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Stringy origin of diboson and dijet excesses at the LHC

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ABSTRACT

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 $Z\gamma$ topology would become a signature consistent only with a stringy origin.

Very recently, searches for narrow resonances at the ATLAS and CMS experiments uncovered various peaks in invariant mass distributions near 2 TeV: *(i)* The ATLAS search for diboson production contains a 3.4σ excess at ~ 2 TeV in boosted jets of WZ [1]. The global significance of the discrepancy above standard model (SM) expectation is 2.5σ. The invariant mass range with significance above 2σ is ~ 1.9 to 2.1 TeV. Because the search is fully hadronic, the capability for distinguishing gauge bosons is narrowed. Therefore, many of the events can also be explained by a ZZ or WW resonance, yielding excesses of 2.9σ and 2.6σ in these channels respectively. *(ii)* The CMS search for diboson production (without distinguishing between the W- and Z-tagged jets) has a 1.4σ excess at ~ 1.9 TeV [2], and the search for diboson production with a leptonically tagged Z yields a 1.5σ excess at invariant mass ~ 1.8 TeV [3]. *(iii)* The CMS search for dijet resonances finds a 2.2σ excess near 1.8 TeV [5]. *(iv)* Around the same invariant mass ATLAS also recorded an excess in the dijet distribution with a 1σ significance [4]. *(v)* The CMS search for resonant HW production yields a 2.1σ excess in the energy bin of 1.8 to 1.9 TeV; here the Higgs boson is highly boosted and decays into $b\bar{b}$, whereas the W decays into charged leptons and neutrinos [6]. Barring the three ATLAS analyses in diboson production, all these excesses are completely independent.

Although none of the excesses is statistical significant yet, it is interesting to entertain the possibility that they correspond to a real new physics signal. On this basis, with the assumption that all resonant channels are consistent with a single resonance energy, a model free analysis of the various excesses has been recently presented [7]. The required cross sections to accommodate the data are quite similar for WZ and ZZ final states, which can be considered as roughly the same measurement. A pure WW signal is

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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 excesses 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 $\sqrt{\hat{s}} \sim$ string scale $\equiv M_s$; 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 $U(N)$ rather than $SU(N)$. (For $N = 2$ the gauge group can be $Sp(1) \cong SU(2)$ 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 herein 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:

$$\sigma(pp \to Z') \times \mathcal{B}(Z' \to ZZ/WW) \sim 5.5^{+5.1}_{-3.7}\ \text{fb} \quad [7], \tag{1}$$

$$\sigma(pp \to Z') \times \mathcal{B}(Z' \to jj) \sim 91^{+53}_{-45}\ \text{fb} \quad [7], \tag{2}$$

$$\sigma(pp \to Z') \times \mathcal{B}(Z' \to e^+e^-) < 0.2\ \text{fb (95\% C.L.)} \quad [19], \tag{3}$$

$$\sigma(pp \to Z') \times \mathcal{B}(Z' \to HZ) < 12.9\ \text{fb (95\% C.L.)} \quad [20]. \tag{4}$$

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 $U(3)_a \times Sp(1)_b \times U(1)_c \times U(1)_d$ [21]. The LHC collisions take place on the (color) $U(3)_a$ 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^a_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^a_μ. 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 $U(1)_c$ and $U(1)_d$ 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].

Table 1
Chiral fermion spectrum of the D-brane model.

Fields	Sector	Representation	Q_a	Q_d	Q_c	Q_Y
U_R	$a \leftrightharpoons d^*$	$(3, 1)$	1	0	1	$\frac{2}{3}$
D_R	$a \leftrightharpoons d$	$(3, 1)$	1	0	-1	$-\frac{1}{3}$
L_L	$c \leftrightharpoons b$	$(1, 2)$	0	1	0	$-\frac{1}{2}$
E_R	$c \leftrightharpoons d$	$(1, 1)$	0	1	-1	-1
Q_L	$a \leftrightharpoons b$	$(3, 2)$	1	0	0	$\frac{1}{6}$
N_R	$c \leftrightharpoons d^*$	$(1, 1)$	0	1	1	0

