Z-pole observables in an effective theory

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(Received 25 June 2013; published 6 January 2014)

There are two Z-peak observables related to the pair production of bottom quarks that show a deviation of about 2.5σ each from Standard Model expectations. While the discrepancy in the forward-backward asymmetry is a long-standing one, the tension for the second observable, namely the ratio of the partial width for a Z decaying to a pair of bottom quarks to the total hadronic decay width of the Z, has recently gone up due to a full two-loop evaluation of the Standard Model contributions. We show how both these discrepancies may be explained in the framework of new physics that couples only to the third generation of quarks. In the paradigm of effective operators, the Wilson coefficients of some of the possible operators are already very tightly constrained by flavor physics data. However, there still remain certain operators, particularly those involving right-chiral quark fields, which can successfully explain the anomalies. We also show how the footprints of such operators may be observed at the upgraded LHC.

DOI: 10.1103/PhysRevD.89.013002 PACS numbers: 13.38.Dg, 13.66.Jn, 13.85.Hd

I. INTRODUCTION

The majority of the electroweak precision observables are in good agreement with the Standard Model (SM) [1]. However, there are two which show a marked tension, albeit not at the level where they can be claimed as incontrovertible evidence for new physics (NP) beyond the SM. One of these is the long-standing anomaly of forward-backward asymmetry in the pair production of b quarks, A_{FB}^b, as measured at the Z peak. The second is the ratio R_b, defined as $R_b = \Gamma(Z \to b\bar{b})/\Gamma(Z \to \text{hadrons})$. Of much interest during the LEP-I and SLC era [2–5], the second tension has resurfaced due to a recent evaluation of R_b in the SM, taking into account all two-loop effects [6].

The Gfitter group [1] has updated the SM fit after the discovery of the Higgs boson at $m_h = 125.7 \pm 0.4$ GeV [7,8]. With the experimental inputs from Ref. [9], the fit [1] shows

$$R_b(\text{exp}) = 0.21629 \pm 0.00066,$$
$$R_b(\text{SM}) = 0.21474 \pm 0.00003, \tag{1}$$

with a pull of -2.35, where for any observable O with a standard deviation σ_{exp}, the pull is defined as[1]

$$\text{Pull} = \frac{O_{\text{SM}} - O_{\text{exp}}}{\sigma_{\text{exp}}}. \tag{2}$$

Note that the pull has increased to -2.35 from -0.8 [as calculated earlier using $R_b(\text{SM}) = 0.21576 \pm 0.00008$] thanks to the recent computation of the full two-loop effects in the SM [6].

The pull for A_{FB}^b is 2.5, computed from [1,9]

$$A_{FB}^b(\text{exp}) = 0.0992 \pm 0.0016,$$
$$A_{FB}^b(\text{SM}) = 0.1032^{+0.0004}_{-0.0006}. \tag{3}$$

Present ever since the LEP-I days, this discrepancy constitutes, perhaps, the most longstanding indicator of NP. Indeed, over the years, numerous attempts have been made to solve this problem in the context of specific NP scenarios. Prominent amongst these are those invoking extra Higgs scalars [10], low-energy supersymmetry [11], or just mixing with exotic quarks [12]. At the same time, the data has a significant constraining power and may be used to rule out certain classes of models particularly in light of the discovery of the 126 GeV scalar.[2]

There is another mild tension in the forward-backward asymmetry of the τ measured at the Z peak. While this has not been updated in Ref. [1] using the m_h data, the value of the asymmetry hardly depends on whether m_h is given as an input or is treated as a free parameter to be determined from the fit. We therefore quote the Particle Data Group (PDG) result [14],

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[1] This definition is consistent with Gfitter but opposite in sign to that used by the Particle Data Group [14].

[2] For example, no supersymmetric model, where the lighter chargino is dominantly a wino, is consistent with both R_b and A_{FB}^b measurements, if we assume the 126 GeV scalar to be the lightest CP-even neutral Higgs boson [13].

$$A^{\tau}_{FB}(\mathrm{exp}) = 0.0188 \pm 0.0017,$$
$$A^{\tau}_{FB}(\mathrm{SM}) = 0.01633 \pm 0.00021, \tag{4}$$

with a pull of -1.5. However, the branching ratio for $Z \to \tau^+\tau^-$ is consistent with that of the other leptons, viz.,

$$\mathrm{Br}(Z \to \tau^+\tau^-) = (3.370 \pm 0.008)\%,$$
$$\mathrm{Br}(Z \to e^+e^-) = (3.363 \pm 0.004)\%. \tag{5}$$

Taking into account the electroweak corrections, $R_\tau \equiv \Gamma(Z \to \mathrm{hadrons})/\Gamma(Z \to \tau^+\tau^-)$ is slightly above the SM predictions, but consistent nevertheless, with a pull of only 0.6,

$$R_\tau(\mathrm{exp}) = 20.764 \pm 0.045,$$
$$R_\tau(\mathrm{SM}) = 20.789 \pm 0.011. \tag{6}$$

