A unified, flavor symmetric explanation for the $t\bar{t}$ asymmetry and $Wjj$ excess at CDF

Ann E. Nelson,1 Takemichi Okui,2 and Tulin S. Roy1

1Department of Physics, University of Washington, Seattle, WA 98195
2Department of Physics, Florida State University, Tallahassee, FL 32306

We present a simple, perturbative, and renormalizable model with a flavor symmetry which can explain both the $t\bar{t}$ forward-backward asymmetry and the bump feature present in the dijet mass distribution of the $W + jj$ sample in the range 120–160 GeV that was recently reported by the CDF collaboration. The flavor symmetry not only ensures the flavor/CP safety of the model, but also relates the two anomalies unambiguously. It predicts a comparable forward-backward asymmetry in $c\bar{c}$. The forward-backward asymmetry in $b\bar{b}$ is, however, small. A bump in the dijet mass distribution in $Z + jj$ sample is also predicted but with a suppressed cross-section.

I. INTRODUCTION

Recently, the CDF collaboration has reported two interesting anomalies — a large $t\bar{t}$ forward-backward (FB) asymmetry [1, 2], and a $3.2\sigma$ excess in the 120–160 GeV range in the dijet mass distribution of the $W + jj$ sample [3] (see, however, Ref. [4] for a DØ analysis). The recent report of the FB asymmetry also confirms the trend suggested by the earlier measurements by CDF [5, 6] and DØ [7, 8].

It is a straightforward exercise to fit these two anomalies by introducing new particles with appropriately chosen masses and couplings. However, the nature of these anomalies suggests that the new physics should couple to standard-model (SM) quarks at tree level with an $O(0.1)$ coupling along with a nontrivial quark flavor structure. Such a new physics typically faces strong constraints from precision flavor and CP constraints, unless the model is equipped with a flavor symmetry.

In this paper, we present a weakly-coupled, renormalizable field theory with a flavor symmetry to explain both anomalies. (For attempts to generate just constraints from precision flavor and CP constraints, un-

In Sec. II, we define our model with an emphasis on the flavor symmetry structure, which keeps flavor/CP violations under control without tuning or an ad hoc choice of couplings. In Sec. III, we go through various potential constraints on the model for the values of parameters necessary for obtaining the $t\bar{t}$ FB asymmetry and $Wjj$ bump. Sec. IV shows our estimation of the asymmetry, while Sec. V shows the details of the bump feature in the dijet mass spectrum as predicted in our model. Our concluding reflections and a brief discussion of various implications of the model are included in Sec. VI.

II. THE MODEL

The flavor symmetry we propose is a subgroup of the $U(3)^3$ quark flavor symmetry of the SM:

$$\prod_{i=1}^{3} U(1)_{q_{Li}} \times U(1)_{u_{Ri}} \times U(3)_d \times Z_3, \quad (1)$$

where $q_{Li}$ and $u_{Ri}$ have charge $+1$ under $U(1)_{q_{Li}}$ and $U(1)_{u_{Ri}}$, respectively, while $d_{Ri}$ is a 3 of $U(3)_d$. $Z_3$ cyclically permutes the flavor indices of $q_{Li}$ and $u_{Ri}$ ($i = 1, 2, 3$), but not of $d_{Ri}$. The lepton sector of our model is identical to that of the SM, and will not be discussed in this paper.

In the SM, one can always go to a basis where $Y_u$ is diagonal. In the limit of neglecting both $Y_u$ and $Y_d$, the SM possesses the flavor symmetry [1]. Turning on the diagonal (but non-degenerate) $Y_u$ breaks the symmetry [1] to its subgroup $U(1)_{B_1} \times U(1)_{B_2} \times U(1)_{B_3} \times U(3)_{d}$, where $q_{Li}$ and $u_{Ri}$ have charge $+1$ under $U(1)_{B_i}$. This subgroup still forbids all flavor violations. Turning on $Y_d$ then breaks $U(1)_{B_1} \times U(1)_{B_2} \times U(1)_{B_3} \times U(3)_{d}$ down to the baryon number $U(1)_B$, thus introducing flavor mixing. However, since $Y_d$ breaks (and only $Y_d$ breaks) $U(3)_{d}$, we can always bring $Y_d$ into the form $Y_d = V_{CKM} diag(y_d, y_s, y_t)$, ensuring that $V_{CKM}$ is the only source of flavor violation. Also, note that $diag(y_d, y_s, y_t)$ and $diag(y_d, y_s, y_t)$ can both be taken to be positive definite, rendering $V_{CKM}$ the only source of CP violation.

