Probing the SM with Dijets at the LHC

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Abstract

The LHC has started to explore the TeV energy regime, probing the SM beyond LEP and Tevatron. We show how the dijet measurements at the LHC are able to test certain sectors of the SM at an unprecedented level. We provide the best bounds on all possible four-quark interactions and translate them into limits on the compositeness scale of the quarks and gluons. We also provide constraints on extra gauge bosons, $Z'$, $W'$ and $G'$, and on new interactions proposed to explain the present measurement of the forward-backward asymmetry of the top.
1 Introduction

The LHC is currently exploring the TeV-energy frontier, searching for new states beyond the Standard Model (SM) [1]. Although not a precision machine, the LHC is also sensitive to deformations of the SM that grow with energy, such as contact-interactions, allowing us then to test sectors of the SM at an unprecedented level.

Among the SM sectors that can be better tested at the LHC than in previous experiments is the quark sector. As we will see, the right-handed quark sector has been poorly tested at LEP and Tevatron, while its left-handed counterpart only electroweak symmetry breaking effects have been scrutinized. There is definitely room for improvement at the LHC.

The purpose of this paper is to use the dijet angular distributions at the LHC [2, 3, 4, 5] to improve our understanding of the quark sector of the SM. By using the observable $F_\chi$, we will put the best bounds on all possible four-quark interactions. These new contact-interactions are present in scenarios in which the quarks are composite states [6], a natural possibility that can arise in models where strong dynamics are expected to be around the TeV-scale, such as composite Higgs models [7, 8] or scenarios with low-energy supersymmetry breaking [9]. Dijets at the LHC provide then the best bounds on some of these scenarios. We will also show how our dijet analysis constrains new heavy gauge sectors beyond the SM, and the degree of compositeness of the SM gauge bosons, most significantly that of the gluon.

The article is organized as follows. We will start showing how well the different sectors of the SM have been tested in the pre-LHC era, and which type of deformations from the SM, parametrized by dimension-six operators, can be much better bounded by the LHC. In section 3 we present the calculation of the dijet angular distributions and the $F_\chi$ parameter that seems to be very well suited to obtain limits on four-quark interactions. By using the LHC data we will obtain bounds on all possible new four-quark interactions (see also Ref. [10]), and translate these bounds into limits on the scale of compositeness of the quarks, on possible new gauge sectors, on the scale of compositeness of the SM gauge bosons, and finally on new physics scenarios responsible for explaining the present experimental discrepancy in the forward-backward asymmetry $A_{FB}$ of the top [11] [12].

2 Tests of the SM sectors before LHC

Let us consider a new sector beyond the SM (BSM) whose physical scale, generically referred by $\Lambda$ (that, for example, can be associated with the mass of the new states), is assumed to be much larger than the momenta $p$ at which we are probing the SM, $\Lambda \gg p$. We can then parametrize the deviations from the SM by higher dimensional operators added to the SM Lagrangian [13]:

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}} + \frac{c_i}{\Lambda^2} \mathcal{O}_i^{d=6} + \cdots,$$  

(1)
where we only keep the dominant contributions corresponding to operators of dimension six, assuming lepton and baryon number conservation. Among these operators it is important to distinguish between two classes:

1. Operators involving extra powers of SM fields:

\[ (\bar{q}_L \gamma_\mu q_L)^2, \ (\bar{q}_L \gamma_\mu q_L)(H^1 D^\mu H), \ ... \]  

2. Operators involving extra (covariant) derivatives:

\[ \bar{q}_L \gamma_\mu q_L D_\nu F^{\mu \nu}, \ \bar{q}_L u_R D_\mu D^\mu H, \ ... \]  

The coefficients \( c_i \) in front of the first class of operators are parametrically proportional to the square of a coupling of the SM fields to the BSM sector, and then they can be as large as \( c_i \lesssim 16\pi^2 \). On the other hand, the coefficients \( c_i \) of the second class of operators should not contain couplings and are expected to be of order one \( c_i \lesssim O(1) \). This distinction is important when considering strongly-coupled BSM with part of the SM fields arising as composite states of this new sector. In this case \( \Lambda \) corresponds to the mass of the heavy resonances of the new strong sector whose couplings, referred as \( g_\rho \), can be as large as \( \sim 4\pi \). Hence operators of the first class with \( c_i \sim g_\rho^2 \) give generically more significant modifications to SM physics than those of the second class [6, 8, 14].

At present we have important constraints on \( c_i/\Lambda^2 \) coming from precision measurements of SM observables. Let us start considering those involving SM fermions. In the Appendix we give the full list of independent operators involving quarks. Neglecting fermion masses (chiral limit), we have that the impact of the dimension-six operators on SM physical processes can generically be parametrized by two new types of interactions:

\[ \frac{\alpha_\psi}{\Lambda^2} (\bar{\psi} \gamma_\mu \psi)^2 + \frac{\beta_\psi v^2}{\Lambda^2} A_\mu \bar{\psi} \gamma^\mu \psi, \]  

where we denote collectively by \( A_\mu = W_\mu, Z_\mu, \ldots \) the SM gauge bosons, by \( \psi = u_L, u_R, \ldots \) the SM fermions of a given chirality, \( v \simeq 246 \text{ GeV} \) is the Higgs vacuum expectation value (VEV), and \( \alpha_\psi \) and \( \beta_\psi \) measure the strength of the interactions. Since both types of interactions in Eq. (4) can arise from operators of the first class (Eq. (2)), one has \( 0 < |\alpha_\psi|, |\beta_\psi| \lesssim 16\pi^2 \). The first term of Eq. (4) gives contributions to four-fermion processes that scale as \( p^2/\Lambda^2 \) where \( p \) characterizes the momenta of the process. The second term corresponds to deviations from the SM gauge interactions at zero-momentum, and therefore these can only arise for the \( W \) and \( Z \) gauge boson and must be proportional to the Higgs VEV.

In principle, very stringent constraints on new interactions of the type of Eq. (4) arise from flavor physics [15]. It is not our purpose here to discuss them; they are very model dependent and
can be avoided if a flavor symmetry is imposed in the BSM sector. For example we can assume a
flavor symmetry for the three left-handed quarks $q_L$, the three right-handed down-quarks $d_R$, and
the two lightest right-handed up-quarks $u_R$, given by

$$G_F \equiv U(3)_q \otimes U(3)_d \otimes U(2)_u,$$

(5)

and similarly for the lepton sector. Due to the absence of important constraints on the flavor physics
for the right-handed top $t_R$, we can consider it a singlet of the flavor symmetry. This allows us
to treat the $t_R$ independently of the other quarks; its physical implications, some of them already
studied in Ref. [14], are left for a future publication. Yukawa couplings break the $G_F$ symmetry,
but it can be shown, by using a spurion’s power counting, that flavor constraints on dimension-six
operators can be easily satisfied for $\Lambda$ slightly above the electroweak scale [15]. From now on, we
will consider BSM that, up to Yukawa couplings, fulfill the flavor symmetry $G_F$.

At LEP the properties of the leptons $\psi = l_L, l_R$ were very well measured, putting bounds at the
per-mille level on deviations from the SM predictions either arising from vertex corrections or new
four-lepton contact interactions. From [16], one gets $\Lambda/|\alpha_{l,L,R}|, \Lambda/|\beta_{l,L,R}| \gtrsim 3 - 4$ TeV. This implies, for example, that the scale of compositeness of the leptons is larger than 40 – 50 TeV for
$\alpha_{l,L,R} \sim g^2_p \sim 16\pi^2$. Thus, the leptonic sector has been very well tested at LEP and recent LHC
data, having only quarks in the initial state, cannot provide better bounds.