The partial width $\Gamma(Z \to b\bar{b})$ is best analyzed by parametrizing the $Zb\bar{b}$ vertex as

$$\frac{g}{\cos\theta_W}\bar{b}\gamma^\mu[(g_L^b + \delta g_L^b)P_L + (g_R^b + \delta g_R^b)P_R]bZ_\mu, \tag{7}$$

where

$$g_L^b = T_3^b - \kappa_b Q_b \sin^2\theta_W, \qquad g_R^b = -\kappa_b Q_b \sin^2\theta_W, \tag{8}$$

with $\kappa_b = 1.0067$. The deviation of κ_b from unity incorporates the electroweak corrections, whereas $\delta g^b_{L,R}$ comprise all possible corrections arising from NP sources. On analyzing all the electroweak data,[3] the best fits are obtained [15] for

$$\begin{aligned} &(i) \quad \delta g_L^b = 0.001 \pm 0.001, \qquad \delta g_R^b = 0.016 \pm 0.005, \\ &(ii) \quad \delta g_L^b = 0.001 \pm 0.001, \qquad \delta g_R^b = -0.170 \pm 0.005, \end{aligned} \tag{9}$$

where both δg_L^b and δg_R^b have been treated as free parameters. Indeed, the χ^2/d.o.f. for the two fits are too close to be considered separately [12,15]. It is easy to see the origin of these two solutions. Apart from some numerical constants,

$$\Gamma(Z \to b\bar{b}) \propto [(g_L^b)^2 + (g_R^b)^2], \qquad A^b_{FB} \propto \frac{(g_R^b)^2 - (g_L^b)^2}{(g_R^b)^2 + (g_L^b)^2}, \tag{10}$$

with $g_R^b = 0.077$ and $g_L^b = -0.423$ within the SM. The partial width $\Gamma(Z \to b\bar{b})$ can be pushed upward by changes in either or both of $g^b_{L,R}$; however, the upward pull on A^b_{FB} preferentially chooses a change in g_R^b. This change must be such that $|g_R^b + \delta g_R^b|^2$ is marginally higher than $(g_R^b)^2$, and so δg_R^b must either be positive and small, or negative and large. It may seem that analogous solutions with large and negative δg_L^b (so that the sign of g_L^b is reversed without changing its magnitude appreciably) should also be admissible. Indeed, this is true as far as the Z-peak observables are concerned. However, away from the Z peak, such a switch would essentially reverse the sign[4] of $A_{FB}(e^+e^- \to b\bar{b})$ and, hence, run afoul of the data [12]. It is intriguing to note that such considerations do not choose between the two solutions of Eq. (9) [12]. It is obvious, though, that if the shifts $\delta g^b_{L,R}$ come only from perturbative quantum corrections, then the first solution would be much easier to achieve than the second.

The strongest phenomenological constraints on NP scenarios arise, typically, from flavor physics, especially from processes involving the first two families. Over the years this has prompted many constructions wherein the coupling of the NP sector to the SM fermions is not flavor democratic, but is preferential to the third generation. Of particular interest in this context are scenarios that proclaim the Higgs to be a condensate effecting a dynamic breaking of electroweak symmetry rather than a fundamental scalar [16], or models with extensions of the gauge group associated with electroweak symmetry [17]. Other examples of models that envisage a special role for heavy fermions include models with extra space-time dimensions [18–20], and models where the electroweak symmetry is broken in a nonlinear way [21], including the Little Higgs models [22]. A still different class of possibilities is afforded by the hypotheses where the SM is augmented by color-triplet or color-sextet scalars that have Yukawa couplings with the third generation [23].

With each such NP scenario being unique in certain respects, it is useful to concentrate on the essential aspects, rather than dwell on the specifics. In particular, if the NP sector is heavy, integrating it out would leave us with new operators in the effective low-energy theory. Moreover, if the NP sector couples preferentially with the third generation, these would primarily be four-fermion operators (and, perhaps, anomalous magnetic moment-like operators) involving third-generation currents with undetermined Wilson coefficients that have to be matched with the full

[3] It should be noted that had we concentrated only on $\Gamma(Z \to b\bar{b})$ and A^b_{FB}, to the exclusion of all else, the fit would have been substantially different, with $\delta g_L^b \sim 0.003$. This, however, would be illogical for such a simple-minded shift would cause the predictions for several other precision variables (such as Γ_Z, Γ_{had}, etc.) to deviate from the measurements.

[4] Away from the Z peak, the domimant contribution to A_{FB} accrues from the interference between the photon and Z-mediated amplitudes.

NP. Reference [24], for example, considered the possibility of such operators explaining certain tensions in the B-physics sector. In this paper, we adopt a similar stance and investigate the implications of such an effective theory for the Z-peak observables, including R_b, A_{FB}^b, and A_{FB}^τ, and whether some of these operators could possibly ameliorate the aforementioned discrepancies. While it might seem that, given the large number of operators available, it would always be possible to find a set that "solves" the problem, it turns out that, in reality, only a subset can play the requisite role. Furthermore, a large Wilson coefficient for any such operator would lead to telltale signatures at the LHC, thereby offering us falsifiability of the ansatz.

This paper is organized as follows. In Sec. II, we introduce new effective dimension-six operators involving only the third-generation fermions. As our aim is to enhance δg_R^b, we might expect that operators involving right-chiral fields would be more suitable for our purpose, and that indeed turns out to be the case. We also delineate the region allowed by the Z-peak observables in the parameter space of the new operators. In Sec. III, we discuss some of the possible signals at the LHC that should show an unambiguous signature of such new physics. We summarize and conclude in the last section. Some calculational details are relegated to the Appendix.

II. NEW OPERATORS

As we are interested essentially in b-sector observables, we begin by introducing generic four-fermion operators involving the b quark, given by

$$\frac{\xi}{\Lambda^2}[\bar{f}\gamma_\mu(v_f + a_f\gamma_5)f][\bar{b}\gamma^\mu(v_b + a_b\gamma_5)b], \tag{11}$$

where ξ is a dimensionless number which is *a priori* undetermined and can only be fixed with a knowledge of the full theory. Λ is the scale up to which the effective theory is valid, and is essentially the scale of NP. The identity of f is undetermined at this point. It is obvious, though, that low-energy constraints on such an operator are the least severe if f is a third-generation fermion. For example, if $m_f < m_b$, we would need $\xi/\Lambda^2 \ll \alpha/M_\Upsilon^2$ so as not to run afoul of $\Upsilon(nS)$ decays. The SM decay is an electromagnetic one, and the width is given by [25]

