_{1111} \\ [\hat{c}^{(3)}_{\ell q}]_{1122} \\ [c_{\ell u}]_{1122} \\ [\hat{c}_{\ell d}]_{1122} \\ [c_{eq}]_{1122} \\ [c_{eu}]_{1122} \\ [\hat{c}_{ed}]_{1122} \\ [\hat{c}^{(3)}_{\ell q}]_{1133} \\ [c_{\ell d}]_{1133} \\ [c_{eq}]_{1133} \\ [c_{ed}]_{1133} \\ [c^{(3)}_{\ell q}]_{2211} \\ [c_{\ell q}]_{2211} \\ [c_{\ell u}]_{2211} \\ [c_{\ell d}]_{2211} \\ [\hat{c}_{eq}]_{2211} \\ [c_{\ell equ}]_{1111} \\ [c_{\ell edq}]_{1111} \\ [c^{(3)}_{\ell equ}]_{1111} \\ \epsilon_P^{d\mu}(2\ \text{GeV})
\end{pmatrix}
=
\begin{pmatrix}
-2.2 \pm 3.2 \\ 100 \pm 180 \\ -5 \pm 11 \\ -5 \pm 23 \\ -1 \pm 12 \\ -4 \pm 21 \\ -61 \pm 32 \\ 2.4 \pm 8.0 \\ -310 \pm 130 \\ -21 \pm 28 \\ -87 \pm 46 \\ 270 \pm 140 \\ -8.6 \pm 8.0 \\ -1.4 \pm 10 \\ -3.2 \pm 5.1 \\ 18 \pm 20 \\ -1.2 \pm 3.9 \\ 1.3 \pm 7.6 \\ 15 \pm 12 \\ 25 \pm 34 \\ 4 \pm 41 \\ -0.080 \pm 0.075 \\ -0.079 \pm 0.074 \\ -0.02 \pm 0.19 \\ -0.02 \pm 0.15
\end{pmatrix}
\times 10^{-2}. \tag{4.8}
$$

The correlation matrix is available in the `Mathematica` notebook attached as a supplemental material [56]. The complete Gaussian likelihood for the Wilson coefficients of dimension-6 SMEFT operators at the scale $\mu = m_Z$ can be reproduced from eq. (4.8) and that correlation matrix. For user's convenience, in the notebook the likelihood is displayed ready-made for cut and paste, and we also provide a translation to the Warsaw basis. That likelihood is relevant to constrain the masses and couplings of any new physics model whose leading effects at the weak scale can be approximated by tree-level contributions of vertex corrections and LLQQ and LLLL operators in the SMEFT.

The model-independent bounds on the vertex corrections are practically the same as the ones obtained from the pole observables only in ref. [32]. This is due to the fact

JHEP08(2017)123

that there are more 4-fermion operators than independent off-pole observables. Hence the latter serve to bound 4-fermion Wilson coefficients but cannot further constrain δg. Nevertheless, there are non-zero correlations between the constraints on vertex corrections and 4-fermion operators that are captured by our analysis. It is worth stressing the CKM-unitarity test Δ_{CKM} of eq. (2.6), which actually provides stronger one-by-one limits on the vertex corrections $\delta g_L^{Wq_1}$ and $\delta g_L^{W\mu}$ than all pole observables combined.

Furthermore, the low-energy flavor observables provide a percent level bound on the W boson coupling to right-handed light quarks $\delta g_R^{Wq_1}$ [55]. Recall that δg_R^{Wq} are not probed by the pole observables at tree level and $\mathcal{O}(\Lambda^{-2})$ in the SMEFT expansion, therefore the model-independent limit in eq. (4.8) (from ref. [55]) is a new result. It is weaker than the one in eq. (4.7) because in the global fit the strong constraints from the CKM-unitarity test of eq. (2.6) are diluted by marginalizing over less precisely probed dimension-6 parameters. Nevertheless, the constraint on $\delta g_R^{Wq_1}$ will typically be stronger in specific new physics scenarios, unless they predict that the particular linear combination on the r.h.s. of eq. (2.6) approximately vanishes at the sub-per-mille level.

The bounds on LLLL operators involving only electrons and/or muons are also similar to the ones previously obtained in ref. [43], with the exception of $[c_{\ell\ell}]_{2222}$ which is now bound due to the inclusion of neutrino trident production data. For the $ee\tau\tau$ operators the bounds are much stronger thanks to including the VENUS $\tau\tau$ polarization data, which resolves the degeneracies present in the fit of ref. [43].

The model-independent bounds on LLQQ operators in eq. (4.8) are new. Previous global SMEFT analyses targeting these operators [9, 10, 39] were carried out assuming some simplifying flavor structure, such as the $U(3)^5$ symmetry [9], which greatly reduces the number of independent Wilson coefficients. On the other hand, previous analyses working in a flavor general setup provided 1-by-1 bounds (see e.g. ref. [15, 22]). Thus, the global bounds applicable for a completely arbitrary flavor structure are obtained for the first time in this paper, and they represent our main result. They are relevant for a large class of new physics scenarios with or without approximate flavor symmetries. In particular, models addressing various flavor anomalies necessarily do not respect the $U(3)^5$ symmetry, and therefore the global likelihood we obtained may provide new constraints on their parameters.

We find several poorly constrained directions in the space of LLQQ operators. As discussed earlier, $[\hat{c}_{eq}]_{1111}$ is currently constrained only by very imprecise measurements of electron neutrino scattering on nucleons, such that the experiments are insensitive to $[\hat{c}_{eq}]_{1111} \lesssim 1$. More surprisingly, another practically unconstrained direction emerges in our fit, which roughly corresponds to the linear combination $[\hat{c}_{ed} + 0.6\,\hat{c}_{\ell d}]_{1122}$. This can be traced to the fact that the LEP-2 collider was scanning a fairly narrow range of $\sqrt{s}$ in e^+e^- collisions. For this reason, not all theoretically available combinations discussed in section 4.2 are resolved in practice. Again, it is should be noted that constraints in typical scenarios generating these LLQQ operators will be stronger unless the operators accidentally align with the flat directions in our fit. We stress that the global likelihood provided in the supplemental material [56] retains the full information about the correlations.