Our fundamental assumption is that this symmetry breaking pattern persists for new physics beyond the SM as well. In other words, new physics should fully respect the flavor symmetry [1] in the limit $Y_u, Y_d \rightarrow 0$, and so $Y_u$ and $Y_d$ remain the only spurious breaking the symmetry [1]. This can be thought of as a variant of minimal flavor violation (MFV) [13, 17], and, in particular, the breaking pattern $U(1)_{B_1} \times U(1)_{B_2} \times U(1)_{B_3} \times U(3)_{d} \rightarrow U(1)_B$ by $Y_d$ has been studied in the context

1 We neglect the QCD vacuum angle. It is straightforward to add an axion to our model to solve the strong CP problem.
of the supersymmetric SM \[18, 19\]. We introduce a \(Z_3\) triplet of complex scalar fields, \(\Phi = (\Phi_1, \Phi_2, \Phi_3)\), where the gauge quantum number of \(\Phi\) is \((1, 2)_{-1/2}\) under \((SU(3)_c, SU(2)_L)_{U(1)_Y}\) representation. \(\Phi_i\) \((i = 1, 2, 3)\) are singlets under \(U(3)_d\), but they are charged under \(U(1)_{qL} \times U(1)_{qL^2} \times U(1)_{qL^3}\) as

\[
\Phi_1 \sim (0, 0, 1), \quad \Phi_2 \sim (1, 0, 0), \quad \Phi_3 \sim (0, 1, 0).
\]

Note that these charge assignments respect \(U(1)_{qL}\) while under \(U(1)_{Y}\), \(\Phi_1, \Phi_2, \Phi_3\) are singlets under \(U(3)_d\). \(\Phi_i\) \((i = 1, 2, 3)\) components of \(\Phi\) as:

\[
L_{\text{tree}} = L_{\text{SM}} + (D_{\mu} \Phi)^\dagger (D^\mu \Phi) - m_{\Phi}^2 \Phi^\dagger \Phi - \lambda (\bar{\tau}_{L1} \Phi_2 \bar{u}_{R3} + \bar{\tau}_{L2} \Phi_3 \bar{u}_{R1} + \bar{\tau}_{L3} \Phi_1 \bar{u}_{R2} + \text{c.c.}) - \xi (H^\dagger \sigma^a H)(\Phi^\dagger \sigma^a \Phi) - \zeta (H^\dagger H)(\Phi^\dagger \Phi) - \frac{\zeta}{4} (\Phi^\dagger \Phi)^2.
\]

The new physics is completely invariant under the flavor symmetry \([19]\). Therefore, the flavor symmetry \([19]\) is broken only by \(Y_u\) and \(Y_d\), just as in the SM. Also, we have chosen \(\lambda\) to be real without loss of generality, by redefining the phase of \(\Phi\). No new CP phase has, therefore, been introduced in this model.

