Stringy origin of diboson and dijet excesses at the LHC

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Very recently, the ATLAS and CMS Collaborations reported diboson and dijet excesses above standard model expectations in the invariant mass region of 1.8–2.0 TeV. Interpreting the diboson excess of events in a model independent fashion suggests that the vector boson pair production searches are best described by WZ or ZZ topologies, because states decaying into W+W− pairs are strongly constrained by semileptonic searches. Under the assumption of a low string scale, we show that both the diboson and dijet excesses can be steered by an anomalous U(1) field with very small coupling to leptons. The Drell–Yan bounds are then readily avoided because of the leptophobic nature of the massive Z′ gauge boson. The non-negligible decay into ZZ required to accommodate the data is a characteristic footprint of intersecting D-brane models, wherein the Landau–Yang theorem can be evaded by anomaly-induced operators involving a longitudinal Z. The model presented herein can be viewed purely field-theoretically, although it is particularly well motivated from string theory. Should the excesses become statistically significant at the LHC13, the associated Z′ topology would become a signature consistent only with a stringy origin.

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disfavored and could only describe the data in combination with another signal. This is because the CMS single lepton analysis sets an upper bound of 6.0 fb at 95% C.L. [3] and a cross section of this magnitude is needed to reproduce the hadronic excesses. Moreover, the CMS dilepton search has a small excess that this channel cannot explain [3].

Several explanations have been proposed to explain the excesses including a new charged massive spin-1 particle coupled to the electroweak sector (which can restore the left-right symmetry) [8], strong dynamics engendering composite models of the bosons [9], dark matter annihilation into right-handed fermions [10], a resonant triboson simulating a diboson through judicious choice of cuts [11], and a heavy scalar [12]. In this Letter we adopt an alternate path. We assume that the source of the excess originates in the decay of a new abelian gauge boson that suffers a mixed anomaly with the SM, but is made self-consistent by the Green–Schwarz (GS) mechanism [13]. Such gauge bosons occur naturally in D-brane TeV-scale string compactifications [14], in which the gauge fields are localized on D-branes wrapping certain compact cycles on an underlying geometry, whose intersection can give rise to chiral fermions [15]. The SM arises from strings stretching between D-branes which belong to the “visible” sector. Additional D-branes are generally required to cancel RR-tadpoles, or to ensure that all space-filling charges cancel. These additional D-branes generate gauge groups beyond the SM which forge the “hidden” sector.

There are two unrivaled phenomenological ramifications for intersecting D-brane models: the emergence of Regge excitations at parton collision energies √s ~ string scale ≡ M_s; and the presence of one or more additional U(1) gauge symmetries, beyond the U(1)Y of the SM. The latter derives from the property that, for N > 2, the gauge theory for open strings terminating on a stack of N identical D-branes is SU(N) rather than SU(N). (For N = 2 the gauge group can be Sp(1) ≃ SU(2) rather than U(2).) In a series of recent publications we have exploited both these ramifications to explore and anticipate new-physics signals that could potentially be revealed at the LHC. Regge excitations most distinctly manifest in the γ + jet [16] and dijet [17] spectra resulting from their decay. The extra U(1) gauge symmetries beyond hypercharge have (in general) triangle anomalies, but are canceled by the GS mechanism and the U(1) gauge bosons get Stückelberg masses. We have used a minimal D-brane construct to show that the massive U(1) field, the Z′, can be tagged at the LHC by its characteristic decay to dijets or dileptons [18]. In the framework of this model wherein we adjust the coupling strengths to be simultaneously consistent with the observed dijet excess and the lack of a significant dilepton excess. Concurrently we show that the model is also consistent with the ATLAS diboson excess as it allows for production of Z-Pairs. At the level of effective Lagrangian, the operator contributing to the Z′ZZ amplitude is induced by the GS anomaly cancellation.