The content of the hypercharge operator is given by

$$Q_Y = \frac{1}{6} Q_a - \frac{1}{2} Q_c + \frac{1}{2} Q_d\,. \tag{5}$$

We also extend the fermion sector by including the right-handed neutrino, with $U(1)$ charges $Q_a = 0$ and $Q_c = Q_d = -1$. The chiral fermion charges of the model are summarized in Table 1. It is straightforward to see that the chiral multiplets yield a $[U(1)_a SU(2)^2_L]$ 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 $Z\gamma$.[1]

The covariant derivative for the $U(1)$ fields in the a, b, c, d basis is found to be

$$\mathcal{D}_\mu = \partial_\mu - i g'_a C_\mu\, Q_a - i g'_c B_\mu\, Q_c - i g'_d X_\mu\, Q_d\,. \tag{6}$$

The fields C_μ, B_μ, X_μ are related to Y_μ, Y'_μ and Y''_μ by the rotation matrix

$$\mathcal{R} = \begin{pmatrix} C_\theta C_\psi & -C_\phi S_\psi + S_\phi S_\theta C_\psi & S_\phi S_\psi + C_\phi S_\theta C_\psi \\ C_\theta S_\psi & C_\phi C_\psi + S_\phi S_\theta S_\psi & -S_\phi C_\psi + C_\phi S_\theta S_\psi \\ -S_\theta & S_\phi C_\theta & C_\phi C_\theta \end{pmatrix}, \tag{7}$$

with Euler angles θ, ψ, and ϕ. Equation (6) can be rewritten in terms of Y_μ, Y'_μ, and Y''_μ as follows

$$\begin{aligned} \mathcal{D}_\mu = \partial_\mu &- iY_\mu \left(-S_\theta g'_d Q_d + C_\theta S_\psi g'_c Q_c + C_\theta C_\psi g'_a Q_a\right) \\ &- iY'_\mu \left[C_\theta S_\phi g'_d Q_d + \left(C_\phi C_\psi + S_\theta S_\phi S_\psi\right) g'_c Q_c \right. \\ &\left. + \left(C_\psi S_\theta S_\phi - C_\phi S_\psi\right) g'_a Q_a\right] \\ &- iY''_\mu \left[C_\theta C_\phi g'_d Q_d + \left(-C_\psi S_\phi + C_\phi S_\theta S_\psi\right) g'_c Q_c \right. \\ &\left. + \left(C_\phi C_\psi S_\theta + S_\phi S_\psi\right) g'_a Q_a\right]. \end{aligned} \tag{8}$$

[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_μ has the hypercharge Q_Y given in (5) we fix the first column of the rotation matrix $\mathcal{R}$

$$\begin{pmatrix} C_\mu \\ B_\mu \\ X_\mu \end{pmatrix} = \begin{pmatrix} Y_\mu \frac{1}{6} g_Y/g'_a & \cdots \\ -Y_\mu \frac{1}{2} g_Y/g'_c & \cdots \\ Y_\mu \frac{1}{2} g_Y/g'_d & \cdots \end{pmatrix}, \tag{9}$$

and we determine the value of the two associated Euler angles

$$\theta = -\arcsin\left[\frac{1}{2} g_Y/g'_d\right] \tag{10}$$

and

$$\psi = \arcsin\left[-\frac{1}{2} g_Y/(g'_c C_\theta)\right]. \tag{11}$$

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

$$\left(-\frac{1}{2g'_c}\right)^2 = \frac{1}{g_Y^2} - \left(\frac{c_1}{6g'_a}\right)^2 - \left(\frac{1}{2g'_d}\right)^2, \tag{12}$$

with g'_a fixed by the relation $g_3(M_s) = \sqrt{6}\, g'_a(M_s)$.[2] In our calculation we take $M_s = 20$ TeV as a reference point for running down to 1.8 TeV the g'_a coupling, ignoring mass threshold effects of stringy states. This yields $g'_a = 0.36$. We have checked that the running of the g'_a coupling does not change significantly for different values of the string scale. The third Euler angle ϕ and the coupling g'_d will be determined by requiring sufficient suppression to leptons to accommodate (3) and a (pre-cut) production rate $\sigma(pp \to Z') \times \mathcal{B}(Z' \to jj)$ in agreement with (2).