For the left-handed quark sector $\psi = q_L$, there are very strong constraints on interactions of the
second type of Eq. (4). The most important ones arise from Kaon and $\beta$-decays [16] which have
allowed to measure very precisely quark-lepton universality of the $W$ interactions. This leads to
bounds on deviations from the $W$ coupling to left-handed quarks as strong as those for leptons,
$\Lambda/|\beta_{q,L}| \gtrsim 3 - 4$ TeV, which we do not expect to be improved substantially at the LHC.
Similar limits are obtained from measurements at LEP of the $Z$ decay to hadrons [16]. Bounds
on four-$q_L$ interactions are weaker, with the main constraint coming from Tevatron and giving
$\Lambda/|\alpha_{q,L}| \gtrsim 1$ TeV [16]. Clearly, the LHC can increase these bounds considerably as we will show
later. While theories of composite Higgs and composite $q_L$, where one expects large $\alpha_{q,L}$ and $\beta_{q,L}$
coefficients (since $\alpha_{q,L} \sim \beta_{q,L} \sim g^2_p \lesssim 16\pi^2)$ [6 8 14], are very constrained by present experimental
data, theories with only $q_L$ composite (and elementary Higgs, as those for example in Ref. [9]) where
only $\alpha_{q,L}$ is expected to grow with $g^2_p$, are not so constrained. LHC dijets can then, as we will see,
probe these scenarios at an unprecedented level.

Regarding right-handed quarks $u_R$ and $d_R$, their couplings to gauge bosons are still poorly
measured, due to their small coupling to $W$ and $Z$. For example, one of the best bounds, arising
from LEP, are on the $Z$ coupling to $b_R$ which reads $0 \lesssim \delta g_{b_R}/g_{b_R} \lesssim 0.2$ [17]. Furthermore these
vertices can be protected by symmetries of the BSM sector [18]. The strongest constraints on
$\alpha_{u_R,d_R}$ are again coming from Tevatron and, as for the left-handed case, LHC can improve them
significantly. Similar conclusions have been recently reached in Ref. [19].

For completeness, we comment on chirality-flip processes that are sensitive to fermion masses. For \( m_\psi = Y_\psi v \neq 0 \), two new types of interactions can be added to Eq. (4):

\[
\gamma_\psi \frac{m_\psi}{\Lambda^2} F^{\mu\nu} \bar{\psi} \sigma_{\mu\nu} \psi + \delta_\psi \frac{Y_\psi^2}{\Lambda^2} (\bar{\psi} \psi)^2, \tag{6}
\]

where \( \gamma_\psi \) and \( \delta_\psi \) are coefficients of order one. Concerning the first one, \( \text{Re}[\gamma_\psi] \) and \( \text{Im}[\gamma_\psi] \) give a contribution to the magnetic and electric dipole moment respectively. Only electric dipole moments give important constraints on BSM sectors, but they can be avoided by demanding CP-invariance in the BSM. Moreover, in most of the BSM they arise at the one-loop level. The second interaction in Eq. (6) corresponds to new four-fermion interactions, but suppressed with respect to those of Eq. (4) by Yukawa couplings, hence we will not consider them in this work.

Let us also mention that bounds on the interactions Eq. (4) can also constrain BSM contributions to the self-energies of the SM gauge bosons. These effects can be parametrized by five quantities \( \hat{S}, \hat{T}, W, Y \) and \( Z \) [20]. The first two, \( \hat{S}, \hat{T} \), are proportional to \( v^2/\Lambda^2 \) and find their best bound from LEP and Tevatron data. The \( W \) and \( Y \) parameters, that measure the compositeness of the \( W_\mu, Z_\mu \) and photons, are also bounded by LEP data at the per-mille level, but since these effects grow with the momenta as \( p^2/\Lambda^2 \), we can expect LHC to improve the bounds. Also at the LHC the best bound on the \( Z \)-parameter, that measures the compositeness of the gluon, can be obtained.

We then conclude that BSM physics generating four-quark interactions are not severely constrained by the pre-LHC data. Especially interesting BSM scenarios that contribute to this type of interactions are theories of composite quarks, either composite \( u_R \) and \( d_R \) (and Higgs), composite \( q_L \) (if the Higgs is elementary), or composite gluons. Below we will show how dijets at the LHC constrain these scenarios.

### 3 Dijets at the LHC

The study at the LHC of the angular distributions of dijets in the process \( pp \rightarrow jj \) has been shown to be a powerful tool to constrain the size of four-quark contact interactions [2, 3, 4, 5]. Here we will follow these analyses to put constraints on all possible four-quark interactions. Out of the complete list of dimension-six operators in \[A.1.1\] only those involving the first family of quarks, up and down, are relevant for our analysis. The reason is the following. In \( pp \rightarrow jj \) the dominant contributions at high dijet invariant-mass \( m_{jj} \) arise from valence-quarks initial states, \( i.e., uu, dd, du \), being \( uu \) or \( uc \) initial state processes very suppressed. For example, in the SM we have

\[
\left( \frac{\sigma(uu \rightarrow u\bar{u})}{\sigma(uu \rightarrow uu)} \right)_{SM}^{m_{jj}>2\text{TeV}} \simeq 0.04 , \quad \left( \frac{\sigma(uu \rightarrow uu)}{\sigma(uu \rightarrow uu)} \right)_{SM}^{m_{jj}>2\text{TeV}} \simeq 0.01 . \tag{7}
\]
Furthermore, processes with other families in the final states but having $u, d$ in the initial state, such as $uu \rightarrow ss, cc$, do not arise from the four-quark operators of Appendix A.1.1 due to the flavor symmetry $G_F$. We are then led to consider partonic processes involving only first family quarks, $uu \rightarrow uu$, $dd \rightarrow dd$ and $ud \rightarrow ud$, that allow us to reduce the set of operators of Eq. (36) to

$$
\mathcal{O}_{uu}^{(1)} = (\bar{u}_R \gamma^\mu u_R)(\bar{u}_R \gamma^\mu u_R) \\
\mathcal{O}_{dd}^{(1)} = (\bar{d}_R \gamma^\mu d_R)(\bar{d}_R \gamma^\mu d_R) \\
\mathcal{O}_{ud}^{(1)} = (\bar{u}_R \gamma^\mu u_R)(\bar{d}_R \gamma^\mu d_R) \\
\mathcal{O}_{ud}^{(8)} = (\bar{u}_R \gamma^\mu T^A u_R)(\bar{d}_R \gamma^\mu T^A d_R) \\
\mathcal{O}_{qq}^{(1)} = (\bar{q}_L \gamma^\mu q_L)(\bar{q}_L \gamma^\mu q_L) \\
\mathcal{O}_{qq}^{(8)} = (\bar{q}_L \gamma^\mu T^A q_L)(\bar{q}_L \gamma^\mu T^A q_L) \\
\mathcal{O}_{q_u}^{(1)} = (\bar{q}_L \gamma^\mu q_L)(\bar{u}_R \gamma^\mu u_R) \\
\mathcal{O}_{q_u}^{(8)} = (\bar{q}_L \gamma^\mu T^A q_L)(\bar{u}_R \gamma^\mu T^A u_R) \\
\mathcal{O}_{q_d}^{(1)} = (\bar{q}_L \gamma^\mu q_L)(\bar{d}_R \gamma^\mu d_R) \\
\mathcal{O}_{q_d}^{(8)} = (\bar{q}_L \gamma^\mu T^A q_L)(\bar{d}_R \gamma^\mu T^A d_R)
$$

(8)

where here we do not sum over flavor indices and from now on $q_L = (u_L, d_L)$. Apart from Eq. (8), there are other dimension-six operators (see the lists of Appendix A.1.2, A.2 and A.3) that can contribute to dijets. Nevertheless, these other operators are either suppressed by $v^2/p^2$ or Yukawa couplings with respect to those of Eq. (8), or can be rewritten, by use of equations of motion, as four-quark operators plus other operators not relevant for dijet physics. Therefore Eq. (8) exhausts the list of all leading operators contributing to dijets.