$$\Gamma_{\Upsilon(1S)\to\ell\ell} = 4\alpha^2 Q_b^2 M_\Upsilon^{-2}|R(0)|^2(1+2x)\sqrt{1-4x},$$

where $x = M_\ell^2/M_\Upsilon^2$ and $R(0)$ is the radial part of the nonrelativistic wave function at the origin. It might be argued that such a decay has nontrivial dependencies on quantities [such as $R(0)$] that can only be calculated in a nonperturbative framework and, thus, the results are model dependent. It is easy to see, though, that apart from the comparable numerical factors, the new physics rate is suppressed by ξ/Λ^2 compared to α/M_Υ^2, and so the bound quoted here is a very conservative one. The terms in Eq. (11) do not exhaust the list of relevant Lorentz-invariant neutral-current four-fermion operators. Scalar (pseudoscalar) and tensor (pseudotensor) structures are also admissible possibilities; however, as would be obvious immediately, the contributions of such operators to the effective $Zb\bar{b}$ vertex are chirality suppressed.[5]

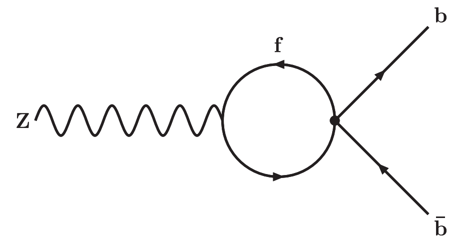


FIG. 1. The effective $Zb\bar{b}$ vertex arising from a single insertion of the operator in Eq. (11).

The operator of Eq. (11) gives rise to a one-loop correction to the $Z\to b\bar{b}$ vertex (see Fig. 1). Formally, this amplitude is quadratically divergent and can be evaluated using a gauge-invariant prescription such as dimensional regularization. While the infinite correction is cancelled by introducing appropriate counterterms[6], the finite part of the correction to the $Zb\bar{b}$ vertex is given by

$$\begin{aligned}\delta g_L^b &= \left(\frac{v_b - a_b}{2}\right)\frac{N_C\xi}{4\pi^2\Lambda^2}\mathcal{J},\\ \delta g_R^b &= \left(\frac{v_b + a_b}{2}\right)\frac{N_C\xi}{4\pi^2\Lambda^2}\mathcal{J},\end{aligned} \tag{12}$$

where $N_C = 3(1)$ if f is a quark (lepton) and $\mathcal{J} \equiv \mathcal{J}(v_f, a_f, m_f, M_Z)$, the expression for which can be found in the Appendix. It should be appreciated that, had we attempted instead to calculate the effective $b\bar{b}\gamma$ vertex, the very form of the corresponding $\mathcal{J}$ would have ensured that the charge radius does not receive any correction. This, of course, is a consequence of gauge invariance and has been ensured by our use of dimensional regularization rather than a naive momentum cutoff.[7] If the scale Λ of new

[5]It might be argued that the contribution of such a (pseudo) scalar term to $\Gamma(\Upsilon(1S)\to\ell^+\ell^-)$ would be chirality true, thereby allowing the corresponding Wilson coefficient to be large. On the other hand, this would be severely constrained by the non-observation of the $\chi_{b0}\to\ell^+\ell^-$ decay.

[6]Although this might seem strange given the higher-dimensional nature of the interaction term, note that the calculation fully conforms to the spirit of effective field theories.

[7]Note that a naive application of a cutoff regularization would have given rise to leading corrections being independent of Λ rather than being suppressed as $(m^2/\Lambda^2)\ln(m^2/\Lambda^2)$, with the consequence that a smaller ξ would be required. Although such a dependence of the corrections would have been expected in a scalar theory, it is clearly not gauge invariant and, hence, inapplicable in the current context.

physics is to be substantially larger than the electroweak symmmetry-breaking scale (as the absence of any new resonances at the LHC seems to suggest), the four-fermion operators need to respect the full $SU(2)_L \otimes U(1)_Y$ symmetry. This is a further restriction on the generic operators of Eq. (11). As we need $\delta g_R^b \gg \delta g_L^b$, it stands to reason that the said operator should involve the b_R field rather than b_L. One of the simplest such operators is given by

$$\mathcal{O}_{RR}^t = \frac{\xi}{\Lambda^2}(\bar{t}_R \gamma_\mu t_R)(\bar{b}_R \gamma^\mu b_R), \tag{13}$$

i.e., with the choice $v_b = a_b = v_t = a_t = \frac{1}{2}$.

In the above, we have deliberately neglected the possibility of quark mixing. Since these operators were presumably generated well above the electroweak scale, it is likely that they were generated in the weak basis instead. If the starting point is indeed so, after the symmetry breaking, the operators need to be reexpressed in terms of mass eigenstates through a Cabibbo-Kobayashi-Maskawa-type rotation [24]. This would then generate a plethora of new operators. The corresponding Wilson coefficients would be constrained by several B-physics observables, such as the mass differences ΔM_d and ΔM_s, and the CP-violating phases β and β_s. We apply the principle of Occam's razor and refrain from considering the entire range of such new operators, restricting ourselves to considering the operator $\mathcal{O}_{RR}^t$ only. Note that, apart from the phenomenological advantages, an operator such as $\mathcal{O}_{RR}^t$ is typically less suppressed than others[8] in scenarios wherein the electroweak symmetry is broken in a nonlinear fashion [21].