The tree-level Lagrangian \([1]\) has three phenomenologically relevant parameters: \(m_{\Phi}, \lambda\), and \(\xi\). \(Z_3\) dictates that all three components of \(\Phi\) have the equal mass \(m_{\Phi}\), and that they all couple to the SM quarks with the same strength \(\lambda\) and likewise to the Higgs via \(\xi\) and \(\xi\). The interesting role of \(\xi\) is that it splits the masses of the neutral \((\Phi^0)\) and charged \((\Phi^-)\) components of \(\Phi\) as:

\[
m_{\Phi^0}^2 = m_{\Phi, \text{eff}}^2 - \xi v^2, \quad m_{\Phi^+}^2 = m_{\Phi, \text{eff}}^2 + \xi v^2,
\]

where \(v = 174\) GeV and \(m_{\Phi, \text{eff}}^2 \equiv m_{\Phi}^2 + \xi v^2\). We choose \(m_{\Phi^0} = 160\) GeV and \(m_{\Phi^-} = 220\) GeV, which corresponds to \(m_{\Phi, \text{eff}} = 192\) GeV and \(\xi = 0.38\). We take \(\lambda\) to be 1.4. This might appear too large to keep \(\lambda\) perturbative up to very high scale. Fortunately, the one-loop RG equation for \(\lambda\) is similar to that of the top Yukawa coupling in the SM and there is a quasi-fixed point near \(\lambda \approx 1.4\).

At loop level, counter-terms \(\delta L\) must be added to the Lagrangian \([1]\) for renormalization. We assume that all terms required to renormalize the theory are present in \(\delta L\), at the minimal level required to avoid fine tuning. This assumption is technically natural, and may be justified by assuming that our Lagrangian arises from a more fundamental theory in which \(Y_u\) and \(Y_d\) are the only parameters breaking the flavor symmetry \([1]\). As an example of terms in \(\delta L\), renormalization requires the counter-terms \(\bar{q}_{L1} Y_{u1} \Phi_2 (Y_d d_R)^2 + (\text{cyclic permutations})\) with a common coefficient \(\sim \lambda/(16\pi^2)\). We so include these operators in \(\delta L\) with a single coefficient \(\sim \lambda/(16\pi^2)\).

This also exemplifies the general principle that all generated operators respect the flavor symmetry \([1]\), broken only by \(Y_u\) and \(Y_d\). Moreover, since no operators are generated with an independent phase, renormalization does not require introducing new phases. The property that \(V_{\text{CKM}}\) is the only source of flavor and CP violations, therefore, remains intact at the quantum level.

### III. CONSTRAINTS

The property that \(V_{\text{CKM}}\) is the only source of flavor/CP violations, and the fact that the mass scale for \(\Phi^0\) and \(\Phi^-\) is similar to or larger than the top quark mass, imply that flavor/CP violations involving \(\Phi\) is at most comparable to those in the SM. For example, consider bounds from \(D^0 - \bar{D}^0\) mixing, that is, 4-fermion operators with two \(c\) and two \(\bar{u}\) fields that arise upon integrating \(\Phi\) out. Recall that, in our flavor symmetry breaking pattern, the only flavor non-diagonal spurion is \(Y_d = V_{\text{CKM}} \text{diag}(y_d, y_s, y_b)\). The flavor symmetry \([1]\) thus dictates that the simplest combination of spurions that can change \(c\) to \(u\) must involve the combination \((Y_d Y_d^\dagger)_{12} = (V_{\text{CKM}} \text{diag}(y_d^2, y_s^2, y_b^2)_{Y_{\text{CKM}}} \sim O(10^{-6})\). Since this combination converts \(c_L\) to \(s_L\), the coefficient of \((i u_L c_L)^2\) is on the order of \([Y_d Y_d^\dagger]_{12}/m_{\Phi}^2 \sim (10^5 \text{ TeV})^{-2}\) multiplied by certain powers of the couplings \(\lambda\) and \(g\) and the loop factor \(1/(16\pi^2)\). This is safely much smaller than the experimental bound \(\sim (10^3 \text{ TeV})^{-2}\) \([20]\). For the \((i \bar{u}_L c_L)/(i u_R c_L)\) operator, the spurions \((Y_u)_{22} \sim 10^{-2}\) and \((Y_u)_{11} \sim 10^{-5}\) have to be further inserted to convert \(c_L\) to \(c_R\) and \(u_L\) to \(u_R\), respectively, rendering the coefficient of the operator way below the bound \(\sim (10^4 \text{ TeV})^{-2}\) \([20]\). Therefore, \(D^0 - \bar{D}^0\) mixing is not an issue at all in our model, thanks to the flavor symmetry.