In our calculations we will adopt as benchmarks:

$$\sigma(pp \to Z') \times B(Z' \to ZZ/WW) \sim 5.5^{+5.1}_{-3.3} \text{ fb}$$ [7].
$$\sigma(pp \to Z') \times B(Z' \to jj) \sim 91^{+52}_{-45} \text{ fb}$$ [7].
$$\sigma(pp \to Z') \times B(Z' \to e^+e^-) < 0.2 \text{ fb} (95\% \text{ C.L.})$$ [19].
$$\sigma(pp \to Z') \times B(Z' \to HH) < 12.9 \text{ fb} (95\% \text{ C.L.})$$ [20].

To develop our program in the simplest way, we will work within the construct of a minimal model with 4 stacks of D-branes in the visible sector. The basic setting of the gauge theory is given by U(3)_c × Sp(1) × U(1)_X × U(1) [21]. The LHC collisions take place on the (color) U(3)_c stack of D-branes. In the bosonic sector the open strings terminating on this stack contain, in addition to the SU(3)_C octet of gluons g_μ, an extra U(1) boson C_μ, most simply the manifestation of a gauged baryon number. The Sp(1)_B stack is a terminus for the SU(2)_L gauge bosons W^μ_L. The U(1)_Y boson Y_μ that gauges the usual electroweak hypercharge symmetry is a linear combination of C_μ and the U(1) bosons B_μ and X_μ terminating on the separate U(1)_X and U(1)_Y branes. Any vector boson orthogonal to the hypercharge, must grow a mass so as to avoid long range forces between baryons other than gravity and Coulomb forces. The anomalous mass growth allows the survival of global baryon number conservation, preventing fast proton decay [22].

The content of the hypercharge operator is given by
$$Q_Y = \frac{1}{6} Q_d - \frac{1}{2} Q_c + \frac{1}{2} Q_d .$$ (5)

We also extend the fermion sector by including the right-handed neutrino, with U(1) charges Q_ν = 0 and Q_ν = −1. The chiral fermion charges of the model are summarized in Table 1. It is straightforward to see that the chiral multiplets yield a U(1)_Y × SU(2)_L^2 mixed anomaly through triangle diagrams with fermions running in the loop. This anomaly is canceled by the GS mechanism, wherein closed string couplings yield classical gauge-variant terms whose gauge variation cancels the anomalous triangle diagrams. The extra abelian gauge field becomes massive by the GS anomaly cancellation, behaving at low energies as a Z′ with a mass in general lower than the string scale by an order of magnitude corresponding to a loop factor. Even though the divergences and anomaly are canceled, the triangle diagrams contribute an univocal finite piece to an effective vertex operator for an interaction between the Z′ and two SU(2)_L vector bosons [23]. This is a distinguishing aspect of the D-brane effective theory, which features a noticeable decay width of the Z′ into WW, ZZ, and ZY.1

The covariant derivative for the U(1) fields in the a, b, c, d basis is found to be
$$D_{aμ} = \partial_μ - ig_μ C_a Q_μ - ig_μ B_μ Q_c - ig_μ X_μ Q_d .$$ (6)

The fields C_μ, B_μ, X_μ are related to Y_μ, Y'_μ, and Y''_μ by the rotation matrix
$$R = \begin{pmatrix}
C_0 & C_ψ & -C_φ S_ψ + S_φ S_ψ & C_φ S_ψ + C_ψ S_φ S_ψ \\
C_φ S_ψ & C_φ C_ψ & S_φ S_C_ψ + C_ψ S_φ S_ψ & -S_ψ C_ψ + C_φ S_ψ + C_ψ S_φ S_ψ \\
-S_φ S_ψ & S_ψ C_φ S_ψ & C_0 & C_ψ C_ψ \\
S_ψ & C_φ S_ψ & C_ψ C_ψ & C_0
\end{pmatrix} ,$$ (7)

with Euler angles θ, ψ, and φ. Equation (6) can be rewritten in terms of Y_μ, Y'_μ, and Y''_μ as follows
$$D_{aμ} = \partial_μ - iY_μ \left(-S_ψ g_μ Q_d + C_φ S_ψ g_μ Q_c + C_ψ S_ψ g_μ Q_d \right)$$
$$- iY'_μ \left(C_φ S_ψ g_μ Q_d + (C_ψ C_ψ + S_ψ S_ψ S_φ) g_μ Q_c + (C_φ S_ψ S_ψ C_φ + S_ψ S_ψ S_ψ g_μ Q_a) \right)$$
$$- iY''_μ \left(C_φ S_ψ g_μ Q_d + (C_ψ S_ψ + C_ψ S_ψ S_ψ) g_μ Q_c + (C_φ C_ψ S_φ + S_ψ S_ψ) g_μ Q_a \right) .$$ (8)