JHEP08(2017)123

4.5 Flavor-universal limit

The general likelihood presented in section 4.4 can be easily restricted to a smaller subspace relevant for any particular scenario. We present here the results for the flavor-universal limit, where dimension-6 operators are invariant under the global flavor symmetry $\mathrm{U}(3)^5$. The symmetry implies that 1) all off-diagonal and chirality-violating operators as well as δg_R^{Wq} are absent, 2) the remaining operators do not carry the flavor index. The only subtlety concerns the $[c_{\ell\ell}]_{IJKL}$ coefficients, since two independent contractions of flavor indices are allowed by the $\mathrm{U}(3)^5$ symmetry. We follow the common practice of parametrizing them in terms of the two $\mathrm{U}(3)^5$-symmetric operators $O_{\ell\ell} \equiv \frac{1}{2}\sum_{I,J}(\bar{\ell}_I\bar{\sigma}_\mu\ell_I)(\bar{\ell}_J\bar{\sigma}_\mu\ell_J)$ and $O_{\ell\ell}^{(3)} \equiv \frac{1}{2}\sum_{I,J}(\bar{\ell}_I\sigma^i\bar{\sigma}_\mu\ell_I)(\bar{\ell}_J\sigma^i\bar{\sigma}_\mu\ell_J)$. All in all, with the parameterization of the dimension-6 space used in this paper, the $\mathrm{U}(3)^5$ symmetry corresponds to the following pattern:

$$\begin{pmatrix} \delta g_L^{We_J} \\ \delta g_L^{Ze_J} \\ \delta g_R^{Ze_J} \\ \delta g_L^{Zu_J} \\ \delta g_R^{Zu_J} \\ \delta g_L^{Zd_J} \\ \delta g_R^{Zd_J} \end{pmatrix} = \begin{pmatrix} \delta g_L^{We} \\ \delta g_L^{Ze} \\ \delta g_R^{Ze} \\ \delta g_L^{Zu} \\ \delta g_R^{Zu} \\ \delta g_L^{Zd} \\ \delta g_R^{Zd} \end{pmatrix}, \quad \begin{pmatrix} [c_{\ell\ell}]_{JJJJ} \\ [c_{\ell\ell}]_{IJJI} \\ [c_{\ell\ell}]_{IIJJ} \\ [c_{\ell e}]_{IIJJ} \\ [c_{ee}]_{IIJJ} \end{pmatrix} = \begin{pmatrix} c_{\ell\ell} + c_{\ell\ell}^{(3)} \\ 2c_{\ell\ell}^{(3)} \\ c_{\ell\ell} - c_{\ell\ell}^{(3)} \\ c_{\ell e} \\ c_{ee} \end{pmatrix}, \quad \begin{pmatrix} [c_{\ell q}^{(3)}]_{IIJJ} \\ [c_{\ell q}]_{IIJJ} \\ [c_{eq}]_{IIJJ} \\ [c_{\ell u}]_{IIJJ} \\ [c_{\ell d}]_{IIJJ} \\ [c_{eu}]_{IIJJ} \\ [c_{ed}]_{IIJJ} \end{pmatrix} = \begin{pmatrix} c_{\ell q}^{(3)} \\ c_{\ell q} \\ c_{eq} \\ c_{\ell u} \\ c_{\ell d} \\ c_{eu} \\ c_{ed} \end{pmatrix}, \tag{4.9}$$

and all the remaining vertex corrections and 4-fermion Wilson coefficients vanish. This setup corresponds to the SMEFT limit studied in the pioneering work of ref. [9].[15]

It turns out that the global likelihood constrains the entire restricted parameter set introduced in eq. (4.9). Thus, unlike in the flavor-generic case, there is no need to define new variables $\hat{c}$ in order to factor out the flat directions. Marginalizing over $\tilde{V}_{ud}$, we find the following constraints:

$$\begin{pmatrix} \delta g_L^{We} \\ \delta g_L^{Ze} \\ \delta g_R^{Ze} \\ \delta g_L^{Zu} \\ \delta g_R^{Zu} \\ \delta g_L^{Zd} \\ \delta g_R^{Zd} \end{pmatrix} = \begin{pmatrix} -1.22 \pm 0.81 \\ -0.10 \pm 0.21 \\ -0.15 \pm 0.23 \\ -1.6 \pm 2.0 \\ -2.1 \pm 4.1 \\ 1.9 \pm 1.4 \\ 15 \pm 7 \end{pmatrix} \times 10^{-3}, \tag{4.10}$$

$$\begin{pmatrix} c_{\ell\ell}^{(3)} \\ c_{\ell\ell} \\ c_{\ell e} \\ c_{ee} \end{pmatrix} = \begin{pmatrix} -3.0 \pm 1.7 \\ 7.2 \pm 3.3 \\ 0.2 \pm 1.3 \\ -2.5 \pm 3.0 \end{pmatrix} \times 10^{-3}, \quad \begin{pmatrix} c_{\ell q}^{(3)} \\ c_{\ell q} \\ c_{eq} \\ c_{\ell u} \\ c_{\ell d} \\ c_{eu} \\ c_{ed} \end{pmatrix} = \begin{pmatrix} -4.8 \pm 2.3 \\ -15.4 \pm 9.1 \\ -14 \pm 23 \\ 4 \pm 24 \\ 6 \pm 42 \\ 4 \pm 11 \\ 26 \pm 18 \end{pmatrix} \times 10^{-3}. \tag{4.11}$$

[15] Let us note that the more recent analysis of ref. [39] corresponds to a more restricted scenario, since the two independent coefficients $c_{\ell\ell}$ and $c_{\ell\ell}^{(3)}$ are controlled by one single coefficient $C_{\ell\ell}$ in that work.