The most stringent flavor bounds on our model arise from the 4-fermion operators that are generated via the tree-level exchange of \(\Phi\). In the gauge basis, these are:

\[
(g_{2u} u_R)(\bar{u}_{R} q_L), \quad (g_{2c} c_R)(\bar{c}_R q_L), \quad (g_{1L} t_R)(\bar{t}_{R} q_L), \quad (g_{1L} t_R)(\bar{t}_{R} q_L),
\]

\[2\] A crucial difference between MFV and our flavor symmetry breaking pattern is in the spurion structure in the up-quark sector. In MFV there are \(9\) complex spurions for the up-quark sector, that is, the \(3 \times 3\) matrix \(Y_u\) transforming as \(3 \times 3\) under \(SU(3)_c \times SU(3)_c\), while in our case there are only three real spurions, \((Y_u)^i_{11}\) \((i = 1, 2, 3)\) carrying the charge \((1, -1, -1)\) for \(U(1)_{qL_1} \times U(1)_{R_3}\). By assumption \(Y_d\) is the only spurion which breaks CP or the \(U(1)_{R_{1,2,3}}\) quantum numbers. Note that the up-type quarks are in the mass basis at the outset and that, unlike in MFV, unitary rotations done on \(q_L\) or \(u_R\) are not approximate symmetries of the Lagrangian.

\[3\] We have the usual “hierarchy problem” for \(m_{\Phi}\) and \(m_{\Phi}\) as well as other dimension-2 operators in \(\delta L\). In this paper, we simply tune \(m_{\Phi}\) and \(m_{\Phi}\) and use dimensional regularization with (modified) minimal subtraction to obtain the natural sizes of dimension-2 operators in \(\delta L\), but it is straightforward to supersymmetrize the model to justify this.
The first and second operators can contribute to hadronic $b$ decays. In the mass basis, they contain

$$V_{ub}^* V_{cd} (\bar{b} L_R)(\bar{d} L_L) + V_{tb}^* V_{tc} (\bar{b} L_R)(\bar{c} L_L).$$

(7)

This comes from the tree-level exchange of $\Phi^-$, so its coefficient is $\lambda^2/m_{\Phi^-}^2$. The first operator contributes to the charmless process $b \to \bar{s} u u$. The particle data book [22] specifies the total inclusive branching fraction of $B$ mesons into charmed modes to be $(95 \pm 5)%$. For $m_{\Phi^-} = 220 \text{ GeV}$ and $\lambda = 1.4$, the leading order spectator decay model gives a branching fraction for the $b \to \bar{s} u u$ mode of $15\%$, which is within the $2\sigma$ margin of error. CP constraints do not pose a problem for the new contribution to decays. The CP phase in the new contribution to $b \to c \bar{c} s$ is almost the same as the phase of the standard model contribution. The phase of the new contribution to the $b \to \bar{s} u u$ mode is the same as the gluonic penguin contribution, which so far, is consistent with experiments. Of greater concern is the nonstandard increase in the hadronic width. However the heavy quark expansion, which is needed for a theoretical computation of the hadronic width [22][24], may have significant uncertainty for this computation [25][26]. Should further experimental tests of $B$ decay modes exclude such a large new contribution to hadronic $B$ decays then the $\Phi^-$ mass would have to be increased.