Equation (12) immediately gives $\delta g_L = 0$ and $\delta g_R \neq 0$. The region in the ξ-Λ plane that generates the required δg_R [as in Eq. (9)] is shown in Fig. 2. Requiring that the coupling ξ be perturbative, at least at the TeV scale, means that only the $\delta g_R > 0$ solution proposed by Ref. [15] is realized.[9] There is a caveat, though. The analysis of Ref. [15] was performed by treating both δg_R^b and δg_L^b as free parameters, whereas invoking $\mathcal{O}_{RR}^t$ necessarily implies that $\delta g_L^b = 0$. In a strict sense, the fit would be different in the two cases. However, quantitatively, the 1σ (or 2σ) allowed regions in the two cases are not too different. Indeed, the required δg_L^b can be generated by positing, in addition, a $\mathcal{O}_{LL}^t$ with a Wilson coefficient much smaller than ξ. This, though, would be tantamount to invoking two new operators to explain two discrepancies, and, hence, we desist from exploring this alternative any further.

[8]Here we discount possible four-top operators as they are not germane to the issue at hand.

[9]One might set the perturbative limit at $\xi \sim \mathcal{O}(10)$, coming from the condition $\xi^2/16\pi^2 < 1$ for higher-order processes in the full theory. This is satisfied for $\Lambda \sim \mathcal{O}(1\ \mathrm{TeV})$, only for the $\delta g_R > 0$ solution and not the other one.

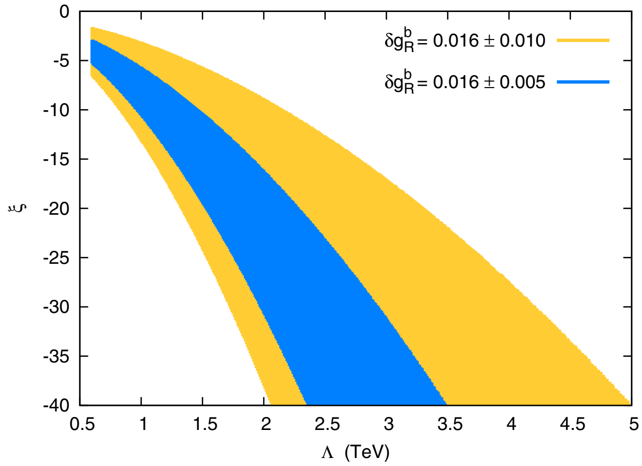


FIG. 2 (color online). The allowed region in the ξ-Λ plane that is consistent with the observed values of R_b and A_{FB}^b.

It is obvious that the operator in Eq. (13) also modifies the $Zt\bar{t}$ coupling, with the b now in the loop. However probing this effect presents a bigger challenge. Even at an e^+e^- collider, $t\bar{t}$ production is dominated by the photon-mediated amplitude with $e^+e^- \to Z^* \to t\bar{t}$ making a small contribution. Hence one needs to consider more complex processes. We shall return to this discussion in the next section.

A. Other operator choices

As discussed in Sec. I, apart from the b sector, some minor discrepancies also exist in the τ sector in the LEP data. One may, therefore, contemplate the introduction of an operator O_{RR}^τ involving τ's and b's analogous to O_{RR}^t above, in the hope that the two sets of discrepancies could perhaps be simultaneously explained. However, note that for O_{RR}^τ, $N_C = 1$ for δg_L^b and δg_R^b, but for the corresponding corrections to g_L^τ and g_R^τ, $N_C = 3$. Thus, in general, the corrections to the $Z\tau^+\tau^-$ couplings will be larger than those to $Zb\bar{b}$ couplings[10]. On the other hand, the disagreements between data and SM predictions are smaller in the case of the τ observables. Hence, with O_{RR}^τ alone, it is not possible to simultaneously generate the requisite corrections to all of g_L^b, g_R^b, g_L^τ, and g_R^τ. If one were to additionally consider O_{RL}^τ, O_{LR}^τ, and O_{LL}^τ as well, it is indeed possible to arrange a conspiracy between the coefficients of the various operators such that the observed values of all the couplings are obtained simultaneously. An easier path to such an explanation is offered by invoking a (set of) $\bar{\tau}\tau\bar{t}t$ operators along with O_{RR}^t. This has the advantage of not upsetting any other low-energy observable to a significant degree. On the other

[10]Although the correction term also carries a dependence on the mass of the fermion in the loop, the difference between m_τ and m_b is small and cannot entirely offset the difference due to the color factor.

hand, it is a construction that is barely testable in current experiments.

A much more intriguing possibility is offered where f [in Eq. (11)] is an exotic fermion. Clearly, few constraints apply to such operators, and it is much easier to arrange for the requisite shifts in $g^b_{L,R}$ as long as f itself does couple to the Z. This is eminently possible, as for example in supersymmetric or extra-dimensional extensions of the SM. While many different choices for f are possible (as long as it is heavy enough not to have been found at the Tevatron or the LHC), a particularly interesting choice is that of f being the dark matter (DM) candidate itself. The tantalizing indications, over the years, for the existence of a DM particle (whether it be from cosmological data fitting, indirect evidence from satellite-based observations, or direct Earth-bound experiments) in the absence of actual discovery has led to much speculation about its nature. It has been realized of late that, quite apart from dedicated DM search experiments, collider experiments can provide substantial information about the DM sector. Indeed, given the complete absence of any information, even dedicated DM searches only parametrize its interactions with matter through effective operators as in Eq. (11). The very same operators would also lead to DM pair production (in association with visible objects) at colliders. Thus, an excess in such channels (with the DM pair providing missing momentum) over the SM expectations would constitute a signal while a lack thereof would constrain the said interactions [26–29].

The situation becomes particularly interesting if the DM particle couples to the SM sector preferentially through the third-generation fermions [30–33]. Direct detection experiments would be rendered rather ineffectual. Even satellite-based indirect detection experiments would have reduced sensitivity. Although collider experiments too would suffer, the suppression in the cross section is not that extreme. Aided by the possibility of tagging heavy flavors, LHC experiments would have the highest sensitivity (amongst all currently operating ones) to such operators [32]. Given this, it is worthwhile to consider this possibility as well. As the formalism is identical to that which we have delineated above, the results would only depend on the choices[11] for the DM couplings to the b current as well as to the Z. And finally, while scalar DM is also a possibility, and may couple to both the Z as well as to a b current, the corresponding corrections to the effective $Zb\bar{b}$ vertex would have a Lorentz structure that does not readily translate to a discernible shift in A^b_{FB}.