At the partonic level the SM differential cross section of $pp \rightarrow jj$ is dominated by QCD interactions [21]:

$$
\frac{s^2}{\pi \alpha_s^2} \frac{d\sigma}{dt} (q_i q_i \rightarrow q_i q_i)_{SM} = \frac{4}{9} \frac{s^2 + \hat{t}^2}{\hat{t}^2} + \frac{4}{9} \frac{s^2 + \hat{u}^2}{\hat{u}^2} - \frac{8}{27} \frac{s^2}{\hat{t} \hat{u}}, \quad (q_i = u, d)
$$

$$
\frac{s^2}{\pi \alpha_s^2} \frac{d\sigma}{dt} (ud \rightarrow ud)_{SM} = \frac{4}{9} \frac{s^2 + \hat{u}^2}{\hat{t}^2},
$$

$$
\frac{s^2}{\pi \alpha_s^2} \frac{d\sigma}{dt} (gq \rightarrow gq)_{SM} = (s^2 + \hat{u}^2) \left( \frac{1}{\hat{t}^2} - \frac{4}{9} \frac{1}{\hat{s} \hat{u}} \right),
$$

$$
\frac{s^2}{\pi \alpha_s^2} \frac{d\sigma}{dt} (gg \rightarrow gg)_{SM} = \frac{9}{2} \left( 3 - \frac{\hat{t} \hat{u}}{s^2} - \frac{\hat{s} \hat{u}}{\hat{t}^2} - \frac{\hat{t} \hat{u}}{\hat{u}^2} \right),
$$

$$
\frac{s^2}{\pi \alpha_s^2} \frac{d\sigma}{dt} (gg \rightarrow g_\tilde{q} \bar{g})_{SM} = \frac{3}{8} \left( \hat{t}^2 + \hat{u}^2 \right) \left( \frac{4}{9} \frac{1}{\hat{s} \hat{u}} - \frac{1}{\hat{s}^2} \right),
$$

(9)

where $s, t$ and $u$ are the partonic Mandelstam variables, and we are working in the massless quark
four-quark operators. They are then always suppressed by an extra \( \sim \) and therefore their coefficients in front are not parametrically larger than those of dimension-six where \( \Lambda \) is a characteristic scale.

Contributions from the operators of Eq. (8) give

\[
\frac{d\sigma}{dt}(q_iq_i \rightarrow q_iq_i)_{\text{BSM}} = -\frac{8\alpha_s}{27\Lambda^2} \left[ A_1^{qq} \frac{s}{t \bar{u}} - A_2^{qq} \left( \frac{\hat{u}^2}{t} + \frac{\hat{t}^2}{\bar{u}} \right) \frac{1}{s^2} \right] + \frac{4}{27\pi\Lambda^4} \left[ B_1^{qq} + B_2^{qq} \frac{\hat{u}^2 + \hat{t}^2}{s^2} \right],
\]

\[
\frac{d\sigma}{dt}(ud \rightarrow ud)_{\text{BSM}} = \frac{2\alpha_s}{9\Lambda^2} \left[ A_3 \frac{1}{\hat{t}} + A_4 \frac{\hat{u}^2}{s^2\hat{t}} \right] + \frac{1}{36\pi\Lambda^4} \left[ B_3 + B_4 \frac{\hat{u}^2}{s^2} \right],
\]

where

\[
A_1^{u,d} = \frac{1}{2} c_{qq}^{(8)} + \frac{3}{2} c_{uu,dd}^{(1)} + c_{qq}^{(1)},
\]

\[
A_2^{u,d} = \frac{3}{4} c_{qq,ud}^{(8)},
\]

\[
A_3 = \frac{1}{2} (2c_{qq}^{(8)} + c_{ud}^{(8)}),
\]

\[
A_4 = \frac{1}{2} (c_{qu}^{(8)} + c_{qd}^{(8)}),
\]

\[
B_1^{u,d} = (A_1^{u,d})^2 + \frac{1}{16} \left[ c_{qq}^{(8)} + 3(c_{qq}^{(1)} - c_{uu,dd}^{(1)}) \right]^2,
\]

\[
B_2^{u,d} = \frac{3}{16} (c_{qq,ud}^{(8)})^2 + \frac{27}{32} (c_{qu,qu}^{(1)})^2,
\]

\[
B_3 = A_3^2 + \frac{1}{4} (2c_{qq}^{(8)} - c_{ud}^{(8)})^2 + \frac{9}{8} (2c_{qq}^{(1)} + c_{ud}^{(1)})^2 + \frac{9}{8} (2c_{qq}^{(1)} - c_{ud}^{(1)})^2,
\]

\[
B_4 = A_4^2 + \frac{1}{4} (c_{qu}^{(8)} - c_{qd}^{(8)})^2 + \frac{9}{8} (c_{qu}^{(1)} + c_{qd}^{(1)})^2 + \frac{9}{8} (c_{qu}^{(1)} - c_{qd}^{(1)})^2,
\]

being the coefficients \( c_i \) defined according to Eq. (1). This extends the results of [22]. It is important to remark that in Eq. (10) we have included terms of order \( c_i^2/\Lambda^4 \); for \( c_i \gg 1 \) these terms are enhanced by an extra \( c_i \) factor with respect to the interference terms (of order \( c_i/\Lambda^2 \)), compensating for the \( \hat{s}/\Lambda^2 \) suppression factor. Therefore they should be considered in the calculations. Contributions from operators of dimension eight or larger are always smaller. For example, dimension-eight operators contributing to dijets involve at most four-fermions and extra derivatives, e.g. \( \partial^\nu \partial_\nu \bar{\psi} \gamma^\mu \psi \bar{\psi} \gamma_\mu \psi \), and therefore their coefficients in front are not parametrically larger than those of dimension-six four-quark operators. They are then always suppressed by an extra \( \sim \hat{s}/\Lambda^2 \).

As compared to the SM contribution Eq. (9), the BSM contribution Eq. (10) is enhanced at large \( \hat{s} \) and large CM scattering angle \( \theta^* \), or equivalently, for large (negative) \( \hat{t} = -\hat{s}(1 - \cos \theta^*)/2 \). It is convenient to define the angular variable \( \chi = (1 + |\cos \theta^*|)/(1 - |\cos \theta^*|) = -(1 + \hat{s}/\hat{t}) \in [1, +\infty) \) that can also be written as \( \chi = e^{y_1 - y_2} \) where \( y_{1,2} \) are the rapidity of the two jets. The QCD contribution to the differential cross section \( d\sigma(pp \rightarrow jj)/d\chi \) is almost flat in \( \chi \), while that of BSM grows for small values of \( \chi \), as can be appreciated in Fig. 1.
Figure 1: Dijet differential cross section as a function of $\chi$ for $m_{jj} > 2$ TeV at the LHC with $\sqrt{s} = 7$ TeV. The QCD contribution is shown in solid red line, while the green dashed line includes the contribution from the operator $O^{(8)}_{qq}$ with $c^{(8)}_{qq} = -0.5$ and $\Lambda = 1$ TeV.

3.1 The $F_\chi$ parameter

To put bounds on BSM four-quark operators, we will follow the method used by the ATLAS collaboration [2, 3]. This is based on the variable $F_\chi$ defined as the quotient of events with $1 < \chi < \chi_c \equiv 3.32$, the central region in the detector, over those with $1 < \chi < \chi_{max} \equiv 30$:

$$
F_{m_{jj}^{\text{cut}}}^{\chi} = \frac{N(\chi < \chi_c, m_{jj} > m_{jj}^{\text{cut}})}{N(\chi < \chi_{max}, m_{jj} > m_{jj}^{\text{cut}})},
$$

where $m_{jj}^{\text{cut}}$ is the cut over the invariant mass of the two-jet pair. Many systematic effects cancel in this ratio, providing an optimal test of QCD and a sensitive probe of hard new physics. It is also useful to write the analytic expression for this observable. Defining the integrated differential cross section over the angular region from 1 to $\chi_0$ as

$$
\sigma_{\chi_0}^{m_{jj}^{\text{cut}}} \equiv \int_1^{\chi_0} d\chi \frac{d\sigma}{d\chi} |_{pp \to jj},
$$