The correction to the electroweak parameter $\alpha T$ is given by

$$\alpha T = \frac{3}{32 \pi^2 v^2} \left[ m_{\Phi^0}^2 + m_{\Phi^-}^2 - \frac{2 m_{\Phi^-} m_{\Phi^0}}{m_{\Phi^-}^2 - m_{\Phi^0}^2} \log \frac{m_{\Phi^-}^2}{m_{\Phi^0}^2} \right].$$

(8)

where $\alpha$ is to be evaluated at the weak scale. For $m_{\Phi^0} = 160 \text{ GeV}$ and $m_{\Phi^-} = 220 \text{ GeV}$, we get $\alpha T = 1.5 \times 10^{-3}$. From the particle data book [22], for a Higgs mass of 117(300) GeV, the $T$ parameter is constrained to be $T = 0.07(0.16) \pm 0.08$. Thus, our model is consistent with precision electroweak constraints without any tuning. However, the $T$ parameter contribution gives an upper bound on the mass of the $\Phi^-$. 

$\Phi$ couplings to leptons are highly loop-suppressed; so precision lepton measurements (e.g. the muon $g - 2$) do not place constraints on our model. Since all components of $\Phi$ are unstable even in the collider time scale, there are no cosmological constraints. Hence, we concentrate on the collider physics constraints below.

- The total $t \bar{t}$ production cross-section is not changed significantly. Also note that the precise value of the theoretical prediction is still open to debate. The NLO+NNL calculations quote $\sim 10\%$ uncertainty [27][30] and the resummations of threshold logs result in a smaller value [31][32]. When compared to the SM leading order result, we find that $d\sigma_{t\bar{t}}/dM_{t\bar{t}}$, where $M_{t\bar{t}}$ is the invariant mass of the $t \bar{t}$ pair, in our model slightly decreases near the threshold but slightly increases at higher values of $M_{t\bar{t}}$. (See Sec. [15] for details).

- Single top production via $\Phi$ is suppressed. The leading single top production comes from $b \bar{c} \to \Phi^- \to W^- \Phi^0 \to W^- t \bar{c}$, which is highly suppressed by the smallness of the $b$ and $c$ parton distribution functions (PDFs). It is also suppressed by the 3-body phase space if the $W^-$ is on-shell and the $\Phi^-$ is off-shell. If the $\Phi^-$ is on-shell, then the $W^-$ must be off-shell, becoming a 4-body process. Either case, it is clearly much smaller than the SM counterpart $u d \to W^+ \to tb$, which is only suppressed by the offshellness of the $W^+$.

- The CDF collaboration has searched for events with same-sign lepton pairs and at least one $b$ jet, and found 3 such events in 2 fb$^{-1}$ of data [33], where they expect $\sim 2$ events from background. Di-top ($tt$) production can give such final states and is, thus, severely constrained. In our model, however, di-top production is extremely suppressed since in the limit of neglecting $Y_d$, baryon numbers for the three up-type quarks are separately conserved as in Eq. [1]. Therefore, this is not a constraint for us.

- The top-quark width is not modified significantly. For $m_{\Phi^0} = 160 \text{ GeV}$ and $\lambda = 1.4$, the total top width is 1.6 GeV, which is well within the experimental limit [31][35].

- There is a sizeable dijet production in our model via $\Phi_3$ exchange in the $s$- or $t$-channel. For masses as low as 160 GeV or 220 GeV, Tevatron has large SM dijet backgrounds due to the gluon PDF’s increasing rapidly at low parton $x$. Thus, there are no constraints from Tevatron [37]. The strongest bound comes from the CERN SPS collider ($pp$ at $\sqrt{s} = 630$ GeV). The UA2 collaboration at the SPS has placed 90% C.L. bounds $\approx 100$ pb on a dijet resonance at 160 GeV, and $\approx 10$ pb at 220 GeV [36]. Among our dijet channels, those from $t$-channel $\Phi_3$ exchange give a smooth dijet mass distribution on top of the smooth huge SM background. So they could not have been picked up by the UA2 search. For the $s$-channel contributions from $\Phi_3^0$ and $\Phi_3^-$, we find the cross-sections to be $\approx 56$ pb for $\Phi_3^0$ and $\approx 15$ pb for $\Phi_3^-$, with $\lambda = 1.4$. The latter may appear to be in conflict with the UA2 90% C.L. bound. However, note that the UA2 bounds are for a narrow resonance such as $W'$ which can distinguish itself from the smooth SM dijet background. Our $\Phi_3^-$ resonance, on the other hand, is not narrow — its width is quite large, $\approx 26$ GeV for $m_{\Phi^-} = 220$ GeV and $\lambda = 1.4$, which is expected to be smoothed out even further once parton showering and detector effects are taken into account. Thus, the search optimized for a very narrow resonance has a reduced sensitivity to our $\Phi_3^-$. Furthermore, note that UA2 performed their analysis in the early days of QCD jet studies. Their answer depends crucially on the quality of the Monte Carlo probes.
and the detector simulation which are primitive by today’s standard. They also use events with two exclusive jets, where jets were constructed using an infrared unsafe jet algorithm \cite{[38]}. We believe, therefore, that there are considerable uncertainties associated with their bounds, and that it is fair to regard the UA2 90\% C.L. bound as an order-of-magnitude limit.