1. The Landau–Yang theorem [24], which is based on simple symmetry arguments, forbids decays of a spin-1 particle into two photons.
Now, by demanding that $Y_\mu$ has the hypercharge $Q_Y$ given in (5) we fix the first column of the rotation matrix $R$

\[ \begin{pmatrix} C_\mu \\ B_\mu \\ X_\mu \end{pmatrix} = \begin{pmatrix} Y_\mu \frac{1}{g_Y} / g'_d & \cdots \\ -Y_\mu \frac{1}{g_Y} / g'_c & \cdots \\ Y_\mu \frac{1}{g_Y} / g'_d & \cdots \end{pmatrix}, \]

and we determine the value of the two associated Euler angles

\[ \theta = -\arcsin \left( \frac{1}{2g_Y / g'_d} \right) \]

and

\[ \psi = \arcsin \left( \frac{1}{2g_Y / (g'_c C_\psi)} \right). \]

The couplings $g'_c$ and $g'_d$ are related through the orthogonality condition,

\[ \left( \frac{-1}{2g'_c} \right)^2 = \frac{1}{g_Y} \left( \frac{c_1}{6g'_d} \right)^2 - \left( \frac{1}{2g'_d} \right)^2, \]

with $g'_d$ fixed by the relation $g_3(M_1) = \sqrt{5} g'_d(M_1)^2$. In our calculation we take $M_t = 20$ TeV as a reference point for running down to 1.8 TeV the $g'_c$ coupling, ignoring mass threshold effects of stringy states. This yields $g'_d = 0.36$. We have checked that the running of the $g'_c$ coupling does not change significantly for different values of the string scale. The third Euler angle $\phi$ and the coupling $g'_d$ will be determined by requiring sufficient suppression to leptons to accommodate (3) and a (pre-cut) production rate $\sigma(pp \to Z') \times B(Z' \to jj)$ in agreement with (2).

The $f \bar{f} Z'$ Lagrangian is of the form

\[ \mathcal{L} = \frac{1}{2} \left( \sum_f \left( \bar{\psi}_{fL} \gamma^\mu \psi_{fL} + \bar{\psi}_{fR} \gamma^\mu \psi_{fR} \right) Z'_\mu \right) \]

where each $\psi_{fL,R}$ is a fermion field with the corresponding $\gamma^\mu$ matrices of the Dirac algebra, and $\bar{\psi}_{fL,R}$ the vector and axial couplings respectively. From (8) and (13) we obtain the explicit form of the chiral couplings in terms of $\phi$ and $g'_d$

\[ \epsilon_{u_L} = \epsilon_{d_L} = \frac{2}{\sqrt{g_Y^2 + g'^2}} \left[ C_\phi S_0 S_\phi - C_\phi S_\phi \right] g'_d, \]

\[ \epsilon_{u_R} = \epsilon_{d_R} = \frac{2}{\sqrt{g_Y^2 + g'^2}} \left[ C_\phi S_0 g'_d + (C_\psi S_0 S_\psi - C_\psi S_\psi) g'_d \right]. \]

The decay width of $Z' \to f \bar{f}$ is given by [25]

\[ \Gamma(Z' \to f \bar{f}) = \frac{G_F e_2^2 M_Z^2}{6\pi \sqrt{2}} N_4 c(M_Z^2) M_{Z'} \sqrt{1 - 4x} \left[ \sqrt{2} (1 + 2x) + a_1^2 (1 - 4x) \right]. \]

