Z′-gauge bosons as harbingers of low-mass strings

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Massive Z'-gauge bosons act as excellent harbingers for string compactifications with a low string scale. In D-brane models they are associated to $U(1)$ gauge symmetries that are either anomalous in four dimensions or exhibit a hidden higher-dimensional anomaly. We discuss the possible signals of massive Z'-gauge bosons at hadron collider machines (Tevatron, LHC) in a minimal D-brane model consisting out of four stacks of D-branes. In this construction, there are two massive gauge bosons, which can be naturally associated with baryon number B and $B - L$ (L being lepton number). Here baryon number is always anomalous in four dimensions, whereas the presence of a four-dimensional $B - L$ anomaly depends on the $U(1)$-charges of the right-handed neutrinos. In case $B - L$ is anomaly-free, a mass hierarchy between the two associated Z'-gauge bosons can be explained. In our phenomenological discussion about the possible discovery of massive Z'-gauge bosons, we take as a benchmark scenario the dijet plus W signal, recently observed by the CDF Collaboration at Tevatron. It reveals an excess in the dijet-mass range 150 GeV/c², 4.1σ beyond standard modelexpectations. We show that in the context of low-mass string theory this excess can be associated with the production and decay of a leptophobic Z', a singlet partner of $SU(3)$ gluons coupled primarily to baryon number. Even if the CDF signal disappears, as indicated by the more recent D0 results, our analysis can still serve as the basis for future experimental search for massive Z'-gauge bosons in low string scale models. We provide the relevant cross sections for the production of Z'-gauge bosons in the TeV region, leading to predictions that are within reach of the present or the next LHC run.

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I. INTRODUCTION

Very recently, the CERN Large Hadron Collider (LHC) has fired mankind into a new era in particle physics. The $SU(3)_C \times SU(2)_L \times U(1)_Y$ standard model (SM) of electroweak and strong interactions was once again severely tested with a dataset corresponding to an integrated luminosity of ~ 4.9 fb^{-1} of pp collisions collected at $\sqrt{s} = 7$ TeV. The SM agrees remarkable well with LHC7 data, but has rather troubling weaknesses and appears to be a somewhat *ad hoc* theory.

It has long been thought that the SM may be a subset of a more fundamental gauge theory. Several models have been proposed, using the fundamental principle of gauge invariance as guidepost. A common thread in most of these proposals is the realization of the SM within the context of D-brane TeV-scale string compactifications [1]. Such D-brane constructions extend the SM with several additional $U(1)$ symmetries [2].[1] The basic unit of gauge invariance for these models is a $U(1)$ field, so that a stack of N identical D-branes eventually generates a $U(N)$ theory

with the associated $U(N)$ gauge group. (For $N = 2$ the gauge group can be $Sp(1) \equiv SU(2)$ rather than $U(2)$.) Gauge interactions emerge as excitations of open strings with endpoints attached to the D-branes, whereas gravitational interactions are described by closed strings that can propagate in all nine spatial dimensions of string theory (these comprise parallel dimensions extended along the D-branes and transverse large extra dimensions).

In this paper we study the main phenomenological aspects of one particular D-brane model that contains two additional $U(1)$ symmetries, which can be chosen to be mostly baryon number B and $B - L$, where L is lepton number. This choice is very natural from the point of view of the SM. Moreover, with this choice of the two additional $U(1)$ gauge symmetries, one can obtain a natural mass gap between the light anomalous $U(1)_B$ gauge boson Z' and the heavier nonanomalous $U(1)_{B-L}$ gauge boson Z''. Our first goal is to survey the basic features of the gauge theory's prediction regarding the new mass sector and couplings. These features lead to new phenomena that can be probed using data from the Tevatron and the LHC. In particular the theory predicts additional gauge bosons that we will show are accessible at the hadron colliders.

The layout of the paper is as follows. In Sec. II we detail some desirable properties which apply to generic models

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with multiple $U(1)$ symmetries. We perform a renormalization group analysis for the running of the gauge couplings, pointing out that the gauge couplings of the two group factors $U(1)_a \times SU(N)_a = U(N)_a$ run differently towards low energies below the string scale. This observation has some interesting phenomenological consequences. Having so identified the general properties of the theory, in Sec. III we outline the basic setting of TeV-scale string compactifications and discuss general aspects of the $U(3)_C \times Sp(1)_L \times U(1)_L \times U(1)_R$ intersecting D-brane configuration that realize the SM by open strings. In Secs. IV and V we discuss the associated phenomenological aspects of anomalous $U(1)$ gauge bosons related to experimental searches for new physics at the Tevatron and at the LHC. Finally, in Sec. VI we explore predictions inherited from properties of the overarching string theory. Concretely, we study the LHC discovery potential for Regge excitations within the D-brane model discussed in this work. Our conclusions are collected in Sec. VII.

II. ABELIAN GAUGE COUPLINGS AT LOW ENERGIES

We begin with the covariant derivative for the $U(1)$ fields in the $1, 2, 3, \ldots$ basis in which it is assumed that the kinetic energy terms containing X^i_μ are canonically normalized

$$\mathcal{D}_\mu = \partial_\mu - i \sum g'_i Q_i X^i_\mu. \tag{1}$$

The relations between the $U(1)$ couplings g'_i and any non-Abelian counterparts are left open for now. We carry out an orthogonal transformation of the fields $X^i_\mu = \sum_j \mathcal{R}_{ij} Y^j_\mu$. The covariant derivative becomes

$$\mathcal{D}_\mu = \partial_\mu - i \sum_i \sum_j g'_i Q_i \mathcal{R}_{ij} Y^j_\mu = \partial_\mu - i \sum_j \bar{g}_j \bar{Q}_j Y^j_\mu, \tag{2}$$

where for each j

$$\bar{g}_j \bar{Q}_j = \sum_i g'_i Q_i \mathcal{R}_{ij}. \tag{3}$$

Next, suppose we are provided with normalization for the hypercharge (taken as $j = 1$)

$$Q_Y = \sum_i c_i Q_i; \tag{4}$$

hereafter we omit the bars for simplicity. Rewriting (3) for the hypercharge

$$g_Y Q_Y = \sum_i g'_i Q_i \mathcal{R}_{i1} \tag{5}$$

and substituting (4) into (5) we obtain

$$g_Y \sum_i Q_i c_i = \sum_i g'_i \mathcal{R}_{i1} Q_i. \tag{6}$$

One can think about the charges $Q_{i,p}$ as vectors with the components labeled by particles p. Assuming that these vectors are linearly independent, Eq. (6) reduces to

$$g_Y c_i = g'_i \mathcal{R}_{i1}, \tag{7}$$

or equivalently

$$\mathcal{R}_{i1} = \frac{g_Y c_i}{g'_i}. \tag{8}$$

Orthogonality of the rotation matrix, $\sum_i \mathcal{R}^2_{i1} = 1$, implies

$$g^2_Y \sum_i \left(\frac{c_i}{g'_i}\right)^2 = 1. \tag{9}$$

Then, the condition

$$P \equiv \frac{1}{g^2_Y} - \sum_i \left(\frac{c_i}{g'_i}\right)^2 = 0 \tag{10}$$

encodes the orthogonality of the mixing matrix connecting the fields coupled to the stack charges $Q_1, Q_2, Q_3, \ldots$ and the fields rotated, so that one of them, Y, couples to the hypercharge Q_Y.

A very important point is that the couplings that are running are those of the $U(1)$ fields; hence the β functions receive contributions from fermions and scalars, but not from gauge bosons. The one-loop correction to the various couplings are

$$\frac{1}{\alpha_Y(Q)} = \frac{1}{\alpha_Y(M_s)} - \frac{b_Y}{2\pi} \ln(Q/M_s), \tag{11}$$

$$\frac{1}{\alpha_i(Q)} = \frac{1}{\alpha_i(M_s)} - \frac{b_i}{2\pi} \ln(Q/M_s), \tag{12}$$

where

$$b_i = \frac{2}{3} \sum_f Q^2_{i,f} + \frac{1}{3} \sum_s Q^2_{i,s}, \tag{13}$$

with f and s indicating contribution from fermion and scalar loops, respectively.

Let us assume that the charges are orthogonal, $\sum_s Q_{i,s} Q_{j,s} = \sum_f Q_{i,f} Q_{j,f} = 0$ for $i \neq j$. Then Eq. (4) implies

$$\sum_s Q^2_{Y,s} = \sum_i c^2_i \sum_s Q^2_{i,s} \tag{14}$$

and the same thing for fermions, hence

$$b_Y = \sum_i c^2_i b_i. \tag{15}$$

On the other, the renormalization group (RG)-induced change of P defined in Eq. (10) reads

$$\Delta P = \Delta\left(\frac{1}{\alpha_Y}\right) - \sum_i c_i^2 \Delta\left(\frac{1}{\alpha_i}\right)$$

$$= \frac{1}{2\pi}\left(b_Y - \sum_i c_i^2 b_i\right)\ln(Q/M_s). \qquad (16)$$

Thus, $P = 0$ stays valid to one loop if the charges are orthogonal. An example of orthogonality is seen in the $U(3)_C \times Sp(1)_L \times U(1)$ D-brane model of [4,5], for which the various $U(1)$ assignments are given in Table I. In the three-stack D-brane models of [6], the charges are linearly independent, but not necessarily orthogonal. If the charges are not orthogonal, the RG equations controlling the running of couplings associated to different charges become coupled. One-loop corrections generate mixed kinetic terms for $U(1)$ gauge fields [7], greatly complicating the analysis.