IV. THE $t\bar{t}$ FORWARD-BACKWARD ASYMMETRY

Neglecting the CKM mixings and non-valence partons in the incoming $p$ and $\bar{p}$, the relevant interactions for this are:

\[
\mathcal{L}_\text{int} = -\lambda (\bar{u}_L \Phi^0_L t_R + \bar{d}_L \Phi^-_L t_R + c.c.) .
\]

The $t\bar{t}$ forward-backward asymmetry arises at the Tevatron from the processes $u\bar{u} \rightarrow t\bar{t}$ with a $t$-channel $\Phi^0_L$ exchange, and $d\bar{d} \rightarrow t\bar{t}$ with a $t$-channel $\Phi^-_L$ exchange.

A dedicated simulation including parton showering and detector effects is beyond the scope of this paper, partly due to the uncertainties in the SM prediction and the experimental measurement of the asymmetry. We simply perform our analysis at the parton level, and show that the asymmetry is generated with the right sign and is of the same order in magnitude.

We define the asymmetry as

\[
A_{t\bar{t}} = \frac{N(\Delta y > 0) - N(\Delta y < 0)}{N(\Delta y > 0) + N(\Delta y < 0)} ,
\]

where $\Delta y$ is the rapidity difference between the $t$ and $\bar{t}$. Our $t\bar{t}$ sample was generated using \texttt{Madgraph v4.4.48} \cite{[39]}. We have imposed the following cuts on the $t\bar{t}$ pairs: $|\eta_t|, |\eta_{\bar{t}}| < 2.0$, and $M_{t\bar{t}} > 450$ GeV. We find the asymmetry to be $A_{t\bar{t}} \approx 0.13$. Note that the asymmetry in our model does not depend linearly on $\cos \theta$ as it is assumed in Ref. \cite{[1]} to extrapolate the asymmetry to the full $4\pi$ solid angle. This is because our asymmetry is generated by $t$-channel exchange of particles similar in mass to the top quark, not from an $s$-channel heavy resonance. It should, thus, be compared to the value $A_{t\bar{t}} \approx 0.212 \pm 0.096$ for the reconstructed $M_{t\bar{t}} > 450$ GeV as actually measured within the CDF detector coverage \cite{[2]}.

We also check the total cross section of $t\bar{t}$ production at tree level. It is found to be within 10% of the value as calculated in the SM at the leading order. Assuming the same $k$-factor as in the case of the SM, we predict the total $t\bar{t}$ cross-section within theoretical uncertainties \cite{[27],[30]}. Deviation is seen when we check the cross-section as a function $M_{t\bar{t}}$. In Fig. 2 we have compared the tree level cross-section, as calculated in our model, to that in the SM. As shown in Fig. 2 and as reported in Refs. \cite{[40]}. The deviation is too small to give clear signal especially in early LHC data. Note that we cannot make direct comparison to the experimental data, since higher order corrections may change the shape of the curve besides the overall cross-section.