we have

$$
F_{m_{jj}^{\text{cut}}}^{\chi} \approx \frac{\sigma_{\chi_c}^{m_{jj}^{\text{cut}}}}{\sigma_{\chi_{max}}^{m_{jj}^{\text{cut}}}} \approx (\frac{\sigma_{\chi_c}^{m_{jj}^{\text{cut}}}}{\sigma_{\chi_{max}}^{m_{jj}^{\text{cut}}}})_{\text{SM}} + (\frac{\sigma_{\chi_c}^{m_{jj}^{\text{cut}}}}{\sigma_{\chi_{max}}^{m_{jj}^{\text{cut}}}})_{\text{BSM}},
$$

where we have split the contribution of the SM from that of the BSM, and considered that the SM contribution, being almost flat, dominates in the denominator. By making this approximation the deviation from the exact value of $F_\chi$ is of order 10%. Using Eqs. (9) and (10), and performing the integration over $\chi$, we obtain the result

$$
F_{m_{jj}^{\text{cut}}}^{\chi} \approx (F_{m_{jj}^{\text{cut}}}^{\chi})_{\text{SM}} - \frac{1}{\Lambda^2} \vec{A} \cdot \vec{P} + \frac{1}{\Lambda^4} \vec{B} \cdot \vec{Q},
$$
where

\[
\vec{A} = (A_1^u, A_2^u, A_1^d, A_2^d, A_3, A_4), \\
\vec{B} = (B_1^u, B_2^u, B_1^d, B_2^d, B_3, B_4),
\]

and

\[
\vec{P} \simeq (0.36P_{uu}, 0.12P_{uu}, 0.36P_{dd}, 0.12P_{dd}, 0.17P_{ud}, 0.074P_{ud}) \text{ TeV}^2, \\
\vec{Q} \simeq (0.013Q_{uu}, 0.0069Q_{uu}, 0.013Q_{dd}, 0.0069Q_{dd}, 0.0024Q_{ud}, 0.00097Q_{ud}) \text{ TeV}^4,
\]

where \((F_{XX}^{m_{jj}^{cut}})^{SM}\) is the SM value of \(F_{XX}^{m_{jj}^{cut}}\) and the coefficients \(P_{q_iq_j}\) and \(Q_{q_iq_j}\) encode the integration over the parton distribution functions (PDF):

\[
P_{q_iq_j} = \frac{1}{(\sigma_{\chi_{max}}^{m_{jj}^{cut}})^{SM}} \int d\tau \int dx f_{q_i}(x)f_{q_j}(\tau/x) \frac{\alpha_s(\tau s)}{x} + (i \leftrightarrow j \text{ for } i \neq j),
\]

\[
Q_{q_iq_j} = \frac{1}{(\sigma_{\chi_{max}}^{m_{jj}^{cut}})^{SM}} \int d\tau \int dx f_{q_i}(x)f_{q_j}(\tau/x) \frac{\tau s}{x} + (i \leftrightarrow j \text{ for } i \neq j),
\]

where \(\hat{s} = \tau s\) is the center of mass energy of the partons \(q_iq_j\) that initiate the collision, with \(\sqrt{s} = 7\) TeV the center of mass energy of the colliding protons. To calculate these coefficients we use MadGraph/MadEvent 4.4.57 [23] and implement the cuts taken in [3]. For this analysis we use the CTEQ6L1 PDF set and fix both the renormalization and factorization scales to \(m_{jj}^{cut} = 2\) TeV. We obtain:

\[
P_{uu} \simeq 0.23 \text{ TeV}^2, \quad P_{dd} \simeq 0.038 \text{ TeV}^2, \quad P_{ud} \simeq 0.28 \text{ TeV}^2, \\
Q_{uu} \simeq 23 \text{ TeV}^4, \quad Q_{dd} \simeq 3.8 \text{ TeV}^4, \quad Q_{ud} \simeq 19 \text{ TeV}^4,
\]

and \((F_{XX}^{2\text{TeV}})^{SM}\) \(\simeq 0.067\), \((\sigma_{\chi_{max}}^{2\text{TeV}})^{SM}\) \(\simeq 0.016\) TeV\(^{-2}\). We have checked the consistency of these results by numerically integrating over the MSTW2008 PDFs [24] using our analytical formulae for the cross sections Eqs. (9) and (10) and implementing the cuts in [3], which translate into the integration limits \(x \in [\sqrt{\tau e^{-y_B^{cut}}}, 1], \tau \in [(m_{jj}^{cut})^2/s, e^{-2y_B^{cut}}]\) and \(x \in [\tau, 1], \tau \in [e^{-2y_B^{cut}}, 1]\), where \(y_B^{cut} = 1.1\) is the cut on the rapidity boost of the partonic center of mass, \(|y_B| = \frac{1}{2}|y_1 + y_2| \leq y_B^{cut}\).

The variation of renormalization and factorization scales (by twice and half) introduces a theoretical uncertainty of the order of 10 – 15%. We have not computed the errors arising from the PDFs, and have not taken into account hadronization or showering effects since it is reasonable to neglect them for high dijet invariant masses [25].

4 Bounds

ATLAS has reported angular distributions of dijets for several \(m_{jj}\) [3]. We are interested in those with the largest invariant masses that correspond to \(m_{jj} > 2\) TeV and are given in Fig. 1. Using
Table 1: Bounds at 95\% CL on the scale suppressing the four-quark interactions. We denote by $\Lambda_{\pm}$ the bound on this scale obtained when taking the coefficient in front of the operator $c_i = \pm 1$, and considering the effects of the operators one by one.

| Operator | $\Lambda_{-}$ (TeV) | $\Lambda_{+}$ (TeV) |
|----------|---------------------|---------------------|
| $O^{(1)}_{uu}$ | 3.2 | 2.1 |
| $O^{(1)}_{dd}$ | 1.8 | 1.5 |
| $O^{(1)}_{ud}$ | 1.5 | 1.5 |
| $O^{(8)}_{ud}$ | 1.3 | 0.8 |
| $O^{(1)}_{qq}$ | 3.5 | 2.4 |
| $O^{(8)}_{qq}$ | 2.5 | 1.3 |
| $O^{(1)}_{qu}$ | 1.7 | 1.7 |
| $O^{(8)}_{qu}$ | 1.4 | 1.0 |
| $O^{(1)}_{qd}$ | 1.3 | 1.3 |
| $O^{(8)}_{qd}$ | 1.0 | 0.8 |

Using this value and the prediction Eq. (30) one can obtain bounds on the scales suppressing the operators Eq. (8). We instead derive the bounds using the coefficients of Eq. (19) but without making any approximation on the denominator of Eq. (14). The results are shown in Table 1. Few comments are in order. These bounds are obtained by taking the coefficient in front of the corresponding operator to be $c_i = \pm 1$. For other values we must rescale the bound by a factor $\sqrt{|c_i|}$. Since we are working in the approximation in which the energy of the physical process is assumed to be smaller than the masses of the BSM states $\Lambda$, we must require $\Lambda > m_{jj}^{cut}$. This implies that our bounds can only strictly be applied if $|c_i| > (2 \, \text{TeV}/\Lambda_{\pm})^2$. We recall that large values of $c_i$ are in principle possible since $c_i \lesssim 16\pi^2$. Also, we notice that for $c_i < 0$ the interference between the BSM contribution and the QCD contribution is constructive (with the exception of $O^{(1)}_{ud,qu,qd}$ where the interference is null), and as a consequence the bound is stronger than for a positive $c_i$.