The correlation matrix reads $\rho =$

$$\begin{pmatrix}
1. & -0.5 & 0.2 & 0.1 & 0.1 & 0. & 0. & 1. & -0.5 & 0. & -0.1 & 0.4 & -0.1 & 0. & 0.1 & 0. & 0.1 & 0. \\
-0.5 & 1. & 0.3 & -0.1 & 0. & -0.2 & -0.2 & -0.5 & 0.2 & 0. & 0.1 & -0.1 & 0.1 & 0. & 0. & 0. & -0.1 & -0.1 \\
0.2 & 0.3 & 1. & 0. & 0. & -0.3 & -0.3 & 0.2 & -0.2 & 0. & 0.1 & 0.3 & 0. & 0.1 & 0.1 & 0.1 & 0. & -0.1 \\
0.1 & -0.1 & 0. & 1. & 0.8 & 0.2 & 0.1 & 0.1 & 0. & 0. & 0. & 0.7 & -0.3 & 0. & 0.1 & 0. & 0.5 & 0.1 \\
0.1 & 0. & 0. & 0.8 & 1. & 0.1 & 0.2 & 0.1 & 0. & 0. & 0. & 0.7 & -0.3 & 0. & 0.1 & 0. & 0.5 & 0.2 \\
0. & -0.2 & -0.3 & 0.2 & 0.1 & 1. & 0.9 & 0. & 0. & 0. & 0. & -0.4 & -0.2 & -0.1 & -0.1 & -0.2 & 0.2 & 0.4 \\
0. & -0.2 & -0.3 & 0.1 & 0.2 & 0.9 & 1. & 0. & 0. & 0. & 0. & -0.5 & -0.2 & -0.1 & -0.1 & -0.2 & 0.2 & 0.4 \\
1. & -0.5 & 0.2 & 0.1 & 0.1 & 0. & 0. & 1. & -0.5 & 0. & -0.1 & 0.4 & -0.1 & 0. & 0.1 & 0. & 0.1 & 0. \\
-0.5 & 0.2 & -0.2 & 0. & 0. & 0. & 0. & -0.5 & 1. & -0.2 & -0.6 & -0.2 & 0. & 0. & 0. & 0. & 0. & 0. \\
0. & 0. & 0. & 0. & 0. & 0. & 0. & 0. & -0.2 & 1. & -0.2 & 0. & 0. & 0. & 0. & 0. & 0. & 0. \\
-0.1 & 0.1 & 0.1 & 0. & 0. & 0. & 0. & -0.1 & -0.6 & -0.2 & 1. & 0. & 0. & 0. & 0. & 0. & 0. & 0. \\
0.4 & -0.1 & 0.3 & 0.7 & 0.7 & -0.4 & -0.5 & 0.4 & -0.2 & 0. & 0. & 1. & -0.1 & 0.1 & 0.2 & 0.1 & 0.3 & -0.1 \\
-0.1 & 0.1 & 0. & -0.3 & -0.3 & -0.2 & -0.2 & -0.1 & 0. & 0. & 0. & -0.1 & 1. & -0.2 & -0.7 & -0.6 & -0.5 & -0.9 \\
0. & 0. & 0.1 & 0. & 0. & -0.1 & -0.1 & 0. & 0. & 0. & 0. & 0.1 & -0.2 & 1. & 0.7 & 0.9 & -0.5 & 0.5 \\
0.1 & 0. & 0.1 & 0.1 & 0.1 & -0.1 & -0.1 & 0.1 & 0. & 0. & 0. & 0.2 & -0.7 & 0.7 & 1. & 0.9 & -0.1 & 0.8 \\
0. & 0. & 0.1 & 0. & 0. & -0.2 & -0.2 & 0. & 0. & 0. & 0. & 0.1 & -0.6 & 0.9 & 0.9 & 1. & -0.2 & 0.7 \\
0.1 & -0.1 & 0. & 0.5 & 0.5 & 0.2 & 0.2 & 0.1 & 0. & 0. & 0. & 0.3 & -0.5 & -0.5 & -0.1 & -0.2 & 1. & 0.3 \\
0. & -0.1 & -0.1 & 0.1 & 0.2 & 0.4 & 0.4 & 0. & 0. & 0. & 0. & -0.1 & -0.9 & 0.5 & 0.8 & 0.7 & 0.3 & 1.
\end{pmatrix} \tag{4.12}$$

where the rows and columns correspond to the ordering of the parameters in eq. (4.10) and eq. (4.11). The correlation matrix with more significant digits (necessary for practical applications) is given in the `Mathematica` notebook attached as supplemental material [56].

Thanks to lifting the exact and approximate flat directions, in the $\mathrm{U}(3)^5$ symmetric limit typical constraints on the dimension-6 parameters are at the per-mille level. We note that the vertex corrections are constrained slightly better than when only the pole observables are used [32], thanks to the precise input from low-energy flavor measurements. Most of the LLQQ operators are constrained at the percent level.

Also working in the flavor-universal limit, ref. [40] obtained bounds on 10 *additional* SMEFT coefficients using Higgs data and WW production at LEP2. The only flavor-universal SMEFT coefficients unconstrained by these two fits are those that are either CP-violating, or contain only quarks, only gluons or only higgses.