Another important element of the RG analysis is that the relation for $U(N)$ unification, $g'_N = g_N/\sqrt{2N}$, holds only at M_s because the $U(1)$ couplings (g'_1, g'_2, g'_3) run differently from the non-Abelian $SU(3)$ (g_3) and $SU(2)$ (g_2).

In this paper we are interested in a minimal four-stack model $U(3)_C \times Sp(1)_L \times U(1)_L \times U(1)_R$, which has the attractive property of elevating the two major global symmetries of the SM (baryon number B and lepton number L) to local gauge symmetries. A schematic representation of the D-brane structure (to be discussed in detail in Sec. III) is shown in Fig. 1. The chiral fermion charges in Table II are not orthogonal as given ($Q_{1L} \cdot Q_{1R} \neq 0$). Orthogonality can be completed by including a right-handed neutrino with charges $Q_3 = 0$, $Q_{1L} = Q_{1R} = \pm1$, $Q_Y = 0$. We turn now to discuss the string origin and the compelling properties of this model.

III. GENERALITIES OF $U(3)_C \times Sp(1)_L \times U(1)_L \times U(1)_R$

The generic perturbative spectrum in intersecting D-brane models consists of products of unitary groups $U(N_i)$ associated to stacks of N_i coincident D-branes and matter in bifundamental representations. In the presence of orientifolds which are necessary for tadpole cancellation, and thus consistency of the theory, open strings become in general nonoriented allowing for orthogonal and
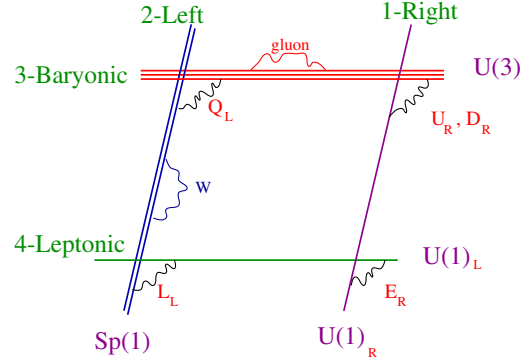

FIG. 1 (color online). Pictorial representation of the $U(1)_C \times Sp(1)_L \times U(1)_L \times U(1)_R$ D-brane model.

symplectic gauge group factors, as well as for symmetric and antisymmetric matter representations.

The minimal embedding of the SM particle spectrum requires at least three brane stacks [6] leading to three distinct models of the type $U(3)_C \times U(2)_L \times U(1)$ that were classified in [4,6]. Only one of them (model C of [4]) has the baryon number as symmetry that guarantees proton stability (in perturbation theory), and can be used in the framework of TeV strings. Moreover, since Q_2 (associated to the $U(1)$ of $U(2)_L$) does not participate in the hypercharge combination, $U(2)_L$ can be replaced by $Sp(1)_L$ leading to a model with one extra $U(1)$, the baryon number, besides hypercharge [5]. The quantum numbers of the chiral SM spectrum are given in Table I. Since the baryon number is anomalous, the extra Abelian gauge field becomes massive by the Green-Schwarz (GS) mechanism, 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 [8]. Given the three SM couplings and the hypercharge combination, this model has no free parameter in the coupling of Z' to the SM fields. Moreover, the lepton number is not a symmetry creating a problem with large neutrino masses through the Weinberg dimension-five operator $LLHH$ suppressed only by the TeV string scale. We therefore proceed to models with four D-brane stacks.

The SM embedding in four D-brane stacks leads to many more models that have been classified in [9,10]. In order to make a phenomenologically interesting choice, we first focus on models where $U(2)_L$ can be reduce to $Sp(1)$.

TABLE I. Quantum numbers of chiral fermions and Higgs doublet for $U(3)_C \times Sp(1)_L \times U(1)$.

Name	Representation	Q_3	Q_1	Q_Y
$\bar{U}_i$	$(\bar{3}, 1)$	-1	1	$-\frac{2}{3}$
$\bar{D}_i$	$(\bar{3}, 1)$	-1	-1	$\frac{1}{3}$
L_i	$(1, 2)$	0	1	$-\frac{1}{2}$
$\bar{E}_i$	$(1, 1)$	0	-2	1
Q_i	$(3, 2)$	1	0	$\frac{1}{6}$
H	$(1, 2)$	0	1	$-\frac{1}{2}$

TABLE II. Chiral fermion spectrum of the $U(3)_C \times Sp(1)_L \times U(1)_L \times U(1)_R$ D-brane model.

Name	Representation	Q_3	Q_{1L}	Q_{1R}	Q_Y
$\bar{U}_i$	$(\bar{3}, 1)$	-1	0	-1	$-\frac{2}{3}$
$\bar{D}_i$	$(\bar{3}, 1)$	-1	0	1	$\frac{1}{3}$
L_i	$(1, 2)$	0	1	0	$-\frac{1}{2}$
$\bar{E}_i$	$(1, 1)$	0	-1	1	1
Q_i	$(3, 2)$	1	0	0	$\frac{1}{6}$

Besides the fact that this reduces the number of extra $U(1)$'s, one avoids the presence of a problematic Peccei-Quinn symmetry, associated in general with the $U(1)$ of $U(2)_L$ under which Higgs doublets are charged [6]. We then impose baryon and lepton number symmetries that determine completely the model $U(3)_C \times Sp(1)_L \times U(1)_L \times U(1)_R$, as described in [10] (see subsection 4.2.4). The corresponding fermion quantum numbers are given in Table II, while the two extra $U(1)$'s are the baryon and lepton number, B and L, respectively; they are given by the following combinations:

$$B = Q_3/3; \qquad L = Q_{1L};$$
$$Q_Y = \frac{1}{6}Q_3 - \frac{1}{2}Q_{1L} + \frac{1}{2}Q_{1R}; \qquad (17)$$

or equivalently by the inverse relations:

$$Q_3 = 3B; \qquad Q_{1L} = L; \qquad Q_{1R} = 2Q_Y - (B - L). \qquad (18)$$

Note that with the "canonical" charges of the right-handed neutrino $Q_{1L} = Q_{1R} = -1$, the combination $B - L$ is anomaly-free, while for $Q_{1L} = Q_{1R} = +1$, both B and $B - L$ are anomalous. Actually, both choices guarantee orthogonality of the charges discussed in the previous section. As mentioned already, anomalous $U(1)$'s become massive necessarily due to the Green-Schwarz anomaly cancellation, but nonanomalous $U(1)$'s can also acquire masses due to effective six-dimensional anomalies associated for instance to sectors preserving $N = 2$ supersymmetry [8].[2] These two-dimensional "bulk" masses become therefore larger than the localized masses associated to four-dimensional anomalies, in the large volume limit of the two extra dimensions. Specifically for Dp-branes with $(p - 3)$-longitudinal compact dimensions the masses of the anomalous and, respectively, the non-anomalous $U(1)$ gauge bosons have the following generic scale behavior:

$$\text{anomalous } U(1)_a: M_{Z'} = g'_a M_s,$$
$$\text{non-anomalous } U(1)_a: M_{Z''} = g'_a M_s^3 V_2. \qquad (19)$$

Here g'_a is the gauge coupling constant associated to the group $U(1)_a$, given by $g'_a \propto g_s/\sqrt{V_\parallel}$ where g_s is the string coupling and $V_\parallel$ is the internal D-brane world-volume along the $(p - 3)$ compact extra dimensions, up to an order-one proportionality constant. Moreover, V_2 is the internal two-dimensional volume associated to the effective six-dimensional anomalies giving mass to the

nonanomalous $U(1)_a$.[3] E.g., for the case of D5-branes, whose common intersection locus is just four-dimensional Minkowski-space, $V_\parallel = V_2$ denotes the volume of the longitudinal, two-dimensional space along the two internal D5-brane directions. Since internal volumes are bigger than one in string units to have effective field theory description, the masses of nonanomalous $U(1)$-gauge bosons are generically larger than the masses of the anomalous gauge bosons. Since we want to identify the light Z' gauge boson with the baryon number, which is always anomalous, a hierarchy compared to the second $U(1)$-gauge boson Z'' can arise, if we identify Z'' with the anomaly-free combination $B - L$, and take the internal world-volume V_2 a bit larger than the string scale.[4] In summary, this model has two free parameters: one coupling and one angle in the two-dimensional space orthogonal to the hypercharge defining the direction of the corresponding Z'. Tuning the later, it can become leptophobic, while the former controls the strength of its interactions to matter. As discussed already, one can distinguish two cases: (i) when B and L have 4d anomalies, the mass ratio of the two extra gauge bosons (Z' and Z'') is fixed by the ratio of their gauge couplings, up to order-one coefficients; (ii) when $B - L$ is anomaly-free and gets a mass from effective six-dimensional anomalies, the mass ratio of the leftover anomalous $U(1)$ compared to the nonanomalous $U(1)$ is suppressed by the two-dimensional volume.