These bounds are subject to a set of theoretical errors. The uncertainty in the parameters Eq. (19) estimated by changing the factorization and renormalization scales results in a $\sim 5\%$ uncertainty in the bounds. Also the NLO QCD correction to $(F_{\chi}^{2\text{TeV}})_{BSM}/(F_{\chi}^{2\text{TeV}})_{SM}$ has shown to be as large as $\sim 30\%$ [26], what amounts to a $\sim 10\%$ uncertainty in the bounds on $\Lambda$. Finally, it has been recently shown in [27] that electroweak corrections reduce the SM prediction of $F_{\chi}$ by a $\sim 2\%$ for large invariant masses $m_{jj} \sim 2\,\text{TeV}$. We therefore expect that our calculations for the bounds on $\Lambda$ can be trusted within a $\sim 10\%$ margin of error.
Table 2: Coefficients of the operators of Eq. (8) induced from integrating out heavy vector resonances for different composite quark scenarios. We have taken $\Lambda = m_\rho$.

| Composite | $c_{uu}^{(1)}/g_\rho^2$ | $c_{ud}^{(1)}/g_\rho^2$ | $c_{dd}^{(1)}/g_\rho^2$ | $c_{ud}^{(8)}/g_\rho^2$ | $c_{uq}^{(1)}/g_\rho^2$ | $c_{qu}^{(8)}/g_\rho^2$ | $c_{qd}^{(1)}/g_\rho^2$ | $c_{qd}^{(8)}/g_\rho^2$ |
|-----------|-------------------------|-------------------------|-------------------------|-------------------------|-------------------------|-------------------------|-------------------------|-------------------------|
| $u_R$     | $-37/72$                | 0                       | 0                       | 0                       | 0                       | 0                       | 0                       | 0                       |
| $d_R$     | 0                       | $-7/18$                 | 0                       | 0                       | 0                       | 0                       | 0                       | 0                       |
| $u_R, d_R$| $-37/72$                | $-7/18$                 | $2/9$                   | $-1$                    | 0                       | 0                       | 0                       | 0                       |
| $q_L$     | 0                       | 0                       | 0                       | $-5/36$                 | $-1$                    | 0                       | 0                       | 0                       |
| $q_L, u_R$| $-37/72$                | 0                       | 0                       | 0                       | $-5/36$                 | $-1$                    | $-1/9$                  | 0                       |
| $q_L, d_R$| 0                       | $-7/18$                 | 0                       | 0                       | $-5/36$                 | $-1$                    | 0                       | $1/18$                  |
| $q_L, u_R, d_R$| $-37/72$| $-7/18$| $2/9$| $-1$| $-5/36$| $-1$| $-1/9$| $1/18$| $-1$

4.1 Bounds on composite quarks

As we mentioned in section 2, previous experiments have not been able to probe the compositeness of quarks beyond the TeV scale. Data from dijets at the LHC can however improve this situation and put stronger constraints on their compositeness scale.

We will focus on models in which quarks arise as composite states of a strong sector whose global symmetry is $\mathcal{G} \equiv SU(3)_c \otimes SU(2)_L \otimes U(1)_Y \otimes G_F$, where $G_F$ is given in Eq. (5). In these theories we expect to have massive vector resonances associated to the current operators of $\mathcal{G}$, and then transforming in the adjoint representation of $\mathcal{G}$. This is in fact the case of the five-dimensional analogs based on the AdS/CFT correspondence [7]. Following Ref. [8], we will assume that all the vector resonances have equal masses and couplings, $m_\rho$ and $g_\rho$ respectively. Let us first consider the case in which only the right-handed up-type quarks $u_R$ are composite states, with charges under the global group $\mathcal{G}$ equal to $(3, 1, 2/3, 1, 1, 1)$. In this type of models, as we said before, the Higgs could also be composite without affecting our conclusions. Now, integrating out the heavy vector resonances at tree level (an approximation valid in the large-$N$ limit or, equivalently, $g_\rho \ll 4\pi$), we find that the four-quark operators of Eq. (8) are induced with coefficients given in Table 2 where we have fixed $\Lambda = m_\rho$. Constraints from dijets give $f \equiv m_\rho/g_\rho \gtrsim 2$ TeV. For $g_\rho \gg 1$, we see that this bound is stronger than that coming from the $S$-parameter that requires $f \gtrsim 4\pi v/g_\rho$ in theories of composite Higgs [8].

Similarly, we can assume a scenario where only the right-handed down-type quarks $d_R$ are composite with quantum numbers under $\mathcal{G}$ equal to $(3, 1, -1/3, 1, 3, 1)$. Again the coefficients of the four-quark operators induced are given in Table 2. We obtain the bound $f \gtrsim 1$ TeV. In the case of both $u_R$ and $d_R$ composite the bound goes up to $f \gtrsim 2.5$ TeV.

For composite left-handed quarks $q_L$ with $\mathcal{G}$-charges $(3, 2, 1/6, 3, 1, 1)$, the bound is $f \gtrsim 3$ TeV. Bounds on other composite quark scenarios are given in Table 3.

For weakly-coupled resonances ($g_\rho \lesssim 1$) with masses close to $m_{jj}^{cut}$ stronger bounds can be
Table 3: 95% CL bounds on the scale $f = m_\rho/g_\rho$ for different composite quark scenarios.

| Composite States | $f$ (TeV) |
|------------------|-----------|
| $d_R$            | 1.1       |
| $u_R$            | 2.3       |
| $u_R, d_R$       | 2.6       |
| $q_L$            | 2.7       |
| $q_L, d_R$       | 2.9       |
| $q_L, u_R$       | 3.5       |
| $q_L, u_R, d_R$  | 3.8       |

obtained from dijet resonance searches at the LHC [19]. This is just a consequence of the resonant enhancement of the cross section for a narrow region of invariant masses, where the resonances sit. This feature however is lost when the resonances are too broad.

4.2 Bounds on heavy gauge bosons

Heavy gauge bosons at the TeV-scale coupled to first family quarks generate four-quark operators that can be constrained by the dijet LHC data. Here we provide some examples. For a gauge boson $Z'$ gauging baryon number or hypercharge we obtain respectively

$$\frac{M_{Z'}}{g_B} \gtrsim 1.2 \text{ TeV}, \quad \frac{M_{Z'}}{g_Y} \gtrsim 1.6 \text{ TeV},$$

while for the gauge bosons $W'$ of a $SU(2)_R$ symmetry, where $q_R = (u_R, d_R)$ is assumed to transform as a doublet, we get

$$\frac{M_{W'}}{g_R} \gtrsim 1.6 \text{ TeV}.$$  \hfill (22)

As mentioned before, the fact that we work within an effective theory Eq. (1), implies that our bounds only apply to resonances with masses above $m_{jj}^{cut} = 2$ TeV. Gluonic resonances $G'^\mu_A$ coupled to first family quarks as $\mathcal{L}_{int} = G'^\mu_A \left[ g_L \bar{q}_L T^A \gamma^\mu q_L + g_R \bar{q}_R T^A \gamma^\mu q_R \right]$, with $T_A = \lambda_A/2$ where $\lambda_A$ are the Gell-Mann matrices, can also be constrained. This kind of resonances have been recently advocated (see for example Ref. [28]) to accommodate the discrepancy in the top forward-backward asymmetry measured at Tevatron. In Fig. 2 we show the excluded region of the parameter space. It can be seen that, for a resonance of mass $M_{G'} = 2.5$ TeV, the allowed range for the couplings is $-1.5 \lesssim g_{L,R} \lesssim 1.5$ at 95% CL. Similar bounds have been also obtained in Ref. [29].