4.6 Oblique parameters

In the literature, precision constraints on new physics are often quoted in the language of *oblique parameters* S, T, W, Y [11, 102]. These correspond to a further restriction of the pattern of the dimension-6 parameters in the $\mathrm{U}(3)^5$ symmetric case [43, 103]:

$$\begin{aligned}
\delta g_{L/R}^{Zf} &= \alpha \left\{ T^3_{f_{L/R}} \frac{T - W - \frac{g_Y^2}{g_L^2} Y}{2} + Q_f \frac{2 g_Y^2 T - (g_L^2 + g_Y^2) S + 2 g_Y^2 W + \frac{2 g_Y^2 (2 g_L^2 - g_Y^2)}{g_L^2} Y}{4(g_L^2 - g_Y^2)} \right\}, \\
\delta g_L^{We} &= \frac{\alpha}{2(g_L^2 - g_Y^2)} \left(-\frac{g_L^2 + g_Y^2}{2} S + g_L^2 T - (g_L^2 - 2 g_Y^2) W + g_Y^2 Y \right), \\
c_{\ell\ell}^{(3)} &= c_{\ell q}^{(3)} = c_{qq}^{(3)} = -\alpha W, \qquad c_{f_1 f_2} = -4 Y_{f_1} Y_{f_2} \frac{g_Y^2}{g_L^2} \alpha Y \,,
\end{aligned} \tag{4.13}$$

where Y_{f_i} is the fermionic hypercharge. With this pattern, all vertex corrections and 4-fermion operators can be redefined away, such that new physics affects only the electroweak gauge boson propagators. Restricting the $\mathrm{U}(3)^5$ symmetric likelihood using eq. (4.13) we find the following constraints on the oblique parameters:

$$\begin{pmatrix} S \\ T \\ Y \\ W \end{pmatrix} = \begin{pmatrix} -0.10 \pm 0.13 \\ 0.02 \pm 0.08 \\ -0.15 \pm 0.11 \\ -0.01 \pm 0.08 \end{pmatrix}, \qquad \rho = \begin{pmatrix} 1. & 0.86 & 0.70 & -0.12 \\ . & 1. & 0.39 & -0.06 \\ . & . & 1. & -0.49 \\ . & . & . & 1. \end{pmatrix}. \tag{4.14}$$

The constraints on the oblique corrections are dominated by the pole-observables and lepton-pair production in LEP-2. The new observables probing LLQQ operators do not affect these constraints significantly. In particular, the low-energy flavor observables do not probe the oblique corrections at all. Compared to the fit in ref. [43], we only observe a small shift of the central values.[16]

5 Comments on LHC reach

Four-fermion LLQQ operators can be probed via the $q\bar{q} \to \ell^+\ell^-$ processes in hadron colliders. Previously several groups set bounds on their Wilson coefficients through the reanalysis within the SMEFT of various ATLAS and CMS exotic searches (see e.g. [22, 104, 105]). In this section we derive analogue bounds using the recently published measurements of the differential Drell-Yan cross sections in the dielectron and dimuon channels [106]. Our main goal here is to present a brief comparison between the sensitivity of the LHC run-1 and of the low-energy observables discussed in this paper.

Precision measurements in hadron collider environments are challenging. Individual observables are typically measured with much worse accuracy than in lepton colliders or very low-energy experiments. However, the effect of 4-fermion operators on scattering amplitudes grows with the collision energy E as $\sim c_{4f}E^2/v^2$. As a consequence, the superior energy reach of the LHC compensates the inferior precision in this case [22, 104]. This message was recently stressed in ref. [107] in the context of the determination of the oblique parameters, which encode new physics corrections to propagators of the electroweak gauge bosons. It turns out that the effect of the oblique parameters W and Y [11] on the high invariant-mass tail of $\frac{d\sigma(pp\to\ell^+\ell^-)}{dm_{\ell\ell}}$ also grows with E (as opposed to that of the more familiar S and T parameters [102]). The current LHC constraint on W and Y are already competitive with those obtained from low-energy precision experiments, and will become more accurate with the full run-2 dataset at $\sqrt{s} \approx$ 13-14 TeV [107]. In the SMEFT framework, W and Y correspond to a particular pattern of vertex corrections and 4-fermion operators [43, 103], cf. eq. (4.13). Therefore we expect that similar arguments apply, and that competitive bounds on the LLQQ operators can be extracted from ATLAS and CMS measurements of $\frac{d\sigma(pp\to\ell^+\ell^-)}{dm_{\ell\ell}}$. Below we present some quantitative illustrations of this message.

[16] The $\mathcal{O}(10\%)$ increase of some errors compared to [43] is due to using different values of the electroweak couplings to evaluate the dimension-6 shifts of the LEP-2 observables.

$(ee)(qq)$

	$[c_{\ell q}^{(3)}]_{1111}$	$[c_{\ell q}]_{1111}$	$[c_{\ell u}]_{1111}$	$[c_{\ell d}]_{1111}$	$[c_{eq}]_{1111}$	$[c_{eu}]_{1111}$	$[c_{ed}]_{1111}$
Low-energy	0.45 ± 0.28	1.6 ± 1.0	2.8 ± 2.1	3.6 ± 2.0	-1.8 ± 1.1	-4.0 ± 2.0	-2.7 ± 2.0
$\mathrm{LHC}_{1.5}$	$-0.70^{+0.66}_{-0.74}$	$2.5^{+1.9}_{-2.5}$	$2.9^{+2.4}_{-2.9}$	$-1.6^{+3.4}_{-3.0}$	$1.6^{+1.8}_{-2.2}$	$1.6^{+2.5}_{-1.5}$	$-3.1^{+3.6}_{-3.0}$
$\mathrm{LHC}_{1.0}$	$-0.84^{+0.85}_{-0.92}$	$3.6^{+3.6}_{-3.7}$	$4.4^{+4.4}_{-4.7}$	$-2.4^{+4.8}_{-4.7}$	$2.4^{+3.0}_{-3.2}$	$1.9^{+2.5}_{-1.9}$	$-4.6^{+5.4}_{-4.1}$
$\mathrm{LHC}_{0.7}$	$-1.0^{+1.4}_{-1.5}$	5.9 ± 7.2	7.4 ± 9.0	-3.6 ± 8.7	3.8 ± 5.9	$2.1^{+3.8}_{-2.9}$	-8 ± 10