To summarize, we will analyze the phenomenology of two D-brane constructions with three mutually orthogonal $U(1)$ charges, in which the combination $B - L$ is either anomalous or anomaly-free. In the next section, we analyze these situations and study the regions of the parameter space where Z' is leptophobic and can accommodate the recent Tevatron data.

IV. LEPTOPHOBIC Z' AT THE TEVATRON

Taken at face value, the disparity between CDF [12,13] and D0 [14] results insinuates a commodious uncertainty

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[2]In fact, also the hypercharge gauge boson of $U(1)_Y$ can acquire a mass through this mechanism. In order to keep it massless, certain topological constraints on the compact space have to be met.

[3]It should be noted that in spite of the proportionality of the $U(1)_a$ masses to the string scale, these are not string excitations but zero modes. The proportionality to the string scale appears because the mass is generated from anomalies, via an analog of the GS anomaly cancellations: either four-dimensional anomalies, in which case the GS term is equivalent to a Stuckelberg mechanism, or from effective six-dimensional anomalies, in which case the mass term is extended in two more (internal) dimensions. The nonanomalous $U(1)_a$ can also grow a mass through a Higgs mechanism. The advantage of the anomaly mechanism versus an explicit vev of a scalar field is that the global symmetry survives in perturbation theory, which is a desired property for the baryon and lepton number, protecting proton stability and small neutrino masses.

[4]In [11] a different (possibly T-dual) scenario with D7-branes was investigated. In this case the masses of the anomalous and nonanomalous $U(1)$'s appear to exhibit a dependence on the entire six-dimensional volume, such that the nonanomalous masses become lighter than the anomalous ones.

as to whether there is an excess of events in the dijet system-invariant mass distribution of the associated production of a W boson with two jets (hereafter Wjj production). The M_{jj} excess showed up in 4.3 fb^{-1} of integrated luminosity collected with the CDF detector as a broad bump between about 120 and 160 GeV [12]. The CDF Collaboration fitted the excess (hundreds of events in the $\ell jj + \not{E}_T$ channel) to a Gaussian and estimated its production cross section times the dijet branching ratio to be 4 pb. This is roughly 300 times the SM Higgs rate $\sigma(p\bar{p} \to WH) \times$ BR$(H \to b\bar{b})$. For a search window of $120 - 200$ GeV, the excess significance above SM background (including systematics uncertainties) has been reported to be 3.2σ [12]. Recently, CDF has included an additional 3 fb^{-1} to their data sample, for a total of 7.3 fb^{-1}, and the statistical significance has grown to $\sim 4.8\sigma$ ($\sim 4.1\sigma$ including systematics) [13]. More recently, the D0 Collaboration released an analysis (which closely follows the CDF analysis) of their Wjj data finding "no evidence for anomalous resonant dijet production" [14]. Using an integrated luminosity of 4.3 fb^{-1} they set a 95% confidence level (C.L.) upper limit of 1.9 pb on a resonant Wjj production cross section.

In a previous work [15] we presented an explanation of the CDF data by identifying the resonance with a Z' inherent to D-brane TeV-scale string compactifications [1]. In this section we repeat our analysis but with two highly significant changes. First, we allow for the experimental uncertainty by focusing on a wide range ($1.6 - 6.0$ pb) of the (precut) Wjj resonant production cross section. This interpolates between the CDF and D0 results. Second, we turn our attention to a different D-brane model which has the attractive property of elevating the two major global symmetries of the SM (baryon number B and lepton number L) to local gauge symmetries.

Related explanations for the CDF anomaly based on an additional leptophobic Z' gauge boson have been offered [16]. Alternative new physics explanations include technicolor, new Higgs sectors, supersymmetry with and without R parity violation, color octect production, quirk exchange, and more [17]. There are also attempts to explain this puzzle within the context of the SM [18].

The suppressed coupling to leptons (or more specifically, to electrons and muons) is required to evade the strong constraints of the Tevatron Z' searches in the dilepton mode [19] and LEP-II measurements of $e^+e^- \to e^+e^-$ above the Z-pole [20]. In complying with the precision demanded of our phenomenological approach it would be sufficient to consider a 1% branching fraction to leptons as consistent with the experimental bound. This approximation is within a factor of a few of *model-independent* published experimental bounds. In addition, the mixing of the Z' with the SM Z boson should be extremely small [21,22] to be compatible with precision measurements at the Z-pole by the LEP experiments [23].

All existing dijet-mass searches via direct production at the Tevatron are limited to $M_{jj} > 200$ GeV [24] and

therefore cannot constrain the existence of a Z' with $M_{Z'} \simeq 150$ GeV. The strongest constraint on a light leptophobic Z' comes from the dijet search by the UA2 Collaboration, which has placed a 90% C.L. upper bound on $\sigma(p\bar{p} \to Z') \times$ BR$(Z' \to jj)$ in this energy range [25]. A comprehensive model-independent analysis incorporating Tevatron and UA2 data to constrain the Z' parameters for predictive purposes at the LHC has been recently presented [26].[5]

In the $U(3)_C \times Sp(1)_L \times U(1)_L \times U(1)_R$ D-brane model the Q_3, Q_{1L}, Q_{1R} content of the hypercharge operator is given by,

$$Q_Y = c_1 Q_{1R} + c_3 Q_3 + c_4 Q_{1L}, \tag{20}$$

with $c_1 = 1/2$, $c_3 = 1/6$, and $c_4 = -1/2$.

The covariant derivative (1) can be rewritten as

$$\mathcal{D}_\mu = \partial_\mu - ig_3' C_\mu Q_3 - ig_4' \tilde{B}_\mu Q_{1L} - ig_1' B_\mu Q_{1R}. \tag{21}$$

The fields $C_\mu, \tilde{B}_\mu, B_\mu$ are related to Y_μ, Y_μ' and Y_μ'' by the rotation matrix

$$\mathcal{R} = \begin{pmatrix} C_\theta C_\psi & -C_\phi S_\psi + S_\phi S_\theta C_\psi & S_\phi S_\psi + C_\phi S_\theta C_\psi \\ C_\theta S_\psi & C_\phi C_\psi + S_\phi S_\theta S_\psi & -S_\phi C_\psi + C_\phi S_\theta S_\psi \\ -S_\theta & S_\phi C_\theta & C_\phi C_\theta \end{pmatrix}, \tag{22}$$

with Euler angles θ, ψ, and ϕ. Equation (21) can be rewritten in terms of Y_μ, Y_μ', and Y_μ'' as follows

$$\begin{aligned} \mathcal{D}_\mu = \partial_\mu &- iY_\mu(-S_\theta g_1' Q_{1R} + C_\theta S_\psi g_4' Q_{1L} \\ &+ C_\theta C_\psi g_3' Q_3) - iY_\mu'[C_\theta S_\phi g_1' Q_{1R} + (C_\phi C_\psi \\ &+ S_\theta S_\phi S_\psi)g_4' Q_{1L} + (C_\psi S_\theta S_\phi - C_\phi S_\psi)g_3' Q_3] \\ &- iY_\mu''[C_\theta C_\phi g_1' Q_{1R} + (-C_\psi S_\phi \\ &+ C_\phi S_\theta S_\psi)g_4' Q_{1L} + (C_\phi C_\psi S_\theta + S_\phi S_\psi)g_3' Q_3]. \end{aligned} \tag{23}$$

Now, by demanding that Y_μ has the hypercharge Q_Y given in Eq. (20) we fix the first column of the rotation matrix $\mathcal{R}$

$$\begin{pmatrix} C_\mu \\ \tilde{B}_\mu \\ B_\mu \end{pmatrix} = \begin{pmatrix} Y_\mu c_3 g_Y / g_3' & \cdots \\ Y_\mu c_4 g_Y / g_4' & \cdots \\ Y_\mu c_1 g_Y / g_1' & \cdots \end{pmatrix}, \tag{24}$$

and we determine the value of the two associated Euler angles

$$\theta = -\arcsin[c_1 g_Y / g_1'] \tag{25}$$

and

$$\psi = \arcsin[c_4 g_Y / (g_4' C_\theta)]. \tag{26}$$

The couplings g_1' and g_4' are related through the orthogonality condition (10)

$$\left(\frac{c_4}{g_4'}\right)^2 = \frac{1}{g_Y^2} - \left(\frac{c_3}{g_3'}\right)^2 - \left(\frac{c_1}{g_1'}\right)^2, \tag{27}$$

with g_3' fixed by the relation $g_3(M_s) = \sqrt{6}g_3'(M_s)$. In what follows, we take $M_s = 5$ TeV as a reference point for running down to 150 GeV the g_3' coupling using (12), that is *ignoring mass threshold effects of stringy states.* This yields $g_3' = 0.383$. We have checked that the running of the g_3' coupling does not change significantly within the LHC range, i.e., 3 TeV $< M_s < 10$ TeV.