4.3 Bounds on oblique parameters $Y$, $W$ and $Z$

The electroweak precision parameters $Y$, $W$ and $Z$ [20] can be regarded as a measure of the compositeness of the transversal components of the $SU(2)_L$, $U(1)_Y$, and $SU(3)_c$ gauge bosons respectively.
They manifest themselves as deviations of the self-energies of such vector bosons, and can be parametrized by the following higher dimensional operators:

\[-Y \frac{Y}{4m_W^2} (\partial_\rho B_{\mu\nu})^2, \quad -W \frac{W}{4m_W^2} (D_\rho W_I^J)^2, \quad -Z \frac{Z}{4m_W^2} (D_\rho G^A_{\mu\nu})^2.\]  

(23)

At large momenta as compared to the masses of the gauge bosons, these operators induce effective four-fermion operators, equivalent to those arising from integrating out a very heavy copy of the corresponding gauge boson. Therefore our dijet analysis can be conveniently used to put bounds on these parameters. We show in Fig. 3 our results in the $W$-$Y$ plane. Although bounds from LEP [20] are still stronger, this analysis shows that LHC will be competitive when running at a higher energy. Regarding the $Z$-parameter our analysis gives the strongest bound up to date:

\[-3 \times 10^{-3} \lesssim Z \lesssim 6 \times 10^{-4}.\]  

(24)

4.4 Bounds on new interactions for the $A_{FB}$ of the top

The recent discrepancy between the measured $A_{FB}$ of the top and its SM prediction [11, 12] has boosted the search for BSM that could explain it. Dijet angular distributions can be useful to constrain these models. As an example, we consider the proposal of Refs. [30, 31] where the measured value of the top asymmetry was explained by the following new interaction:

\[\mathcal{L}_{eff} = \frac{c_A^{(8)}}{\Lambda^2} O_A^{(8)} = \frac{c_A^{(8)}}{\Lambda^2} (\bar{u} T^A \gamma_\mu \gamma^5 u)(\bar{t} T^A \gamma_\mu \gamma^5 t).\]  

(25)
In terms of chirality eigenstates the operator $O^{(8)}_A$ reads
\[
O^{(8)}_A = (\bar{u}_R \gamma^\mu T^A u_R)(\bar{t}_R \gamma_\mu T^A t_R) - (\bar{u}_L \gamma^\mu T^A u_L)(\bar{t}_R \gamma_\mu T^A t_R) \\
- (\bar{u}_R \gamma^\mu T^A u_R)(\bar{t}_L \gamma_\mu T^A t_L) + (\bar{u}_L \gamma^\mu T^A u_L)(\bar{t}_L \gamma_\mu T^A t_L). \tag{26}
\]
If these operators arise from BSM that are invariant under the SM gauge group and $G_F$ (up to small effects $v^2/\Lambda^2$ and Yukawa couplings), the presence of $c^{(8)}_A \neq 0$ requires, in the basis A.1.1,
\[
c^{(8)}_{ut} = -c^{(8)}_{qt} = -c^{(8)}_{qu} = c^{(8)}_{qq} = c^{(8)}_A. \tag{27}
\]
In other words, the flavor symmetry requires that if the operator $O^{(8)}_A$ is generated, also operators involving four up-quarks must be present. Bounds from our dijet analysis (mostly from the bounds on $c^{(8)}_{qu}$ and $c^{(8)}_{qq}$) lead then to
\[
\frac{c^{(8)}_A}{\Lambda^2} \lesssim \frac{0.4}{\text{TeV}^2}, \tag{28}
\]
excluding the possibility to fit the recent top asymmetry measurement which requires $c^{(8)}_A/\Lambda^2 \sim 2 \text{ TeV}^{-2}$ \[31\].

If we relax the assumption of flavor invariance of the BSM sector, an operator involving four up-quarks can still be generated from $O^{(8)}_A$ at the one-loop level. The one-loop contribution, involving tops, is divergent and therefore sensitive to physics at the BSM scale $\Lambda$. We can get an estimate by regulating the divergence with a hard cut-off taken to be $\Lambda$. We obtain
\[
c^{(8)}_{qq} = \frac{1}{3} c^{(8)}_{uu} = -2 c^{(8)}_{qu} \simeq -\frac{(c^{(8)}_A)^2}{4\pi^2}. \tag{29}
\]
In Fig. 4 we show the region of the parameter space that fits the $A_{FB}$ of the top with the region excluded by dijets using Eq. (29). One can see that dijets with $m_{jj} > 2$ TeV exclude a large region of the parameter space, although, as we mentioned before, these results cannot be strictly applied if $\Lambda < m_{jj}$. For this reason we also show the exclusion region arising from dijets with smaller invariant masses, $m_{jj} > 1.2$ TeV.

5 Conclusions

We have shown that the present dijet LHC data is already testing all the quark sector of the SM at almost the same level of accuracy as leptons were tested at LEP after a decade of data collection. This is due to the very high-energy of the dijet scattering at the LHC that enhances the effects from four-quark interactions. We have used the $F_{\chi}$ parameter, defined in Eq. (12), to put bounds on all possible four-quark operators. Our results, presented in Table 1, show bounds on the scales suppressing these operators $\Lambda / \sqrt{|c_i|}$ ranging between $1 - 3$ TeV.

Among the most interesting BSM scenarios to be tested by dijet angular distributions are theories in which strong dynamics are postulated to solve the hierarchy problem. These theories demand a composite scale $\Lambda$ around the TeV-scale. We have seen that if the SM quarks are composite states arising from a new strong sector, $\Lambda \gtrsim 50$ TeV $\times (g_{\rho}/4\pi)$. Other possibilities are also significantly constrained, as can be seen from Table 3. We also derive the best bounds on the $Z$-parameter, Eq. (24), that measures the degree of compositeness of the gluons.

We also show that extra gauge bosons with sizable couplings to quarks are constrained to lie
above the TeV scale, limiting then their possible contribution to the $A_{FB}$ of the top.

Finally, we would like to stress that these results are based on the 2010 LHC data corresponding to 36 pb$^{-1}$ of integrated luminosity \cite{3}. It is expected that the 2011 LHC data set, containing more luminosity, will significantly improve all the bounds derived throughout this analysis.

**Note Added:** The 2011 data set for dijet events at CMS has been recently reported in Ref. \cite{33}, corresponding to a luminosity of 2.2 fb$^{-1}$ and a cut in the dijet invariant mass of $m_{jj} > 3$ TeV. The analysis made throughout this article can be applied in this case; by doing so we obtain more stringent bounds in all our results.

In this new analysis we define the observable $F_{\chi}$ as the quotient of events with $1 < \chi < \chi_c \equiv 3$, the first bin in the experimental analysis, over those with $1 < \chi < \chi_{max} \equiv 16$:

$$F_{\chi}^{m_{jj}} = \frac{N(\chi < \chi_c, m_{jj} > m_{jj}^{cut})}{N(\chi < \chi_{max}, m_{jj} > m_{jj}^{cut})}. \quad (30)$$

The experimental value for this observable is $F_{\chi}^{(m_{jj}>3 \text{ TeV})} \simeq 0.09$ with a 2$\sigma$ interval,

$$0.003 \lesssim F_{\chi}^{(m_{jj}>3 \text{ TeV})} \lesssim 0.15 \quad \text{at 95\% C.L.}, \quad (31)$$

while the SM prediction is $F_{\chi}^{(m_{jj}>3 \text{ TeV})} \simeq 0.12$. The new definition of $F_{\chi}$ together with the new cut in the dijet invariant mass will modify the numerical results in Sec. 3.1 in the following way: first of all the values in Eq. (17) and Eq. (19) have to be replaced by the following ones,

$$\vec{P} \simeq \frac{1}{(\sigma_{3 \text{ TeV}}^{\chi_{max}})^{\lambda_{SM}}}(0.33P_{uu}, 0.10P_{uu}, 0.33P_{dd}, 0.10P_{dd}, 0.15P_{ud}, 0.064P_{ud}) \text{ TeV}^2,$$

$$\vec{Q} \simeq \frac{1}{(\sigma_{3 \text{ TeV}}^{\chi_{max}})^{\lambda_{SM}}}(0.012Q_{uu}, 0.0064Q_{uu}, 0.012Q_{dd}, 0.0064Q_{dd}, 0.0022Q_{ud}, 0.00087Q_{ud}) \text{ TeV}^4,$$

and

$$P_{uu} \simeq 0.013 \quad , \quad P_{dd} \simeq 0.0019 \quad , \quad P_{ud} \simeq 0.015 ,$$

$$Q_{uu} \simeq 2.8 \text{ TeV}^2 \quad , \quad Q_{dd} \simeq 0.37 \text{ TeV}^2 \quad , \quad Q_{ud} \simeq 2.5 \text{ TeV}^2,$$

where $(\sigma_{3 \text{ TeV}}^{\chi_{max}})^{\lambda_{SM}} \simeq 0.0131 \text{ TeV}^{-2}$.