$(\mu\mu)(qq)$

	$[c_{\ell q}^{(3)}]_{2211}$	$[c_{\ell q}]_{2211}$	$[c_{\ell u}]_{2211}$	$[c_{\ell d}]_{2211}$	$[c_{eq}]_{2211}$	$[c_{eu}]_{2211}$	$[c_{ed}]_{2211}$
Low-energy	-0.2 ± 1.2	4 ± 21	18 ± 19	-20 ± 37	40 ± 390	-20 ± 190	40 ± 390
$\mathrm{LHC}_{1.5}$	$-1.22^{+0.62}_{-0.70}$	1.8 ± 1.3	2.0 ± 1.6	-1.1 ± 2.0	1.1 ± 1.2	$2.5^{+1.8}_{-1.4}$	-2.2 ± 2.0
$\mathrm{LHC}_{1.0}$	$-0.72^{+0.81}_{-0.87}$	$3.2^{+4.0}_{-3.5}$	$3.9^{+4.8}_{-4.4}$	$-2.3^{+4.9}_{-4.7}$	$2.3^{+3.1}_{-3.2}$	$1.6^{+2.3}_{-1.8}$	-4.4 ± 5.3
$\mathrm{LHC}_{0.7}$	$-0.7^{+1.3}_{-1.4}$	$3.2^{+10.3}_{-4.8}$	$4.3^{+12.5}_{-6.4}$	-3.6 ± 9.0	3.8 ± 6.2	$1.6^{+3.4}_{-2.7}$	-8 ± 11

Chirality-violating operators ($\mu = 1\,\mathrm{TeV}$)

	$[c_{\ell equ}]_{1111}$	$[c_{\ell edq}]_{1111}$	$[c_{\ell equ}^{(3)}]_{1111}$	$[c_{\ell equ}]_{2211}$	$[c_{\ell edq}]_{2211}$	$[c_{\ell equ}^{(3)}]_{2211}$
Low-energy	$(-0.6 \pm 2.4)10^{-4}$	$(0.6 \pm 2.4)10^{-4}$	$(0.4 \pm 1.4)10^{-3}$	$0.014(49)$	$-0.014(49)$	$-0.09(29)$
$\mathrm{LHC}_{1.5}$	0 ± 2.0	0 ± 2.6	0 ± 0.91	0 ± 1.2	0 ± 1.6	0 ± 0.56
$\mathrm{LHC}_{1.0}$	0 ± 2.9	0 ± 3.7	0 ± 1.4	0 ± 2.9	0 ± 3.7	0 ± 1.4
$\mathrm{LHC}_{0.7}$	0 ± 5.3	0 ± 6.6	0 ± 2.6	0 ± 5.5	0 ± 6.9	0 ± 2.6

Table 6. Comparison of low-energy and LHC constraints (in units of 10^{-3}) on the Wilson coefficients of the chirality-conserving $(ee)(qq)$ and $(\mu\mu)(qq)$ and chirality-violating operators defined at the scale $\mu = 1\,\mathrm{TeV}$. The 68% CL bounds are derived assuming only one 4-fermion operator is present at a time, and that the vertex corrections and $[c_{\ell\ell}]_{1221}$ are absent. The low-energy constraints combine all experimental input summarized in table 4. The $\mathrm{LHC}_{1.5}$ constraints use the $m_{\ell\ell} \in [0.5\text{-}1.5]\,\mathrm{TeV}$ bins of the measured differential e^+e^- and $\mu^+\mu^-$ cross sections at the 8 TeV LHC [106]. We also separately show the constraints obtained when the $m_{\ell\ell} \in [0.5\text{-}1.0]\,\mathrm{TeV}$ ($\mathrm{LHC}_{1.0}$) and $m_{\ell\ell} \in [0.5\text{-}0.7]\,\mathrm{TeV}$ ($\mathrm{LHC}_{0.7}$) data range is used.

In the situation when only one LLQQ operator is present at a time and all other dimension-6 operators are absent, the sensitivity of the LHC run-1 and of the low-energy observables is contrasted in table 6. To estimate the LHC reach we use 3 bins in the range $m_{\ell\ell} \in [0.5\text{-}1.5]\,\mathrm{TeV}$ of the ATLAS measurement of the differential e^+e^- and $\mu^+\mu^-$ cross sections at the 8 TeV LHC ($20.3\,\mathrm{fb}^{-1}$) [106]. This is shown under the label of $\mathrm{LHC}_{1.5}$ constraints in table 4, and it is compared to the combined constraints using the low-energy input. For the chirality conserving $(ee)(qq)$ operators the two are indeed similarly sensitive. For the chirality conserving $(\mu\mu)(qq)$ operators the low-energy bounds are relatively weaker, especially for the operators that do not affect the muon neutrino couplings. With the exception of $[O_{\ell q}^{(3)}]_{2211}$ probed by the flavor observables, the LHC sensitivity is superior by at least an order of magnitude. Therefore in these directions in the parameter space

JHEP08(2017)123

of dimension-6 SMEFT the LHC is in a completely uncharted territory. The situation is quite opposite for the chirality-violating $(ee)(qq)$ and $(\mu\mu)(qq)$ operators. There the light quark transitions offer a superior sensitivity with which the LHC cannot compete in most cases. The exception is the $[O^{(3)}_{\ell equ}]_{2211}$ operator where the LHC reach is comparable.