The phenomenological analysis thus far has been formulated in terms of the mass-diagonal basis set of gauge fields (Y, Y', Y''). As a result of the electroweak phase transition, the coupling of this set with the Higgses will induce mixing, resulting in a new mass-diagonal basis set (Z, Z', Z''). It will suffice to analyze only the 2×2 system (Y, Y') to see that the effects of this mixing are totally negligible. We consider simplified zeroth and first-order (mass)^2 matrices

$$(M^2)^{(0)} = \begin{pmatrix} 0 & 0 \\ 0 & M'^2 \end{pmatrix} \quad (M^2)^{(1)} = \begin{pmatrix} \bar{M}_Z^2 & \epsilon \\ \epsilon & m'^2 \end{pmatrix}, \tag{28}$$

where M' is the mass of the Y' gauge field, $\bar{M}_Z = \sqrt{(g_2^2 v^2 + g_Y^2 v^2)/2}$ is the usual tree-level formula for the mass of the Z particle in the electroweak theory (before mixing), g_2 is the electroweak coupling constant, v is the vacuum expectation value of the Higgs field, and ϵ, m'^2 are of $\mathcal{O}(\bar{M}_Z^2)$.

Standard Rayleigh-Schrodinger perturbation theory then provides the (mass)^2 (to second order in $\bar{M}_Z^2$) and wave functions (to first order) of the mass-diagonal eigenfields (Z, Z') corresponding to (Y, Y')

$$M_Z^2 = \bar{M}_Z^2 - \left(\frac{\epsilon^2}{M'^2}\right), \quad M_{Z'}^2 = M'^2 + m'^2 + \left(\frac{\epsilon^2}{M'^2}\right), \tag{29}$$

and

$$Z = Y - \left(\frac{\epsilon}{M'^2}\right)Y', \quad Z' = Y' + \left(\frac{\epsilon}{M'^2}\right)Y. \tag{30}$$

From Eqs. (29) and (30) the shift in the mass of the Z is given by $\delta M_Z^2 = (\epsilon/M')^2$, so that $\epsilon = M'\sqrt{2M_Z\delta M_Z}$. The admixture of Y in the mass-diagonal field Z' is

$$\theta = \frac{\epsilon}{M'^2} = \frac{M_Z}{M'}\sqrt{\frac{2\delta M_Z}{M_Z}} \simeq 0.004, \tag{31}$$

where we have taken $\delta M_Z = 0.0021$ GeV [28]. Interference effects which are proportional to θ are present in processes with fermions (e.g., Drell-Yan). However, these vanish at the peak of the resonance. Because of the smallness of θ, modifications of SM partial decay rates of the Z are negligible. (See, e.g., [21], for an analysis of such effects.) Remaining effects are order $\theta^2 \simeq 1.6 \times 10^{-5}$, and therefore all further discussion will be, with negligible

error, in terms of Z'. By the same token, the admixture of Y' in the eigenfield Z is negligible, so that the discussion henceforth will reflect $Z \simeq Y$ and $\bar{M}_Z^2 \simeq M_Z^2$.

The third Euler angle ϕ and the coupling g_1' are determined by requiring sufficient suppression ($\lesssim 1\%$) to leptons, a (precut) production rate $1.6 \lesssim \sigma(p\bar{p} \rightarrow WZ') \times \text{BR}(Z' \rightarrow jj) \lesssim 6.0$ pb at $\sqrt{s} = 1.96$ TeV, and compatibility with the 90% C.L. upper limit reported by the UA2 Collaboration on $\sigma(p\bar{p} \rightarrow Z') \times \text{BR}(Z' \rightarrow jj)$ at $\sqrt{s} = 630$ GeV [25].

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

$$\mathcal{L} = \frac{1}{2}\sqrt{g_Y^2 + g_2^2}\sum_f (\epsilon_{f_L}\bar{\psi}_{f_L}\gamma^\mu\psi_{f_L} + \epsilon_{f_R}\bar{\psi}_{f_R}\gamma^\mu\psi_{f_R})Z'_\mu$$
$$= \sum_f ((g_{Y'}Q_{Y'})_{f_L}\bar{\psi}_{f_L}\gamma^\mu\psi_{f_L} + (g_{Y'}Q_{Y'})_{f_R}\bar{\psi}_{f_R}\gamma^\mu\psi_{f_R})Z'_\mu, \tag{32}$$

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. The (precut) Wjj production rate at the Tevatron $\sqrt{s} = 1.96$ pb, for arbitrary couplings and $M_{Z'} \simeq 150$ GeV, is found to be [26]

$$\sigma(p\bar{p} \rightarrow WZ') \times \text{BR}(Z' \rightarrow jj)$$
$$\simeq [0.719(\epsilon_{u_L}^2 + \epsilon_{d_L}^2) + 5.083\epsilon_{u_L}\epsilon_{d_L}]$$
$$\times \Gamma(\phi, g_1')_{Z' \rightarrow q\bar{q}}\text{pb}, \tag{33}$$

where $\Gamma(\phi, g_1')_{Z' \rightarrow q\bar{q}}$ is the hadronic branching fraction. The presence of a W in the process shown in Fig. 2 restricts the contribution of the quarks to be purely left-handed. The dijet production rate at the UA2 $\sqrt{s} = 630$ GeV can be parametrized as follows [26]

$$\sigma(p\bar{p} \rightarrow Z') \times \text{BR}(Z' \rightarrow jj)$$
$$\simeq \frac{1}{2}[773(\epsilon_{u_L}^2 + \epsilon_{u_R}^2) + 138(\epsilon_{d_L}^2 + \epsilon_{d_R}^2)]$$
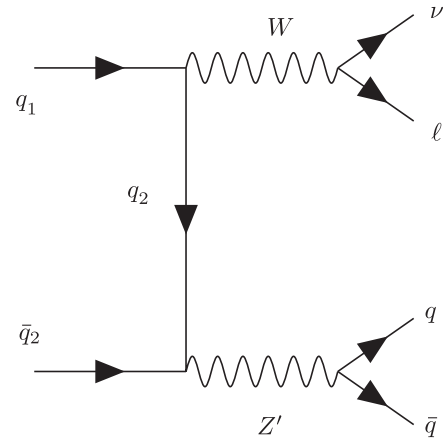$$\times \Gamma(\phi, g_1')_{Z' \rightarrow q\bar{q}}\text{pb}. \tag{34}$$

FIG. 2. Feynman diagram for $q\bar{q} \rightarrow WZ' \rightarrow \nu\ell jj$.

(Our numerical calculation using CTEQ6 [29] agrees within 5% with the result of [26].) The dilepton production rate at UA2 energies is given by

$$\sigma(p\bar{p} \to Z') \times \mathrm{BR}(Z' \to \ell\bar{\ell})$$

$$\simeq \frac{1}{2}[773(\epsilon_{u_L}^2 + \epsilon_{u_R}^2) + 138(\epsilon_{d_L}^2 + \epsilon_{d_R}^2)]$$

$$\times \Gamma(\phi, g_1')_{Z' \to \ell\bar{\ell}} \mathrm{pb}, \qquad (35)$$

where $\Gamma(\phi, g_1')_{Z' \to \ell\bar{\ell}}$ is the leptonic branching fraction. From (23) and (32) we obtain the explicit form of the chiral couplings in terms of ϕ and g_1'

$$\epsilon_{u_L} = \epsilon_{d_L} = \frac{2}{\sqrt{g_Y^2 + g_2^2}}(C_\psi S_\theta S_\psi - C_\phi S_\psi)g_3',$$

$$\epsilon_{u_R} = -\frac{2}{\sqrt{g_Y^2 + g_2^2}}[C_\theta S_\phi g_1' + (C_\psi S_\theta S_\psi - C_\phi S_\psi)g_3'],$$

$$\epsilon_{d_R} = \frac{2}{\sqrt{g_Y^2 + g_2^2}}[C_\theta S_\phi g_1' - (C_\psi S_\theta S_\psi - C_\phi S_\psi)g_3']. \quad (36)$$

Using (33)–(36) the ratio of branching ratios of electrons to quarks is minimized within the $\phi - g_1'$ parameter space, subject to sufficient Wjj production and saturation of the 90% C.L. upper limit. For a (precut) Wjj production varying between $1.6 - 6.0$ pb, one possible allowed region of the $\phi - g_1'$ parameter space is found to be $-1.16 \lesssim \phi \lesssim 2.12$ and $0.20 \lesssim g_1' \lesssim 0.27$.

A. Anomalous $B - L$

Let us first consider a reference point of the $\phi - g_1'$ parameter space consistent with the recent D0 limit [14]. For $\phi = -1.16$ and $g_1' = 0.27$, corresponding to a suppression $\Gamma_{Z' \to e^+ e^-}/\Gamma_{Z' \to q\bar{q}} \sim 1\%$, we obtain $\sigma(p\bar{p} \to WZ') \times \mathrm{BR}(Z' \to jj) \simeq 1.6$ pb at $\sqrt{s} = 1.96$ TeV. From Eqs. (25) and (26), this also corresponds to $\theta = -0.722$, $\psi = -1.37$.[6] All the couplings of the Y' (or equivalently Z') gauge boson are now determined and contained in Eq. (23). Numerical values are given in Table III under the heading of $g_{Y'}Q_{Y'}$.

In Fig. 3 we show a comparison of $\sigma(p\bar{p} \to Z') \times \mathrm{BR}(Z' \to jj)$ at $\sqrt{s} = 630$ GeV and the UA2 90% C.L. upper limit on the production of a gauge boson decaying into two jets. One can see that for our fiducial values, $\phi = -1.16$ and $g_1' = 0.27$, the single $Z' \to jj$ production cross section saturates the UA2 90% C.L. upper limit.