V. THE CDF EXCESS IN $Wjj$  

The relevant interactions for these processes are:

\[
\mathcal{L}_\text{int} = -\lambda (\tilde{c}_L \Phi^0_L u_R + \tilde{s}_L \Phi^-_L u_R + c.c.) .
\]

$Wjj$ final states via an intermediate $\Phi$ then dominantly arise at the Tevatron from the processes $u\bar{s} \rightarrow W^+\Phi^0_L \rightarrow W^+u\bar{c}$ and its charge-conjugated process.

For $\lambda = 1.4$, the $Wjj$ signal cross-section is found to be 2.1 pb. We use \texttt{Madgraph v4.4.48} to generate the signal events, which are subsequently decayed, showered, and finally hadronized by \texttt{PYTHIA v6.4} \cite{[41]}. We group the hadronic output of \texttt{PYTHIA} into “cells” of size $\Delta \eta \times \Delta \phi = 0.1 \times 0.1$. We sum the four momentum of all particles in
The dijet invariant mass distribution. The red curve is due to the resonances in our model at 160 GeV. The black curve shows the same for SM diboson events. Note that the center of the peak is shifted to 150 GeV.

VI. CONCLUSIONS

The $t\bar{t}$ forward-backward asymmetry and the $3.2\sigma$ excess in the 120–160 GeV range of the dijet mass distribution in the $Wjj$ sample at the Tevatron, if real, signify the existence of new physics at the electroweak scale. We constructed a simple, weakly-coupled renormalizable theory with one multiplet of scalar particles obeying the $(\prod_{i=1}^{3}U(1)_{1}\times U(1)_{u_{R_{i}}})\times U(3)_{d}\times Z_{3}$ flavor symmetry, which ensures that the only source of flavor/CP violations is $V_{\text{CKM}}$. We showed that the model can explain the two anomalies in terms of a single mass and a single coupling constant, without conflicting with existing bounds.

The flavor symmetry of the model also makes definite predictions on the amount of forward-backward asymmetries in $c\bar{c}$ and $b\bar{b}$. The $c\bar{c}$ asymmetry arises predominantly from the $uis$ initial state via the $t$-channel exchange of $\Phi_{d}^{0}$, while the $bb$ asymmetry is dominated by $c\bar{c}$ via $\Phi_{0}^{+}$ exchange. Therefore, the $c\bar{c}$ asymmetry is predicted to be comparable to the $t\bar{t}$ asymmetry, while the $bb$ asymmetry is expected to be suppressed due to the smallness of the $c$ parton distribution function.

As mentioned in Sec. V, we also have a $Zjj$ production via $us\rightarrow Z\Phi_{3}^{0}\rightarrow Zjj$, with smaller cross-section than $Wjj$. For the LHC with $\sqrt{s} = 7$ TeV, the production cross-sections of $Wjj$ and $Zjj$ are 47 pb and 12 pb respectively for the 160 GeV resonances, and are 20 pb and 8 pb respectively for the 220 GeV resonances.

Finally, note that we get a sizeable forward-backward asymmetry because of $O(1)$ coupling of the top quark with scalars of masses comparable to top mass, and as a result, we expect to see larger $t\bar{t}$ production cross section that in the SM at high values of the invariant mass of the $t\bar{t}$ pairs. This would certainly be an interesting feature to observe at the LHC.

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[Note added] While this manuscript was in progress, Refs. [43]–[47] discussing the CDF $Wjj$ excess (but not the $t\bar{t}$ asymmetry) appeared in arXiv.