An important difference between the LHC and low-energy constraints should be emphasized. The latter are obtained in the energy regime where it is very plausible to assume the validity of the EFT. Here, by validity we mean that the SMEFT with dimension-6 operators adequately describes the physics of the underlying UV completion. First of all, if such completion contains new states at $\sim 1\,\mathrm{TeV}$ then clearly the LHC bounds in table 6 cannot be applied and a model-dependent approach becomes necessary. This is however not the case for the SMEFT bounds derived from low-energy data in the previous section, which are still valid. On the other hand, even in the absence of such "light" states one should analyze the sensitivity to $\mathcal{O}(\Lambda^{-4})$ terms. The precisely measured low-energy observables are dominated by $\mathcal{O}(\Lambda^{-2})$ contributions of dimension-6 operators, whereas the quadratic terms in the Wilson coefficients, formally $\mathcal{O}(\Lambda^{-4})$, are negligible. In contrast, the one-by-one LHC constraints on 4-fermion operators in table 6 have in general a similar sensitivity to linear and quadratic terms.[17] Notice that this problem becomes much more severe in a global fit and that in the particular case of the chirality-violating operators there is no interference at all. This may undermine the SMEFT $1/\Lambda^2$ expansion for generic UV completions, and it is not clear whether the dimension-8 and higher operators can be neglected in the analysis. As discussed in ref. [108], in such a case the EFT is still valid for strongly coupled UV completions, where the dimension-6 squared terms are parametrically enhanced with respect to the dimension-8 contributions by a large new physics coupling. On the other hand, for weakly coupled UV completion one should use weaker LHC bounds obtained by truncating the $\sqrt{s}$ range of the analyzed data at some M_{cut} above which the SMEFT is no longer valid. For illustration, in table 4 we show the analogous LHC constraints with $M_{\mathrm{cut}} = 1\,\mathrm{TeV}$ ($\mathrm{LHC}_{1.0}$) and $M_{\mathrm{cut}} = 0.7\,\mathrm{TeV}$ ($\mathrm{LHC}_{0.7}$).

Another practical consequence of the quadratic terms domination at the LHC is that the likelihood for the Wilson coefficients is not approximately Gaussian. That means it is not fully characterized by the central values, $1\,\sigma$ errors, and the correlation matrix, as is the case for the low-energy observables. This makes the presentation of the global fit results more cumbersome.

Last, let us notice that the dilepton-production cross section is also sensitive to SMEFT coefficients that are flavor non-diagonal in the quark bilinear if we go beyond the $V = 1$ approximation at order Λ^{-2}. This was exploited in ref. [55] to set bounds on the Wilson coefficients of chirality-violating $\ell\ell 21$ operators.

6 Conclusions

This paper compiles information from a number of experiments sensitive to flavor-conserving LLQQ operators. The main focus is on experiments probing physics well below the weak scale, such as neutrino scattering on nucleon targets, atomic parity violation,

[17] In fact, in a few $\mathrm{LHC}_{0.7}$ entries in table 6 there is an additional (not shown) second solution far from the origin.

parity-violating electron scattering on nuclei, and so on. Information from e^+e^- collisions at the center-of-mass energies around the weak scale is also included. This is combined with previous analyses studying 4-lepton operators and the strength of the Z and W boson couplings to matter. The ensemble of data is interpreted as constraints on heavy new physics encoded in tree-level effects of dimension-6 operators in the SMEFT. The main strength of this analysis is that we allow all independent operators to be simultaneously present with an arbitrary flavor structure. Another novelty is the inclusion of low-energy flavor constraints from pion, neutron, and nuclear decays, recently summarized in ref. [55]. The leading renormalization group running effects from low energies to the weak scale are taken into account.

We obtain simultaneous constraints on 61 linear combinations of Wilson coefficients in the SMEFT. The results are presented as a multi-dimensional likelihood function, which is provided in a `Mathematica` notebook attached as supplemental material [56]. The likelihood can easily be projected onto more restricted new physics scenarios. As an illustration, we provide constraints on the SMEFT operators in the $\mathrm{U}(3)^5$-symmetric scenario, and on the oblique parameters S, T, W, Y. The likelihood can be used to place limits on masses and couplings in a large class of theories beyond the SM when the mapping between these theories and the SMEFT is known.

Finally, a brief comparison of the sensitivity of low-energy experiments to LLQQ operators with that of the LHC is provided. For many directions in the SMEFT parameters space, dilepton production at the LHC is exploring virgin territories not constrained by previous experiments. This is especially true for the chirality-conserving $2\mu 2q$ operators, where q are light quarks, while for the chirality-conserving $2e2q$ operators the LHC and low-energy probes are similarly sensitive. It would be beneficial to recast the LHC dilepton results in a model-independent form of a global likelihood on the SMEFT Wilson coefficients. We leave this task for future publications.

The SMEFT constraints summarized in this paper should be improved in the near future. Measurements of the differential Drell-Yan production cross sections at the LHC run-2 will provide a more powerful probe of LLQQ operators, thanks to the increased center-of-mass energy of the collisions and higher luminosities.[18] Progress is imminent on the low-energy front as well, e.g. thanks to more precise measurements of low-energy electron scattering in the Q-weak, MOLLER and P2 experiments. In this paper we have stressed the importance of probing new physics in multiple low- and high-energy experiments. The huge number of independent SMEFT operators requires a rich and diverse set of observables in order to lift flat directions in the global likelihood. In fact, several poorly or not-at-all constrained directions in the SMEFT parameter space persist, as is evident from eq. (4.8). This is especially true for operators involving the second and third generation quarks or the third generation leptons, but some flat directions involve the first generation fermions. The existence of these unexplored directions could be an inspiration to design new experiments and observables.

[18] As the recent recast of 13-TeV ATLAS data carried out in ref. [105] shows, this is already the case with the currently available luminosity ($36.1\,\mathrm{fb}^{-1}$). Expected bounds with $3000\,\mathrm{fb}^{-1}$ of data can also be found in that work.