The Tevatron rate for the associated production channels [26]

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[6]The UA2 data has a dijet-mass resolution $\Delta M_{jj}/M_{jj} \sim 10\%$ [25]. Therefore, at 150 GeV the dijet-mass resolution is about 15 GeV. This is much larger than the resonance width, which is calculated to be $\Gamma(Z' \to f\bar{f}) \simeq 5$ GeV [30].

TABLE III. Chiral couplings of Y' and Y'' gauge bosons for $\phi = -1.16$ and $g_1' = 0.27$.

Name	$g_{Y'}Q_{Y'}$	$g_{Y''}Q_{Y''}$
$\bar{U}_i$	-0.013	-0.411
$\bar{D}_i$	-0.386	-0.251
L_i	-0.125	-0.125
$\bar{E}_i$	-0.061	-0.027
Q_i	0.199	0.331

$$\sigma(p\bar{p} \to ZZ') \times \mathrm{BR}(Z' \to jj)$$

$$\simeq \frac{1}{4}[381.5\epsilon_{u_L}^2 + 221\epsilon_{u_R}^2 + 1323\epsilon_{d_L}^2$$

$$+ 44.1\epsilon_{d_R}^2] \times \Gamma_{Z' \to q\bar{q}} \mathrm{fb} \qquad (37)$$

and

$$\sigma(p\bar{p} \to \gamma Z') \times \mathrm{BR}(Z' \to jj)$$

$$\simeq \frac{1}{2}[767(\epsilon_{u_L}^2 + \epsilon_{u_R}^2) + 72.7(\epsilon_{d_L}^2 + \epsilon_{d_R}^2)] \times \Gamma_{Z' \to q\bar{q}} \mathrm{fb} \qquad (38)$$

is always substantially smaller. (In (37) the SM leptonic branching fractions have been included to ease comparison with the experiment.) It is straightforward to see that these processes should not yet have been observed at the Tevatron.

The second strong constraint on the model derives from the mixing of the Z and the Y' through their coupling to the two Higgs doublets H_1 and H_2. The criteria we adopt here
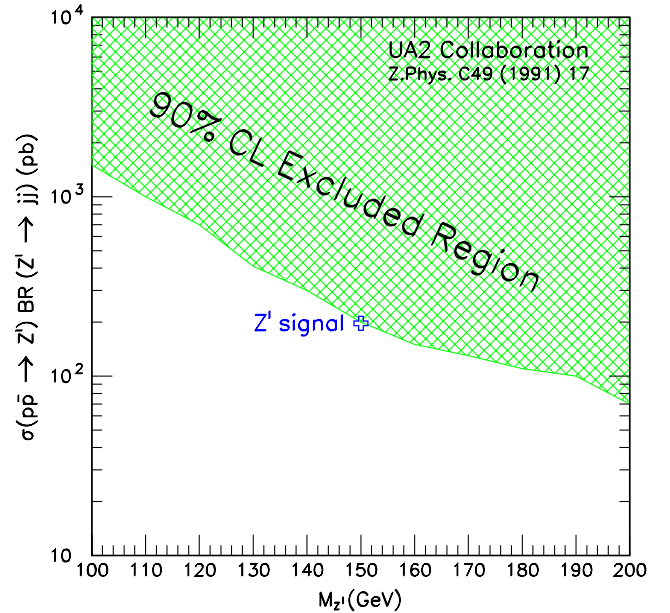

FIG. 3 (color online). Comparison of the total cross section for the production of $p\bar{p} \to Z' \to jj$ at $\sqrt{s} = 630$ GeV and the UA2 90% C.L. upper limit on the production of a gauge boson decaying into two jets [25]. We have taken $\phi = -1.16$ and $g_1' = 0.27$.

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. From Table II, $\bar{u} q$ has the charges $(0, 0, -1)$ and $\bar{d} q$ has $(0, 0, 1)$. So the Higgs H_u has $Q_3 = Q_{1L} = 0$, $Q_{1R} = 1$, $Q_Y = 1/2$, whereas H_d has opposite charges $Q_3 = Q_{1L} = 0$, $Q_{1R} = -1$, $Q_Y = -1/2$.

Let us consider first the case in which the right-handed neutrino has the following $U(1)$ charges $(0, 1, 1)$. As explained before, $B - L$ is then anomalous, and there is no hierarchy among the masses of Z' and Z''. For simplicity we can assume that $H_u \equiv H_1$ and $H_d = H_1^*$, with

$$\langle H_1 \rangle = \begin{pmatrix} 0 \\ v_1 \end{pmatrix}.$$

For the second Higgs field $H_\nu \equiv H_2$ the charges are $Q_3 = 0$, $Q_{1L} = -2$, $Q_{1R} = -1$, $Q_Y = 1/2$.[7] Here,

$$\langle H_2 \rangle = \begin{pmatrix} 0 \\ v_2 \end{pmatrix},$$

$v = \sqrt{v_1^2 + v_2^2} = 174$ GeV, and $\tan\delta \equiv v_1/v_2$. The last two terms in the covariant derivative,

$$\mathcal{D}_\mu = \partial_\mu - i \frac{1}{\sqrt{g_2^2 + g_Y^2}} Z_\mu (g_2^2 T^3 - g_Y^2 Q_Y)$$

$$- i g_{Y'} Y'_\mu Q_{Y'} - i g_{Y''} Y''_\mu Q_{Y''}, \tag{39}$$

are conveniently written as

$$- i \frac{x_{H_i}}{v_i} \bar{M}_Z Y'_\mu - i \frac{y_{H_i}}{v_i} \bar{M}_Z Y''_\mu \tag{40}$$

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

$$x_{H_1} = 1.9 \sqrt{g_1'^2 - 0.032} S_\phi, \tag{41}$$

$$x_{H_2} = -x_{H_1}$$
$$- 1.9 \left[0.054 g_1'^2 \sqrt{\frac{1}{(g_1'^2 - 0.032)(g_1'^2 - 0.033)}} C_\phi \right.$$
$$\left. + \frac{0.064 S_\phi}{\sqrt{g_1'^2 - 0.032}} \right], \tag{42}$$

$$y_{H_1} = 1.9 \sqrt{g_1'^2 - 0.032} C_\phi, \tag{43}$$

and

$$y_{H_2} = -y_{H_1}$$
$$- 1.9 \left[0.054 g_1'^2 \sqrt{\frac{1}{(g_1'^2 - 0.032)(g_1'^2 - 0.033)}} S_\phi \right.$$
$$\left. + \frac{0.064 C_\phi}{\sqrt{g_1'^2 - 0.032}} \right]. \tag{44}$$

For our fiducial values of ϕ and g_1' we obtain $x_{H_1} = -0.351$, $x_{H_2} = 0.822$, $y_{H_1} = 0.151$, and $y_{H_2} = -0.556$.

The Higgs field kinetic term together with the Green-Schwarz mass terms ($-\frac{1}{2} M'^2 Y'_\mu Y'^\mu - \frac{1}{2} M''^2 Y''_\mu Y''^\mu$, see Appendix) yield the following mass square matrix

$$\begin{pmatrix} \bar{M}_Z^2 & \bar{M}_Z^2 (x_{H_1} C_\delta^2 + x_{H_2} S_\delta^2) & \bar{M}_Z^2 (y_{H_1} C_\delta^2 + y_{H_2} S_\delta^2) \\ \bar{M}_Z^2 (x_{H_1} C_\delta^2 + x_{H_2} S_\delta^2) & \bar{M}_Z^2 (C_\delta^2 x_{H_1}^2 + S_\delta^2 x_{H_2}^2) + M'^2 & \bar{M}_Z^2 (C_\delta^2 x_{H_1} y_{H_1} + S_\delta^2 x_{H_2} y_{H_2}) \\ \bar{M}_Z^2 (y_{H_1} C_\delta^2 + y_{H_2} S_\delta^2) & \bar{M}_Z^2 (C_\delta^2 x_{H_1} y_{H_1} + S_\delta^2 x_{H_2} y_{H_2}) & \bar{M}_Z^2 (y_{H_1}^2 C_\delta^2 + y_{H_2}^2 S_\delta^2) + M''^2 \end{pmatrix}.$$

The free parameters are $\tan\delta$, $M_{Z'}$, and $M_{Z''}$ which will be fixed by requiring the shift of the Z mass to lie within 1 standard deviation of the experimental value and $M_{Z'} = 150 \pm 5$ GeV. We are also minimizing $M_{Z''}$ to ascertain whether it can be detected at existing colliders. This leads to $\tan\delta = 0.65$ and $M_{Z''} \simeq M'' \geq 0.90$ TeV.