From these values we can obtain the results corresponding to the new data set. All the bounds derived in Sec. 4 have to be replaced by the ones given in this added note. Concerning the 4-quark operators the results are shown in Table 4, while for the composite-quark scenarios the updated bounds are given in Table 5.

The results given in Sec. 4.2 have to be replaced by the following ones,

$$\frac{M_{W'}}{g_R} \gtrsim 2.3 \text{ TeV} \quad , \quad \frac{M_{Z_B}}{g_B} \gtrsim 1.6 \text{ TeV} \quad , \quad \frac{M_{Z_Y}}{g_Y} \gtrsim 2.3 \text{ TeV \quad \text{ at 95\% C.L.}}, \quad (33)$$
Table 4: Bounds at 95% CL on the scale suppressing the four-quark interactions corresponding to the 2011 dijet data set by CMS. We denote by $\Lambda_{\pm}$ the bound on this scale obtained when taking the coefficient in front of the operator $c_i = \pm 1$, and considering the effects of the operators one by one.

| Operator $\mathcal{O}$ | $\Lambda_-/\sqrt{c_i}$ | $\Lambda_+/\sqrt{c_i}$ (TeV) |
|------------------------|------------------------|-------------------------------|
| $O_{uu}^{(1)}$         | 4.5                    | 3.0                           |
| $O_{dd}^{(1)}$         | 2.4                    | 2.0                           |
| $O_{ud}^{(1)}$         | 2.2                    | 2.2                           |
| $O_{ud}^{(8)}$         | 1.8                    | 1.3                           |
| $O_{qq}^{(1)}$         | 5.0                    | 3.5                           |
| $O_{qq}^{(8)}$         | 3.4                    | 2.0                           |
| $O_{qu}^{(1)}$         | 2.5                    | 2.5                           |
| $O_{qu}^{(8)}$         | 1.9                    | 1.5                           |
| $O_{qd}^{(1)}$         | 1.9                    | 1.9                           |
| $O_{qd}^{(8)}$         | 1.4                    | 1.2                           |

Table 5: 95% CL bounds on the scale $f = m_{\rho}/g_{\rho}$ for different composite quark scenarios corresponding to the 2011 dijet data set by CMS.

| Composite States | $f$ (TeV) |
|------------------|-----------|
| $d_R$            | 1.5       |
| $u_R$            | 3.2       |
| $u_R \cdot d_R$  | 3.6       |
| $q_L$            | 3.8       |
| $q_L \cdot d_R$  | 4.0       |
| $q_L, u_R$       | 4.9       |
| $q_L, u_R, d_R$  | 5.2       |
including Fig. 5 for a gluonic resonance coupled to both LH and RH quarks.

In the case of Sec. 4.3 the new bound for the $Z$ parameter is,

$$-9 \times 10^{-4} \lesssim Z \lesssim 3 \times 10^{-4}.$$  \hspace{1cm} (34)

While the bounds for the $W$ and the $Y$ parameters are shown in Fig. 6.

Finally, concerning Sec. 4.4 and the flavor invariant case we have to replace Eq. (28) by,

$$\frac{e_A^{(8)}}{\Lambda^2} \lesssim \frac{0.2}{\text{TeV}^2},$$  \hspace{1cm} (35)
while for the case that there is no flavor invariance the best bounds on the parameter space are
given by the previous analysis, since with the current data we can just exclude masses above 3 TeV
in the parameter space.

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A Dimension six operators involving quarks

Here we list the set of independent higher-dimensional operators involving SM quarks. As explained
in section 2 we assume a flavor symmetry for the three left-handed quarks $q_L$, the three right-handed
down-quarks $d_R$, and the two lightest right-handed up-quarks $u_R$, given by $U(3)_q \otimes U(3)_d \otimes U(2)_u$. The top right-handed quark $t_R$ will be considered a singlet of the flavor symmetry. We use the
following notation. We label with $A$, $I$ and $F$ the color, electroweak and flavor index respectively
in the adjoint representation. The contraction of the indices in the fundamental representation
of these symmetries is understood within the fields in parenthesis, and flavor indices can also be
contracted with Yukawa matrices $Y_{u,d}$. We identify $T_A = \lambda_A/2$, being $\lambda_A$ the Gell-Mann matrices,
and $\tau_I = \sigma_I/2$, where $\sigma_I$ are the Pauli matrices.

We classify the operators according to their expected suppression. First, we show the list of
independent operators unsuppressed by Yukawa couplings (those generated in the massless quark
limit). Following the discussion of section 2 we separate these operators as those of first class and
second class, Eq. (2) and Eq. (3) respectively. We finally show the list of independent operators
suppressed by Yuwaka couplings.
A.1 First class operators

A.1.1 Four-quark operators

\[
\begin{align*}
\mathcal{O}_{dd}^{(1)} &= (\bar{d}_R \gamma^\mu d_R)(\bar{d}_R \gamma_\mu d_R) \\
\mathcal{O}_{ud}^{(1)} &= (\bar{u}_R \gamma^\mu u_R)(\bar{d}_R \gamma_\mu d_R) \\
\mathcal{O}_{uu}^{(1)} &= (\bar{u}_R \gamma^\mu u_R)(\bar{u}_R \gamma_\mu u_R) \\
\mathcal{O}_{qd}^{(1)} &= (\bar{q}_L \gamma^\mu q_L)(\bar{d}_R \gamma_\mu d_R) \\
\mathcal{O}_{qq}^{(1)} &= (\bar{q}_L \gamma^\mu q_L)(\bar{q}_L \gamma_\mu q_L) \\
\mathcal{O}_{qq}^{(3W)} &= (\bar{q}_L \gamma^\mu q_L)(\bar{q}_L \gamma_\mu q_L) \\
\mathcal{O}_{qq}^{(8P)} &= (\bar{q}_L \gamma^\mu q_L)(\bar{q}_L \gamma_\mu T^P q_L) \\
\mathcal{O}_{uu}^{(8)} &= (\bar{u}_R \gamma^\mu T^A u_R)(\bar{u}_R \gamma_\mu T^A u_R) \\
\mathcal{O}_{ud}^{(8)} &= (\bar{d}_R \gamma^\mu T^A d_R)(\bar{d}_R \gamma_\mu T^A d_R) \\
\mathcal{O}_{ud}^{(8C)} &= (\bar{q}_L \gamma^\mu T^A q_L)(\bar{q}_L \gamma_\mu T^A q_L) \\
\mathcal{O}_{qq}^{(8)} &= (\bar{q}_L \gamma^\mu q_L)(\bar{q}_L \gamma_\mu T^A q_L) \\
\mathcal{O}_{eq}^{(8)} &= (\bar{q}_L \gamma^\mu T^A q_L)(\bar{q}_L \gamma_\mu T^A q_L) \\
\mathcal{O}_{eq}^{(3W)} &= (\bar{q}_L \gamma^\mu q_L)(\bar{q}_L \gamma_\mu q_L) \\
\mathcal{O}_{eq}^{(8P)} &= (\bar{q}_L \gamma^\mu q_L)(\bar{q}_L \gamma_\mu T^P q_L) \\
\mathcal{O}_{eq}^{(8C)} &= (\bar{q}_L \gamma^\mu q_L)(\bar{q}_L \gamma_\mu T^C q_L)
\end{align*}
\]

For physics involving only the first family quarks, that as explained in section 3 mainly corresponds to the LHC dijet data \( pp \rightarrow jj \), we can reduce the above set of four-quark operators to the set of Eq. (8). In this reduction, we have

\[
\begin{align*}
c_{uu}^{(1)} + \frac{1}{3} c_{uu}^{(8)} &\rightarrow c_{uu}^{(1)} \\
c_{dd}^{(1)} + \frac{1}{3} c_{dd}^{(8)} &\rightarrow c_{dd}^{(1)} \\
c_{qq}^{(1)} - \frac{1}{12} c_{qq}^{(3W)} + \frac{1}{3} c_{qq}^{(8P)} &\rightarrow c_{qq}^{(1)} \\
c_{qq}^{(8C)} + c_{qq}^{(3W)} &\rightarrow c_{qq}^{(8)}
\end{align*}
\]