JHEP08(2017)123

A Translation to Warsaw basis

In this paper we parametrize the relevant part of the space of dimension-6 operators using an independent set of vertex corrections δg and Wilson coefficients of 4-fermion operators. The latter are directly inherited from the Warsaw basis, such that the translation is trivial. The former are related to the Wilson coefficients of dimension-6 operators in the Warsaw basis by the following linear transformation:

$$\begin{aligned}
\delta g_L^{We} &= c_{H\ell}^{(3)} + f(1/2,0) - f(-1/2,-1), \\
\delta g_L^{Ze} &= -\frac{1}{2}c_{H\ell}^{(3)} - \frac{1}{2}c_{H\ell} + f(-1/2,-1), \\
\delta g_R^{Ze} &= -\frac{1}{2}c_{He} + f(0,-1), \\
\delta g_R^{Wq} &= -\frac{1}{2}c_{Hud}, \\
\delta g_L^{Zu} &= \frac{1}{2}c_{Hq}^{(3)} - \frac{1}{2}c_{Hq} + f(1/2,2/3), \\
\delta g_L^{Zd} &= -\frac{1}{2}V^\dagger c_{Hq}^{(3)} V - \frac{1}{2}V^\dagger c_{Hq} V + f(-1/2,-1/3), \\
\delta g_R^{Zu} &= -\frac{1}{2}c_{Hu} + f(0,2/3), \\
\delta g_R^{Zd} &= -\frac{1}{2}c_{Hd} + f(0,-1/3),
\end{aligned} \tag{A.1}$$

where

$$\begin{aligned}
f(T^3,Q) = &-I_3 Q \frac{g_L g_Y}{g_L^2 - g_Y^2} c_{HWB} \\
&+ I_3 \left(\frac{1}{4}[c_{\ell\ell}]_{1221} - \frac{1}{2}[c_{H\ell}^{(3)}]_{11} - \frac{1}{2}[c_{H\ell}^{(3)}]_{22} - \frac{1}{4}c_{HD}\right)\left(T^3 + Q\frac{g_Y^2}{g_L^2 - g_Y^2}\right),
\end{aligned} \tag{A.2}$$

and I_3 is the 3×3 identity matrix in the generation space. Using eq. (A.1) one can easily recast the results of this paper as a likelihood for the Wilson coefficients in the Warsaw basis. See ref. [59] for the dictionary between δg and the Wilson coefficients in the SILH basis.

B More general approach to low-energy flavor observables

The low-energy flavor observables discussed in ref. [55] also probe precisely 4-fermion operators with a strange quark. In the framework of the SMEFT the corresponding observables receive contributions from flavor off-diagonal dimension-6 operators, and in this paper we marginalized our likelihood over them. We also approximated the CKM matrix as $V = 1$ when acting on $\mathcal{O}(\Lambda^{-2})$ terms in the Lagrangian. For completeness, in this appendix we provide the formalism that allows one to take into account the constraints from strange observables and retrieve the terms suppressed by off-diagonal elements of the CKM matrix.

JHEP08(2017)123

First, the effective low-energy Lagrangian in eq. (2.4) is generalized to

$$
\begin{aligned}
\mathcal{L}_{\text{eff}} \supset - \sum_{I,J=1,2} \frac{2\tilde{V}_{uI}}{v^2} \Bigg[& \left(1+\bar{\epsilon}_L^{d_I e_J}\right)(\bar{e}_J \bar{\sigma}_\mu \nu_J)(\bar{u}\bar{\sigma}^\mu d_I) + \epsilon_R^{d_I e}(\bar{e}_J \bar{\sigma}_\mu \nu_J)(u^c \sigma^\mu \bar{d}_I^c) \\
& + \frac{\epsilon_S^{d_I e_J} + \epsilon_P^{d_I e_J}}{2}(e_J^c \nu_J)(u^c d_I) + \frac{\epsilon_S^{d_I e_J} - \epsilon_P^{d_I e_J}}{2}(e_J^c \nu_J)(\bar{u}\bar{d}_I^c) \\
& + \epsilon_T^{d_I e_J}(e_J^c \sigma_{\mu\nu}\nu_J)(u^c \sigma_{\mu\nu} d_I) + \text{h.c.} \Bigg],
\end{aligned} \tag{B.1}
$$

such that it also includes charged currents with the strange quark ($s \to u\ell\nu_\ell$). At tree level, the low-energy parameters are related to the SMEFT parameters as

$$
\begin{aligned}
\epsilon_R^{de} = -\bar{\epsilon}_L^{de} &= \frac{1}{V_{ud}} \delta g_R^{Wq_1}, \\
\epsilon_R^{se} = -\bar{\epsilon}_L^{se} &= \frac{1}{V_{us}} [\delta g_R^{Wq}]_{12}, \\
\bar{\epsilon}_L^{d\mu} &= -\frac{1}{V_{ud}} \delta g_R^{Wq_1} + \delta g_L^{W\mu} - \delta g_L^{We} + \left([c_{lq}^{(3)}]_{111J} - [c_{lq}^{(3)}]_{221J}\right) \frac{V_{Jd}}{V_{ud}}, \\
\bar{\epsilon}_L^{s\mu} &= -\frac{1}{V_{us}} [\delta g_R^{Wq}]_{12} + \delta g_L^{W\mu} - \delta g_L^{We} + \left([c_{lq}^{(3)}]_{111J} - [c_{lq}^{(3)}]_{221J}\right) \frac{V_{Js}}{V_{us}},
\end{aligned} \tag{B.2}
$$