where $G_F$ is the Fermi coupling constant, $c(M_Z^2) = 1 + \alpha_s / \pi + 1.409(\alpha_s / \pi)^2 - 12.77(\alpha_s / \pi)^3$, $\alpha_s = \alpha_s(M_Z^2)$ is the strong coupling constant at the scale $M_Z$, $x = m_f^2 / M_{Z'}^2$, and $N_4 = 3$ or 1 if $f$ is a quark or a lepton, respectively. The couplings of the $Z'$ to the electroweak gauge bosons are model dependent, and are strongly dependent on the spectrum of the hidden sector. Following [26] we parametrize the model-dependence of the decay width in terms of two dimensionless coefficients,

\[ \Gamma(Z' \to ZZ) = \frac{c_1^2 \sin^2 \theta \mu M_Z^2}{192\pi M_Z^2} \left( 1 - \frac{4 M_Z^2}{M_Z^2} \right)^{5/2} \]

\[ \approx \frac{c_1^2 (45 \text{ GeV})}{M_Z^2 / \text{TeV}} + \cdots, \]

(16)

\[ \Gamma(Z' \to W^+ W^-) = \frac{c_2^2 M_Z^2}{48\pi M_W^2} \left( 1 - \frac{4 M_W^2}{M_Z^2} \right)^{5/2} \]

\[ \approx \frac{c_2^2 (1.03 \text{ TeV})}{M_Z^2 / \text{TeV}} + \cdots, \]

(17)

\[ \Gamma(Z' \to Y Y') = \frac{c_3^2 \cos \theta \mu M_Z^2}{96\pi M_Z^2} \left( 1 - \frac{M_Z^2}{M_Z^2} \right)^{3/2} \left( 1 + \frac{M_L^2}{M_Z^2} \right) \]

\[ \approx \frac{c_3^2 (307 \text{ GeV})}{M_Z^2 / \text{TeV}} + \cdots. \]

(18)

The $Z'$ production cross section at the LHC8 is found to be [8]

\[ \sigma(pp \to Z') \approx 5.2 \left( \frac{2 \Gamma(Z' \to \mu \mu) + \Gamma(Z' \to d \bar{d})}{\text{GeV}} \right) \text{ fb}. \]

(19)

Next, we scan the parameter space to obtain agreement with (1) to (4). In Fig. 1 we show contour plots, in the $(g'_d, \phi)$ plane, for constant $\sigma(pp \to Z') \times B(Z' \to jj)$, $\sigma(pp \to Z') \times B(Z' \to e^+ e^-)$, and $\sigma(pp \to Z') \times B(Z' \to W^+ W^-)$. To accommodate (1), (2), and (3) the ratio of branching fractions of electrons to quarks must be minimized subject to sufficient dijet and diboson production. It is easily seen in Fig. 2 that $\phi = 0.96$ and $g'_d(M_t) = 0.29$, $c_1 = 0.08$, and $c_2 = 0.02$ yield $\sigma(pp \to Z') \approx 228 \text{ nb}$, $\sigma(B(Z' \to jj)) \approx 0.54$, $\sigma(B(Z' \to e^+ e^-)) \approx 8.9 \times 10^{-4}$, $\sigma(B(Z' \to W^+ W^-)) \approx 3.4 \times 10^{-4}$, which are consistent with (1), (2), and (3) at the 1σ level. In addition, $\sigma(B(Z' \to HZ) \approx 7.4 \times 10^{-3}$. Thus, the upper limit set by (4) is also satisfied by our fiducial values of $\phi$, $g'_d$, $c_1$, and $c_2$. The chiral couplings of $Z'$ and $Z''$ are given in Table 2. All fields in a given set have a common $g_Y Q_Y$, $g_Y Q_{Y'}$ couplings.

The second constraint on the model derives from the mixing of the $Z$ and the $Y'$ through their coupling to the two Higgs doublets. The criteria we adopt here to define the Higgs charges is to make the Yukawa couplings $(h_u d Q, h_d d Q, h_d \psi \epsilon, h_\epsilon \psi \epsilon)$ invariant under all three $U(1)'s$ [27]. Two “supersymmetric” Higgses $H_u \equiv H_v$ and $H_d$ (with charges $Q_u = Q_c = 0$, $Q_d = 1$, $Q_Y = 1/2$) and $Q_u = Q_c = 0$, $Q_d = -1$, $Q_Y = -1/2$ are sufficient to give masses to all the chiral fermions. Here, $(H_d) = (u) = (0)$.