Next, we scan the parameter space to obtain a larger Wjj production cross section at the Tevatron. For $\phi = 2.12$ and $g_1' = 0.26$, corresponding to a suppression $\Gamma_{Z' \to e^+ e^-}/\Gamma_{Z' \to q\bar{q}} \sim 0.6\%$, $\theta = -0.76$, and $\psi = -1.36$, one obtains $\sigma(p\bar{p} \to WZ') \times \text{BR}(Z' \to jj) = 2.9$ pb. The associated $g_{Y'} Q_{Y'}$ and $g_{Y''} Q_{Y''}$ couplings are given in Table IV. It is straightforward to see that for $x_{H_1} = 0.303$, $x_{H_2} = -0.740$, $y_{H_1} = -0.187$, and $y_{H_2} = 0.690$ the shift of the Z mass would lie within 1 standard deviation of the experimental value if $\tan\delta \simeq 0.64$, $M_{Z'} = 150$ GeV, and $M_{Z''} \geq 0.92$ TeV.

B. Nonanomalous $B - L$

We now turn to discuss the alternative framework where the $U(1)$ right-handed neutrino charges are $Q_3 = 0$, $Q_{1L} = Q_{1R} = -1$, which means that $B - L$ is anomaly-free. Therefore this case is somewhat preferred compared to the previous case, since there can be a natural hierarchy among Z' and Z''. In fact, as we will show below, the

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[7]Note that H_2 cannot correspond to an elementary open string excitation, since it has $Q_{1L} = -2$. One possibility is to regard H_2 as a composite scalar field, built from two elementary open string scalars, a SM singlet ϕ and another Higgs doublet H_2', $H_2 \sim \phi H_2'$, with the following $U(1)$-charges: ϕ: $(0, -1, -1)$ and H_2': $(0, -1, 0)$. In case H_2 is a composite scalar field so that the corresponding Yukawa coupling arises from a dimension-5 effective operator, one expects that its vacuum expectation value is somewhat suppressed compared to the vev of H_1, i.e., $\tan\delta \equiv v_1/v_2 > 1$.

TABLE IV.   Chiral couplings of Y' and Y'' gauge bosons for $\phi = 2.12$ and $g_1' = 0.26$.

Name	$g_{Y'}Q_{Y'}$	$g_{Y''}Q_{Y''}$
$\bar{U}_i$	0.088	0.395
$\bar{D}_i$	0.410	0.197
L_i	0.116	0.116
$\bar{E}_i$	0.045	0.034
Q_i	−0.249	−0.296

mixing angle ϕ is small and therefore Z' and Z'' become essentially B and $B - L$, respectively. In such a case, two "supersymmetric" Higgses $H_u \equiv H_\nu$ and H_d with charges $Q_3 = Q_{1L} = 0$, $Q_{1R} = 1$, $Q_Y = 1/2$ and $Q_3 = Q_{1L} = 0$,

$Q_{1R} = -1$, $Q_Y = -1/2$ would be sufficient to give masses to all the chiral fermions. Here,

$$\langle H_u \rangle = \begin{pmatrix} 0 \\ v_u \end{pmatrix},$$

$$\langle H_d \rangle = \begin{pmatrix} v_d \\ 0 \end{pmatrix},$$

$v = \sqrt{v_u^2 + v_d^2} = 174$ GeV, and $\tan\beta \equiv v_u/v_d$. For this particular selection of $U(1)$ charges $x_{H_u} = -x_{H_d} = x_{H_1}$ and $y_{H_u} = -y_{H_d} = y_{H_1}$. Therefore, it is straightforward to see that the corresponding mass square matrix for the $Z - Z'$ mixing,

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

does not impose any constraint on the $\tan\beta$ parameter. We then use the 2 degrees of freedom of the model (g_1', ϕ) to demand the shift of the Z mass to lie within 1 standard deviation of the experimental value and leptophobia. Taking $M_{Z'} = 150$ GeV, with $g_1' = 0.20$, $\phi = 0.0028$, and $M_{Z''} = 5$ TeV, we find that $\Gamma_{Z''\to e^+e^-}/\Gamma_{Z''\to q\bar{q}} \simeq 1\%$.

Recall that for this particular $U(1)$ charge selection of the right-handed neutrino the combination $B - L$ is anomaly-free. Therefore, the mass ratio of the anomalous and the nonanomalous $U(1)$ can be ascribed to a suppression induced by the large two-dimensional volume. The $g_{Y'}Q_{Y'}$ and $g_{Y''}Q_{Y''}$ couplings to the chiral fields are fixed and given in Table V. The Z' couplings to quarks leads to a large (precut) Wjj production ($\simeq 6$ pb) at the Tevatron, and at $\sqrt{s} = 630$ GeV, a direct (precut) $Z' \to jj$ production ($\simeq 700$ pb) in the region excluded by UA2 data. However, it is worthwhile to point out that the UA2 Collaboration performed their analysis in the early days of QCD jet studies. Their upper bound depends crucially on the quality of the Monte Carlo and 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 [31]. In view of the considerable uncertainties associated with the UA2

analysis we remain skeptical of drawing negative conclusions. Instead we argue that our supersymmetric D-brane construct could provide an explanation of the CDF anomaly if acceptance and pseudorapidity cuts reduce the Wjj production rate by about 35% and the UA2 90% C.L. bound is taken as an order-of-magnitude limit [32].

The $U(1)$ vector bosons couple to currents

$J_Y = 1.8 \times 10^{-1}Q_{1R} + 5.9 \times 10^{-2}Q_3 - 1.8 \times 10^{-1}Q_{1L}$

$J_{Y'} = 2.5 \times 10^{-4}Q_{1R} + 3.7 \times 10^{-1}Q_3 + 1.4 \times 10^{-1}Q_{1L}$

$J_{Y''} = 9.0 \times 10^{-2}Q_{1R} - 1.2 \times 10^{-1}Q_3$

$\qquad\qquad + 3.5 \times 10^{-1}Q_{1L}. \qquad (45)$

Using Eq. (18), we rewrite $J_{Y'}$ and $J_{Y''}$ as

$J_{Y'} = 2.5 \times 10^{-4}Q_{1R} + 1.11B + 1.4 \times 10^{-1}L$

$J_{Y''} = 9.0 \times 10^{-2}Q_{1R} - 2.5 \times 10^{-3}(B + L)$

$\qquad\qquad - 3.55 \times 10^{-1}(B - L). \qquad (46)$

Since $\text{Tr}[Q_{1R}B] = \text{Tr}[Q_{1R}L] = 0$, the decay widths are given by

$\Gamma_{Y'} = \Gamma_{Y'\to Q_{1R}} + \Gamma_{Y'\to B} + \Gamma_{Y'\to L}$

$\qquad \propto (2.5 \times 10^{-4})^2\text{Tr}[Q_{1R}^2] + (1.11)^2\text{Tr}[B^2]$

$\qquad\quad + (1.4 \times 10^{-1})^2\text{Tr}[L^2]$

$\qquad = 8 \times (2.5 \times 10^{-4})^2 + \dfrac{4}{3} \times (1.11)^2$

$\qquad\quad + 4 \times (1.4 \times 10^{-1})^2$

$\qquad = 5 \times 10^{-7} + 1.64 + 7.84 \times 10^{-2}, \qquad (47)$

TABLE V.   Chiral couplings of Y' and Y'' gauge bosons for $\phi = 0.0028$ and $g_1' = 0.20$.

Name	$g_{Y'}Q_{Y'}$	$g_{Y''}Q_{Y''}$
$\bar{U}_i$	−0.368	0.028
$\bar{D}_i$	−0.368	0.209
L_i	0.143	0.143
$\bar{E}_i$	−0.142	−0.262
Q_i	0.368	−0.119

$$\Gamma_{Y''} = \Gamma_{Y'' \to Q_{1R}} + \Gamma_{Y' \to B+L} + \Gamma_{Y' \to B-L}$$
$$\propto (9.0 \times 10^{-2})^2 \text{Tr}[Q_{1R}^2] + (2.5 \times 10^{-3})^2 \text{Tr}[(B$$
$$+ L)^2] + (3.55 \times 10^{-1})^2 \text{Tr}[(B-L)^2]$$
$$= 6.48 \times 10^{-2} + 3.33 \times 10^{-5} + 6.7 \times 10^{-1}. \quad (48)$$

Thus, the corresponding branching fractions are

$$\text{BR}\, Y' \to Q_{1R} : \text{BR}\, Y' \to B : \text{BR}\, Y' \to L \qquad (49)$$
$$2.9 \times 10^{-7} : 0.95 : 0.046$$

and

$$\text{BR}\, Y'' \to Q_{1R} : \text{BR}\, Y'' \to B + L : \text{BR}\, Y'' \to B - L \qquad (50)$$
$$0.09 : 4.5 \times 10^{-5} : 0.91.$$

Of course, since the quiver construction has each particle straddling two adjacent branes, there can be considerable variation in decay channels particle by particle. The dominance of B for the Y' decay channel and $B - L$ for the Y'' decay channel is valid after averaging over decay channels.[8] It is important to note that a 100% coupling of the Y' and Y'' to B and $B - L$, respectively, is possible only if the $U(1)$ gauge coupling constants are equal, see Appendix B.

V. LEPTOPHOBIC Z' AT THE LHC

Since the CDF signal is in dispute, it is of interest to study the predictions of the model for energies not obtainable at the Tevatron, but within the range of the LHC. To illustrate the LHC phenomenology of our D-brane construct, we consider the model in which the $U(1)$ right-handed neutrino has charges $Q_3 = 0$, $Q_{1L} = Q_{1R} = -1$, i.e., $B - L$ is nonanomalous.