III. O^T_{RR} AT THE LHC

In the last section, we saw that the low-energy constraints on the operator O^t_{RR} (or analogous ones) are not

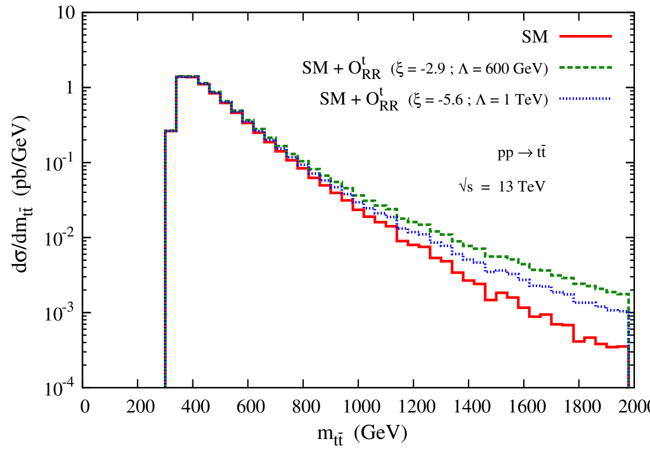


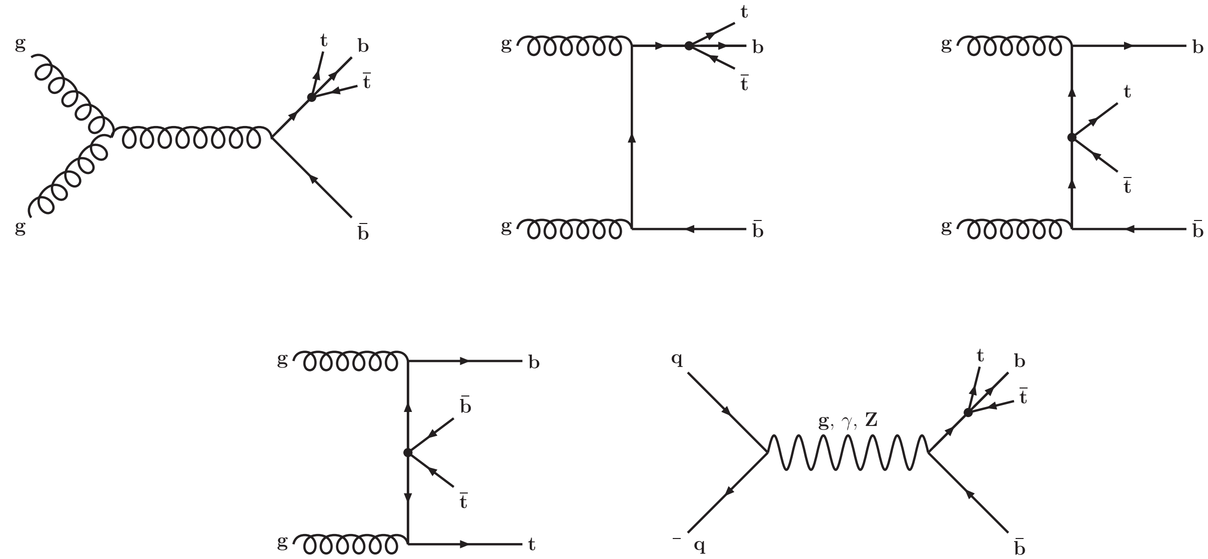
FIG. 3 (color online). The m_{tt} distribution of the cross section for $pp \to t\bar{t}$ at $\sqrt{s} = 13$ TeV, in the presence of an anomalous $b\bar{b} \to t\bar{t}$ contribution driven by O^t_{RR}. Included is only the leading-order cross section computed with the CTEQ6L distributions. The upper (lower) dashed line is for lower (higher) value of Λ.

strong enough to call into question a possible role for it in the explanation of the anomaly in the $Zb\bar{b}$ vertex. Thus, the only theatre for studying such an operator is provided by colliders. Although O^t_{RR} also engenders changes in the $Zt\bar{t}$ vertex analogous to those wrought for the $Zb\bar{b}$ one, such a change is of little relevance either at the LHC or even at a linear collider.[12] And as we have already argued, loops induced by such operators do not generate any corrections to the electric or color charge radii of the fermions. Although anomalous (chromo) magnetic moments are indeed generated, once again, these are of little immediate concern as the change in $gg \to t\bar{t}$ is hardly discernible.

However, there is a tree-level subprocess that could receive a large contribution from O^t_{RR}, namely $b\bar{b} \to t\bar{t}$. Despite the smallness of the b flux within the proton, the additional contribution to the cross section, at $\sqrt{s} = 13$ TeV, can be as large as $\sim 10\%$ for values of ξ/Λ^2 required to reproduce the correct δg^b_R (see Fig. 3). While this might seem very promising in view of the accuracy in the $t\bar{t}$ cross-section measurement (especially in the dilepton channel), note that the theoretical errors due to higher-order corrections and parton distribution function ambiguities are much larger. The latter is of particular relevance here as the b flux is relatively poorly known. One might attempt to exploit the fact that owing to the higher-dimensional nature of the interaction term, the corresponding amplitude grows with energy. While this is certainly true at the subprocess level, the growth of the anomalous cross section is muted owing to the rapid fall of the b flux with Bjorken x. Moreover, the reconstruction of m_{tt} is less efficient in the dilepton channel, whereas the use of the hadronic channels

[11] It must be remembered though that if the DM is a Majorana fermion, it may not have a vector-like coupling to the Z, whereas an axial coupling is allowed.