The last two terms in the covariant derivative $\partial_\mu = \partial_\mu - i \frac{1}{\sqrt{g_Y^2 + g'^2}} Z_\mu (g'^2 + g_Y^2) Y_Y$ \[ \times -ig_Y Y'_\mu Q_Y - ig_Y Y'_\mu Q_{Y'}, \]

(20)
are conveniently written as
\[ -i \frac{X_{H_i}}{V_i} \overline{M}_Z Y_{\mu}^\nu = -i \frac{Y_{H_i}}{V_i} \overline{M}_Z Y_{\mu}^\nu \]  \hspace{1cm} (21)
for each Higgs $H_i$, with $T^3 = \sigma^3/2$, where for the two Higgs doublets
\[ x_{H_u} = -x_{H_d} = 1.9 \sqrt{g_d^2 - 0.032 S_\phi} \]  \hspace{1cm} (22)
and
\[ y_{H_u} = -y_{H_d} = 1.9 \sqrt{g_d^2 - 0.032 C_\phi}. \]  \hspace{1cm} (23)

The Higgs field kinetic term together with the GS mass terms
\[ \left( -\frac{1}{2} M_{d,\mu}^2 Y_{\mu}^\nu + \frac{1}{2} M_{u,\mu}^2 Y_{\mu}^\nu \right) \] yield the following mass square
matrix for the $Z - Z'$ mixing,
\[
\begin{pmatrix}
M_0^2 & M_0^2 (c_{\phi}^2 - s_{\phi}^2) & M_0^2 (y_{H_u} c_{\mu}^2 - y_{H_d} S_{\mu} S_{\phi}^2) \\
M_0^2 (c_{\phi}^2 - s_{\phi}^2) & M_0^2 (c_{\mu}^2 - S_{\mu}^2) + M^2 & M_0^2 (c_{\phi}^2 c_{\mu}^2 + s_{\phi}^2 S_{\mu}^2 + S_{\phi}^2 S_{\mu}^2) \\
M_0^2 (y_{H_u} c_{\mu}^2 - y_{H_d} S_{\mu} S_{\phi}^2) & M_0^2 (c_{\phi}^2 c_{\mu}^2 + s_{\phi}^2 S_{\mu}^2 + S_{\phi}^2 S_{\mu}^2) & M_0^2 (y_{H_u} c_{\mu}^2 + y_{H_d} S_{\mu} S_{\phi}^2) + M^2
\end{pmatrix},
\]
which does not impose any constraint on the $\tan \beta$ parameter. We have verified that, for our fiducial values of $\phi$ and $g_d$, if $M_{Z'} \gg M_Z$, the shift of the $Z$ mass would lie within 1 standard deviation of the experimental value.

In summary, we have shown that recent results by ATLAS and CMS searching for heavy gauge bosons decaying into $W W / ZZ$ and $jj$ final states could be a first hint of string physics. In D-brane string compactifications the gauge symmetry arises from a product of $U(N)$ groups, guaranteeing extra $U(1)$ gauge bosons in the spectrum. The weak hypercharge is identified with a linear combination of anomalous $U(1)'s$ which itself is anomaly free. The extra anomalous $U(1)$ gauge bosons generically obtain a St"uckelberg mass. Under the assumption of a low string scale, we have shown that the diboson and dijet excesses can be steered by an anomalous $U(1)$ field with very small coupling to leptons. The Drell–Yan bounds are then readily avoided because of the leptophobic nature of the massive $Z'$ gauge boson. The resulting loop diagrams, along with tree-level higher-dimension couplings arising from the GS anomaly cancellation mechanism, generate an effective vertex that couples the anomalous $U(1)$ field to two electroweak gauge bosons. The effective vertex renders viable the decay of the $Z'$ into $Z$-pairs, which is necessary to fit the data. Should the excesses become statistically significant at the LHC13, the associated $Z\gamma$ topology would become a signature consistent only with a stringy origin.

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