A.1.2 Higgs-quark operators

\[
\begin{align*}
\mathcal{O}_{Hu} &= i(H^\dagger D^\mu H)(\bar{u}_R \gamma_\mu u_R) \\
\mathcal{O}_{Hd} &= i(H^\dagger D^\mu H)(\bar{d}_R \gamma_\mu d_R) \\
\mathcal{O}_{Hq}^{(1)} &= i(H^\dagger D^\mu H)(\bar{q}_L \gamma_\mu q_L) \\
\mathcal{O}_{Hq}^{(3)} &= i(H^\dagger \tau^i D^\mu H)(\bar{q}_L \gamma_\mu \tau^i q_L)
\end{align*}
\]

\[(38)\]
Notice that we just consider the antisymmetric part of the corresponding operators in [13]. This is so because the symmetric part of the above operators can be put in the form \( \partial_{\mu}(|H|^2)(\bar{\psi}\gamma^{\mu}\psi) \), where \( \psi = (q_L, d_R, u_R, t_R) \). It can be shown that this kind of operators can be expressed in terms of operators appearing in A.3.

### A.2 Second class operators

\[
\begin{align*}
O_{uB} &= (\bar{u}_R\gamma^\mu u_R)\partial^\nu B_{\mu\nu} & O_{tB} &= (\bar{t}_R\gamma^\mu t_R)\partial^\nu B_{\mu\nu} \\
O_{dB} &= (\bar{d}_R\gamma^\mu d_R)\partial^\nu B_{\mu\nu} & O_{qB} &= (\bar{q}_L\gamma^\mu q_L)\partial^\nu B_{\mu\nu} \\
O_{qW} &= (\bar{q}_L\tau^I\gamma^\mu q_L)D^\nu W^{\mu I} & O_{qW} &= (\bar{q}_L\tau^I\gamma^\mu q_L)D^\nu W^{\mu I} \\
O_{tB} &= (\bar{t}_R\gamma^\mu t_R)\partial^\nu B_{\mu\nu} & O_{tB} &= (\bar{t}_R\gamma^\mu t_R)\partial^\nu B_{\mu\nu} \\
O_{dH} &= (\bar{q}_L\gamma^\mu q_L)\partial^\nu H^{\mu I} & O_{dH} &= (\bar{q}_L\gamma^\mu q_L)\partial^\nu H^{\mu I} \\
O_{dB} &= (\bar{d}_L\gamma^\mu d_L)\partial^\nu B_{\mu\nu} & O_{dB} &= (\bar{d}_L\gamma^\mu d_L)\partial^\nu B_{\mu\nu} \\
O_{dW} &= (\bar{d}_L\gamma^\mu d_L)\partial^\nu H^{\mu I} & O_{dW} &= (\bar{d}_L\gamma^\mu d_L)\partial^\nu H^{\mu I} \\
O_{dW} &= (\bar{d}_L\gamma^\mu d_L)\partial^\nu B_{\mu\nu} & O_{dW} &= (\bar{d}_L\gamma^\mu d_L)\partial^\nu B_{\mu\nu}
\end{align*}
\]

As explained in [32] this kind of operators can be rewritten, by using the equations of motion (EOM) for the field strengths, as four-fermion operators involving quarks and leptons. Also notice that operators involving gluons are not included since they can be rewritten as four-quark operators.

### A.3 Yukawa-suppressed operators

The Yukawa couplings break the flavor symmetry and generate extra dimension-six operators. Assigning to the \( 3 \times 3 \) Yukawa matrices, \( Y_d \) and \( Y_u \), the following quantum numbers under \( G_F \): \( Y_d \in (\mathbf{3}, \overline{\mathbf{3}}, \mathbf{1}), \tilde{Y}_u \equiv (Y_u)_{ik} \in (\mathbf{3}, \mathbf{1}, 2) \) (i = 1, 2, 3; \( \tilde{k} = 1, 2 \)) and \( \tilde{Y}_t \equiv (Y_u)_{i3} \in (\mathbf{3}, \mathbf{1}, 1) \), we can write the following \( G_F \)-invariant operators \((\bar{H} = i\sigma_2 H^*)\):

\[ (i) \]

\[
\begin{align*}
O_{uH} &= (H^\dagger H)(\bar{q}_L\tilde{Y}_u\tilde{H}u_R) & O_{tH} &= (H^\dagger H)(\bar{q}_L\tilde{Y}_t\tilde{H}t_R) \\
O_{dH} &= (H^\dagger H)(\bar{q}_LY_dHd_R)
\end{align*}
\]

\[ (ii) \]

\[
\begin{align*}
O_{uGH} &= (\bar{q}_L\tilde{Y}_u\sigma^{\mu\nu}T^Au_R)\tilde{H}G^A_{\mu\nu} & O_{tGH} &= (\bar{q}_L\tilde{Y}_t\sigma^{\mu\nu}T^At_R)\tilde{H}G^A_{\mu\nu} \\
O_{uWH} &= (\bar{q}_L\tilde{Y}_u\sigma^{\mu\nu}\tau^Iu_R)\tilde{H}W^I_{\mu\nu} & O_{tWH} &= (\bar{q}_L\tilde{Y}_t\sigma^{\mu\nu}\tau^It_R)\tilde{H}W^I_{\mu\nu} \\
O_{uBH} &= (\bar{q}_L\tilde{Y}_u\sigma^{\mu\nu}u_R)\tilde{H}B^A_{\mu\nu} & O_{tBH} &= (\bar{q}_L\tilde{Y}_t\sigma^{\mu\nu}t_R)\tilde{H}B^A_{\mu\nu} \\
O_{dGH} &= (\bar{q}_LY_d\sigma^{\mu\nu}T^Ad_R)HG^A_{\mu\nu} & O_{dWH} &= (\bar{q}_LY_d\sigma^{\mu\nu}\tau^Id_R)HW^I_{\mu\nu} \\
O_{dBH} &= (\bar{q}_LY_d\sigma^{\mu\nu}d_R)HB_{\mu\nu}
\end{align*}
\]
(iii)\[ \mathcal{O}_{uDH} = (\bar{q}_L \tilde{Y}_u u_R) D_\mu D^\mu \tilde{H} \quad \mathcal{O}_{tDH} = (\bar{q}_L \tilde{Y}_t t_R) D_\mu D^\mu \tilde{H} \]
\[ \mathcal{O}_{dDH} = (\bar{q}_L \tilde{Y}_d d_R) D_\mu D^\mu H \quad \mathcal{O}_{tDH} = (\bar{q}_L \tilde{Y}_t t_R) D_\mu D^\mu \tilde{H} \]

By applying the EOM of $H$ these operators could be rewritten as other operators of the list plus four-fermion operators involving leptons.

(iv)\[ \mathcal{O}_{Hud} = i(\bar{H}^T \tilde{D}_\mu H)(\bar{u}_R \tilde{Y}_u u_L) \quad \mathcal{O}_{Htd} = i(\bar{H}^T \tilde{D}_\mu H)(\bar{t}_R \tilde{Y}_t t_L) \]

(v)\[ \mathcal{O}_{qud}^{(1)} = (\bar{q}_L \tilde{Y}_u u_R)(\bar{q}_L \tilde{Y}_d d_R) \quad \mathcal{O}_{qtd}^{(1)} = (\bar{q}_L \tilde{Y}_t t_R)(\bar{q}_L \tilde{Y}_d d_R) \]
\[ \mathcal{O}_{qud}^{(8)} = (\bar{q}_L \tilde{Y}_u T^A u_R)(\bar{q}_L \tilde{Y}_d T^A d_R) \quad \mathcal{O}_{qtd}^{(8)} = (\bar{q}_L \tilde{Y}_t T^A t_R)(\bar{q}_L \tilde{Y}_d T^A d_R) \]

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