The ATLAS Collaboration has searched for narrow resonances in the invariant mass spectrum of dimuon and dielectron final states in event samples corresponding to an integrated luminosity of 1.21 fb^{-1} and 1.08 fb^{-1}, respectively [33]. The spectra are consistent with SM expectations and thus upper limits on the cross section times branching fraction for Z' into lepton pairs have been set.

Using a data set with an integrated luminosity of 1 fb^{-1}, the CMS Collaboration has searched for narrow resonances in the dijet-invariant mass spectrum [34]. For $M_{Z'} \simeq$ 1 TeV, the CMS experiment has excluded production rates $\sigma(pp) \times \text{BR}(Z' \to jj) \times \mathcal{A} > 1$ pb at the 95% C.L. Each event in the search is required to have its two highest-p_T jets with (pseudorapidity) $|\eta_j| < 2.5$ and the leading jet must satisfy $p_t^{j_1} > 150$ GeV, with $|\Delta \eta_{jj}| < 1.3$. The acceptance $\mathcal{A}$ of selection requirements is reported to be ≈ 0.6.

To compare our predictions with LHC experimental searches in dilepton and dijets it is sufficient to consider the production cross section in the narrow Z' width approximation,

$$\hat{\sigma}(q\bar{q} \to Z') = K \frac{2\pi}{3} \frac{G_F M_Z^2}{\sqrt{2}} [v_q^2(\phi, g_1')$$
$$+ a_q^2(\phi, g_1')]\delta(\hat{s} - M_{Z'}^2), \qquad (51)$$

where G_F is the Fermi coupling constant and the K-factor represents the enhancement from higher-order QCD processes estimated to be $K \simeq 1.3$ [35]. After folding $\hat{\sigma}$ with the CTEQ6 parton distribution functions [29], we determine (at the parton level) the resonant production cross section. In Fig. 4 we compare the predicted $\sigma(p\bar{p} \to Z') \times \text{BR}(Z' \to jj)$ (left panel) and $\sigma(p\bar{p} \to Z') \times \text{BR}(Z' \to \ell\ell)$ (right panel) production rates with 95% C.L. upper limits recently reported by the CMS [34] and ATLAS [33] collaborations. Selection cuts will probably reduce event rates by factors of 20%. Keeping this in mind, we conclude that the 2012 LHC7 run will probe $M_{Z'} \sim 3$ TeV, whereas future runs from LHC14 will provide a generous discovery potential of up to about $M_{Z'} \sim 8$ TeV.

VI. REGGE EXCITATIONS

In TeV-scale gravity scenarios where the SM is realized on the world-volume of D-branes the presence of fundamental strings can also be unearthed by searching for the effects of their vibrations. The particles that appear as the quanta of oscillating string modes are called Regge excitations and have squared masses quantized in units of $M_s = 1/\sqrt{\alpha'}$, where α' is the Regge slope parameter [36]. The leading contributions of Regge recurrences to certain processes at hadron colliders are *universal*. This is because the full-fledged string amplitudes which describe $2 \to 2$ parton scattering subprocesses involving four gauge bosons as well as those with two gauge bosons and two chiral matter fields are (to leading order in string coupling, but all orders in α') independent of the D-brane configuration, the geometry of the extra dimensions, and whether supersymmetry is broken or not. Therefore, the s-channel pole terms of the average square amplitudes contributing to dijet production can be obtained independently of the details of the compactification scheme [37]. For phenomenological purposes, the poles need to be softened to a Breit-Wigner form by obtaining and utilizing the correct *total* widths of the resonances [38]. After this is done, it is feasible to compute genuine string corrections to dijet signals at the LHC [39].[9] The CMS Collaboration has

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[8]An analogue is in the SM. The Z couples to a current $J_Z \propto T_3 - \tan^2\theta_W \frac{Y}{2}$, where $Q = T_3 - \frac{Y}{2}$. In this case, $\sum(\frac{Y}{2})^2 = \frac{17}{6}$ and $\text{Tr}[T_3^2] = 2$; we have $\text{BR}Z \to T_3 : \text{BR}Z \to \frac{Y}{2} = 2 : \frac{17}{6}\tan^4\theta_W = 2 : 0.25 = 8 : 1$. However, this certainly does not hold particle by particle; e.g., for the neutrino electron doublet: $\Gamma_{Z\to\nu} \propto (1 + \tan^2\theta_W)^2 \sim 1.7$, whereas $\Gamma_{Z\to e} \propto (1 - \tan^2\theta_W)^2 \sim 0.5$.

[9]Phenomenological studies of Regge excitations and associated collider signatures, based on simple toy model embedding parts of the SM into string theory, have been carried out in [40]. The discovery potential of string resonances via top quark pair production in the context of cannonical D-brane constructions has been recently established [41].
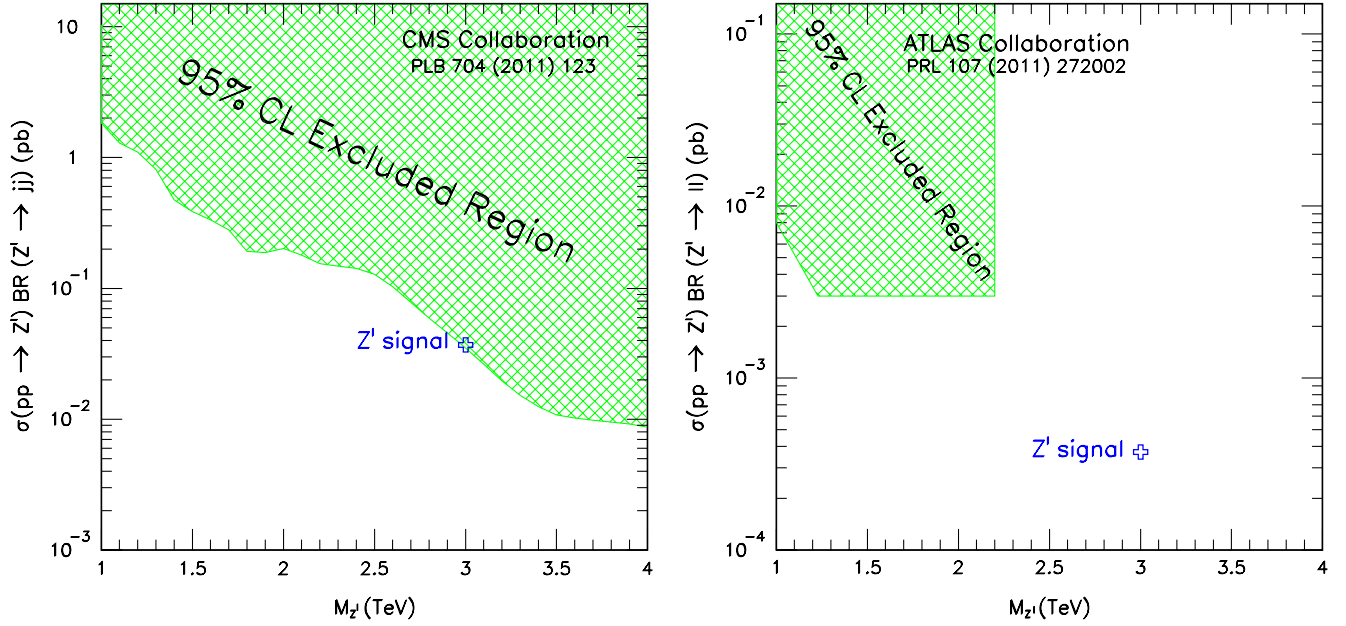
FIG. 4 (color online).   Comparison of the (precut) total cross section for the production of $pp \to Z' \to jj$ (left) and $pp \to Z' \to \ell\ell$ (right) with the 95% C.L. upper limits on the production of a gauge boson decaying into two jets (left) and two leptons (right), as reported by the CMS (corrected by acceptance) [34] and ATLAS [33] collaborations, respectively. We have taken $\phi = 0.0028$, $g_1' = 0.20$. For isotropic decays (independently of the resonance), the acceptance for the CMS detector has been reporetd to be $\mathcal{A} \approx 0.6$ [34]. The predicted Z' production rates for $\sqrt{s} = 7$ TeV and $M_{Z'} \simeq 3$ TeV saturate the current limits.

searched for such narrow resonances in their dijet-mass spectrum [42]. After operating for only few months, with merely 2.9 inverse picobarns of integrated luminosity, the LHC CMS experiment has ruled out $M_s < 2.5$ TeV. The LHC7 has recently delivered an integrated luminosity in excess of 1 fb^{-1}. This extends considerably the search territory for new physics in events containing dijets. The new data exclude string resonances with $M_s < 4$ TeV [34]. In fact, as shown in Fig. 5, the LHC has the capacity of discovering strongly interacting resonances via dijet final states in practically all range up to $\frac{1}{2}\sqrt{s_{\text{LHC}}}$. Of particular interest here, for the $U(3)_C \times Sp(1)_L \times U(1)_L \times U(1)_R$ D-brane model, the anomaly cancellation fixes the projection of the hypercharge into the color stack at the string scale: $\kappa = c_3\sqrt{6}g_Y/g_3$ [43]. Therefore one can also cleanly extract the leading string corrections to γ + jet signals at the LHC [44].[10] The precise predictions for the branching fraction of two different topologies (dijet and γ + jet) can be used as a powerful discriminator of low-mass string excitations from other beyond SM scenarios.