[12] Even the best sensitivity, provided by a high-luminosity $t\bar{t}$ threshold scan at the linear collider, is not adequate to probe the required values of ξ.



FIG`. 4. Some of the new Feynman diagrams that come into play when O^t_{RR} is introduced.

typically leads to larger experimental uncertainties. Given this situation, we desist from further consideration of this channel.

Instead, we consider the process $pp \to t\bar{t}b\bar{b}$. As such, this final state is of interest as an SM background for analyses concerning Higgs production in association with a top pair where the Higgs then decays into a bottom pair. With the introduction of O^t_{RR}, several new diagrams come into play. Rather than listing all of them, we illustrate some representative topological classes in Fig. 4. At the LHC, the gluon-initiated contribution is, understandably, the dominant one. At first, it might seem that, owing to a different color structure, the O^t_{RR} diagrams cannot interfere with the pure QCD ones. This argument, though, holds only for those pairs of diagrams wherein the O^t_{RR} vertex is replaced by a gluon propagator, and not in general. Similarly, the new diagrams do interfere with the majority of the mixed QCD-electroweak diagrams in the SM. We incorporate all such potential contributions (including the subdominant ones) and calculate the cross section through a simple modification of the CalcHEP [34] software.

For a quantitative assessment, one must impose a minimal set of acceptance cuts on the final-state particles. To this end, we require that the transverse momentum and the rapidity of the two primary b jets (i.e., the b jets emanating from the primary hard process, rather than the decays of the top) satisfy

$$p_T(b) > 50 \text{ GeV}, \qquad |\eta(b)| < 2.5. \tag{14}$$

To veto Z and Higgs events (for example, as occasioned from $t\bar{t}Z$ or $t\bar{t}h$ production), we impose, in addition,

$$M(b,\bar{b}) \notin [75, 135] \text{ GeV}. \tag{15}$$

For a pp collider operating at a center-of-mass energy of 13 TeV, the SM prediction for the cross section for this process as calculated using CalcHEP [34] is ~60 fb. This could be enhanced by as much as an order of magnitude for (Λ, ξ) values consistent with the $Z \to b\bar{b}$ measurements (see Fig. 2). Owing to the higher-dimensional nature of the coupling, the excess would, typically, be concentrated in phase-space regions corresponding to large momentum transfers. In Figs. 5 and 6, we show some such kinematic distributions. We find that rather than requiring individual particles to be harder or more central, as in Eq. (14), it is more profitable to impose stronger cuts on variables such as those appearing in Figs. 5 and 6. Apart from the simplistic observables considered here, in the actual experimental setup, the use of more advanced analysis techniques will offer additional means for the extraction of signal from the background [35].

Note that the QCD cross section for the production of a $t\bar{t}$ pair along with two well-separated and hard jets is much larger than the $\sigma(t\bar{t}b\bar{b})$ that is quoted here. Thus, b tagging is of prime importance. The corresponding efficiency has a strong dependence on $p_T(b)$, and thus requiring it to be very large would lead to a drastic reduction in signal sizes. On the other hand, the typical values of $p_T(t/\bar{t})$ are not so large as to warrant worries pertaining to the identification of highly boosted tops. Thus, stiffening the cuts on the top momenta would seem to be called for. Reconstructing a top, however, is associated with certain limitations. With the additional bottom pair introducing further combinatoric ambiguities, the errors would be amplified to an extent. Note, though, that owing to its four-fermion nature, the signal events would tend to concentrate at higher values of $M(b,\bar{b})$ where the b jets emanate from the hard process. Thus, requiring that for at least one pairing $M(b,\bar{b})$ is much larger than the cut of Eq. (15) stipulates would enhance the signal-to-noise ratio [32,36]. Indeed, given that the NP cross sections are significantly large, the nominal

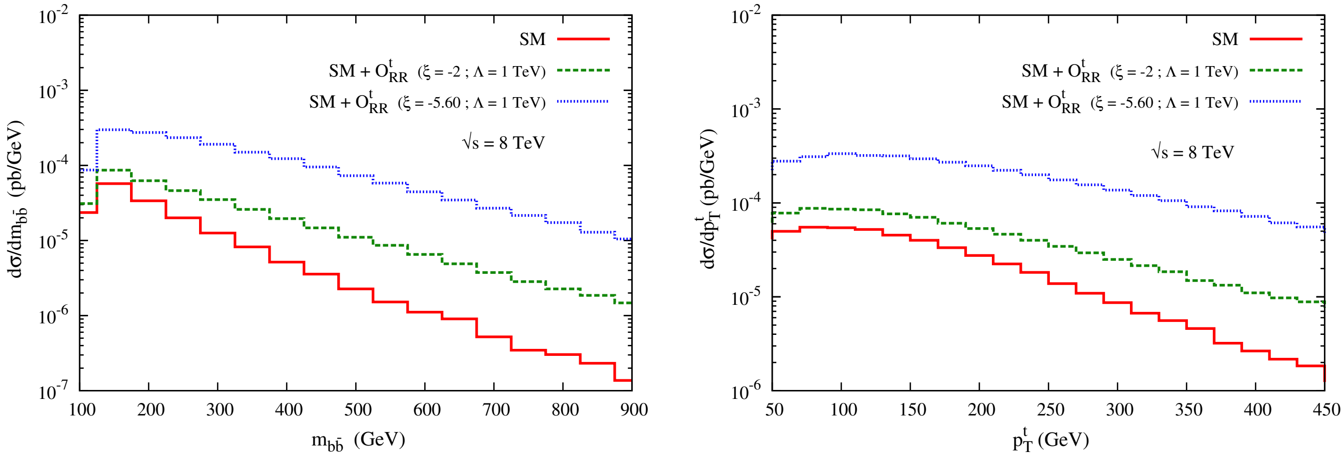


FIG. 5 (color online). $pp \to b\bar{b}t\bar{t}$ at $\sqrt{s} = 8$ TeV. Left panel: Invariant mass of the $b\bar{b}$ system. Right panel: Transverse momentum of the top. The upper (lower) dashed line is for $\xi = -5.6(-2)$.