[Note added 2] A week after the first version of our preprint appeared on arXiv, we have received Ref. [48], which claims that the $\phi_{-}$ mass has to be heavier than 540 GeV to avoid too large a rate for $B^{+}\rightarrow K^{+}\pi^{0}$, and that this would make our $Wjj$ excess signal go away. These claims, however, are incorrect. First, even granting 540 GeV for the $\Phi^{-}$ mass, it is not true that the signal would disappear, since it comes from the production of...
the scalar $\Phi^0$ \textit{directly from SM quarks}, whereas Ref. [18] misunderstands that the signal is from an $s$-channel $\Phi^-$ decaying to a $\Phi^0$ and a $W^-$. Second, the $B^+ \to K^+\pi^0$ (i.e. $b \to d\bar{s}$) calculation of Ref. [18] is actually technically incorrect. The essential error is the misidentification of the first term of 4-fermion operators [7] as the standard QCD penguin operator $O_b$, where the latter is summed over all (active) quark flavors while the former is \textit{not}. As a result, dangerous decays such as $b \to d\bar{s}$ are only generated via renormalization group running and are small. A proper calculation analyzing $\overline{B^0} \to \pi^+K^-$ (i.e. $b \to u\bar{u}s$) was recently performed by Ref. [29], which found that the rate for this process is enhanced by two orders of magnitude in our model.

We propose two ways to avoid this constraint while keeping our signals intact. The first is simply to make $\Phi^-$ heavier by a factor of a few to suppress operators [7]. This will enhance the $T$ parameter [8] but it can be tuned to be small by adding, for example, an electroweak-triplet scalar with a nonzero vacuum expectation value which contributes negatively to $T$.

The second way, which would require no tuning, is to enlarge the $Z_3$ symmetry of the flavor symmetry [1] to the full permutation group $S_3$. This amounts to the following. To keep track of $S_3$ more simply, we rename $\Phi_{1,2,3}$ as $\Phi_{32,13,21}$, respectively, and introduce $\Phi_{23,31,12}$ with the same gauge quantum numbers as $\Phi_{32,13,21}$ and the $\prod_{i=1}^3 U(1)_{\bar{q}iL} \times U(1)_{uiR}$ charges in the manner obvious from the notation (e.g., $\Phi_{23}$ has charges $+1$ and $-1$ under $U(1)_{\bar{q}2L}$ and $U(1)_{uiR}$, respectively, which is the opposite of $\Phi_{32}$). The six $\Phi$ fields form a sextet of $S_3$, and thus have a common mass $m_\Phi$ and a $\Phi^0,\Phi^-$ mass splitting parameter $\xi$ in the lagrangian [4]. The Yukawa couplings in Eq. 4 has to be generalized to be $S_3$ invariant, i.e.,

$$\mathcal{L}_{\text{Yukawa}} = \lambda \sum_{\text{all permutations}} \overline{q}_{L1} \Phi_{12} u_{R2} + \text{c.c.}. \quad (12)$$

Then, instead of the dangerous operator $(\overline{q}_{L2} u_{R})(\overline{u}_{R} q_{L2})$ (the first one of Eq. 9), we generate

$$\overline{q}_{L2} u_{R}(\overline{u}_{R} q_{L2}) + (\overline{q}_{L3} u_{R})(\overline{u}_{R} q_{L3})$$

$$= - (\overline{q}_{L1} u_{R})(\overline{u}_{R} q_{L1}) + \sum_{i=1}^{3} (\overline{q}_{Li} u_{R})(\overline{u}_{R} q_{Li}), \quad (13)$$

in the gauge basis. Going to the mass basis, the 2nd term will remain flavor-diagonal by the unitarity of $V$, hence not contributing to $b \to u\bar{u}s$, while the 1st term gives the $b$-dependent operator

$$V_{u\bar{u}} V_{u\bar{u}}(\overline{b}_{L1} u_{R})(\overline{u}_{R} d_{L1}). \quad (14)$$

Remarkably, this is much smaller than the dangerous operator in Eq. 9, which is only suppressed by $V_{cb}^* V_{cb}$. Therefore, the $b \to u\bar{u}s$ transition rate in this new model is suppressed by the same CKM factors as the tree level standard model contribution, and is sufficiently small even with a fairly light $\Phi^-$. The doubling of $\Phi$ will increase the dijet cross-section but still within the uncertainties discussed in section III especially due to the fact that $\lambda$ in this model is smaller for the same values of the $t\bar{t}$ asymmetry. A detailed study of this model is under investigation.

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