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

$$\begin{aligned}\mathcal{L} &= \frac{1}{2}\sqrt{g_Y^2 + g_2^2} \sum_f \left(\epsilon_{f_L} \bar{\psi}_{f_L} \gamma^\mu \psi_{f_L} + \epsilon_{f_R} \bar{\psi}_{f_R} \gamma^\mu \psi_{f_R}\right) Z'_\mu \\ &= \sum_f \left((g_{Y'} Q_{Y'})_{f_L} \bar{\psi}_{f_L} \gamma^\mu \psi_{f_L} + (g_{Y'} Q_{Y'})_{f_R} \bar{\psi}_{f_R} \gamma^\mu \psi_{f_R}\right) Z'_\mu,\end{aligned} \tag{13}$$

where each $\psi_{f_{L\,(R)}}$ is a fermion field with the corresponding γ^μ matrices of the Dirac algebra, and $\epsilon_{f_L, f_R} = v_q \pm a_q$, with v_q and a_q the vector and axial couplings respectively. From (8) and (13) we obtain the explicit form of the chiral couplings in terms of ϕ and g'_d

$$\begin{aligned}\epsilon_{u_L} = \epsilon_{d_L} &= \frac{2}{\sqrt{g_Y^2 + g_2^2}} (C_\psi S_\theta S_\psi - C_\phi S_\psi) g'_a, \\ \epsilon_{u_R} &= -\frac{2}{\sqrt{g_Y^2 + g_2^2}} [C_\theta S_\phi g'_d + (C_\psi S_\theta S_\psi - C_\phi S_\psi) g'_a], \\ \epsilon_{d_R} &= \frac{2}{\sqrt{g_Y^2 + g_2^2}} [C_\theta S_\phi g'_d - (C_\psi S_\theta S_\psi - C_\phi S_\psi) g'_a].\end{aligned} \tag{14}$$

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

$$\begin{aligned}&\Gamma(Z' \to f\bar{f}) \\ &= \frac{G_F M_Z^2}{6\pi\sqrt{2}} N_c C(M_{Z'}^2) M_{Z'} \sqrt{1-4x}\left[v_f^2(1+2x) + a_f^2(1-4x)\right],\end{aligned} \tag{15}$$

[2] Throughout g_3 and g_2 are the strong and weak gauge coupling constants.

where G_F is the Fermi coupling constant, $C(M_{Z'}^2) = 1 + \alpha_s/\pi + 1.409(\alpha_s/\pi)^2 - 12.77(\alpha_s/\pi)^3$, $\alpha_s = \alpha_s(M_{Z'})$ is the strong coupling constant at the scale $M_{Z'}$, $x = m_f^2/M_{Z'}^2$, and $N_c = 3$ 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,

$$\begin{aligned}\Gamma(Z' \to ZZ) &= \frac{c_1^2 \sin^2\theta_W M_{Z'}^3}{192\pi M_Z^2}\left(1 - \frac{4M_Z^2}{M_{Z'}^2}\right)^{5/2} \\ &\approx c_1^2\,(45\text{ GeV})\left(\frac{M_{Z'}}{\text{TeV}}\right)^3 + \cdots,\end{aligned} \tag{16}$$

$$\begin{aligned}\Gamma(Z' \to W^+W^-) &= \frac{c_2^2 M_{Z'}^3}{48\pi M_W^2}\left(1 - \frac{4M_W^2}{M_{Z'}^2}\right)^{5/2} \\ &\approx c_2^2\,(1.03\text{ TeV})\left(\frac{M_{Z'}}{\text{TeV}}\right)^3 + \cdots,\end{aligned} \tag{17}$$

$$\begin{aligned}\Gamma(Z' \to Z\gamma) &= \frac{c_1^2 \cos^2\theta_W M_{Z'}^3}{96\pi M_Z^2}\left(1 - \frac{M_Z^2}{M_{Z'}^2}\right)^3 \left(1 + \frac{M_Z^2}{M_{Z'}^2}\right) \\ &\approx c_1^2\,(307\text{ GeV})\left(\frac{M_{Z'}}{\text{TeV}}\right)^3 + \cdots.\end{aligned} \tag{18}$$

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

$$\sigma(pp \to Z') \simeq 5.2\left(\frac{2\Gamma(Z' \to u\bar{u}) + \Gamma(Z' \to d\bar{d})}{\text{GeV}}\right)\text{ fb}. \tag{19}$$