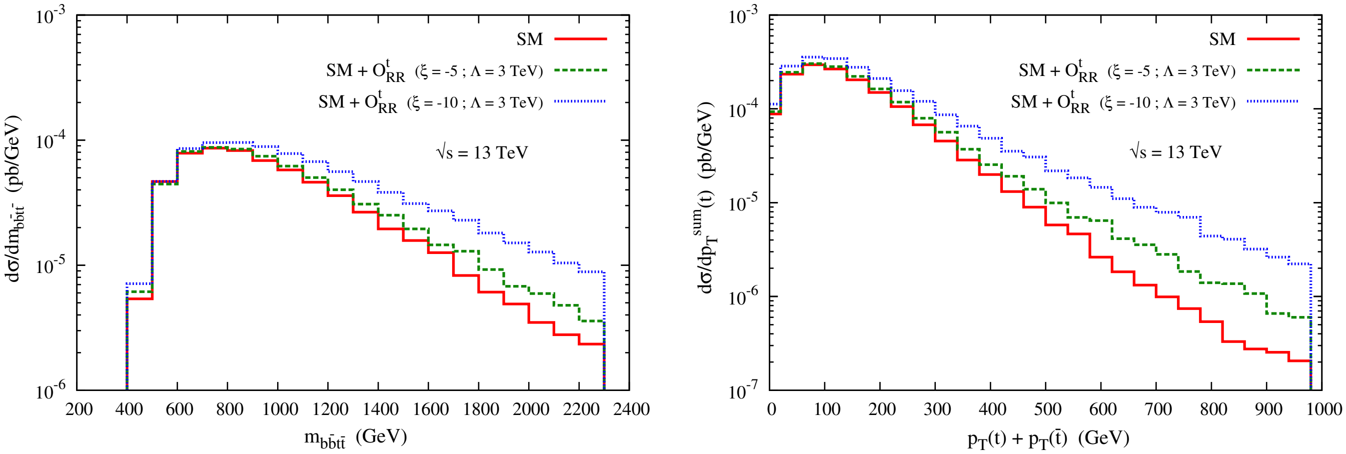


FIG. 6 (color online). $pp \to b\bar{b}t\bar{t}$ at $\sqrt{s} = 13$ TeV. Left panel: Invariant mass of the $b\bar{b}t\bar{t}$ system. Right panel: Transverse momentum of the $t\bar{t}$ system. The upper (lower) dashed line is for $\xi = -10(-5)$.

luminosity expected for the 13 TeV run of the LHC would be enough for a discovery even after accounting for the branching fractions, b-tagging efficiencies, and combinatoric ambiguities, as well as detector acceptance and efficiencies for a Λ near 3 TeV. This contention is supported by the detailed simulation of Ref. [32], where the production of dark matter particles in association with a top pair was considered. Although the final state is different ($t\bar{t} + \not{E}_T$), the analysis is similar; the absence of the missing transverse momentum is amply compensated for by the two hard b jets. Were one to admit smaller values of $\Lambda \sim 1$ TeV, large deviations from the SM would be expected even in the 8 TeV LHC data (see Fig. 5). Thus this mode is potentially the best bet for a direct confirmation of such an ansatz as that presented here.

We refrain, though, from using this study to extract information on ξ/Λ. For one, we have not taken into account the complexities of event reconstruction for this final state in the LHC environment. Furthermore, the theoretical predictions are only the leading-order ones. Nonetheless, it is suggestive of a method that could be used to further investigate a LEP/SLD anomaly at the LHC, where a direct repetition of the measurement is not possible.

IV. CONCLUSION

We have tried to gain some insight into the possible structure of NP at the TeV scale that might successfully address the mismatch between measurements and theoretical predictions of R_b and A^b_{FB}. We have used a bottom-up approach, not being confined to any specific model, with the sole assumption being that the NP couples only to the third-generation fermions. While there can be several such operators with different fermion fields and Lorentz structures, electroweak precision data and B-physics observables already put severe constraints on the Wilson coefficients of most of these operators. The quest for an operator that can resolve the anomalies while being relatively unconstrained has motivated us to work with one involving right-chiral top- and bottom-quark fields. At the same time, other choices are also possible, e.g., one with b quarks and dark matter particles that couple to the Z.

The four-fermion operators arise from a more fundamental theory at the higher scale. We performed our analysis in the spirit of an effective theory, with a high cutoff at the TeV scale (possibly indicative of the NP masses). The shifts in the $Zb\bar{b}$ couplings are caused by the parameters of the full theory, and we can only make the leading-order estimate of these in the effective theory. It turns out that there is a significant region in the parameter space that is consistent with the R_b and A_{FB}^b data, without being in contradiction with other observables.

Finally, we looked for the possible signals of this operator at the LHC. Altough $b\bar{b} \to t\bar{t}$ is the lowest-order process that features the new coupling, given the experimental as well as theoretical uncertainties the sensitivity is likely to be low. On the other hand, $pp \to t\bar{t}b\bar{b}$ is well suited for this task. We found that several observables would show a clear deviation from the SM, thus opening up clear channels to investigate such interactions. The results will be eagerly anticipated.

ACKNOWLEDGMENTS

A. K. acknowledges support from CSIR, India, and the DRS programme of UGC, India. P. S. would like to acknowledge financial support from NSERC, Canada.