$$
\begin{aligned}
\epsilon_S^{de_J} &= -\frac{1}{2V_{ud}} \left(V_{Kd} [c_{lequ}]^*_{JJK1} + [c_{ledq}]^*_{JJ11} \right), \\
\epsilon_P^{de_J} &= -\frac{1}{2V_{ud}} \left(V_{Kd} [c_{lequ}]^*_{JJK1} - [c_{ledq}]^*_{JJ11} \right), \\
\epsilon_S^{se_J} &= -\frac{1}{2V_{us}} \left(V_{Ks} [c_{lequ}]^*_{JJK1} + [c_{ledq}]^*_{JJ12} \right), \\
\epsilon_P^{se_J} &= -\frac{1}{2V_{us}} \left(V_{Ks} [c_{lequ}]^*_{JJK1} - [c_{ledq}]^*_{JJ12} \right), \\
\epsilon_T^{de_J} &= -\frac{V_{Kd}}{2V_{ud}} [c_{lequ}^{(3)}]^*_{JJK1}, \\
\epsilon_T^{se_J} &= -\frac{V_{Ks}}{2V_{us}} [c_{lequ}^{(3)}]^*_{JJK1}.
\end{aligned} \tag{B.3}
$$

In addition to $\tilde{V}_{ud}$ we also introduce the rescaled CKM matrix element parameter $\tilde{V}_{us}$. Both are distinct from the elements of the unitary matrix V, to which they are related by $V_{ud} = \tilde{V}_{ud}(1+\delta V_{ud})$, $V_{us} = \tilde{V}_{us}(1+\delta V_{us})$, where

$$
\begin{aligned}
\delta V_{ud} &= -\frac{1}{V_{ud}} \delta g_L^{Wq_1} - \frac{1}{V_{ud}} \delta g_R^{Wq_1} + \delta g_L^{W\mu} - \frac{1}{2}[c_{\ell\ell}]_{1221} + [c_{lq}^{(3)}]_{111J} \frac{V_{Jd}}{V_{ud}}, \\
\delta V_{us} &= -\frac{1}{V_{us}} [\delta g_L^{Wq}]_{12} - \frac{1}{V_{us}} [\delta g_R^{Wq}]_{12} + \delta g_L^{W\mu} - \frac{1}{2}[c_{\ell\ell}]_{1221} + [c_{lq}^{(3)}]_{111J} \frac{V_{Js}}{V_{us}}.
\end{aligned} \tag{B.4}
$$

The purpose of this rescaling is to impose the relation $\bar{\epsilon}_L^{d_I e} = -\epsilon_R^{d_I e}$ in eq. (B.1). After the rescaling, $\tilde{V}_{ud}$ and $\tilde{V}_{us}$ are no longer related by the standard unitarity equation. In the limit where the mixing with the 3rd generation is neglected we have $|\tilde{V}_{ud}|^2 + |\tilde{V}_{us}|^2 = 1 + \Delta_{\text{CKM}}$, where

$$
\begin{aligned}
\Delta_{\text{CKM}} &= -2V_{ud}\delta V_{ud} - 2V_{us}\delta V_{us} \\
&= 2V_{ud}\left(\delta g_L^{Wq_1} + \delta g_R^{Wq_1} - [c_{lq}^{(3)}]_{111J}V_{Jd}\right) \\
&\quad +2V_{us}\left([\delta g_L^{Wq}]_{12} + [\delta g_R^{Wq}]_{12} - [c_{lq}^{(3)}]_{111J}V_{Js}\right) \\
&\quad -2\delta g_L^{W\mu} + [c_{\ell\ell}]_{1221}.
\end{aligned}
\tag{B.5}
$$

As before, $\tilde{V}_{ud}$ may be affected by new physics contributing to ϵ_S^{de} and should be treated as a free parameter in the fit. Ref. [55] obtained the following constraints on the low-energy parameters

$$
\begin{pmatrix} \tilde{V}_{ud}^e \\ \Delta_{\text{CKM}} \\ \Delta_L^s \\ \Delta_{LP}^d \\ \epsilon_P^{de} \\ \epsilon_R^{de} \\ \epsilon_P^{se} \\ \epsilon_P^{s\mu} \\ \epsilon_R^s \\ \epsilon_S^{s\mu} \\ \epsilon_T^{s\mu} \\ \epsilon_S^{de} \\ \epsilon_T^{de} \\ \epsilon_S^{se} \\ \epsilon_T^{se} \end{pmatrix}
=
\begin{pmatrix} 0.97451 \pm 0.00038 \\ -1.2 \pm 8.4 \\ 1.0 \pm 2.5 \\ 1.9 \pm 3.8 \\ 4.0 \pm 7.8 \\ -1.3 \pm 1.7 \\ -0.4 \pm 2.1 \\ -0.7 \pm 4.3 \\ 0.1 \pm 5.0 \\ -3.9 \pm 4.9 \\ 0.5 \pm 5.2 \\ 1.4 \pm 1.3 \\ 1.0 \pm 8.0 \\ -1.6 \pm 3.3 \\ 0.9 \pm 1.8 \end{pmatrix}
\times 10^{\wedge}
\begin{pmatrix} 0 \\ -4 \\ -3 \\ -2 \\ -6 \\ -2 \\ -5 \\ -3 \\ -2 \\ -4 \\ -3 \\ -3 \\ -4 \\ -3 \\ -2 \end{pmatrix},
\tag{B.6}
$$

in the $\overline{MS}$ scheme at $\mu = 2\,\text{GeV}$. Here $\Delta_L^s = \bar{\epsilon}_L^{s\mu} - \bar{\epsilon}_L^{se}$ and $\Delta_{LP}^d \approx \bar{\epsilon}_L^{de} - \bar{\epsilon}_L^{d\mu} + 24\epsilon_P^{d\mu}$. The associated correlation matrix is given in ref. [55]. We note that some entries in this matrix are very close to one, so it is crucial to take it into account.

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JHEP08(2017)123

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JHEP08(2017)123

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JHEP08(2017)123