Next, we scan the parameter space to obtain agreement with (1) to (4). In Fig. 1 we show contour plots, in the (g'_d, ϕ) plane, for constant $\sigma(pp \to Z') \times \mathcal{B}(Z' \to jj)$, $\sigma(pp \to Z') \times \mathcal{B}(Z' \to e^+e^-)$, and $\sigma(pp \to Z') \times \mathcal{B}(Z' \to W^+W^-)$. 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 $g'_d(M_s) = 0.29$, $c_1 = 0.08$, and $c_2 = 0.02$ yield $\sigma(pp \to Z') \simeq 228$ fb, $\mathcal{B}(Z' \to jj) \simeq 0.54$, $\mathcal{B}(Z' \to e^+e^-) \simeq 8.9 \times 10^{-4}$, $\mathcal{B}(Z' \to WW/ZZ) \simeq 3.4 \times 10^{-2}$, which are consistent with (1), (2), and (3) at the 1σ level. In addition, $\mathcal{B}(Z' \to HZ) \simeq 7.4 \times 10^{-4}$. Thus, the upper limit set by (4) is also satisfied by our fiducial values of ϕ, 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 ($H_u\bar{u}q$, $H_d\bar{d}q$, $H_d\bar{e}\ell$, $H_\nu\bar{\nu}\ell$) invariant under all three $U(1)$'s [27]. Two "supersymmetric" Higgses $H_u \equiv H_\nu$ and H_d (with charges $Q_a = Q_c = 0$, $Q_d = 1$, $Q_Y = 1/2$ and $Q_a = Q_c = 0$, $Q_d = -1$, $Q_Y = -1/2$) are sufficient to give masses to all the chiral fermions. Here, $\langle H_u\rangle = \binom{0}{v_u}$, $\langle H_d\rangle = \binom{v_d}{0}$, $v = \sqrt{v_u^2 + v_d^2} = 174$ GeV, and $\tan\beta \equiv v_u/v_d$.

The last two terms in the covariant derivative

$$\begin{aligned}\mathcal{D}_\mu = \partial_\mu &- i\frac{1}{\sqrt{g_2^2 + g_Y^2}} Z_\mu (g_2^2 T^3 - g_Y^2 Q_Y) \\ &\times -i g_{Y'} Y_\mu{}' Q_{Y'} - i g_{Y''} Y_\mu{}'' Q_{Y''},\end{aligned} \tag{20}$$

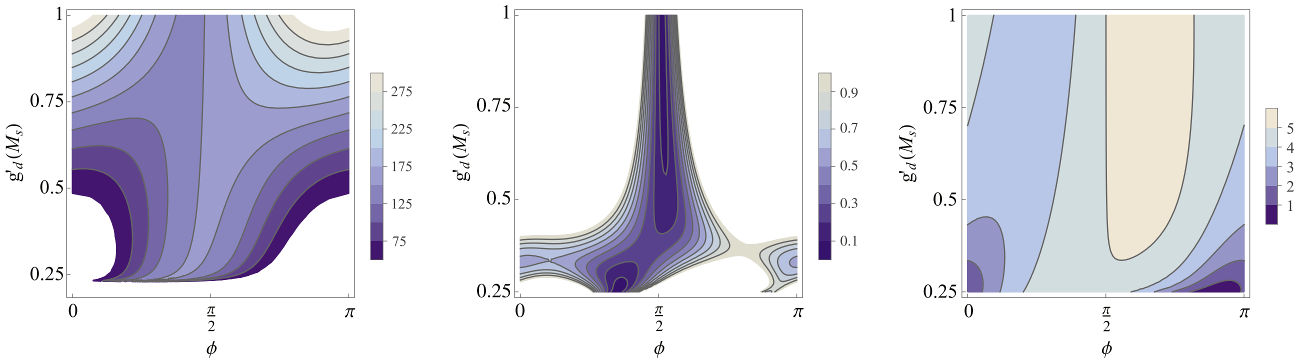

Fig. 1. Contours of constant cross section $pp \to Z'$ times branching into dijet (left), e^+e^- (middle), and diboson (right), for $M_{Z'} \simeq 1.8$ TeV and $\sqrt{s} = 8$ TeV. The color encoded scales are in fb. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