Note added:—Recently Freitas and Huang [6] have revised their two-loop calculation of R_b and $\sin^2\theta_{\rm eff}^{b\bar{b}}$. As a result the discrepancy between the SM expectations and the experimental value has again come down to 1.2σ. If it stands, this result would serve to restrict the parameter space for the higher-dimensional effective operators. However, even the remaining part of the parameter space would still be of great interest in the context of the LHC as well as the paradigm of nonlinear realization of the electroweak symmetry. Moreover, if the anomaly in R_b indeed disappears after the inclusion of two-loop effects, then it leaves the heavy-quark sector in a rather intriguing position. It appears now that, for third-generation quarks, production cross sections agree with SM predictions but A_{FB} measurements do not. This throws up interesting possibilities and is sure to spur further activity in this area in the near future.

APPENDIX: ANALYTIC EXPRESSIONS

We parametrize the $Zb\bar{b}$ vertex within the Standard Model by

$$\frac{ig}{2\cos\theta_W}\bar{b}\gamma^\mu(v_Z^b + a_Z^b\gamma_5)b. \tag{A1}$$

The one-loop correction to this vertex on account of the interaction of Eq. (11) is given by the diagram in Fig. 7. The expression for the corresponding correction is given by

$$\frac{gN_C\xi}{2\cos\theta_W\Lambda^2}[\bar{b}\gamma_\alpha(v_b + a_b\gamma_5)b]\cdot\Gamma^{\mu\alpha}, \tag{A2}$$

where

$$\Gamma^{\mu\alpha} = -\int\frac{d^4k}{(2\pi)^4}\frac{\mathrm{Tr}[\gamma^\mu(v_Z^f + a_Z^f\gamma_5)(\not{k} + m_f)\gamma^\alpha(v_f + a_f\gamma_5)(\not{k} + \not{p}_1 + m_f)]}{(k^2 - m_f^2)[(k+p_1)^2 - m_f^2]}. \tag{A3}$$

The evaluation of this integral is best done by ignoring the higher-dimensional nature of the coupling and the possible role of Λ as a cutoff. Treating (ξ/Λ^2) as just a dimensionful parameter in the theory, we employ dimensional regularization and the finite part of the correction is given by (note that $\chi < 0$ denotes the presence of a threshold)

$$\Gamma^{\mu\alpha} = \frac{ig^{\mu\alpha}}{4\pi^2}\mathcal{J}, \tag{A4}$$

with

$$\begin{aligned}\mathcal{J} = {}& \frac{2}{3}\mathcal{A}_+ p_1^2\left[\frac{1}{2}\ln\left(\frac{m_f^2}{\Lambda^2}\right) + \frac{1}{6} - 2\chi - \frac{3}{2} + \sqrt{\chi}(3 + 4\chi)\tan^{-1}\left(\frac{1}{2\sqrt{\chi}}\right)\right] \\ & - m_f^2(\mathcal{A}_+ - \mathcal{A}_-)\left[\ln\left(\frac{m_f^2}{\Lambda^2}\right) - 2 + 4\sqrt{\chi}\tan^{-1}\left(\frac{1}{2\sqrt{\chi}}\right)\right],\end{aligned} \tag{A5}$$

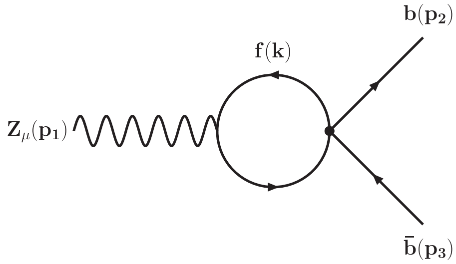


FIG. 7. One-loop correction to the $Zb\bar{b}$ vertex owing to NP interactions.

where

$$\chi = \frac{m_f^2}{p_1^2} - \frac{1}{4}, \qquad \mathcal{A}_\pm = v_Z^f v_f \pm a_Z^f a_f. \tag{A6}$$

In other words, on the inclusion of NP,

$$\begin{aligned} v_Z^b &\to v_Z^b + v_b \frac{N_C \xi}{4\pi^2 \Lambda^2} \mathcal{J}, \\ a_Z^b &\to a_Z^b + a_b \frac{N_C \xi}{4\pi^2 \Lambda^2} \mathcal{J}, \end{aligned} \tag{A7}$$

or, in terms of g_L^b and g_R^b,

$$\begin{aligned} g_L^b &\to g_L^b + \left(\frac{v_b - a_b}{2}\right) \frac{N_C \xi}{4\pi^2 \Lambda^2} \mathcal{J} \quad \text{and} \\ g_R^b &\to g_R^b + \left(\frac{v_b + a_b}{2}\right) \frac{N_C \xi}{4\pi^2 \Lambda^2} \mathcal{J}. \end{aligned} \tag{A8}$$

A couple of points need to be noted here. Had we employed a naive cutoff regularization instead, we would have encountered a quadratic divergence instead of the logarithmic one present in $\mathcal{J}$. This, however, would have been a spurious one occasioned by the facts that the loop integral is a tensorial one and that the naive cutoff regularization does not respect the symmetries of the theory [37]. Indeed, the adoption of such a regularization would have induced anomalous corrections to the (chromo) electric charge radius of the b and the t, thereby violating gauge invariance. On the other hand, had we used a gauge- and Lorentz-invariant prescription such as the Pauli-Villars scheme, we would have obtained a term exactly analogous to that which we already have, albeit achieved after a much more tedious calculation. A further issue relates to our implicit equalization of the renormalization scale μ_R with Λ. While this choice is a natural one, it is by no means the only possible one. Note, though, that the additional term introduced by using $\mu_R \neq \Lambda$ is a subdominant one and of little consequence here.

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