FIG. 5 (color online).   Signal-to-noise ratio of $pp \to$ dijet and $pp \to \gamma$ + jet, for $\sqrt{s} = 14$ TeV, $\mathcal{L} = 100$ fb^{-1}, and $\kappa^2 \simeq 0.02$. The approximate equality of the background due to misidentified π^0's and the QCD background, across a range of large p_T^γ as implemented in [44], is maintained as an approximate equality over a range of γ-jet-invariant masses with the rapidity cuts imposed ($|y_{\text{max}}^j| < 1.0$ and $|y_{\text{max}}^\gamma| < 2.4$). Details of the signal and background calculations have been given elsewhere [39].

[10]E.g., the tree-level amplitude for gluon fusion into γ + jet, $\mathcal{M}(gg \to \gamma g) = \cos\theta_W \mathcal{M}(gg \to Yg) = \kappa \cos\theta_W \mathcal{M}(gg \to Cg)$, has a unique free parameter that is the string scale M_s. Here, θ_W is the weak angle and C is the extra $U(1)$ boson tied to the color stack. For details see [44].
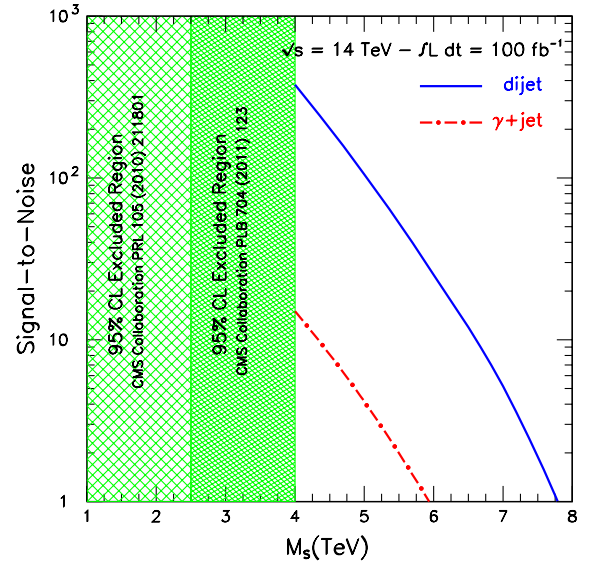
VII. CONCLUSIONS

We have shown that a Z' that can explain the CDF Wjj excess and is in full agreement with existing limits on Z' couplings to quarks and leptons can materialize in the context of D-brane TeV-scale string compactifications. The existence of additional, largely leptophobic, Z's with anomalous masses somewhat less than the string scale, is generic to the D-brane models discussed in some detail in this paper. Thus, even if the CDF anomaly does not survive additional scrutiny, there may exist such Z's with masses $\gtrsim 1$ TeV, whose discovery is out of reach of Tevatron, but open to such at LHC. In that case the analysis presented here can be directly applied to the higher-energy realm, with a view toward identifying the precise makeup of the various Abelian sectors, and pursuing with strong confidence a signal at LHC for the Regge excitations of the string. We have long imagined strings to be minuscule objects which could only be experimentally observed in the far distant future. It is conceivable that this future has already arrived.

APPENDIX A: PROPERTIES OF THE ANOMALOUS MASS SECTOR

Outside of the Higgs couplings, the relevant parts of the Lagrangian are the gauge couplings generated by the $U(1)$ covariant derivatives acting on the matter fields, and the (mass)^2 matrix of the anomalous sector

$$\mathcal{L} = Q^T g X + \frac{1}{2} X^T \mathbb{M}^2 X, \qquad (A1)$$

where X_i are the three $U(1)$ gauge fields in the D-brane basis $(B_\mu, C_\mu, \tilde{B}_\mu)$, g is a diagonal coupling matrix (g'_1, g'_3, g'_4), and Q are the 3 charge matrices.

As in Sec. II, perform a rotation $X = \mathcal{R} Y$ and require that one of the Y's (say Y_μ) couple to hypercharge. We then obtain the constraint on the first column of $\mathcal{R}$ given in Eq. (8). However, there is now an additional constraint: *the field Y_μ is an eigenstate of $\mathbb{M}^2$ with zero eigenvalue.* Under the $\mathcal{R}$ rotation, the mass term becomes

$$\frac{1}{2} X^T \mathbb{M}^2 X = \frac{1}{2} Y^T \bar{\mathbb{M}}^2 Y, \qquad (A2)$$

with $\bar{\mathbb{M}}^2 = \mathcal{R}^T \mathbb{M}^2 \mathcal{R}$. We know that at least Y_μ is an eigenstate with eigenvalue 0. We also know that Poincare invariance requires the complete diagonalization of the mass matrix in order to deal with observables. However, further similarity transformations will undo the coupling of the zero eigenstate to hypercharge. There seems no way of eventually fulfilling all these conditions except to require that the same $\mathcal{R}$ which rotates to couple Y_μ to hypercharge simultaneously diagonalizes $\mathbb{M}^2$ so that

$$\bar{\mathbb{M}}^2 = \mathrm{diag}(0, M'^2, M''^2). \qquad (A3)$$

This implies that the original $\mathbb{M}^2$ in the D-brane basis is given by

$$\mathbb{M}^2 = \mathcal{R}\,\mathrm{diag}(0, M'^2, M''^2)\mathcal{R}^T, \qquad (A4)$$

which results in the following baroque matrix:

$$\mathbb{M}^2 = \begin{pmatrix} a & b & c \\ b & d & e \\ c & e & f \end{pmatrix}, \qquad (A5)$$

where

$$a = M'^2(C_\psi S_\theta S_\phi - C_\phi S_\psi)^2 + M''^2(C_\phi C_\psi S_\theta + S_\phi S_\psi)^2,$$

$$b = (M'^2 - M''^2)C_\phi C_{2\psi} S_\theta S_\phi + C_\phi^2 C_\psi(-M'^2 + M''^2 S_\theta^2)S_\psi$$
$$+ C_\psi(-M''^2 + M'^2 S_\theta^2)S_\phi^2 S_\psi,$$

$$c = C_\theta[M''^2 C_\phi^2 C_\psi S_\theta + M'^2 C_\psi S_\theta S_\phi^2$$
$$- (M'^2 - M''^2)C_\phi S_\phi S_\psi],$$

$$d = M''^2(C_\psi S_\phi - C_\phi S_\theta S_\psi)^2 + M'^2(C_\phi C_\psi + S_\theta S_\phi S_\psi)^2,$$

$$e = C_\theta[(M'^2 - M''^2)C_\phi C_\psi S_\phi + M''^2 C_\phi^2 S_\theta S_\psi$$
$$+ M'^2 S_\theta S_\phi^2 S_\psi],$$

$$f = C_\theta^2(M''^2 C_\phi^2 + M'^2 S_\phi^2). \qquad (A6)$$

APPENDIX B: B AND $B - L$ COUPLINGS ON THE ROTATED BASIS

For given a set of $U(1)$ fields with orthogonal charges in the 1, 2, 3, ... basis, an obvious question is whether each of the fields on the rotated basis couples to a single charge $\bar{Q}_i$. Let

$$\mathcal{L} = X^T g Q, \qquad (B1)$$

be the Lagrangian in the 1, 2, 3, ... basis, with X_μ^i and Q_i vectors and g a diagonal matrix in N-dimensional 'flavor' space. Now rotate to new orthogonal basis $(\bar{Q})$ for Q

$$Q = \mathcal{O}\bar{Q}; \tag{B2}$$

(B1) becomes

$$\mathcal{L} = X^T g \mathcal{O}\bar{Q}. \tag{B3}$$

As it stands, each X^i_μ does not couple to a unique charge $\bar{Q}_i$; hence we rotate X,

$$X = \mathcal{R}\bar{Y}, \tag{B4}$$

to obtain

$$\mathcal{L} = \bar{Y}^T \mathcal{R}^T g \mathcal{O}\bar{Q}. \tag{B5}$$

We wish to see if, for given $\mathcal{R}$ and g, we can find an $\mathcal{O}$ so that

$$\mathcal{R}^T g \mathcal{O} = \bar{g}(\text{diagonal}). \tag{B6}$$

This allows each $\bar{Y}^i_\mu$ to couple to a unique charge $\bar{Q}_i$ with strength $\bar{g}_i$. To see the problem with this, we rewrite (B6) in terms of components

$$(\mathcal{R}^T)_{ij} g_j \mathcal{O}_{jk} = \bar{g}_i \delta_{ik}; \tag{B7}$$

for $i \neq k$, (B7) leads to

$$(\mathcal{R}^T)_{ij} g_j \mathcal{O}_{jk} = 0. \tag{B8}$$

In general, in Eq. (B8) there are $N(N-1)$ equations, but only $N(N-1)/2$ independent $\mathcal{R}_{ij}$ generators in $SO(N)$; therefore the system is overdetermined. Of course, if $g = g\mathbb{1}$, the equation becomes

$$\mathcal{R}^T \mathcal{O} = \mathbb{1}, \tag{B9}$$

and so $\mathcal{O} = \mathcal{R}$.

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