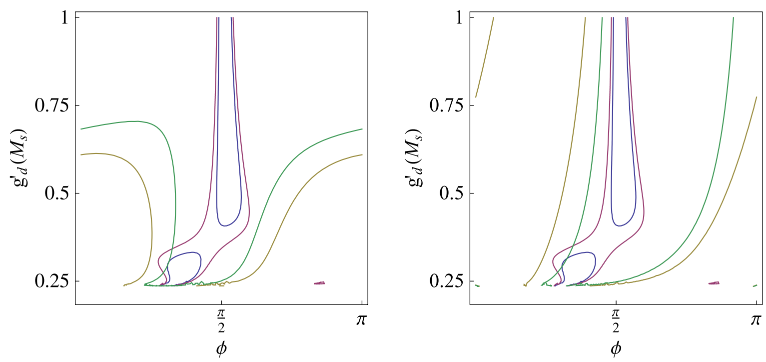

Fig. 2. Best fit of cross section contours $pp \to Z'$ times branching into dijet/leptons (left) and $pp \to Z'$ times branching into dibosons (right), for $M_{Z'} \simeq 1.8$ TeV and $\sqrt{s} = 8$ TeV. The blue and red contours correspond to $\sigma(pp \to Z') \times \mathcal{B}(e^+e^-) = 0.2$ and 0.3 fb, respectively. The yellow and green contours on the left correspond to $\sigma(pp \to Z') \times \mathcal{B}(jj) = 91$ and 123 fb, respectively. The yellow and green contours on the right correspond to $\sigma(pp \to Z') \times \mathcal{B}(W^+W^-) = 4$ and 4.5 fb, respectively. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

Table 2
Chiral couplings of Z' and Z'' gauge bosons.

Fields	$g_{Y'} Q_{Y'}$	$g_{Y''} Q_{Y''}$
U_R	−0.24	0.26
D_R	−0.41	0.02
L_L	0.06	0.23
E_R	−0.02	0.10
Q_L	−0.33	0.14
N_R	0.15	0.35

are conveniently written as

$$-i\frac{x_{H_i}}{v_i}\overline{M}_Z Y'_\mu - i\frac{y_{H_i}}{v_i}\overline{M}_Z Y''_\mu \tag{21}$$

for each Higgs H_i, with $T^3 = \sigma^3/2$, where for the two Higgs doublets

$$x_{H_u} = -x_{H_d} = 1.9\sqrt{{g'_d}^2 - 0.032}\, S_\phi \tag{22}$$

and

$$y_{H_u} = -y_{H_d} = 1.9\sqrt{{g'_d}^2 - 0.032}\, C_\phi\,. \tag{23}$$

The Higgs field kinetic term together with the GS mass terms $(-\frac{1}{2}M'^2 Y'_\mu Y'^\mu - \frac{1}{2}M''^2 Y''_\mu Y''^\mu)$ yield the following mass square matrix for the $Z - Z'$ mixing,

$$\begin{pmatrix} \overline{M}_Z^2 & \overline{M}_Z^2(x_{H_u}C_\beta^2 - x_{H_d}S_\beta^2) & \overline{M}_Z^2(y_{H_u}C_\beta^2 - y_{H_d}S_\beta^2) \\ \overline{M}_Z^2(x_{H_u}C_\beta^2 - x_{H_d}S_\beta^2) & \overline{M}_Z^2(C_\beta^2 x_{H_u}^2 + S_\beta^2 x_{H_d}^2) + M'^2 & \overline{M}_Z^2(C_\beta^2 x_{H_u}y_{H_u} + S_\beta^2 x_{H_d}y_{H_d}) \\ \overline{M}_Z^2(y_{H_u}C_\beta^2 - y_{H_d}S_\beta^2) & \overline{M}_Z^2(C_\beta^2 x_{H_u}y_{H_u} + S_\beta^2 x_{H_d}y_{H_d}) & \overline{M}_Z^2(y_{H_u}^2 C_\beta^2 + y_{H_d}^2 S_\beta^2) + M''^2 \end{pmatrix},$$

which does not impose any constraint on the $\tan\beta$ parameter. We have verified that, for our fiducial values of ϕ and g'_d, if $M_{Z''} \gtrsim M_{Z'}$ 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 WW/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 Stückelberg 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 $U(1)$ fields 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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